

# **JAC: Jena Atomic Calculator**

## **— User Guide, Compendium & Theoretical Background —**

<http://www.atomic-theory.uni-jena.de/>

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# 1. Overview about JAC. Structure of this User Guide

## 1.1. Goals of the JAC toolbox

### Purpose of the JAC module:

- The Jena Atomic Calculator (JAC) provides tools for performing atomic (structure) calculations at various degrees of complexity and sophistication. This toolbox has been designed to calculate not only atomic level structures and properties [g-factors, hyperfine and isotope-shift parameters, etc.] or transition amplitudes between bound-state levels [dipole operator, Schiff moment, parity non-conservation, etc.] but, in particular, also (atomic) transition probabilities, Auger rates, photoionization cross sections, radiative and dielectronic recombination rates as well as cross sections and parameters for many other (elementary) processes.
- JAC also facilitates interactive computations, the simulation of atomic cascades and atomic responses, the time-evolution of statistical tensors as well as various semi-empirical estimates of atomic properties. It provides a diverse and wide-ranging, yet consistent set of methods which can be applied in different fields of atomic physics and elsewhere.
- In addition, the JAC module has been designed to readily support the display of level energies, electron and photon spectra, radial orbitals and several others entities.
- To find (the details about) individual features of the JAC program, see the index or search for keywords/phrases in this text.
- In practice, the design of JAC has been based on an analysis of typical user requirements and a hierarchical structure of the code.
- Since the theoretical background and data, implemented in JAC, have been extracted from quite many sources, we also hope to develop JAC as a repository of previous experience with electronic structure calculations of atoms in different environments, and that is to be further refined, expanded and developed here.
- The source code, an extensive documentation as well as a number of tutorials and examples are available from our Web site <https://www.github.com/sfritzsche/JAC.jl>.
- In order to support all these goals, several types of computation are distinguished within the JAC toolbox.

## Types of computations:

- **Atomic computations, based on explicitly specified electron configurations:** A typical computation, that is based on explicitly specified electron configurations, refers to the level energies, atomic states and to either **one (or several) atomic properties** for levels from a given multiplet or to the rates and cross sections of just **one selected atomic process**. For further details, see the supported amplitudes, properties and processes in Sections 5–7 below.
- **Restricted active-space computations (RAS):** A RAS computation refers to systematically-enlarged calculations of atomic states and level energies due to a specified (and usually restricted) active space of orbitals as well as due to the number and/or kind of virtual excitations to be included. Such **RAS computations are internally performed *stepwise*** in JAC by utilizing the self-consistent field (orbitals) from some prior step. This type has **not yet been properly implemented so far**.
- **Interactive computations:** In an interactive computation, the **functions/methods of the JAC program are applied interactively, either directly within the REPL or by just a short JULIA script**, in order to compute energies, expansion coefficients, transition matrices, rates, cross sections, etc. An interactive computation typically first prepares and generates (instances of) of different data structures of JAC, such as orbitals, (configuration-state) bases, multiplets, and later applies these computed data to obtain the desired information. More generally, **all methods from JAC and its submodules can be utilized also interactively**, although some specialized methods are often available and may facilitate the computations. Like for other JULIA functions, the functions and methods provided by JAC can be seen as (high-level) language elements in order to perform atomic computations at various degrees of sophistication.
- **Atomic cascade computations:** A cascade computation typically refers to **three or more charge states of an atom**, and which are connected to each other by different atomic processes, such as photoionization, dielectronic recombination, Auger decay, radiative transitions, etc. Different (cascade) approaches have been predefined in order to deal with atomic cascades. The particular atomic processes that are to be taken into account for the individual steps of the cascade need to be specified explicitly; these cascade computations have been **only partly implemented so far**.
- **Atomic responses:** Atomic response computations will support **simulations of how atoms *respond* to an incident (beam of) light pulses and particles**, such as field-induced ionization processes, high-harmonic generation and others. For these responses, the detailed atomic structure has often not been considered in much detail in the past, though it will become relevant as more elaborate and accurate measurements are carried out. These response computations have **not yet been implemented so far**.
- **Time evolution of statistical tensors in (intense) light pulses:** A time evolution of statistical tensors always proceeds within a pre-specified set of sublevels  $\{ |\alpha JM\rangle \}$ , i.e. subspace of the many-electron Hilbert space; all further (decay) processes that lead the system *out of this*

*subspace* must be treated by loss rates. Although such a time evolution can deal with pulses of different shape, strength and duration, it is assumed that they are *weak enough not to substantially disturb the level structure and level sequence of the atoms in their neutral or ionic stage*, i.e. that every sublevel can still be characterized by its (total) energy and symmetry. This time-evolution has *not yet been implemented in detail so far*. No attempt is made in JAC to solve the time-dependent (many-electron) Schrödinger equation explicitly.

- **Semi-empirical estimates of atomic properties, cross sections, asymptotic behaviour, etc.:** A semi-empirical ‘estimate’ of atomic data refers to some simple model computation or to the evaluation of fit functions in order to *provide such atomic data, that cannot be generated so easily by *ab-initio* computations*. These semi-empirical estimates are typically built on — more or less — sophisticated models and external parameter optimizations, although *only a very few of such estimates have been implemented so far*.

## 1.2. Notations

### Atomic shells & subshells:

- **Atomic shell model:** This model, in which electrons fill a (more or less) regular list of atomic shells or make transitions (quantum jumps) between different shells, is *key and guidance for calculating the electronic structure of atoms, ions and molecules, and most of their properties*.
- **Electron configuration:** Describes the occupation of shells within the atomic shell model, for instance,  $1s^2 2s^2 2p^6 3s^2$ . In JAC, closed-shell configurations can often be abbreviated, for instance, by [Ne]  $3s^2$ .
- **Shells and subshells:** Shells and subshells are the building blocks for the atomic shell model. In the relativistic theory, each non-relativistic  $n\ell$ -shell (apart from the  $ns$ -shells) splits into two relativistic subshells due to  $j = \ell \pm 1/2$ . In JAC, the shell and subshell notations are therefore used in order to denote the electron configurations, configuration state functions (CSF) and the orbitals of equivalent electrons. In JAC, there are special data struct’s available to easily deal and communicate the (sub-) shell specification of levels and wave functions.
- **Relativistic angular-momentum quantum number:**  $\kappa = \pm(j + 1/2)$  for  $\ell = j \pm 1/2$  carries information about both, the total angular momentum  $j$  and the parity  $(-1)^\ell$  of the (single-electron) wavefunction.

## Atomic level notations:

- **Bound (many-electron) levels and states:**  $|\alpha\mathbb{J}\rangle \equiv |\alpha J^P\rangle$ ;  $|\alpha\mathbb{J}M\rangle \equiv |\alpha J^P M\rangle$ .

Here, the multi-index  $\alpha$  refers to all additional quantum numbers that are needed for a unique specification of some many-electron level or state. — For instance, we shall often use below  $|\alpha_i\mathbb{J}_i\rangle$  and  $|\alpha_f\mathbb{J}_f\rangle$  in order to denote the initial and final-ionic bound states of some atomic amplitude and process.

- **(Many-electron) resonances or scattering states with a single free electron:**  $|(\alpha\mathbb{J}, \varepsilon\kappa)\mathbb{J}_t\rangle \equiv |(\alpha J^P, \varepsilon\kappa) J_t^{P_t}\rangle$

describes a many-electron scattering wave with a single free electron in the partial wave  $|\varepsilon\kappa\rangle$ , and where  $\mathbb{J}_t$  denotes the overall symmetry of the many-electron scattering state (level).

## Atomic multipoles:

- The multipole components (fields) of the radiation field occur at many different places in atomic theory due to the electron-photon interaction, though often within slightly different contexts. We here use the standard notations for the E1 (electric-dipole), M1 (magnetic-dipole), E2 (electric-quadrupole), etc. components and refer to them briefly as *multipoles*  $\mathbb{M}$  of the radiation field.
- These multipoles or multipole components of the radiation field frequently occurs also in terms of *multipole transition operators*  $\mathbb{O}^{(\mathbb{M})}$  and *multipole moment operators*  $\mathbb{Q}^{(\mathbb{M})}$  as well as in summations  $\sum_{\mathbb{M}} \dots$  over such operators and/or the associated transition amplitudes (matrix elements).
- A multipole  $\mathbb{M}$  is internally characterized by its multipolarity  $L$  and its (boolean) character *electric* (= `true/false`).

## Frequent technical terms used in the manual:

- **Amplitude:** Transition amplitudes are the many-electron matrix elements upon which atomic structure and collision theory is usually built on.
- **Basis:** In JAC, a basis refers to a many-electron basis that is specified in terms of a CSF list and the radial orbitals of all (equivalent) electrons. Each CSF is specified uniquely by a proper set of quantum numbers; here, the (so-called) seniority scheme is applied for the unique classification of the CSF and the evaluation of all matrix elements.



- **Channel:** Sublines of a given line that need to be distinguished due to the symmetry properties of the multipoles or partial waves in the decomposition of the many-electron levels or matrix elements.
- **Grid:** In JAC, all radial orbital functions are always represented on a grid. Only their grid representation is applied in the evaluation of all (single- and many-electron) matrix elements. Predefined grids refer to a 'exponential grid', suitable for bound-state computations, as well as a log-lin grid, which increases exponentially in the inner part and linearly in the outer part. The latter grid is suitable for collision processes and for dealing with the electron continuum.
- **Lines and transitions:** A line refers to an atomic transition that is characterized in terms of a well-defined initial and final level; it frequently occurs in the computation of different properties, such as cross sections or rates, angular distribution parameters. Typically, a line contains various channels (sublines), for instance, due to occurrence of multipoles or partial waves in the decomposition of the many-electron matrix elements
- **Multiplet:** Atomic levels are naturally *grouped together* into (so-called) multiplets; most often, this term just stands for all fine-structure levels of one or several given configurations. More generally, multiplets may refer to any group of levels, for instance, to groups levels with the same total angular momentum  $J$  and/or parity  $P$ , or fine-structure levels of closely related configurations, etc.
- **Orbital:** In atomic physics, an orbital typically refers to a one-electron function in an radial-spherical representation. Often, only the radial function(s) are meant. In the relativistic theory, of course, one needs to distinguish between the large and small components of the orbital, following Dirac's theory.
- **Pathway:** In contrast to an (atomic) line, that is characterized by an initial- and a final-level (and the corresponding multiplets), a pathway describe a sequence of three or more levels, and which often correspond to different atomic processes. These levels are usually referred to as initial, (one or several) intermediate and final level. Pathways occur naturally in dielectronic recombination as well as in various excitation-ionization of excitation-autoionization processes.
- **Settings:** In JAC, the control of most, if not all, computations is made by **Settings** that are associated to particular amplitudes, properties and processes. These settings are used to specify all details about the requested computation and enable one, for instance, to select individual levels or lines as well as various physical and technical parameters, such as the multipoles, gauges, etc.

## 1.3. A quick overview about amplitudes, level properties and processes handled in JAC

### In Jac (partly) implemented amplitudes:

- Selected many-electron (reduced) amplitudes that are accessible within the JAC program. Further details about the call of these amplitudes can be found below in this manual or interactively by `?<module>.amplitude`.

Amplitude	Call within JAC	Brief explanation
$\langle \alpha \mathbb{J} \parallel \mathbb{T}^{(1)} \parallel \beta \mathbb{J}' \rangle$ , $\langle \alpha \mathbb{J} \parallel \mathbb{T}^{(2)} \parallel \beta \mathbb{J}' \rangle$	Hfs.amplitude	Amplitude for the hyperfine interaction with the magnetic-dipole and electric-quadrupole field of the nucleus.
$\langle \alpha \mathbb{J} \parallel \mathbb{N}^{(1)} \parallel \beta \mathbb{J}' \rangle$	LandeZeeman.amplitude	Amplitude for the interaction with an external magnetic field.
$\langle \alpha_f \mathbb{J}_f \parallel \mathbb{O}^{(\mathbb{M}, \text{emission})} \parallel \alpha_i \mathbb{J}_i \rangle$	Radiative.amplitude	Transition amplitude for the emission of a multipole ( $\mathbb{M}$ ) photon.
$\langle \alpha_f \mathbb{J}_f \parallel \mathbb{O}^{(\mathbb{M}, \text{absorption})} \parallel \alpha_i \mathbb{J}_i \rangle$	Radiative.amplitude	Transition amplitude for the absorption of a multipole ( $\mathbb{M}$ ) photon.
$\langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J}_t \parallel \mathbb{O}^{(\mathbb{M}, \text{photoionization})} \parallel \alpha_i \mathbb{J}_i \rangle$	PhotoIonization.amplitude	Photoionization amplitude for the absorption of a multipole ( $\mathbb{M}$ ) photon and the release of an electron into the partial wave $ \varepsilon \kappa\rangle$ .
$\langle \alpha_f \mathbb{J}_f \parallel \mathbb{O}^{(\mathbb{M}, \text{recombination})} \parallel (\alpha_i \mathbb{J}_i, \varepsilon \kappa) \mathbb{J}_t \rangle$	PhotoRecombination.amplitude	Photorecombination amplitude for the emission of a multipole ( $\mathbb{M}$ ) photon and the capture of an electron that comes <i>in</i> the partial wave $ \varepsilon \kappa\rangle$ .
$\langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J}_t \parallel \mathbb{V}^{(\text{Auger})} \parallel \alpha_i \mathbb{J}_i \rangle$	Auger.amplitude	Auger transition amplitude due to the electron-electron interaction and the release of an electron into the partial wave $ \varepsilon \kappa\rangle$ .
$\langle \alpha_f \mathbb{J}_f \parallel \sum \exp i \mathbf{q} \cdot \mathbf{r}_i \parallel \alpha_i \mathbb{J}_i \rangle$	FormFactor.amplitude	Amplitude for a momentum transfer $\mathbf{q}$ with an external particle of photon field.
$\langle \alpha_f \mathbb{J}_f \parallel \mathbb{H}^{(\text{weak} - \text{charge})} \parallel \alpha_i \mathbb{J}_i \rangle$	PNC.weakChargeAmplitude	Amplitude for the nuclear-spin independent Hamiltonian of the ( $P$ -odd, $T$ -even) interaction.
$\langle \alpha_f \mathbb{J}_f \parallel \mathbb{H}^{(\text{Schiff} - \text{moment})} \parallel \alpha_i \mathbb{J}_i \rangle$	PNC.schiffMomentAmplitude	Amplitude for the nuclear Schiff moment of the ( $P$ -odd, $T$ -odd) interaction.
$\langle \alpha_f \mathbb{J}_f \parallel \mathbb{H}^{(\text{scalar} - \text{pseudo} - \text{scalar})} \parallel \alpha_i \mathbb{J}_i \rangle$	PNC.scalarPseudoScalarAmplitude	Amplitude for the scalar-pseudo-scalar ( $P$ -odd, $T$ -odd) interaction.

## In Jac (partly) implemented atomic level properties:

- In JAC implemented or partly-implemented atomic properties. For these properties, different parameters (observables) can generally be obtained by performing an `Atomic.Computation(..., properties=[id1, id2, ...])`, if one or more of the given identifiers are specified. For each of these properties, moreover, the corresponding (default) `Settings` can be overwritten by the user in order to control the detailed computations.

Property	id	Brief explanation.
$ \alpha\mathbb{J}\rangle \longrightarrow  \alpha(J)\mathbb{F}\rangle$	HFS	Hyperfine splitting of an atomic level into hyperfine (sub-) levels with total angular momentum $F =  I - J , \dots, I + J - 1, I + J$ ; hyperfine $A$ and $B$ coefficients; hyperfine energies and interaction constants; representation of atomic hyperfine levels in a $IJF$ -coupled basis.
$ \alpha\mathbb{J}\rangle \longrightarrow  \alpha\mathbb{J}M\rangle$	LandeJ	Zeeman splitting of an atomic level into Zeeman (sub-) levels; Lande $g_J \equiv g(\alpha\mathbb{J})$ and $g_F \equiv g(\alpha\mathbb{F})$ factors for the atomic and hyperfine levels.
$K^{(\text{MS})}, F$	Isotope	Isotope shift of an atomic level for two isotopes with mass numbers $A, A'$ : $\Delta E^{AA'} = E(\alpha\mathbb{J}; A) - E(\alpha\mathbb{J}; A')$ ; mass-shift parameter $K^{(\text{MS})}$ and field-shift parameter $F$ .
$\alpha$ -variations	AlphaX	Differential sensitivity of an atomic level $ \beta\mathbb{J}\rangle$ with regard to variation of the fine-structure constant; $\Delta E(\delta\alpha; \beta\mathbb{J}), \Delta q(\delta\alpha; \beta\mathbb{J}), K(\beta\mathbb{J})$ .
$F(q; \alpha\mathbb{J})$	FormF	Standard and modified atomic form factor of an atomic level $ \alpha\mathbb{J}\rangle$ with a spherical-symmetric charge distribution.
$\omega(\alpha\mathbb{J}) + a(\alpha\mathbb{J}) = 1$	Yields	Fluorescence & Auger decay yields of an atomic level, or averaged over an electron configuration.
$\alpha^{(\mathbb{M})}(\omega)$		Static and dynamic (ac, multipolar) polarizabilities.
$E(\alpha\mathbb{J}; \text{plasma model})$	Plasma	Plasma shift of an atomic level as obtained for different but still rather simple plasma models.
$ \alpha_i\mathbb{J}_i\rangle \longrightarrow  \alpha_f\mathbb{J}_f\rangle + \hbar\omega$	EinsteinX <sup>a</sup>	Photon emission from an atom or ion; Einstein $A$ and $B$ coefficients and oscillator strength between levels $ \alpha_i\mathbb{J}_i\rangle \rightarrow  \alpha_f\mathbb{J}_f\rangle$ that belong to a single multiplet (representation).

<sup>a</sup> Although the Einstein coefficients are not the property of a single level, we here still support a quick computation of these coefficients by means of the `Einstein` module for pairs of levels that are represented within a single CSF basis.

## In Jac (partly) implemented atomic processes:

- In JAC implemented or partly-implemented atomic processes. For *one* process at a time, different parameters (observables) can generally be obtained by performing an `Atomic.Computation(..., process=id)`, if the corresponding identifier is specified. For this selected property, moreover, the corresponding (default) `Settings` can be overwritten by the user in order to control the detailed computations.

Process	id	Brief explanation
$A^* \rightarrow A^{(*)} + \hbar\omega$	RadiativeX	Photon emission from an atom or ion; transition probabilities; oscillator strengths; angular distributions.
$A + \hbar\omega \rightarrow A^*$	PhotoExc	Photoexcitation of an atom or ion; alignment parameters; statistical tensors.
$A + \hbar\omega \rightarrow A^{++} + e_p^-$	PhotoIon	Photoionization of an atom or ion; cross sections; angular parameters; statistical tensors.
$A^{q+} + e^- \rightarrow A^{(q-1)+} + \hbar\omega$	Rec	Photorecombination of an atom or ion; recombination cross sections; angular parameters.
$A^{q+*} \rightarrow A^{(q+1)+(*)} + e_a^-$	AugerX	Auger emission (autoionization) of an atom or ion; rates; angular and polarization parameters.
$A^{q+} + e^- \rightarrow A^{(q-1)+*} \rightarrow A^{(q-1)+(*)} + \hbar\omega$	Dierec	Dielectronic recombination (DR) of an atom or ion; resonance strengths.
$A + \hbar\omega_i \rightarrow A^* \rightarrow A^{(*)} + \hbar\omega_f$	PhotoExcFluor	Photoexcitation of an atom or ion with subsequent fluorescence emission.
$A + \hbar\omega \rightarrow A^* \rightarrow A^{(*)} + e_a^-$	PhotoExcAuto	Photoexcitation & autoionization of an atom or ion.
$A + \hbar\omega_i \rightarrow A^{(*)} + \hbar\omega_f$	Compton	Rayleigh or Compton scattering of photons at an atom or ion; angle-differential and total cross sections.
$A + n\hbar\omega \rightarrow A^*$ or $A^* \rightarrow A^* + n\hbar\omega$	MultiPhoton	Multi-photon (de-) excitation of an atom or ion, including two-photon decay, etc.
$A + Z_p \rightarrow A^* + Z_p$	CoulExc	Coulomb excitation of an atom or ion by fast, heavy ions; energy-differential, partial and total Coulomb excitation cross sections.
$A + \hbar\omega \rightarrow A^* + e_p^- \rightarrow A^{(*)} + e_p^- + \hbar\omega'$	PhotoIonFluor	Photoionization of an atom or ion with subsequent fluorescence emission.
$A + \hbar\omega \rightarrow A^* + e_p^- \rightarrow A^{(*)} + e_p^- + e_a^-$	PhotoIonAuto	Photoionization of an atom or ion with subsequent autoionization.
$A^{q+} + e^- \rightarrow A^{(q-1)+*} \rightarrow A^{(q-1)+(*)} + \hbar\omega$ $\rightarrow A^{(q-1)+} + \hbar\omega + \hbar\omega'$	DierecFluor	Dielectronic recombination of an atom or ion with subsequent fluorescence.

Process	id	Brief explanation
$e_s^- + A \longrightarrow A^* + e_s^{-'}$	Eimex	Electron-impact excitation of an atom or ion; collision strength.
$A + e_s^- \rightarrow A^* + e_s^{-'} \rightarrow A^{+(*)} + e_s^{-'} + e_a^-$	EimexAuto	Electron-impact excitation and subsequent autoionization of an atom or ion.
$A^{q+*} \longrightarrow A^{(q+1)+(*)} + (e_a^- + \hbar\omega)$	RadAuger	Radiative-Auger (autoionization) of an atom or ion.
$A + n \hbar\omega \longrightarrow A^{(*)} + e_p^-$	MultiIon	Multi-photon ionization of an atom or ion.
$A + n \hbar\omega \longrightarrow A^{(*)} + e_{p_1}^- + e_{p_1}^-$	MultiDoubleIon	Multi-photon double ionization of an atom or ion.
$A^{q+}[\text{nucleus}^*] \longrightarrow A^{(q+1)+*} + e_c^-$	Conversion	Internal conversion, i.e. electron emission due to nuclear de-excitation.

## 1.4. Short comparison of JAC with other existing codes

We here compile some (incomplete) information about other existing atomic structure codes for the computation of level energies, transition rates, cross sections, etc. We remind to some of their special features and briefly summarize how these codes differ from the implementation of the JAC toolbox.

### CATS (Cowan: 'Theory of Atomic Spectra', 1980):

- CATS relies on a semi-relativistic self-consistent potential and by using either a nonlocal Hartree-Fock (HF) or local Hartree-Fock-Slater (HFS) approach in order to deal with the exchange interaction.
- **Level energies & ASF:** Since the late 1960s, Cowan's HFX code has set some standard for many experimentalists and has, together with his well-known textbook, helped many (atomic) physicist to understand and make use of atomic structure theory. While these earlier developments are highly appreciated (and are still utilized for various applications), **CATS has several severe limitations in the layout and implementation of the code, which are hard to overcome.** The same applies also for the computation of various atomic properties, such as transition probabilities, photo-excitation and ionization cross sections and Auger rates.
- Cowan's code has been found a mature tool for identifying new lines, especially if additional information is available from experimental observations to support empirical adjustments.

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- **Display of data & spectra:** Several tools and facility-programs have been developed for CATS in order to display the computed data and to compare them with each other and with experiment. The success of some of these tools has or will motivate us for developing some graphical interfaces for the display of (radial) orbitals, line spectra, etc. also for the JAC toolbox.
- **Selected advantages of JAC:** JAC provides methods to display, for instance, radial functions and line spectra; cf. `JAC.display()`.

## GRASP (Grant *et al.*, 1980; Jönsson *et al.*, 2013; Fischer *et al.*, 2018):

- **Level energies & ASF:** GRASP has originally been developed since the late 1960s in order to provide level energies and eigenvectors for quite general open-shell atoms. Much emphasis during the last two decades was placed upon the systematic improvement of these energies and representations. While we also provide such level energies and atomic state functions by the JAC toolbox, we *do not* intend to facilitate such extensive wave function expansions. Instead, **approximate level energies and ASF are mainly considered as the *technical preposition for describing further atomic properties and processes***. With the restricted-active space (RAS) computations, however, we shall provide in JAC useful features for a (more or less) systematic improvement of such *ab-initio* computations.
- **Transition probabilities & oscillator strength:** Apart from the level energies, GRASP has been extensively applied in order to compute and tabulate transition probabilities for many atoms, ions and isoelectronic sequences throughout the periodic table of elements. **With JAC, we provide analogue or even simpler tools for such computations**. Moreover, (the many-electron amplitudes that arise from) the coupling of the radiation field provides the natural *building blocks* for a large number of other atomic processes, cf. section 7 on atomic processes below.
- **Selected advantages of JAC:**
  - JAC supports larger flexibility in handling the output data and, in particular, does not know practically-relevant limitations with regard to the length of filenames (in contrast, for instance, to 24 letters in grasp2K).
  - SCF fields can be generated in JAC at different levels of complexity, including several (local) mean-field potentials as well as, in the future, the average-level and extended average-level schemes.
  - JAC enables one to handle a much larger number of atomic properties and processes as well as atomic cascades and several other types of computation.
  - The use of the JULIA language clearly facilitates the coding and maintenance of the JAC code, when compared to previous Fortran codes.

### RATIP (Fritzsche, 2001, 2012):

- **Relativistic CI (RELCI):** While RATIP has always used the SCF computations and orbitals from the GRASP code, it also supports relativistic CI computations. For several years, it helped define a new standard for performing the angular intergration (angular coefficients; cf. Gaigalas *et al.*, 2002). These angular coefficients are also utilized in JAC by an interface to the Fortran modules of RATIP.
- **Atomic properties and processes:** RATIP was (one of) the first codes that made use of GRASP's systematically improved wave functions in order to compute a good number of atomic properties and processes, such as relaxed-orbital transition probabilities (REOS; Fritzsche and Froese Fischer 1999), Auger rates, photoionization cross sections and angular parameters, radiative and dielectronic recombination rates, electron-impact excitation cross sections, and several others. *The experience with RATIP has been found central for the development of JAC and has find its continuation here.*

### HULLAC (Bar-Shalom *et al.*, 2001):

- HULLAC has been developed as an integrated code for calculating atomic structure and cross sections for collisional and radiative atomic processes, based on the relativistic configuration interaction method.
- All collisional cross sections are calculated in the distorted wave approximation with special emphasis on efficiency.
- A **parametric potential method** is applied for the generation of both, the bound and free orbitals, while a (so-called) factorization-interpolation method is utilized in order to derive all collisional rates.

### LADW (Bar-Shalom *et al.*, 2001):

- LADW, the Los Alamos Distorted-Wave code, has been developed by Sampson and co-workers; it has been further utilized and partly incorporated into the LASER code.

### ELSEPA (Salvat *et al.*, 2005):

- ELSEPA facilitates the computation of differential cross sections, total cross sections and transport cross sections for (single) elastic scattering of electrons and positrons by neutral atoms, positive ions as well as randomly oriented molecules. Moreover, this code help compute scattering amplitudes and spin polarization functions for projectiles with kinetic energies  $E_{\text{kin}} \lesssim 5$  MeV by means of a conventional relativistic Dirac partial-wave analysis.

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- For high kinetic energies, however, the convergence of the partial-wave series becomes slow and is then replaced by some approximate factorization method. ELSEPA has been found useful to explore the sensitivity of DCS for different choices of the (static) interaction potential.

### GEANT4 (Amako *et al.*, 2005):

- GEANT4 is an object-oriented toolkit for analyzing and simulating the passage of particles through matter that provides a **variety of semi-empirical models to describe the underlying electromagnetic and hadronic interactions**. GEANT4 combines theoretical models with experimental data or parameterizations of such data.
- GEANT4 is especially based on a number of separate packages to deal with the electromagnetic interactions of (either) electrons, muons, positrons, photons, hadrons and ions as well as for specific energy range of the processes considered.
- **Implemented processes:** The electromagnetic packages of this code include: multiple scattering, ionization, Bremsstrahlung, positron annihilation, photoelectric effect, Compton and Rayleigh scattering, pair production, synchrotron and transition radiation, Cherenkov effect, refraction, reflection, absorption, scintillation, fluorescence as well as Auger electron emission (Amako *et al.*, 2005). **Less attention has been placed however on the electronic structure of atoms and ions.**

### FAC (Gu 2008):

- FAC is a well-known relativistic atomic structure code based on the fit of free parameters in order to define the atomic potentials and is mainly based on the (standard) Dirac-Fock-Slater method.
- **Level energies & wave functions:** The simplified and more object-oriented treatment of wave functions in the FAC code, when compared with GRASP, has stimulated the development of the JAC program. While JAC will enable the user to perform also systematic improvements on the underlying computational models, the support of rather simple approximations is crucial and need to be supported, for example, for dealing with cascades and time-evolutions.

### AUTOSTRUCTURE (Badnell, 2011):

- AUTOSTRUCTURE is a rather general atomic code for the description of free-bound electron and photon collision processes, based on the original SUPERSTRUCTURE code by Eissner and coworkers.
- AUTOSTRUCTURE supports efficient computations of dielectronic recombination cross sections and rates, especially if large numbers of highly-excited states are involved in the radiative stabilization of an atom or ion.



- The code applies the Breit-Pauli distorted wave method for the electron-impact excitation of atomic ions in order to support problems that are impractical or even impossible for more sophisticated methods.
- AUTOSTRUCTURE mainly computes (Maxwell-averaged) effective collision strengths at temperatures of broad ionic abundance, rather than the detailed collision strength at all the energies.

### LASER (Fontes *et al.*, 2015):

- LASER, i.e. the Los Alamos suite of relativistic atomic physics code, comprises various codes for fundamental atomic structure calculations as well as for various processes, such as photoexcitation, electron-impact excitation and ionization, photoionization and autoionization, within a consistent framework. It may help develop atomic physics models in either configuration- average and fine-structure modes, and by including a proper self-consistency. This suite has been developed for more than 20 years.
- **Applications of the code:** The LASER code has been applied to the collisional-radiative modeling of plasmas, for line identifications in plasma spectroscopy and for testing relativistic atomic and quantum electro-dynamics (QED) theories. The code has been applied also for feasibility studies of the collisional-radiative modeling of non-LTE (optically thin) gold plasmas for and to questions from inertial confinement fusion.
- **Approximations:** LASER mainly employs the semi-relativistic theory, similar and often by directly applying Cowan's atomic structure code (CATS). The bound-electron wavefunctions are obtained from the semi-relativistic approach in CATS, while the continuum-electron wavefunctions are obtained as solutions of the Schrödinger equation from some specialized routines. Typically, the bound and continuum radial wave functions are single-component type wavefunctions associated with the Schrödinger equation, rather than the four-component spinors associated with the Dirac equation.
- **Features:** Cross sections and other properties can be calculated for five fundamental processes: photo-excitation, photo-ionization, electron-impact excitation and ionization and autoionization. The code supports the IPCRESS (Independent of Platform and Can be Read by Existing Software Subroutines) random-access binary file format that is used to store large amounts of data, and which can be ported to any platform.

### QEDMOD (Shabaev *et al.*, 2015):

- **Self energy,  $\mathfrak{h}^{(\text{QED})}$ :** Computes a model QED operator  $\mathfrak{h}^{(\text{QED})}$  that accounts for the Lamb shift in accurate atomic-structure calculations. However, there are various difficulties with QEDMOD which make a direct application of the code and its combination with GRASP or JAC rather cumbersome.

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- **Effective QED Hamiltonian:** The QEDMOD code provides one-electron matrix elements that, in principle could be directly added to any CI matrix and, hence, to the computation of level energies and multiplets. For the vacuum polarization, it includes automatically both, the Uehling and Wichmann-Kroll terms.
- **Selected advantages of JAC:** A simplified version of the effective QED Hamiltonian from QEDMOD has been implemented also in JAC.

## 1.5. To-do's, next steps & desired features of the JAC program

### Encouragement for external users and developers:

- While we (will further) develop JAC for those applications, which are requested frequently by the users, here I shall compile a number of *desired features which will make JAC even more powerful and/or easy to use*. For these *additional features*, I wish to encourage collaboration with external developers. We welcome in particular all help from outside if the overall style of the program is maintained, and if some prior consensus exist how to add and implement additional features.
- New code developments may concern incremental improvements or also multiple approaches for algorithms and modules in order to provide well-defined alternatives, for instance, if some particular approach does not work properly.
- Emphasis will be placed first usually on those applications that receive enough attention by the community.

### 1.5.a. To-do lists

#### Urgent to-do's:

- Visualize a (given) continuum orbital and its normalization as obtained at sufficiently large  $r$ -values.
- *Use of JAC on remote clusters:* Work out some prototype example how 'job scripts' (similar to those from examples) can be exported and handled at remote cluster computers and how to re-import the results later on.
- *Documenter.jl:* How to establish a documentation of JAC.

### Short-term to-do's:

- Implement ... `Jac.modify("level energies: interactive", multiplet::Multiplet)`
- Implement ... `Jac.display("level energies: HFS", multiplets::HFSMultiplet[..])`
- Implement ... `Jac.display("level energies", multiplets::Multiplet[..])`
- Implement ... `Jac.display("configuration list: from basis", basis::Basis)`

### Medium-term to-do's:

- Implement ... `Jac.apply("restrictions: CSF list", csfs::CsfR[..], basis::Basis)` ... to apply a number of restrictions interactively to a list of CSF. The procedure proceeds in three steps: (i) by taking and applying a restriction to a given CSF list; (ii) showing the number of CSF to be deleted from the given list; (iii) making this restrictions explicit. The user is requested to enter one restriction after the other, and until the reduction process is terminated by the user. A `csfList::CsfR[]` is returned; ..... Here we might adopt the form of restrictions from the RATIP program.
- Implement ... `Jac.apply("biorthogonal transformation", mpltA::Multiplet, mpltB::Multiplet, grid::Radial.Grid)`
- **Atoms in plasma environments:** Implement 2-3 plasma models in order to deal with atoms in a few (averaged but) different plasma environments. This usually works via some effective `Jac.InteractionStrength.XL.BreitXL.plasma_ionSphere(L::Int64, a::Orbital, b::Orbital, c::Orbital, d::Orbital, lambda::Float64)` that depends on some particular model and plasma parameters; cf. Saha and Fritzsche, PRE (2004). This will likely require also the set-up of a corresponding Hamiltonian matrix: `Jac.compute("matrix: plasma, ion-sphere model", settings::Plasma.Settings, basis::Basis)`.
- **Interpolation:** We might need (from time to time) a proper interpolation of functions from one grid to another, for instance, for using GRASP-type orbitals that have been generated on a different grid. Implement some function `Jac.interpolate("function: for new radial grid, trapez rule", from::Tuple(rOld::Float64[..],gridOld::Jac.Radial.Grid), to::Tuple(rNew::Float64[..],gridNew::Jac.Radial.Grid) )`.

## 1.5.b. Discussion about (further) implementations

### Issues that need to be discussed:

- **Further documentation of the code:** How can the internal *doc-strings* be readily combined with the JAC websites ??
- **Parallelization and performance of the code:** How can one make the code parallel without that the user need to know and provide much information about the cluster that is used for the computations.
- **Modern input forms in scientific computing:** Which modern formats do exist ?? Which simple graphical (applet) features exist ?  
tomel.jl ??

## 1.5.c. Desired medium- and long-term features of the JAC program

### Plotting and visualization:

- Implement ... `Jac.plot("radial orbitals", orbitals::Orbital[...])`
- Implement ... `Jac.plot("spectrum: oscillator strength over energy, emission", lines::RadiativeLine[...]; widths=value::Float64)` and `Jac.plot("spectrum: oscillator strength over energy, absorption", lines::RadiativeLine[...]; widths=value::Float64)`
- Implement ... `Jac.plot("spectrum: transition rates over energy, Gaussian", lines::RadiativeLine[...]; widths=value::Float64)` and `Jac.plot("spectrum: transition rates over energy, Lorentzian", lines::RadiativeLine[...]; widths=value::Float64)`
- Visualize the convergence of energies or other results as function of the size of the computation and/or model space.
- Visualize the level structure of a given multiplet, for instance, by displaying the level energies in different colors for different (leading) configurations or groups of such configurations.
- Visualize the level structure of a given multiplet together with further level or transition properties, such as lifetimes, HFS parameters, isotope parameters, etc.

## Excitation and decay cascades:

- Analyze (and report) the fine-structure level population following the decay of an inner-shell hole state.
- Compute and extract the Fano parameters and line-shapes for a given set of autoionizing resonances.
- For a given cascade (data), evaluate the ion yields, electron spectra, (fine-structure) level population, fluorescence spectra, etc.

## More physics in Jac ?

- **Atomic spectra:** Evaluate and display the photoabsorption spectra from calculated photoexcitation and photoionization cross sections.
- **Computation of approximate single-electron properties:**
  - Subshell-dependent differential and total photoionization cross sections; cf. Eichler and Meyerhof (1995, Eqs. 9.34 and 9.47).
  - Non-relativistic total K-shell or subshell radiative recombination cross sections by using the Stobbe cross section; cf. Eichler & Meyerhof (1995, Eqs. 9.49, 9.50)
  - Dirac energy (subshell); Dirac  $r^k$  expectation values; Dirac-matrices; `Dirac.Omega(subshell,theta,phi)`.
  - Coulomb-Greens functions for some given hydrogenic orbital.
- **Collisional-radiative models:** Such collisional-radiative models have been frequently applied to describe the evolution of plasma and to derive information for plasma diagnostics. JAC provides many, if not all, the rates and cross sections to built-up such models for selected (plasma) environments.
- **Electron-momentum distributions:** Provide the expectation values  $\langle p^k \rangle$ ,  $k = -2, \dots, 4$  of the single-electron radial momentals, i.e. the radial orbitals in momentum space. These expectation values are frequently applied in crystallography and in studying Compton profiles; cf. Koga and Thakkar (1996), Eq. (10-11). In the first instance, these expectation values could be readily provided as semi-empirical values by following the work above.

## 1.6. Remarks on the implementation of JAC

### Why Julia ?

- Here, we just recall a few remarks from the literature as well as some own experience **why JULIA have been found helpful for developing the JAC program**. Some of these arguments are directly taken from the work of Bezanson *et al.* (2017, 2018).
- **JULIA is a language for scientific computing that offers many of the features of productivity languages**, namely rapid development cycles; exploratory programming without having to worry about types or memory management; reflective and meta-programming; and language extensibility via multiple dispatch (Bezanson *et al.*, 2018).
- **Productivity vs. performance:** JULIA is often said to stand for the **combination of productivity and performance** through a careful language design and carefully chosen technologies; it never forces the user to resort to C or Fortran for fast computations. — JULIA's design allows for gradual learning of modern concepts in scientific computing; from a manner familiar to many users and towards well-structured and high-performance code.
- **JULIA's productivity features include:** dynamic typing, automatic memory management, rich type annotations, and multiple dispatch. JULIA also supports some control of the memory layout and just-in-time compilation in order to eliminate much of the overhead of these features above (Bezanson *et al.*, 2018).
- Julia promises scientific programmers the ease of a productivity language at the speed of a performance language (Bezanson *et al.*, 2018).
- **High-level languages:** Most traditional high-level languages are hampered by the overhead from the interpreter, and which typically results into more run-time processing of what is strictly necessary. One of these hindrances is **(missing) type information**, and which then results in the request for supporting vectorization. **JULIA is a 'verb'-based language in contrast to most object-oriented 'noun'-based language**, in which the generic functions play a more important role than the datatypes.
- **Language design:** Julia includes a number of (modern) features that are common to many productivity languages, namely dynamic types, optional type annotations, reflection, dynamic code loading, and garbage collection (Bezanson *et al.*, 2018).
- **Multiple dispatch:** This concept refers to the dynamically selected implementation and to running the right code at the right time. This is achieved by overloading by means of multiple-argument function, a very powerful abstraction. Multiple dispatch makes it easier to structure the programs close to the underlying science.
- Multiple dispatch is perhaps the most prominent feature of JULIA's design and is crucial for the performance of the language and its ability to inline code efficiently. Another promise of multiple dispatch is that it can be used to extend existing behavior with new features.

## 1. Overview about JAC. Structure of this User Guide

- **Multiple dispatch:** At run-time, a function call is dispatched to the most specific method applicable to the types of its arguments. JULIA's type annotations can also be attached to datatype declarations so that they can be checked whenever typed fields are assigned to. Multiple dispatch also help the programmers to extend the core languages functionalities to their particular needs.
- Multiple dispatch also reduces the needs for argument checking at the begin of a function. The overloading of functions by multiple dispatch is also called ad-hoc polymorphism. Instead of encapsulating methods inside classes, JULIA's multiple dispatch is a paradigm in which methods are defined on combinations of data types (classes). JULIA shows that this is remarkably well-suited for numerical computing.
- **JULIA's type system:** JULIA's expressive type system allows optional type annotations; this type system supports an aggressive code specialization against run-time types. To a large extent, however, JULIA code can be used without any mentioning of types (in contrast to C and Fortran); this is achieved by data-flow interference. — User's own types are also first class in JULIA, that is there is no meaningful distinction between built-in and user-defined types. There are mutable and (default: immutable) composite types.
- **Optimization in JULIA:** The JULIA compiler is built on three strategies that are performed on a high-level intermediate representation, while all native code generation is later delegated to the LLVM compiler infrastructure. These optimization strategies are: (1) method inlining which devirtualizes multi-dispatched calls and inline the call target; (2) object unboxing to avoid heap allocation; and (3) method specialization where code is special cased to its actual argument types (Bezanson *et al.*, 2018).
- **Performance:** There are helpful macros, such as `@timing function_call(parameters)` or `@benchmark function_call(parameters)` to analyze the performance of the program and to find (and resolve) bottlenecks.
- **Data types:** JULIA distinguishes between concrete data types, that can be instantiated, and abstract types, that can (only) be extended by subtypes to built up an hierarchy of such types.
- In JULIA, users are always encouraged to make their programs, whenever possible, type stable. Much of the efficiency of a JULIA code relies on being type stable and on devirtualization and inlining.
- **LAPACK:** All of LAPACK is available in JULIA, not just the most common functions. LAPACK wrappers are fully implemented by `ccall` and can be called directly from the Julia prompt.

### Requests in building large software packages:

- These and further requests have been summarized by Post and Kendall (2004).
- **Physical models:** In general, better physics is more important than better computer science. It is recommended to use modern but well-proven computer-science techniques, and a 'physics code' should not be a computer-science research project. Instead, one should use best



engineering practices to improve quality rather than processes. Emphasis should be given to improvements of the physics capabilities. Do not use the latest computer-science features; let the new ideas mature first. **Better physics is the most important product of the code.**

- **Code development and evolution:** The scale of code-development can become truly immense; a good overview/quantitative database about (previously) successful software projects in some given field is typically required for good estimation for resources and schedules. **It is easy to loose motivation on a project that last years and which has few incremental deliveries. Continues replacement of code modules is recommended as better tools and techniques are developed.** Every code development typically proceeds in steps: First develop a core capability (with a small team) and let this small core be tested by users and, if successful, add further capabilities (so-called incremental delivery).
- **Success criteria:** One of the important success criteria is the costumer focus. — **What do the user really need ?**
- **Code specification:** Some flexibility in the requirement specification phase is essential because it is difficult to predict when (or if) a new algorithm/approach will be available. **There is a need to pursue multiple approaches for algorithms and modules near to the critical path.** If one approach is not feasible, another one can be used.
- **External users and developers:** The validation of physics models is a critical issue for all extended software packages and need to be done with regard to authoritative reference data.
- Documentation should be intended to cover all needs of the beginner through to the expert user who wishes to expand the capabilities of JAC.



## 2. Dirac's hydrogenic atom

### 2.1. Energies and wave functions

#### Notations & application in Jac:

- In JAC, non-relativistic hydrogenic orbitals can be easily generated by a call to `JAC.HydrogenicIon.radialOrbital()`.
- JAC also provides approximate radial (relativistic bound-state) orbitals by just applying the kinetic-balance condition to the corresponding non-relativistic radial orbital and by re-normalizing it afterwards.

#### Dirac's (one-electron) energies:

- **Relativistic bound-state energy spectrum:** For electrons with principal quantum number  $n$  and angular-momentum quantum number  $\kappa$ , the relativistic bound-state energy spectrum is given by

$$\varepsilon_{n\kappa} = mc^2 W_{n\kappa} = \frac{mc^2}{\left[1 + \left(\frac{\alpha Z}{n' + s}\right)^2\right]}$$

where  $\alpha$  is the fine-structure constant,  $n' = n - |\kappa| = 0, 1, 2, \dots$  is the number of nodes, and  $s = \sqrt{\kappa^2 - (\alpha Z)^2}$ , respectively.

- Since  $W_{n\kappa} < 1$ , it can be expanded in terms of  $(\alpha Z)$  and written as

$$W_{n\kappa} = 1 - \frac{1}{2} \frac{(\alpha Z)^2}{n^2} - \frac{1}{2} \frac{(\alpha Z)^4}{n^3} \left( \frac{1}{j + 1/2} - \frac{3}{4n} \right) - \dots$$

where the second term, multiplied with  $mc^2$ , represents the nonrelativistic binding energy of a hydrogenic atom.

## 2. Dirac's hydrogenic atom

- For a given (principal) shell  $n$ , the eigenvalues are the same for equal values of  $j$  but are different for equal values of  $\ell$ . For a given value of  $\ell$ , the spin-orbit splitting between states with  $j = \ell + 1/2$  and  $j = \ell - 1/2$  gives rise to the **fine-structure in the spectrum of hydrogen-like atoms**. The equation above shows that the relativistic corrections to the one-electron energies decreases rapidly with  $n$ . The relativistic correction to the non-relativistic energy is therefore important, especially for highly-charged ions with rather large  $Z$ .

### Non-relativistic radial orbitals:

- **Radial orbital functions  $P(r)$ :** These functions in the non-relativistic orbital  $\psi(r, \vartheta, \varphi) = \frac{P(r)}{r} Y_{\ell m}(\vartheta, \varphi)$  is known to obey the *radial* Schrödinger equation and belongs to either the discrete part of the spectrum with the (negative) energies  $E_n = -\frac{Z^2}{2n^2}$  a.u.  $< 0$  (bound states), or to the continuous part for all energies  $E > 0$  (the so-called *free-electron* or *continuum* states).
- **Radial orbital functions  $P(r)$ :** An analytic solution of the radial functions  $P(r)$  in terms of the **confluent hypergeometric function  $F(\alpha, \beta; x)$**  are known for both, the bound-states

$$P_{n\ell}(r) = r^{l+1} \frac{1}{(2\ell+1)!} \sqrt{\frac{(n+\ell)!}{(n-\ell-1)! 2n}} \left(\frac{2Z}{n}\right)^{3/2+\ell} e^{(-\frac{Zr}{n})} F\left(-(n-\ell-1), 2\ell+2; \frac{2Zr}{n}\right)$$

as well as for the *continuum* ( $E > 0$ ) with  $k = \sqrt{2E}$  and  $n' = Z/k$

$$P_{E\ell}(r) = \frac{2\sqrt{Z}}{\sqrt{1-e^{-2\pi n'}}} \left[ \prod_{s=1}^{\ell} \sqrt{s^2 + n'^2} \right] \frac{(2kr)^{\ell}}{(2\ell+1)!} e^{-ikr} F(in' + \ell + 1, 2\ell + 2, 2ikr).$$

### Relativistic radial orbitals:

- **Radial orbital functions  $P(r)$ ,  $Q(r)$ :** Analogue to the non-relativistic case, the solution  $\psi(\mathbf{r})$  of the Dirac equation can be separated for a spherical potential into a radial and angular part

$$\psi_{\kappa m}(r, \vartheta, \varphi) = \frac{1}{r} \begin{pmatrix} P(r) \Omega_{\kappa m}(\vartheta, \varphi) \\ i Q(r) \Omega_{-\kappa m}(\vartheta, \varphi) \end{pmatrix},$$

where  $\Omega_{\kappa m}(\vartheta, \varphi)$  denotes a **standard Dirac spin-orbital function**, and where  $\kappa = \pm(j + 1/2)$  for  $\ell = j \pm 1/2$  is called the *relativistic* angular-momentum quantum number.

- **Relativistic angular-momentum quantum number:** Owing to the definition of  $\kappa = \pm 1, \pm 2, \dots$ , this relativistic quantum number carries information about both, the total angular momentum  $j$  and the parity  $(-1)^\ell$  of the orbital function.
- **Large & small component:** The radial part of a relativistic wavefunctions is given by the two functions  $P(r)$  and  $Q(r)$ , and which are often called the *large* and *small* components.
- **Dirac spin-orbital:** As usual, the Dirac spin-orbitals can be written in terms of the spin-1/2 Pauli spinors  $\chi_\pm$  and the spherical harmonics by

$$\Omega_{\kappa m}(\vartheta, \varphi) = \sum_{m_\ell} \langle \ell m_\ell, 1/2, m - m_\ell | jm \rangle Y_{\ell m_\ell}(\vartheta, \varphi) \chi_{m-m_\ell}.$$

$$Q(r) = \frac{1}{2m c} \left( \frac{d}{dr} + \frac{\kappa}{r} \right) P(r).$$

## 2.2. Coulomb-Green function

### Relativistic radial Coulomb-Greens function:

- **Radial Green function of the Dirac equation:** Following Yerokhin and Shabaev (1999, appendix D), the radial Green function of the Dirac equation for electrons with the angular-momentum symmetry  $\kappa$  can be written in the form:

$$G_\kappa(\omega, r_1, r_2) = -\frac{1}{W_\kappa(\omega)} \left[ \phi_\kappa^\infty(\omega, r_1) \phi_\kappa^{0T}(\omega, r_2) \Theta(r_1 - r_2) + \phi_\kappa^0(\omega, r_1) \phi_\kappa^{\infty T}(\omega, r_2) \Theta(r_2 - r_1) \right]$$

where  $\phi_\kappa^0(\omega, r)$  and  $\phi_\kappa^\infty(\omega, r)$  are solutions of the radial Dirac equation, bounded at the origin and at the infinity, and where  $W_\kappa(\omega)$  is the Wronskian.

- The functions  $\phi_\kappa^0(\omega, r)$  and  $\phi_\kappa^\infty(\omega, r)$  can be calculated in terms of the Whittaker functions of the first and second kind, although special care has to be taken.

## 2.3. Matrix elements with Dirac orbitals

### 2.3.a. Matrix elements with radial orbitals

#### Matrix elements with non-relativistic radial orbitals:

➤ **Special non-relativistic  $r^k$  expectation values:** The following expectation values are displayed by Marxer (1991)

$$\langle r^{-1} \rangle = \frac{Z}{a_o n^2}$$

$$\langle r^{-6} \rangle = \left[ \frac{Z}{a_o} \right]^6 \frac{35n^4 - n^2 [30\ell(\ell+1) - 25] + 3(\ell-1)\ell(\ell+1)(\ell+2)}{8n^7 (\ell-3/2)(\ell-1)(\ell-1/2)\ell(\ell+1/2)(\ell+1)(\ell+3/2)(\ell+2)(\ell+5/2)}$$

$$\langle r^{-k-2} \rangle = \left[ \frac{Z}{a_o} \right]^{k+2} \frac{1}{n^{k+3} \ell^{k+1}} \mathcal{P}_{k,\ell} \left( \frac{n}{\ell} \right), \quad \mathcal{P}_{k,\ell}(x) = \frac{(2\ell)^{k+1} (2\ell-k)!}{(2\ell+1)!} {}_3F_2(-k, k+1, \ell+1-\ell x; 1, 2\ell+2; 1)$$

➤ **Pasternack-Kramers rekursion relation:** For  $k \leq 0$ , the  $r^{-k}$  expectation values of any  $(n\ell)$  level fulfill the relation (Marxer 1991)

$$\langle r^{-k-2} \rangle = \frac{4}{(2\ell+1)^2 - k^2} \left[ \left( \frac{2k-1}{k} \right) \frac{Z}{a_o} \langle r^{-k-1} \rangle - \left( \frac{k-1}{k} \right) \frac{Z^2}{n^2 a_o^2} \langle r^{-k} \rangle \right]$$

$$\langle r^k \rangle = \frac{(2\ell+k+2)!}{(2\ell-k-1)!} \left[ \frac{n a_o}{2Z} \right]^{2k+3} \langle r^{-k-3} \rangle$$

From this relation, explicit formulas for low principal quantum number  $n$  can be computed rather easily.

#### Matrix elements with relativistic Dirac wavefunctions:

➤ An analytical expression of the matrix elements with operators  $r^k e^{-sr}$  can be found (also) for Dirac's relativistic wavefunctions. In order to derive such expressions, it is typically more convenient first to re-write the standard representation of the hydrogenic functions in terms of a series expansion in  $r$  as it was suggested originally by Rose (1961).

➤ An expression for the corresponding matrix element is given by:

$$\begin{aligned}
\langle n \kappa | r^k e^{-\sigma r} | n' \kappa' \rangle &= \int_0^\infty dr e^{-\sigma r} (P_{n\kappa}(r) P_{n'\kappa'}(r) + Q_{n\kappa}(r) Q_{n'\kappa'}(r)) \\
&= N_{n\kappa} N_{n'\kappa'} \sum_{\nu=0}^{n-|\kappa|} \sum_{\nu'=0}^{n'-|\kappa'|} (c_{n\kappa,\nu}^+ c_{n'\kappa',\nu'}^+ + c_{n\kappa,\nu}^- c_{n'\kappa',\nu'}^-) \frac{q^{s+\nu-1} q'^{s'+\nu'-1}}{(q+q'+\sigma)^{1+\nu+\nu'+s+s'+k}} \Gamma(\nu + \nu' + k + s + s' + 1), \\
c_{n\kappa,\nu}^\pm &= \left(1 \pm \sqrt{1-q^2}\right)^{1/2} \frac{((-n+|\kappa|)_\nu 2^\nu)}{\nu! (2s+1)_\nu} [(\nu - n + |\kappa|) \pm (\alpha Z/q - \kappa)]
\end{aligned}$$

with  $(a)_\nu$  being the Pochhammer symbol and where the parameters  $q$ ,  $s$  and the normalization factor  $N_{n\kappa}$  were defined above.

## 2.3.b. Matrix elements including the angular part of Dirac orbitals

### Matrix elements of spherical tensors:

➤ **Matrix elements of the  $\mathbb{C}^K$  tensors:** Gaidamauskas *et al.* (2011) use the following definition:

$$\langle \kappa_a || \mathbb{C}^{(K)} || \kappa_b \rangle = (-1)^{j_a + 1/2} [j_a, j_b]^{1/2} \begin{pmatrix} j_a & K & j_b \\ 1/2 & 0 & -1/2 \end{pmatrix} \delta_{\ell_a + \ell_b + 1, \text{even}}$$

## 2.4. Frequently applied expansions and identities in atomic theory

### 2.4.a. Partial-wave expansions of free electrons

#### Partial-wave expansion:

- **Partial-wave components of a plane-wave electron:** The computation of partial (and total) ionization cross sections often requires an integration over all possible angles  $\Omega_p = (\vartheta_p, \varphi_p)$  of the photoelectrons, emitted in  $4\pi$ . In practice, this integration over  $\Omega_p$  can be carried out rather easily by making use of the *decomposition* of the free-electron wavefunction  $|\mathbf{p} m_s\rangle$  into *partial-wave* components.
- In practice, however, this expansion generally depends on the choice of the quantization axis.
- **If the quantization axis is taken along  $\mathbf{e}_z$ ,** this expansion is given by

$$|\mathbf{p} m_s\rangle = \sum_{\kappa m} i^\ell e^{-i\Delta_\kappa} \langle lm_\ell, 1/2 m_s | jm\rangle Y_{\ell m_\ell}^*(\vartheta_p, \varphi_p) |\varepsilon \kappa m\rangle ,$$

and where the summation runs over all partial waves,  $\kappa = \pm 1, \pm 2, \dots$  as well as  $m = -j, \dots, j$ . In this expansion, moreover,  $\Delta_\kappa$  is the **Coulomb phase shift**, and the (nonrelativistic orbital angular momentum) quantum number  $\ell$  just distinguishes the parity of the partial waves.

### 2.4.b. Expansions including spherical harmonis

#### Expansions including the spherical harmonics:

- **Expansion of a plane-wave in terms of spherical Bessel functions:**

$$e^{i\mathbf{k}\cdot\mathbf{r}} = 4\pi \sum_{\ell m} i^\ell j_\ell(kr) Y_{\ell m}^*(\hat{k}) Y_{\ell m}(\hat{r}) .$$

- **Multipole expansion of the (transverse-gauge) plane-wave vector potential:**

$$\mathbf{A}^{(\text{transverse})}(\mathbf{r}, \omega) = 4\pi \sum_{LM p=0,1} i^{L-p} \left( \mathbf{Y}_{LM}^{(p)}(\hat{k}) \cdot \boldsymbol{\epsilon} \right) \mathbf{a}_{LM}^{(p)}(\mathbf{r}),$$



where the vector functions  $\mathbf{a}_{LM}^{(p)}(\mathbf{r})$  are referred to as **multipole potentials**, and where only terms with  $p = 0$  (magnetic) and  $p = 1$  (electric) contribute, since  $\mathbf{Y}_{LM}^{(-1)}(\hat{k})$  is orthogonal to  $\boldsymbol{\epsilon}$ .

### 2.4.c. Useful identities

#### Useful identities:

➤ **Derivative of the spherical Bessel functions:**

$$j_{n-1}(z) = \frac{n+1}{z} j_n(z) + j'_n(z),$$

$$j_{n+1}(z) = \frac{n}{z} j_n(z) - j'_n(z)$$

## 2.5. Frequently occuring radial integrals

### Notation & applications in Jac:

- In JAC, all these radial integrals can be directly computed by a call to some function from the `RadialIntegrals` module.

### Grant's radial integrals for the coupling of the radiation field:

- Grant (1988) defines the following radial integrals that frequently occur in the coupling of the radiation field

$$I_L^o(q; ab) = \int_0^\infty dr j_L(qr) [P_a Q_b] = \int_0^\infty dr j_L(qr) [P_a(r) Q_b(r)]$$

$$I_L^\pm(q; ab) = \int_0^\infty dr j_L(qr) [P_a Q_b \pm Q_a P_b]$$

$$J_L(q; ab) = \int_0^\infty dr j_L(qr) [P_a P_b + Q_a Q_b]$$

- These integrals are useful also if the spherical Bessel function occurs in the integrant as  $j_L(qr)/qr$ , and where the following recursion relations can be used:

$$\frac{j_L(x)}{x} = \frac{1}{2L+1} [j_{L-1}(x) + j_{L+1}(x)].$$

### Further radial integrals for different one-electron spherical tensor operators:

- In the treatment of the electron nucleus (hyperfine) interaction and the interaction with an external magnetic field, some radial integrals

occur frequently:

$$I_L^{[r]}(ab) \equiv [r^L]_{ab} = \int_0^\infty dr r^L [P_a Q_b + Q_a P_b]$$

$$J_L^{<r>}(ab) = \langle r^L \rangle_{ab} \equiv \int_0^\infty dr r^L [P_a P_b + Q_a Q_b]$$

➤ **Vinti-Integral:** In the treatment of the (relativistic) mass shift, the (so-called) Vinti-Integrals occur

$$R^{(\text{Vinti})}(a, b) = \int_0^\infty dr P_a \left[ \frac{d}{dr} - \frac{\kappa_a(\kappa_a + 1) - \kappa_b(\kappa_b + 1)}{2r} \right] P_b + \int_0^\infty dr Q_a \left[ \frac{d}{dr} - \frac{-\kappa_a(-\kappa_a + 1) + \kappa_b(-\kappa_b + 1)}{2r} \right] Q_b .$$

## 2.6. B-splines

### Properties of B-splines:

- A B-spline (basis spline) is a spline function with a minimal support and which can be used to express all spline function of any given as a finite linear combination.
- In physics, B-splines are often utilized for curve-fitting and in order to fit derivatives to experimental data.
- **Knot sequence:** A sequence of (radial) grid points  $\{t_1, t_2, \dots, t_m\}$  with  $t_1 \leq t_2 \leq \dots \leq t_m$ .
- **Set of B-splines:** Set of piecewise polynomial functions of order  $k$ , the so-called B-splines:

$$\left\{ B_1^{(k)}(x), B_2^{(k)}(x), \dots, B_n^{(k)}(x) \right\}, \quad B_i^{(1)}(x) = \begin{cases} 1 & \text{for } t_i \leq x \leq t_{i+1} \\ 0 & \text{otherwise,} \end{cases} \quad i = 1, 2, \dots, m-1.$$

➤ **Recursion relation for the generation of B-splines:**

$$B_i^{(k)}(x) = \frac{x - t_i}{t_{i+k-1} - t_i} B_i^{(k-1)}(x) + \frac{t_{i+k} - x}{t_{i+k} - t_{i+1}} B_{i+1}^{(k-1)}(x) \quad i = 1, 2, \dots, m-k$$

## 2. Dirac's hydrogenic atom

The number of B-splines  $n = m - k$  ( $k < m$ ) is determined by the order  $k$  and the number of knots  $m$ . For B-splines of a given order, it is therefore sufficient to use the notation  $B_i^{(k)}(x)$

- Each B-spline is non-negative and is non-zero only for:  $B_i(x) > 0$  for  $t_i < x < t_{i+k}$  and  $B_i(x) = 0$  for  $x < t_i$ ,  $x > t_{i+k}$ . — Essentially, the B-spline  $B_i(x)$  starts at  $t_i$  and ends at  $t_{i+k}$  with only positive values in between.

### Application of B-splines:

- Application of B splines for constructing single- and many-electron basis sets are nowadays widely employed in computational atomic and molecular physics.
- **B-splines:** For a B-spline representation of single-electron orbitals, only the low-lying orbitals will usually contribute significantly to the many-electron wave function. More generally, the configuration mixing between two configuration states is inversely proportional to the energy interval between them. It is therefore said sometimes that a B-spline representation of the single-electron orbitals provide an effective and natural way in order to adopt the basis to the problem of interest.
- **B-splines in relativistic theory:** In contrast to the nonrelativistic theory, the use of B-splines in relativistic computations is often accompanied by the occurrence of spurious states (Shabaev *et al.*, 2004). In a Coulomb potential, spurious states appear especially for  $\kappa > 0$  as the lowest bound states but with rather nonphysical energies.

## 2.7. Generation of continuum orbitals

### Simple approximations in Jac:

- For a first estimate of continuum processes, we have implemented a number of (rather) simple approximations/methods for the continuum orbitals that occur frequently at various places in the program. In JAC, the particular method can be selected by (re-) defining the global variable `JAC_CONT_SOLUTION` by a proper call to `Jac.define("method: continuum; ...)`.
- **Approximations implemented in JAC:**
  - (a) spherical Bessel functions;
  - (b) non-relativistic Coulomb waves for the large component of the continuum orbital, together with the kinetic-balance condition to obtain a small component;

- (c) asymptotically correct, relativistic Coulomb orbital;
- (d) Galerkin method to solve for continuum orbitals within a B-spline basis and for a given atomic potential.
- In JAC, spherical Bessel orbitals are generated if the global constant `JAC_CONT_SOLUTION = ContBessel`; cf. `JAC.define()` and section 2.7.a.
- In JAC, a free nonrelativistic Coulomb orbitals are generated if the global constant is `JAC_CONT_SOLUTION = NonrelativisticCoulomb`; cf. `JAC.define()` and section 2.7.b. Not available at present.
- In JAC, asymptotically-correct relativistic Coulomb orbitals are generated for a potential with a Coulombic tail, if the global constant `JAC_CONT_SOLUTION = AsymptoticCoulomb`; cf. `JAC.define()` and section 2.7.c.
- In JAC, a B-spline representation of the relativistic orbital  $\varepsilon\kappa$  in a given local potential are generated if the global constant `JAC_CONT_SOLUTION = BsplineGalerkin`; cf. `JAC.define()` and section 2.7.d.
- In JAC, different normalization methods can be selected for all continuum orbitals; a pure-sine normalization for `JAC_CONT_NORMALIZATION = PureSine`, an asymptotic Coulomb normalization for `JAC_CONT_NORMALIZATION = CoulombSine`, or a WKB-type normalization for multiply and highly-charged ions for `JAC_CONT_NORMALIZATION = OngRussek`; cf. `JAC.define()`.
- At present, no attempt has yet been made to incorporate the exchange interaction of the continuum and bound-state orbitals.

### 2.7.a. Spherical Bessel orbitals

#### Implementation and numerical details:

- **Generating equation and solutions:** The spherical Bessel functions  $j_\ell(|\mathbf{k}|r)$  obey the (non-relativistic) *free* radial Schrödinger equation for positive energies  $\varepsilon > 0$ , and the corresponding wave number  $k = \sqrt{2\varepsilon}$ ,  $\rho = kr$  and for well-defined orbital angular momentum  $\ell$

$$R_{\varepsilon\ell}(r) = \frac{P_{\varepsilon\ell}}{r} = N j_\ell(|\mathbf{k}|r) = N j_\ell(\sqrt{2\varepsilon} r) \quad \Longleftrightarrow \quad j_\ell''(\rho) + 2 \frac{j_\ell'(\rho)}{\rho} + \left(1 - \frac{\ell(\ell+1)}{\rho^2}\right) j_\ell(\rho) = 0.$$

- While the large component of such a continuum (Bessel) orbital  $|\varepsilon\kappa\rangle$  are represented by the Bessel function  $P = N r j_\ell(|\mathbf{k}|r)$  above, the **small components are here obtained from the kinetic-balance condition.**

## 2.7.b. Non-relativistic Coulomb orbitals

### Implementation and numerical details:

- The non-relativistic Coulomb orbitals  $|\varepsilon\ell\rangle$  are known analytically for  $\varepsilon > 0$  and can be utilized to represent the large components of a relativistic orbital with this (given kinetic) energy, while the **small components are here obtained again from the kinetic-balance condition**.
- **Non-relativistic free Coulomb waves:** The free solutions with  $\varepsilon > 0$  to the radial Schrödinger equation with a central-field potential  $V(r)$  with Coulombic tail satisfy the boundary condition:

$$P_{\varepsilon\ell}(0) = 0; \quad P_{\varepsilon\ell}(r \rightarrow \infty) \simeq \sqrt{\frac{2}{\pi k}} \sin \left( kr + \frac{\bar{Z}}{k} \ln(2kr) - \frac{\ell\pi}{2} + \sigma_{\ell}^{(\text{Coulomb})} + \delta_{\ell} \right),$$

where  $k^2 = 2\varepsilon$ ,  $\bar{Z} = Z - N$  is the effective charge as seen by the free electron at large distance and  $\sigma_{\ell}^{(\text{Coulomb})} = \arg \Gamma(\ell + 1 - i\bar{Z}/k)$  is the (well-known) **Coulomb phase**. Note that this solution is normalized on the energy scale.

- See section 2.1 for a representation of a non-relativistic Coulomb orbital with positive energy  $\varepsilon > 0$ . **This approach does not yet work properly since there is no hypergeometric function with complex arguments available in JULIA.**

## 2.7.c. Asymptotically-correct relativistic Coulomb orbitals

### Implementation and numerical details:

- **Relativistic free-Coulomb waves:** The free solutions to the radial Dirac equation in a spherical potential with Coulombic tail satisfy the

boundary conditions (Åberg and Howat, 1982; Eqs. (19.29-30))

$$\begin{aligned}
 P_{\varepsilon\kappa}(0) &= Q_{\varepsilon\kappa}(0) = 0 \\
 P_{\varepsilon\kappa}(r \rightarrow \infty) &\simeq \sqrt{\frac{\varepsilon + 2c^2}{\pi c^2 k}} \cos [\theta_\kappa(\varepsilon) + \delta_\kappa(\varepsilon)], & Q_{\varepsilon\kappa}(r \rightarrow \infty) &\simeq -\sqrt{\frac{\varepsilon}{\pi c^2 k}} \sin [\theta_\kappa(\varepsilon) + \delta_\kappa(\varepsilon)] \\
 \theta_\kappa(\varepsilon) &= k r + y \ln(2k r) - \arg \Gamma(\bar{\gamma} + i y) - \frac{1}{2} \pi \bar{\gamma} + \eta, & k^2 &= 2\varepsilon + \frac{\varepsilon}{c^2} \\
 \bar{\gamma} &= + \left( \kappa^2 - \frac{\bar{Z}^2}{c^2} \right)^{1/2}, & y &= \frac{\bar{Z} (\varepsilon + c^2)}{c^2 k} & \exp(2i\eta) &= -\frac{\kappa - i y / (\varepsilon + c^2)}{\bar{\gamma} + i y}.
 \end{aligned}$$

➤ The phase-shift  $\delta_\kappa$  is the non-Coulombic phase shift and is set to  $\delta_\kappa = 0$  for a pure Coulombic potential.

## 2.7.d. Continuum orbitals in an atomic potential: Galerkin method

### Implementation and numerical details:

- One advantages of using B-splines is that solutions of the Schrödinger or Dirac equation can be found for any energy  $\varepsilon > 0$  by properly adopting  $r_{\max}$ .
- While both the Schrödinger and Dirac equations have formally a solution for any positive energy  $\varepsilon > 0$ , a diagonalization in a B-spline representation provides only a discrete set of solutions due to the boundary condition  $P(r_{\max}) = Q(r_{\max}) = 0$ , and where  $r_{\max}$  here denotes the size of the *numerical box*.
- **Galerkin method:** This method determines an approximate solution of the generalized eigenvalue equations

$$\sum_i \langle B_i | (\mathbb{H} - \varepsilon) | B_j \rangle c_j = \sum_i A_{ij} c_j = 0, \quad A_{ij} = \langle B_i | (\mathbb{H} - \varepsilon) | B_j \rangle = H_{ij} - \varepsilon S_{ij}.$$

- **Galerkin method:** A possible approximation to this generalized eigenvalue problem is to compute the smallest eigenvalue of  $\mathbf{A} \cdot \mathbf{c} = \lambda \mathbf{c}$ , a more accurate solution follows from a minimization of  $|A\psi|^2$  under the constraint  $|\psi|^2 = 1$ . This variational problem is mathematically

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equivalent to find the smallest eigenvalue of the matrix  $A^+ A$ :  $\mathbf{A}^+ \mathbf{A} \cdot \mathbf{c} = a \mathbf{c}$ . — The eigenvector of the smallest eigenvalue  $a_i > 0$  is then the requested (approximate) solution of the eigenvalue problem above.

### 2.7.e. Normalization and phase of continuum orbitals

#### Pure sine behaviour at large $r$ :

- **Pure sine behaviour:** The free radial Schrödinger equation without an external potential always satisfies the asymptotic pure sine behaviour  $\sim \sin(kr + \varphi)$ . A very similar asymptotic form  $\sim \sin(kr - \ell\pi/2 + \varphi)$  also applies for all partial-wave solutions  $\ell \geq 0$  and positive energy  $\varepsilon > 0$  for the standard radial Schrödinger equation with just a centrifugal potential  $\sim 1/r^2$ .
- **Normalization for a pure sine behaviour:** For this normalization, the large component is assumed to be given with  $k = \sqrt{2\varepsilon}$  by

$$P(r \approx r_{\max}) = \sqrt{\frac{2}{\pi k}} \begin{cases} \sin(kr + \varphi) & \text{for } \ell = 0 \\ \sin(kr - \frac{\ell\pi}{2} + \varphi) & \text{for } \ell > 0. \end{cases}$$

With this behaviour, the normalization constant  $N$  and the phase  $\varphi$  can be obtained by means of the logarithmic derivative  $P/P'$ .

#### Asymptotic Coulomb behaviour at large $r$ :

- **Normalization for an asymptotic Coulomb behaviour:** For this normalization, the large and small components are assumed to be described for all  $r$ -values by the asymptotic behaviour as shown in subsection 2.7.c. With this assumption, the normalization constant  $N$  and the phase  $\varphi$  can again be obtained by means of the logarithmic derivative  $P'/P$ .
- Instead of the logarithmic derivative  $P'/P$ , one can of course also apply the ratio  $Q/P$ ; this option is currently hard-coded in JAC but need to be further tested.

#### Normalization for multiply-charged ions due to Ong and Russek (1973):

- **This normalization scheme has not yet been implemented and tested so far.**
- **Normalization on energy scale:** Continuum orbital functions are often **normalized on an energy scale**,  $\langle \varepsilon \kappa | \varepsilon' \kappa' \rangle = \delta(\varepsilon - \varepsilon') = \delta(E - E')$ , and where  $E, E'$  are the free-electron energies *inclusive* the rest mass energy, i.e.  $E > mc^2$ .



- **WKB ansatz for the large component  $P_{\varepsilon\kappa}(r)$ :** If  $V(r)$  is an atomic central-field potential of the ionic core and if we make use of the ansatz below for the large component in the Dirac equation, a second-order ODE can be derived for the function  $\phi(r)$ , cf. Ong and Russek (1978)

$$P_{\varepsilon\kappa}(r) = A \sqrt{\frac{E - V(r) + c^2}{c (d\phi/dr)}} \cos \phi(r),$$

$$- \left( \frac{d\phi}{dr} \right)^2 + \frac{(E - V)^2 - c^4}{c^2} - \frac{\kappa(\kappa + 1)}{r^2} + \left( \frac{d\phi}{dr} \right)^{1/2} \frac{d^2}{dr^2} \left( \frac{d\phi}{dr} \right)^{-1/2} - (E - V + c^2)^{1/2} \frac{d^2}{dr^2} (E - V + c^2)^{-1/2} + \frac{\kappa}{r} \frac{dV/dr}{E - V + c^2} = 0.$$

- **WKB approximation:** If the last three terms are neglected in the equation above, an approximate solution to this ODE is given by

$$\left( \frac{d\phi}{dr} \right)^2 \approx \frac{(E - V)^2 - c^4}{c^2} - \frac{\kappa(\kappa + 1)}{r^2}, \quad P_{\varepsilon\kappa}(r) \approx \frac{A \sqrt{E - V + c^2}}{[(E - V)^2 - c^4 - c^2 \kappa(\kappa + 1)/r^2]^{1/4}} \cos \phi(r).$$

- **Normalization and phase of the WKB solution at some finite value  $r_o$ :** Following Ong and Russek (1978, Eq. 12), one can introduce an auxiliary potential  $U(r)$ , from which the normalization constant  $A(r_o)$  and the phase  $\phi(r_o)$  at some finite radius  $r_o$  is obtained by

$$U(r) = \frac{c}{2} \frac{(E - V) \frac{dV}{dr} - c^2 \kappa(\kappa + 1)/r^3}{[(E - V)^2 - c^4 - c^2 \kappa(\kappa + 1)/r^2]^{3/2}} P_{\varepsilon\kappa}(r) - c \frac{\frac{dP_{\varepsilon\kappa}}{dr} + \frac{dV/dr}{2(E - V + 2c^2)} P_{\varepsilon\kappa}(r)}{[(E - V)^2 - c^4 - c^2 \kappa(\kappa + 1)/r^2]^{1/2}}$$

$$A(r_o) = \sqrt{\frac{[(E - V)^2 - c^4 - c^2 \kappa(\kappa + 1)/r^2]^{1/2}}{E - V + c^2}} (P_{\varepsilon\kappa}^2 + U^2) |_{r=r_o}, \quad \phi(r_o) = \text{atan2}(U, P_{\varepsilon\kappa}) |_{r=r_o}.$$

- **Normalization and phase of the asymptotic wave function  $P_{\varepsilon\kappa}(r)$ :** The normalization of  $A(r_o)$  can be used to normalize also the large (and small) component  $P_{\varepsilon\kappa}(r)$ . For  $V(r_o) \approx 0$  and  $\phi(r) \approx kr + \delta_\kappa(\varepsilon)$  and if  $C$  is the normalization constant so that  $CP_{\varepsilon\kappa}$  has the correct asymptotic amplitude, one finally obtains

$$C = \frac{1}{A} \sqrt{\frac{(E + c^2) c d\phi/dr}{\pi c^2 k (E - V + c^2)}} \approx \frac{1}{A \sqrt{\pi c}}, \quad \delta_\kappa(\varepsilon) \approx \phi(r) - kr.$$

## 2.8. Nuclear models and potentials

### 2.8.a. Uniform nuclear model

#### Uniform charge distribution and potential:

- **Uniform nuclear-charge distribution:** For a nucleus with radius  $R$ , charge  $Z$  and mass number  $A$ , the nuclear charge distribution and potential is given for  $r \leq R$  by

$$\rho^{(\text{uniform})}(r) = \frac{3Z}{4\pi R^3}, \quad \phi^{(\text{uniform})}(r) = \begin{cases} \frac{Ze}{2R} \left[ 3 - \left( \frac{r}{R} \right)^2 \right] & \text{if } r \leq R \\ \frac{Ze}{R} & \text{if } r > R \end{cases}$$

$$R = R_n \sqrt{\frac{1 + (5s^2/2R_n^2)}{1 + (3s^2/4R_n^2)}}, \quad R_n = 1.07 \times 10^{-15} A^{1/3} \text{ m}, \quad s = 2.0 \times 10^{-15} \text{ m}.$$

- **Form factor for nuclei with uniform nuclear-charge distribution:**

$$F^{(\text{uniform})}(q; R) \equiv 4\pi \int_0^\infty dr r^2 \frac{\sin(qr)}{qr} \rho^{(\text{uniform})}(r) = \frac{3Z}{(qR)^3} [\sin(qR) - qR \cos(qR)].$$

### 2.8.b. Helm's uniform-uniform nuclear model

#### Uniform-uniform charge distribution and potential:

- **Helm's uniform-uniform nuclear-charge distribution:** An alternative realistic nuclear-charge distribution can be expressed as convolution of two uniform distributions  $\rho_1^{(\text{uniform})}$ ,  $\rho_2^{(\text{uniform})}$  with nuclear radii  $R_1 > R_2$  and normalization constant  $N$  (Salvat *et al.*, 2005)

$$\rho^{(\text{Helm})} = Z N \int d^3 \mathbf{r}' \rho_1^{(\text{uniform})}(\mathbf{r}') \rho_2^{(\text{uniform})}(\mathbf{r} - \mathbf{r}')$$

$$\phi^{(\text{Helm})} = \begin{cases} N & \text{if } r \leq R_1 - R_2 \\ N \frac{[(r + R_1)^2 - R_2^2 + 2r R_1][(r - R_1)^2 - R_2^2]^2}{32 r^3 R_1 - 2^3} + \frac{[(r + R_2)^2 - R_1^2 + 2r R_2][(r - R_2)^2 - R_1^2]^2}{32 r^3 R_2 - 2^3} & \text{if } R_1 - R_2 \leq r \leq R_1 + R_2 \\ 0 & \text{if } r > R_1 + R_2 \end{cases}$$

- The normalization constant  $N$  of Helm's uniform-uniform nuclear-charge distribution and the corresponding electrostatic potential  $\phi^{(\text{Helm})}$  need to be determined numerically.

### 2.8.c. Fermi nuclear model

#### Fermi charge distribution and potential:

- **Fermi model:** For a nucleus with *root-mean-square* (rms) radius  $R = \sqrt{\langle r^2 \rangle}$

$$\rho(r, R) = \frac{N}{1 + \exp[(r - c)/a]}, \quad \int dr r^2 \rho(r, R) = 1$$

where the **thickness parameter**  $a$  is often chosen as  $a = 2.3/4 \ln 3$ .

- **Fermi nuclear-charge distribution:** For a nucleus with *half-density* density  $R_n$ , the Fermi distribution is uniform in the center and falls off smoothly at the surface of the nucleus from 0.9 to 0.1 of the central density within a skin of thickness  $t$

$$\rho^{(\text{Fermi})} = \frac{\rho_o}{\exp[(r - R_n)/z] + 1}, \quad z = \frac{t}{(4 \ln 3)} = 0.546 \times 10^{-15} \text{ m}$$

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- The normalization constant  $\rho_o$  of the Fermi nuclear-charge distribution and the corresponding electrostatic potential  $\phi^{(\text{Fermi})}$  need to be determined numerically.
- With high accuracy, the parameters  $N$  and  $c$  can be obtained from the analytical formulas

$$N = \frac{3}{4\pi c^3} \left(1 + \frac{\pi^2 a^2}{c^2}\right)^{-1}, \quad c = \sqrt{\frac{5}{3} \langle r^2 \rangle - \frac{7}{3} \pi^2 a^2}.$$

### Nuclear potential:

- Potential of extended nucleus:

$$\mathbb{V}^{(\text{nuc})}(r; R) = -4\pi \alpha Z \int_0^\infty dr' r'^2 \frac{\rho(r', R)}{r_{>}} = -\frac{4\pi \alpha Z}{r} \int_0^r dr' r'^2 \rho(r', R) - 4\pi \alpha Z \int_r^\infty dr' r' \rho(r', R), \quad r_{>} = \max(r, r').$$

# 3. Many-electron atomic interactions, state functions, density operators and statistical tensors

## 3.1. Electron-electron interaction

### Motivation:

- In the (non-relativistic) limit of a rather slow motion of the electrons, each electron pair just interact by the instantaneous Coulomb repulsion

$$\frac{1}{r_{12}} \equiv \frac{1}{|\mathbf{r}_1 - \mathbf{r}_2|}.$$

Already this quite simple (interaction) operator results mathematically in various complications and is generally the major source of electron-electron correlations in all many-electron computations.

- For moving charges, Maxwell's theory also predicts both, current-current and retarded interactions, to which we briefly refer as Breit interactions in atomic physics. — These relativistic contributions to the electron-electron interactions must be added to the electro-static Coulomb repulsion.
- An effective (relativistic) operator for the electron-electron interaction can be formally derived from QED in perturbation theory with regard to the number of virtually exchanged photons (Sapirstein 1987, Grant und Quiney 1988), while the same expressions can be obtained also from heuristic arguments as given, for example, by Bethe and Salpeter (1957) und Johnson (1995).
- In practice, the computation of the electron-electron interaction usually requires the largest effort in all electronic-structure codes and, therefore, deserves special consideration. In atomic physics, a considerable simplification is usually achieved if all electron-electron interaction operators are represented as spherical tensors.

### 3. Many-electron atomic interactions, state functions, density operators and statistical tensors

- For two electrons in orbitals  $a, b$  and with energies  $\varepsilon_a, \varepsilon_b$ , the **interaction energy can be written as difference of a direct and exchange matrix element** of the potential

$$\mathbb{V}^{(\text{tr})} = -\frac{1}{2\pi} \int \frac{d^3k}{k^2 - \omega^2} e^{i\mathbf{k} \cdot (\mathbf{r}_1 - \mathbf{r}_2)} \sum_{i,j=1}^3 \left[ \delta_{ij} - \frac{k_i k_j}{k^2} \right] \alpha_{1i} \alpha_{2j}$$

where  $\alpha_{1i}$  is the  $i$ -th component of the  $\boldsymbol{\alpha}$  matrix (vector) of the electron at position  $\mathbf{r}_1$  and

$$\omega = \begin{cases} 0 & \text{direct matrix element} \\ \frac{|\varepsilon_a - \varepsilon_b|}{c} & \text{exchange matrix element} \end{cases}$$

refers to the momentum of the exchanged photon.

- **The exchange matrix elements of the potential  $\mathbb{V}^{(\text{tr})}$  is generally complex**; while the *real* part describes an energy shift due to this transverse interaction, the *imaginary* part represents a rate for the decay  $a \rightarrow b$  for  $\varepsilon_a > \varepsilon_b$  or  $b \rightarrow a$  *vice versa*.
- **Frequency-independent Breit interaction**: For  $\omega \rightarrow 0$ , the transversal Breit interaction gives rise to the **frequency-independent Breit interaction**

$$b_{12}^o = -\frac{1}{2r_{12}} \left[ \boldsymbol{\alpha}_1 \cdot \boldsymbol{\alpha}_2 + \frac{(\boldsymbol{\alpha}_1 \cdot \mathbf{r}_{12})(\boldsymbol{\alpha}_2 \cdot \mathbf{r}_{12})}{r_{12}^2} \right],$$

which neglects all contributions  $\sim \alpha^4 Z^3$  (as well as of higher order in  $\alpha Z$ ).

- Typically, the explicit energy-dependence of the transverse interaction [upon the energy difference of the two electrons] gives rather tiny corrections, so that, in practice, **often no difference is made between the transversal and zero-frequency Breit interaction**.
- **Gaunt interaction**: The first term of the zero-frequency Breit interaction (operator) is known also as Gaunt interaction (Gaunt 1929). This term describes the current-current interaction due to the motion of the electrons and contributes about 90 % of the total Breit interaction for closed-shell atoms.
- For many elements of the periodic table, the corrections due to the Breit interaction are usually small, when compared with the dominant Coulomb repulsion. This applies especially for the Breit contributions to the electron-electron correlation.
- For describing most processes, it is **sufficient, if at all, to include the Breit interaction perturbatively in lowest order**. Formally, however, both terms

$$v_{12} = \frac{1}{r_{12}} + b_{12} \quad \Longleftrightarrow \quad \mathbb{V}^{(\text{e-e})} = \mathbb{V}^{(\text{Coulomb})} + \mathbb{V}^{(\text{Breit})}$$

are equivalent, and they should be treated simply together for all (super-) heavy elements.

➤ In JAC, only the frequency-independent Breit interaction is currently implemented.

### Effective interaction strength $X^{(L)}(abcd)$ of scalar electron-electron interactions:

➤ All operators of the (scalar) electron-electron interaction can be generally represented as spherical tensors in the form:

$$g_{12} \equiv g(\mathbf{r}_1, \mathbf{r}_2) = \sum_L g_L(r_1, r_2) (\mathbb{T}^{(L)}(\vartheta_1, \varphi_1) \cdot \mathbb{T}^{(L)}(\vartheta_2, \varphi_2))$$

$$\langle n_a \kappa_a m_a(1) n_b \kappa_b m_b(2) | \mathbb{G}_{12} | n_c \kappa_c m_c(1) n_d \kappa_d m_d(2) \rangle = \sum_{LM} (-1)^{L-M+j_a-m_a+j_b-m_b} \begin{pmatrix} j_a & L & j_c \\ -m_a & M & m_c \end{pmatrix} \begin{pmatrix} j_b & L & j_d \\ -m_b & -M & m_d \end{pmatrix} X^{(L)}(abcd),$$

leading to a simple **factorization of the two-electron matrix elements**, and where the magnetic quantum numbers (i.e. the angular dependence) of the matrix elements only occurs in the phase and the Wigner 3-j symbols.

➤ **Effective interaction strength  $X^{(L)}(abcd)$  of order  $L$ :** The strengths  $X^{(L)}(abcd)$  describe the physical interaction and are specific for every operator. They are often used as the **building blocks in order to handle the electron-electron interactions efficiently in electronic structure computations**.

### Effective interaction strength $X^{(L, \text{Coulomb})}(abcd)$ of the Coulomb repulsion:

➤ **Coulomb repulsion:** The decomposition of the Coulomb operator is given by:

$$\frac{1}{r_{12}} = \sum_{L=0}^{\infty} U_L(r, s) P_L(\cos(\vartheta)) = \sum_{L=0}^{\infty} U_L(r, s) (\mathbb{C}^{(L)}(\vartheta_1, \varphi_1) \cdot \mathbb{C}^{(L)}(\vartheta_2, \varphi_2)), \quad U_L(r, s) = \begin{cases} \frac{r^L}{s^{L+1}} & r < s \\ \frac{s^L}{r^{L+1}} & r > s. \end{cases}$$

In this expansion,  $P_L(x)$  denotes a Legendre-polynomial and  $\vartheta = \mathbf{r} \cdot \mathbf{s}/rs$  the angle between the two electron coordinates.

### 3. Many-electron atomic interactions, state functions, density operators and statistical tensors

#### ➤ Effective interaction strength of the Coulomb repulsion:

$$X^{(L, \text{Coulomb})}(abcd) = \delta(j_a, j_c, L) \delta(j_b, j_d, L) \Pi^e(\kappa_a, \kappa_c, L) \Pi^e(\kappa_b, \kappa_d, L) (-1)^L \langle \kappa_a \parallel \mathbb{C}^{(L)} \parallel \kappa_c \rangle \langle \kappa_b \parallel \mathbb{C}^{(L)} \parallel \kappa_d \rangle R^L(abcd),$$

$$\delta(j_a, j_b, j_c) = \begin{cases} 1 & |j_a - j_b| \leq j_c \leq |j_a + j_b| \text{ and cyclic interchanged} \\ 0 & \text{otherwise,} \end{cases}$$

and where  $\Pi^e(\kappa, \kappa', L)$  represent the angular momentum and parity selection rules.

#### Effective interaction strength $X^{(L, \text{Breit})}(abcd)$ of the Breit interaction:

#### ➤ Effective interaction strength of the zero-frequency Breit: For the zero-frequency interaction, we find

$$X^{(L, \text{Breit})}(abcd) = (-1)^L \langle \kappa_a \parallel \mathbb{C}^{(L)} \parallel \kappa_c \rangle \langle \kappa_b \parallel \mathbb{C}^{(L)} \parallel \kappa_d \rangle \left[ \Pi^o(\kappa_a, \kappa_c, L-1) \Pi^o(\kappa_b, \kappa_d, L+1) \sum_{\mu=1}^8 s_{\mu}^L(abcd) S_{\mu}^L(abcd) + \right. \\ \left. + \sum_{\nu=L-1}^{L+1} \Pi^o(\kappa_a, \kappa_c, \nu) \Pi^o(\kappa_b, \kappa_d, \nu) \sum_{\mu=1}^4 t_{\mu}^{\nu, L}(abcd) T_{\mu}^{\nu}(abcd) \right].$$

$$\Pi^o(\kappa_a, \kappa_b, L) = \begin{cases} 1 & \text{if } l_a + l_b + L \text{ odd} \\ 0 & \text{otherwise.} \end{cases}$$

#### ➤ The constant coefficients $s_{\mu}^L$ and $t_{\mu}^{\nu, L}$ just depend on the angular quantum numbers $\{\kappa_a, \kappa_b, \kappa_c, \kappa_d, \nu, L\}$ and are displayed, for example, by Grant (1988; tables 3 and 4, and where we here use $t_{\mu}^{\nu, L} \equiv r_{\mu}^{\nu, L}$ ).



Table 3.1.: Definition of the integrales  $S_\mu^\nu(abcd)$  und  $T_\mu^\nu(abcd)$ 

$S_1^\nu(abcd) = S^\nu[ac    bd]$	$S_2^\nu(abcd) = S^\nu[bd    ac]$	$T_1^\nu(abcd) = T^\nu[ac    bd]$
$S_3^\nu(abcd) = S^\nu[ca    db]$	$S_4^\nu(abcd) = S^\nu[db    ca]$	$T_2^\nu(abcd) = T^\nu[ca    db]$
$S_5^\nu(abcd) = S^\nu[ac    db]$	$S_6^\nu(abcd) = S^\nu[db    ac]$	$T_3^\nu(abcd) = T^\nu[ac    db]$
$S_7^\nu(abcd) = S^\nu[ca    bd]$	$S_8^\nu(abcd) = S^\nu[bd    ca]$	$T_4^\nu(abcd) = T^\nu[ca    bd]$

➤ The radial integrals  $S_\mu^\nu(abcd)$  and  $T_\mu^\nu(abcd)$  are different (linear) combinations of the two integrals [cf. table 3.1]

$$S^\nu[ac || bd] = \int_0^\infty dr \int_0^\infty ds P_a(r) Q_c(r) \frac{1}{2} [W_{\nu-1, \nu+1, \nu}(r, s; \omega_{ac}) + W_{\nu-1, \nu+1, \nu}(r, s; \omega_{bd})] P_b(s) Q_d(s)$$

$$T^\nu[ac || bd] = \int_0^\infty dr \int_0^\infty ds P_a(r) Q_c(r) \frac{1}{2} [V_\nu(r, s; \omega_{ac}) + V_\nu(r, s; \omega_{bd})] P_b(s) Q_d(s)$$

$$V_\nu(r, s; \omega) = \begin{cases} -(2\nu+1) j_\nu(\omega r) n_\nu(\omega s) & r < s \\ -(2\nu+1) j_\nu(\omega s) n_\nu(\omega r) & r > s \end{cases}$$

$$W_{\nu-1, \nu+1, \nu}(r, s; \omega) = \begin{cases} -(2\nu+1) j_{\nu-1}(\omega r) n_{\nu+1}(\omega s) + \left(\frac{2\nu+1}{\omega}\right)^2 \frac{r^{\nu-1}}{s^{\nu+2}} & r < s \\ -(2\nu+1) \omega j_{\nu-1}(\omega s) n_{\nu+1}(\omega r) & r > s \end{cases}.$$

➤ In these expressions,  $j_\nu(\omega r)$  and  $n_\nu(\omega r)$  denote the spherical Bessel functions, and  $\omega_{ac} = |\epsilon_a - \epsilon_c|$  the difference of the corresponding single-particle energies.

## 3.2. Atomic potentials

### Parametrized many-electron densities:

- **Thomas-Fermi-Moliere (TFM) electron density:** A useful analytical model for the spherical symmetric electron density follows from a simple parametrization of the Thomas-Fermi (TF) screening function  $\chi^{(\text{TFM})}(r)$  and its substitution into the Poisson equation (Salvat *et al.*, 2005)

$$\chi^{(\text{TFM})}(r) = \sum_{i=1}^3 A_i \exp\left(\frac{a_i r}{b}\right), \quad A_1 = 0.1, \quad A_2 = 0.55, \quad A_3 = 0.35, \quad a_1 = 6, \quad a_2 = 1.2, \quad a_3 = 0.3$$

$$b = \frac{(3\pi)^{2/3}}{2^{7/3}} \frac{\hbar^2}{m e^2 Z^{1/3}} = \frac{0.88534 a_o}{Z^{1/3}}$$

$$\rho^{(\text{TFM, electrons})}(r) = \frac{Z}{4\pi r} \sum_{i=1}^3 A_i a_i^2 \exp(a_i r).$$

- **Thomas-Fermi-Dirac (TFD) electron density:** The TFD theory also accounts for the exchange term of the electron gas and leads again to a differential equation that need to be solved numerically for each individual atom or ion. A parametrized form for an approximated TFD screening function  $\chi^{(\text{TFD})}$  has been suggested in the literature but results in a rather lengthy and cumbersome expansion.

### 3.2.a. In JAC implemented potentials

#### Atomic potentials:

- **Core-Hartree potential:**

$$V^{(\text{core-Hartree})}(r) = \int_0^\infty dr' \frac{\rho_c(r')}{r_{>}}, \quad r_{>} = \max(r, r'), \quad \rho_c(r) = \sum_a (P_a^2(r) + Q_a^2(r)),$$

and where the summation runs over all *core* orbitals of electron configurations. This potential is applied if just a single valence electron occurs (alkali atoms). In JAC, this potential can be obtained for a given level from `Jac.computePotentialCoreHartree(grid::Radial.Grid, level::Level)`.

➤ Dirac-Fock-Slater potential:

$$V^{(\text{DFS})}(r) = \int_0^\infty dr' \frac{\rho_t(r')}{r_{>}} - \left( \frac{3}{4\pi^2 r^2} \rho_t(r) \right)^{1/3}, \quad r_{>} = \max(r, r'), \quad \rho_t(r) = \sum_a (P_a^2(r) + Q_a^2(r)),$$

and where the summation runs over *all* orbitals (electrons). In JAC, this potential can be obtained for a given level from `Jac.computePotentialDFS(grid::Radial.Grid, level::Level)`.

➤ Hartree potential:

$$V^{(\text{Hartree})}(r) = - \sum_a \bar{q}_a r Y_{aa}^0(r)$$

where  $\bar{q}_a$  is the generalized occupation number, and where the summation runs over all orbitals here. In JAC, this potential can be obtained for a given level from `Jac.computePotentialHartree(grid::Radial.Grid, level::Level)`.

➤ Hartree-Slater potential:

$$V^{(\text{HS})}(r) = - \sum_a \bar{q}_a Y_{aa}^0(r) + \frac{3}{2} \left( \frac{3}{4\pi^2 r^2} \rho(r) \right)^{1/3} \frac{r}{2}, \quad \rho(r) = \sum_a \bar{q}_a (P_a^2(r) + Q_a^2(r)),$$

and where the summation runs over *all* orbitals. In JAC, this potential can be obtained for a given level from `Jac.computePotentialDFS(grid::Radial.Grid, level::Level)`.

➤ Kohn-Sham potential:

$$V^{(\text{Kohn-Sham})}(r) = \int_0^\infty dr' \frac{\rho_t(r')}{r_{>}} - \frac{2}{3r} \left( \frac{81}{32\pi^2} r \rho_t(r) \right)^{1/3}, \quad r_{>} = \max(r, r'), \quad \rho_t(r) = \sum_a (P_a^2(r) + Q_a^2(r)),$$

and where the summation runs over *all* orbitals (electrons). In JAC, this potential can be obtained for a given level from `Jac.computePotentialKohnSham(grid::Radial.Grid, level::Level)`.

### 3.2.b. Further atomic potentials, not yet considered in JAC

#### Atomic potentials:

➤ Tietz (1954) potential:

$$V^{(\text{Tietz})}(r) = -\frac{\alpha Z}{r} \phi(x), \quad \phi(x) = \frac{1}{(1 + (\pi/8)^{2/3} x)^2}, \quad x = \frac{r Z^{1/3}}{\mu}, \quad \mu = \left(\frac{3\pi}{4}\right)^{2/3} \approx 0.8853.$$

➤ Modified-Tietz potential: (Neuffer and Commins, 1977)

$$V^{(\text{modified-Tietz})}(r) = -\frac{\alpha}{r} \left(1 + \frac{(Z-1)}{(1+tr)^2} e^{-\gamma r}\right)$$

where  $\gamma$  and  $t$  are parameters adjusted to fit the low-lying spectra of the atoms under consideration. Johnson *et al.* (1985) tabulate these parameters for rubidium, cesium, gold and thallium.

### 3.3. Construction of symmetry-adapted CSF basis

#### Excitation schemes for the constructing a many-electron CSF basis:

- Several (de-) excitation schemes for the construction of non-relativistic configuration lists are defined within JAC in order to facilitate the set-up of (lists of) CSF bases with selected total symmetry  $\mathbb{J}$ . These bases are required specially for the computation of cascades and/or approximate many-electron Green functions.
- In JAC, various excitation schemes are distinguished by the (abstract) `scheme::Basics.AbstractExcitationScheme`, and which is used for the construction of systematically enlarged basis sets, the computation of approximate Green functions as well as for the generation of atomic cascades.

- (De-) excitation of a single electron from a given set of (non-relativistic) configurations (**DeExciteSingleElectron**): This schemes includes all excitations and de-excitations of a single electron from a given list of (nonrelativistic bound-electron) configurations. Here, of course, the number of electrons of the newly generated configurations is the *same* as given by the configuration, while up to *one* free electron occur in each CSF for representing  $N$ -electron scattering levels.

Example:  $1s^2 2s^2 2p^6 \longrightarrow \left\{ \begin{array}{l} 1s 2s^2 2p^6 (ns + np + nd + \dots) \\ 1s^2 2s 2p^6 (ns + np + nd + \dots) \\ 1s^2 2s^2 2p^5 (ns + np + nd + \dots) \end{array} \right\}$

- (De-) excitation of two electrons from a given set of (non-relativistic) configurations (**DeExciteTwoElectrons**): Similar as above, this schemes includes all excitations and de-excitations of up to *two* electrons from the given list of (bound-electron) configurations and, hence, allows up to *two* free electrons for representing the  $N$ -electron scattering levels.

Example:  $1s^2 2s^2 2p^6 \longrightarrow \left\{ \begin{array}{l} 1s 2s^2 2p^6 (ns + np + nd + \dots) + 2s^2 2p^6 nl n'l' \\ 1s^2 2s 2p^6 (ns + np + nd + \dots) + 1s^2 2p^6 nl n'l' \\ 1s 2s 2p^6 nl n'l' \\ \dots \end{array} \right\}$

- Add a single electron to a given set of configurations (**AddSingleElectronWithoutHoles**): This scheme generates all configurations with one additonal electron to the given set of configurations but without any replacement from the occupied orbitals.

Example:  $1s^2 2s^2 2p^6 \longrightarrow \{ 1s^2 2s^2 2p^6 (ns + np + nd + \dots) \}$

## 3.4. Multi-configuration Dirac-Hartree-Fock (MCDHF) method

### Motivation:

- The multi-configuration Dirac-Hartree-Fock (MCDHF) method has been found a versatile tool for calculating the level structure and approximate wave functions for atoms and ions of all elements across the periodic table. Not much need to be said about this method, which has been presented at several places elsewhere (Grant 1989; Parpia *et al.* 1996; Fritzsche 2002).
- **Atomic state function (ASF):** Similar to the well-known Hartree-Fock (HF) method from non-relativistic quantum theory, in which the state of a quantum system is approximated by a single (symmetry-adapted) Slater determinant, an **atomic state function (ASF)** is written in the MCDHF model as linear combination of configuration state functions (CSF) with well adapted symmetry

$$|\alpha \mathbb{J}M\rangle \equiv \psi_{\alpha}(\mathbb{J}M) = \sum_{r=1}^{n_c} c_r(\alpha) |\gamma_r P J M\rangle .$$

In this ansatz,  $n_c$  is the number of CSF,  $\mathbb{J} \equiv J^P$  the total angular momentum and parity of the state, and  $\alpha$  refers to all (further) quantum numbers that are needed to specify the state uniquely. Moreover,  $\{c_r(\alpha)\}$  denotes the representation of the atomic state in the given basis.

- In most standard computations, the set of CSF  $\{|\gamma_r P J M\rangle\}$  are constructed as antisymmetrized products of a common set of orthonormal orbitals. In contrast to the HF method, **both the radial (one-electron) functions as well as the expansion coefficients  $\{c_r(\alpha), r = 1, \dots, n_c\}$  are optimized simultaneously** on the basis of the Dirac-Coulomb Hamiltonian. All **further relativistic contributions to the energies and representation of the atomic states**, which arise for instance from the (relativistic) interaction among the electrons or with the radiation field, are then treated in first-order perturbation theory by diagonalizing the Dirac-Coulomb-Breit Hamiltonian matrix.
- Standard MCDHF calculations give rise to the energies and wave functions  $(E_{\alpha}, |\alpha \mathbb{J}M\rangle)$  for one or several selected bound states. From the symmetry of atoms and ions with regard to a rotation or inversion of the coordinates it is clear, that the Hamiltonian matrix is always block-diagonal in the total parity and angular momentum quantum numbers,  $\mathbb{J} \equiv J^P$  and  $M$ , respectively.
- Based on the (one-electron) Dirac Hamiltonian, the **MCDHF method enables one to describe consistently the dominant effects of relativity and correlations within the same framework.**
- Owing to the symmetry of the Hamiltonian, the computational effort of the MCDHF method can be reduced by a rather large factor ( $> 10$ ), if the various (total) angular momenta and parities  $\mathbb{J} = J^P$  of the atomic states of interest are considered independently.

- The decision about the atomic Hamiltonian operator, that is utilized in a particular MCDHF computation, is often made on the basis of (1) the nuclear charge, (2) the charge state of the atom or ion, (3) its particular **shell structure** as well as (4) the atomic property under consideration. **Further relativistic and radiative corrections** to the (total energies of the) atomic levels can be added in various (effective) approximations.
- **Steps for practical MCDHF computations:** In the MCDHF method and, hence, in JAC, the atomic bound states are typically generated by a series of steps:
  - Definition of the nuclear parameters and the angular structure (and extent) of the CSF basis  $\{|\gamma P J M\rangle\}$  as used in the MCDHF ansatz above;
  - Algebraic evaluation and **computation of the spin-angular integrals (the so-called angular coefficients)**, based on standard techniques from Racah's algebra;
  - Initial estimates of all radial orbital functions and self-consistent field (SCF) calculations, based on the Dirac-Coulomb Hamiltonian;
  - **Configuration interaction (CI) calculations** in order to incorporate further relativistic contributions into the Hamiltonian matrix and/or to enlarge the CSF basis beyond the given SCF model.

Often, **these four steps are repeated at several stages by systematically enlarging the one-electron (active) orbital space.**

### Restricted active spaces:

- **Romplete active space (RAS) method:** To include the **correlated motion of the bound electrons** in a more systematic fashion, the RAS method has been found useful for simple shell structures and has been applied along various isoelectronic sequences. The idea of this method is to account for an **excitation of the active electrons from the outer shells (into a number of predefined unoccupied orbitals)**, while some 'electronic core' remains usually fixed in the set-up of the RAS.
- For a restricted active space, it is generally supposed that all possible excitations (distributions) of the valence electrons can be taken into account among the set of active orbitals. In practice, **even the concept of a restricted active space is often not only unfeasible but also unnecessary for predicting many atomic properties**, since only those CSF with a "total excitation energy" in the same region as the reference configurations need often to be taken into account in the MCDHF expansion.
- The necessary CSF expansions often increases by one or even several orders of magnitude for open *d*- or *f*-shells, and especially if **reference configurations with nearly half-filled shells occur**. — With the JAC tools, here we aim for a new generation of atomic structure programs which enables one to compute the properties of such open *d*- and *f*-shell elements with an accuracy comparable to that of simpler shell structures.

## 3.5. Atomic self-consistent-field and configuration-interaction calculations

### Hamiltonian matrix elements:

- For a given CSF basis, the representation of an atomic state in the MCDHF ansatz above (i.e. the mixing coefficients  $\mathbf{c}(\alpha) \equiv (c_1(\alpha), c_2(\alpha), \dots, c_{n_c}(\alpha))$ ) is obtained by solving the **secular equation**

$$\det(\mathbf{H} - E_\alpha^{(n_c)} \mathbf{I}) = 0, \quad \mathbf{H} = (H_{rs}) = (\langle \gamma_r P J M | \mathbb{H} | \gamma_s \bar{P} \bar{J} \bar{M} \rangle \delta_{P\bar{P}} \delta_{J\bar{J}} \delta_{M\bar{M}})$$

where  $E_\alpha^{(n_c)}(PJ)$  denotes the eigenvalue.

- The **Hamiltonian matrix is block-diagonal in the total parity and angular momentum of the atom**, and is independent of the particular choice of the electron-electron interaction in the Hamiltonian.
- Since the **Hamiltonian matrix is real and symmetric**, all atomic states are orthogonal for  $E_\alpha \neq E_\beta$  or can be chosen in this way for  $E_\alpha = E_\beta$ . In practice, an **efficient decomposition of the many-electron matrix elements  $H_{rs}$  in the Hamiltonian matrix above is central to every implementation of the MCDHF method**.
- **Spin-angular integrals:** The integration over the spin-angular variables of all  $N$  electrons can be performed algebraically and enables one to write the matrix elements always in the form (Grant, 1989)

$$H_{rs} = \sum_t u_{rs}(a_t b_t) \langle a_t | \mathbb{H}_D | b_t \rangle + \sum_t v_{L;rs}(a_t b_t c_t d_t) X^{(L_t)}(a_t b_t c_t d_t).$$

In this expansion,  $u_{rs}(ab)$  and  $v_{L;rs}(abcd)$  are one- and two-particle (scalar) angular coefficients, while the  $X^{(L)}(abcd)$  describe the effective interaction strengths of (formally) the four electrons  $a, b, c, d$  that are involved in the interaction. As usual, we here abbreviate the one-particle quantum numbers  $a = (n_a, \kappa_a)$ ,  $b = (n_b, \kappa_b)$ , ... to have a compact notation for the subshells of equivalent electrons.

- **Reduced one-electron matrix elements of the Dirac Hamiltonian:** The reduced matrix element in the first term of the Hamiltonian matrix elements

$$\langle a | \mathbb{H}_D | b \rangle = \delta_{\kappa_a \kappa_b} \int_0^\infty dr \left[ c Q_a \left( \frac{d}{dr} + \frac{\kappa_a}{r} \right) P_b + c P_a \left( -\frac{d}{dr} + \frac{\kappa_a}{r} \right) Q_b - 2c^2 Q_a Q_b + V_{\text{nuc}}(r) (P_a P_b + Q_a Q_b) \right],$$

contains the kinetic and potential energy of an electron in subshell  $a$ , and can be expressed in terms of the *large* and *small* (radial) components,  $P_a(r)$  and  $Q_a(r)$ , of the corresponding one-electron orbital functions.



### Combined CI-MBPT method:

- **All-order techniques:** From many-body and coupled-cluster (CC) computations, it is known that certain classes of correlation diagrams need to be incorporated to all orders to obtain accurate results. One of the most popular all-order technique is the **CC approach in which correlation contributions are classified due to the number of excited particles rather than by the order of the perturbation**. In practice, most CC computations include single and pair correlations to all orders but neglect triple and higher correlations, though selected triple and quadrupole excitations might be included (Dzuba and Johnson, 1998).
- **Combined CI-MBPT method:** The **diagonalization of the many-electron Hamiltonian matrix** is perhaps the simplest way to include electronic correlations, though only very few core excitations can typically be taken into account by the this configuration interaction method. More often, correlations between the valence and core electrons are calculated by means of many-body perturbation theory (MBPT), and later included into the CI matrix. This method appears however very sensitive to the particular choice and representation of the valence electrons. Of course, no correlations between the core and valence electrons are included if excitations from the core are prohibited.
- **Combined CI-MBPT method:** If, instead of the standard MCDF expansion above, we consider an **extended expansion with single and/or double excitations from the core** ( $\sum_k$ )

$$\psi_\alpha(PJM) = \sum_r c_r(\alpha) |\gamma_r PJM\rangle + \sum_t c_t(\alpha) |\gamma_t PJM\rangle$$

and if the second term is treated as a perturbation with regard to a variation of the  $\{c_r\}$  coefficients, one obtains a generalized eigenvalue problem (Dzuba and Johnson, 1998)

$$\sum_r \left( \langle \Phi_r | \mathbb{H} | \Phi_s \rangle + \sum_t \frac{\langle \Phi_r | \mathbb{U} | \Phi_t \rangle \langle \Phi_r | \mathbb{U} | \Phi_t \rangle}{E_{\text{val}} - E_k} \right) c_s(\alpha) = E_s^{(\text{valence})} c_s.$$

Here,  $E_t = \sum_{r=\text{particles}} \varepsilon_r - \sum_{s=\text{holes}} \varepsilon_s$  is the energy of the virtual configuration with single and/or double excitation from the core and

$$\mathbb{U} = \sum_{i < j}^N \frac{1}{r_{ij}} - \sum_i V^N(r_i)$$

refers to the residual Coulomb interaction. **The generalized eigenvalue equation above is the main equation of the combined CI-MBPT method** and just differs from the standard CI equations by the second-order term with regard to the residual Coulomb interaction  $\mathbb{U}$ . **This**

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combined CI-MBPT method can be applied if the residual Coulomb interaction is small when compared to the excitation energies of the core.

- **Effective Hamiltonian:** The combined (second-order) CI+MBPT methods coincides with the standard CI method if an effective Hamiltonian

$$\mathbb{H} \rightarrow \mathbb{H} + \sum_t \frac{\mathbb{U} |t\rangle \langle t| \mathbb{U}}{E_s^{(\text{valence})}}$$

is used, while the standard expression for the wave function ( $\sum_i$ ) is employed. Note that this combined (second-order) CI+MBPT means that only single and double excitations from the core can be included. Together with excitations of the valence electrons, however, the combined CI+MBPT method goes far beyond pair correlations.

## 3.6. Atomic estimates of quantum-electrodynamic (QED) corrections

### 3.6.a. QED model operators & model potentials

#### Motivation & notations:

- Various proposals have been made in the literature to incorporate the (radiative) quantum-electrodynamic corrections by different model potentials into correlated many-electron methods, such as the MCDHF, many-body perturbation or coupled-cluster theories.
- **Effective single-electron QED Hamiltonian:** To a good approximation, these QED corrections can be incorporated into the (many-electron) representation  $\{c_r(\alpha\mathbb{J})\}$  of some level  $|\alpha\mathbb{J}\rangle$  by means of a local single-electron QED Hamiltonian. This Hamiltonian can be separated into two parts

$$\mathbb{h}^{(\text{QED})} = \mathbb{h}^{(\text{SE})} + \mathbb{h}^{(\text{VP})},$$

the self-energy (SE) and vacuum-polarization (VP) Hamiltonians, and which are often written in terms of (effective) potentials.

- **Vacuum-polarization:** The vacuum-polarization is given by a *local* potential and can be further splitted into an Uehling and a Wichmann-Kroll potential; this local potential is rather straightforward to calculate and has already been applied in many applications in the past.

- **Self-energy Hamiltonian  $\mathfrak{h}^{(\text{SE})}$**  : Such a Hamiltonian can be generally represented as sum of a local and non-local potential. These potentials are typically based on *ab initio* calculations of the diagonal and nondiagonal matrix elements of the one-loop QED operator with H-like wave functions (Shabaev *et al.*, 2013).
- Shabaev and coworkers (2013) have shown that the **QED corrections can be systematically incorporated into an effective Hamiltonian that acts in the space of Slater determinants**, if these determinants are built from one-electron positive-energy states only and if the total (many-electron binding) energies of the system are smaller than the pair-creation energy. This effective QED Hamiltonian can then be added to the Dirac-Coulomb-Breit Hamiltonian.
- **Beyond QED model Hamiltonians**: Explicit calculations of radiative corrections are extremely sophisticated and time-consuming for many-electron systems, if the standard QED perturbation theory is to be applied. **Until the present, such detailed QED computations can be carried out only for highly-charged, few-electron ions and by still making use of an effective screening potential** (instead of the complete perturbation expansion for all electron-electron pairs).

### Vacuum polarization potentials:

- **Uehling potential**: This potential is the dominant part of the VP and can be obtained from the direct numerical integration of the well-known formula

$$V^{(\text{Uehling})}(r) = -\alpha Z \frac{2\alpha}{3\pi} \int_0^\infty dr' 4\pi r' \rho(r') \int_1^\infty dt \left(1 + \frac{1}{2t^2}\right) \frac{\sqrt{t^2 - 1}}{t^2} \frac{e^{-2m|r-r'|t} - e^{-2m(r+r')t}}{4mrt},$$

and where  $Z\rho(r)$  denotes the density of the nuclear charge distribution, if normalized to  $\int dV \rho(r) = 1$ .

- **Approximate Uehling potential**: The Uehling potential can be calculated also by an approximate formula for a point-like nucleus (Flambaum and Ginges, 2005)

$$V^{(\text{simplified Uehling})}(r) = \frac{2\alpha}{3\pi} V(r) \int_1^\infty dt \left(1 + \frac{1}{2t^2}\right) \frac{\sqrt{t^2 - 1}}{t^2} e^{-2mrt}.$$

- **Wichmann-Kroll potential**: A detailed evaluation of the Wichmann-Kroll potential is (much) more sophisticated but can be estimated again to a rather good accuracy by the approximate formula for a point-like nucleus and with the classical radius  $r_o$  (Flambaum and

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Ginges, 2005)

$$V^{(\text{simplified WK})}(r) = -\frac{2\alpha}{3\pi} V(r) \frac{0.092 Z^2 \alpha^2}{1 + (1.62 r/r_c)^4}.$$

- The Wichmann-Kroll correction is generally small and gives a  $\sim 1\%$  contribution to the VP shifts only for  $Z \gtrsim 80$  (Flaumbaum and Ginges, 2005).

#### Local self-energy potentials:

- Shabaev *et al.* (2013) suggest a **local but independent part of the self-energy potential for each symmetry block  $\kappa$**  of the one-electron orbitals of the form

$$V_{\kappa}^{(\text{local SE})}(r) = A_{\kappa} \exp\left(-\frac{2\pi r}{\lambda_C}\right) = A_{\kappa} \exp\left(-\frac{r}{\alpha}\right),$$

where the constants  $A_{\kappa}$  are chosen to reproduce the SE shift for the lowest (one-electron, H-like) energy of each given  $\kappa$  symmetry, and where  $\lambda_C = h/mc$  is the Compton wavelength of the electron.

- These local contributions can also be cast be into the form

$$V^{(\text{local SE})}(r) = \sum_{\kappa} A_{\kappa} P_{\kappa} \exp\left(-\frac{2\pi r}{\lambda_c}\right), \quad \langle a | P_{\kappa} | b \rangle = \delta_{\kappa, \kappa_a} \delta_{\kappa, \kappa_b}.$$

- **Self-energy interaction strength:** Following the discussion with Volotka (2019), a (single-electron) **local self-energy interaction strength** for the (local) potential above can be written

$$\langle a || \mathbb{h}^{(\text{local SE})} || b \rangle = \left[ \frac{\langle n_g \kappa || \mathbb{h}^{(\text{local SE})} || n_g \kappa \rangle}{\langle n_g \kappa || \exp\left(-\frac{r}{\alpha}\right) || n_g \kappa \rangle} \right]_{\text{hydrogenic}} \langle a || \exp\left(-\frac{r}{\alpha}\right) || b \rangle$$

where  $|n_g \kappa\rangle = \{1s_{1/2}, 2p_{1/2}, 2p_{3/2}, 3d_{3/2}, \dots\}$  refers to the lowest hydrogenic orbitals of symmetry  $\kappa$  and nuclear charge  $Z$ .

- **Self-energy interaction strength:** Shabaev *et al.* (2013) have tabulated values for  $\langle n_g \kappa || \mathbb{h}^{(\text{local SE})} || n_g \kappa \rangle_{\text{hydrogenic}}$  for selected nuclear charges, which can be readily applied in order to compute the self-energy interactions strengths for general many-electron atoms and ions.

### Radiative potentials:

- **Radiative potential:** Flambaum and Ginges (2005) suggest an approximate expression for a (so-called) **radiative potential** which can be utilized in order to estimate the radiative corrections in strong Coulomb fields to the energies and electric-dipole transition amplitudes in atomic (many-electron) computation. This (alternative) potential is said to give good QED estimates with an accuracy of just a few percent but has been mainly tested and applied to (neutral) cesium only.
- **Radiative potential:** Flambaum and Ginges (2005) decompose this radiative potential into five terms which are explained either above or below:

$$V^{(\text{total QED})} = V^{(\text{simplified Uehling})} + V^{(\text{magnetic form})} + V^{(\text{electric form})} + V^{(\text{low frequency})} + V^{(\text{simplified WK})}.$$

- **Magnetic form-factor potential:** For the magnetic form-factor contribution, Flambaum and Ginges (2005) give

$$V^{(\text{magnetic form})} = \frac{\alpha}{4\pi m} i \boldsymbol{\gamma} \cdot \boldsymbol{\nabla} \left[ V(r) \left( \int_1^\infty dt \frac{\exp(-2m r t)}{t^2 \sqrt{t^2 - 1}} - 1 \right) \right],$$

and where  $\boldsymbol{\gamma} = (\gamma_1, \gamma_2, \gamma_3)$  is the vector of gamma matrices.

- **Low-frequency potential to the electric form-factor:** Flambaum and Ginges (2005) suggest the (low-frequency) expression

$$V^{(\text{low frequency})} = -\frac{B(Z)}{e} Z^4 \alpha^5 m c^2 \exp\left(\frac{Z r}{a_o}\right), \quad B(Z) = 0.074 + 0.35 Z \alpha,$$

with the Bohr radius  $a_o$  and where the function  $B(Z)$  has been fitted in order to reproduce the radiative energy shifts for  $p$ -levels of selected heavy elements. This term should be used with some care however.

- **Electric form-factor potential (including high-frequency contributions):** Flambaum and Ginges (2005) display a final expression for the electric form-factor contribution

$$V^{(\text{electric form})}(r) = -A(Z, r) \frac{\alpha}{\pi} V(r) \int_1^\infty dt \frac{\exp(-2m r t)}{\sqrt{t^2 - 1}} \left[ \left(1 - \frac{1}{2t^2}\right) [\ln(t^2 - 1) + 4 \ln(1/Z\alpha + 0.5)] - \frac{3}{2} + \frac{1}{t^2} \right].$$

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Here, the function  $A(Z, r) = (1.071 - 1.976 x^2 - 2.128 x^3 + 0.169 x^4) m r / (m r + 0.07 Z^2 \alpha^2)$ , has been suggested with  $x = (Z - 80) \alpha$ . The detailed form of  $A(Z, r)$  was found by fitting the radiative shifts for the high Coulomb  $s$  levels from the literature. **This term should be used with good care again.**

- Note that this local potential by Flambaum and Ginges (2005) was optimized for weakly bound valence states of heavy neutral atoms and may be less accurate for strongly bound ionic or even core-hole states.

#### 3.6.b. In JAC implemented QED estimates

##### Models & applications:

- In JAC, the radiative QED corrections can be incorporated into the computation and representation of the wave functions if `AsfSettings(..)` is initialized with `qedModel = QedPetersburg()` or `qedModel = QedSydney()`. This singleton datatype tells JAC how the estimates for the vacuum polarization and the (local) self-energy contributions are incorporated into the Hamiltonian matrix, cf. below.
- In JAC, use `qedModel = NoneQed()` in `AsfSettings(..)` if no QED estimates are to be included; this is also the default.

##### Implementation:

- **QedPetersburg model:** Here, we simply add to all matrix elements  $\langle a || \mathbb{h}^{(1\text{-particle})} || b \rangle$  in the many-electron Hamiltonian matrix the (single-electron) matrix elements

$$\langle a || \mathbb{h}^{(\text{QED})} || b \rangle = \langle a || \mathbb{h}^{(\text{local SE})} || b \rangle + \langle a || \mathbb{h}^{(\text{simplified Uehling})} || b \rangle.$$

- **QedSydney model:** Analogous, here we add to all matrix elements  $\langle a || \mathbb{h}^{(1\text{-particle})} || b \rangle$  in the many-electron Hamiltonian matrix the (single-electron) matrix elements

$$\langle a || \mathbb{h}^{(\text{QED})} || b \rangle = \langle a || \mathbb{h}^{(\text{magnetic form})} || b \rangle + \langle a || \mathbb{h}^{(\text{electric form})} || b \rangle + \langle a || \mathbb{h}^{(\text{low frequency})} || b \rangle + \langle a || \mathbb{h}^{(\text{simplified Uehling})} || b \rangle.$$

- **Effective QED operator** With the definitions from above, the total QED operator for a many-electron atom includes of course a summation over all atomic electrons:

$$\mathbb{H}^{(\text{QED})} = \mathbb{H}^{(\text{SE})} + \mathbb{H}^{(\text{VP})} = \sum_j \mathbb{h}_j^{(\text{QED})} = \sum_j \left( \mathbb{h}_j^{(\text{SE})} + \mathbb{h}_j^{(\text{VP})} \right),$$

where the  $\mathbb{h}_j^{(\text{SE})}$  and  $\mathbb{h}_j^{(\text{VP})}$  refers to the one-electron operators for electron  $j$ .

## 3.7. Unitary $jjJ - LSJ$ transformation of atomic states

### 3.7.a. Transformation matrices from $jjJ$ - to $LSJ$ -coupling

#### Atomic coupling schemes:

- **Coupling schemes:** Atomic and ionic levels are usually identified and labeled by means of quantum numbers of an appropriate coupling scheme. However, these quantum numbers are well-conserved only for a pure coupling of symmetry-adapted many-electron state, and which is realized only (rather) approximately in Nature.
- $LSJ$ - and  $jjJ$ -coupling are nowadays the two most frequently applied coupling schemes in atomic theory.
- As usual in quantum mechanics, the transformation from one to another orthonormal basis (i.e. the representation of atomic levels in a particular coupling scheme) is simply given by a unitary matrix, although it is generally not easy to evaluate this transformation matrix for arbitrary shell structures of atoms and ions.
- **$LSJ$ -notation:** In atomic spectroscopy, the standard  $LSJ$ -notation of the levels is frequently applied for classifying the low-lying level structures of atoms or ions.
- **Labeling of levels:** Atomic energy levels are often labeled by means of their leading  $LSJ$ -coupled CSF in a wave function expansion. However, special care need to be taken, if the same CSF occurs largest in the expansion of two levels, since then, the classification is no longer unique.
- The lack of providing a fast and proper spectroscopic notation in relativistic computations may hamper the spectroscopic level classification of medium and heavy elements as well as the interpretation and analysis of inner-shell processes.

#### $jjJ - LSJ$ transformation of subshell states:

- **Unique subshell order:** Since each non-relativistic  $n\ell$ -shell (apart from the  $ns$  shells) split into two relativistic subshells with  $j = \ell \pm 1/2$ , the order of these subshell is relevant for the transformation of the many-electron basis and is always fixed in JAC to  $n\ell_{j=\ell-1/2}$ ,  $n\ell_{j=\ell+1/2}$ . For the  $np$ - and  $nd$ -shells, for example, we always assume the sequence  $np_{1/2}$ ,  $np_{3/2}$  and  $nd_{3/2}$ ,  $nd_{5/2}$  in all formulas below.
- **Non-relativistic subshell states:** In the transformation of the non-relativistic shell states  $|n\ell^N \alpha \nu LS\rangle$ , we generally have to consider the product of two (relativistic) subshell states  $\left| \begin{smallmatrix} \bar{N} & \bar{\nu} \\ n\bar{\kappa} & \bar{\nu}J \end{smallmatrix} ; \begin{smallmatrix} N & \nu \\ n\kappa & \nu J \end{smallmatrix} \right\rangle$ , and with  $\bar{N} + N = N$ . Here, the  $+$  and  $-$  on top of the quantum numbers just refer to the relativistic subshells with  $j = \ell \pm 1/2$ .



- The transformation between the (sub-) shell states in  $jjJ$ - and  $LSJ$ -coupling can be expressed with real Fourier coefficients as

$$\left| (n\bar{\kappa}^{\bar{N}} \bar{\nu} \bar{J}, n\kappa^{\dagger N} \dagger \nu \dagger J) J \right\rangle = \sum_{\alpha \nu LS} |n\ell^N \alpha \nu LSJ\rangle \left\langle n\ell^N \alpha \nu LSJ \left| (n\bar{\kappa}^{\bar{N}} \bar{\nu} \bar{J}, n\kappa^{\dagger N} \dagger \nu \dagger J) J \right. \right\rangle, \quad N \equiv \bar{N} + \dagger N$$

$$\left\langle n\ell^N \alpha \nu LSJ \left| (n\bar{\kappa}^{\bar{N}} \bar{\nu} \bar{J}, n\kappa^{\dagger N} \dagger \nu \dagger J) J \right. \right\rangle = \left\langle (n\bar{\kappa}^{\bar{N}} \bar{\nu} \bar{J}, n\kappa^{\dagger N} \dagger \nu \dagger J) J \left| n\ell^N \alpha \nu LSJ \right. \right\rangle.$$

- The evaluation of the (unitary) transformation does generally not involve the radial (orbital) functions but depends only on the spin-angular functions.
- Although the notation of the non-relativistic and relativistic (sub-) shell states is quite similar in  $jjJ$ - and  $LSJ$ -coupling, these states generally belong to different irreducible representations of the  $SO_3$  rotation group. This can be seen already from the splitting of each non-relativistic shell into two subshells as outlined above.

### Transformation of atomic states:

- *LSJ-coupled, non-relativistic CSF:*

$$|\gamma^{(\text{NR})} LSJ\rangle = |((\cdot ((n_1 \ell_1^{N_1} \alpha_1 \nu_1 L_1 S_1), (n_2 \ell_2^{N_2} \alpha_2 \nu_2 L_2 S_2)) L_{12} S_{12}, (n_3 \ell_3^{N_3} \alpha_3 \nu_3 L_3 S_3)) L_{123} S_{123}, \dots) LS) J\rangle$$

- *jjJ-coupled, relativistic CSF:*

$$|\gamma^{(\text{R})} J\rangle = \left| \left( \cdot \left( \left( (n_1 \bar{\kappa}_1^{\bar{N}_1} \bar{\nu}_1 \bar{J}_1), (n_1 \kappa_1^{\dagger N_1} \dagger \nu_1 \dagger J_1) \right) \bar{X}_1, (n_2 \bar{\kappa}_2^{\bar{N}_2} \bar{\nu}_2 \bar{J}_2) \right) \bar{X}_2, (n_2 \kappa_2^{\dagger N_2} \dagger \nu_2 \dagger J_2) \right) \bar{X}_2, \dots \right) J \right\rangle$$

- To determine the explicit transformation matrix for going from one to another (many-electron) basis, of course, all these quantum numbers above play an important role in practice. For the sake of simplicity, however, it is sufficient to just re-call that the  $jj$ -coupled CSF basis spans a (finite) part  $\mathcal{H}^{(\text{finite})}$  of the  $N$ -electron Hilbert space and that the matrix  $(\langle \gamma^{(\text{NR})} LSJ | \gamma^{(\text{R})} J \rangle)$  just represent the (usual) Fourier coefficients for the expansion of an atomic state with respect to the basis  $\{|\gamma^{(\text{NR})} LSJ\rangle\}$  in  $\mathcal{H}^{(\text{finite})}$ .
- *Representation of atomic state functions:*

$$|\psi_\alpha\rangle = \sum_r c_r^{(\text{R})}(\alpha) |\gamma_r^{(\text{R})} J\rangle = \sum_t c_t^{(\text{NR})}(\alpha) |\gamma_t^{(\text{NR})}(L_t, S_t) J\rangle$$

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- **Formal construction of the  $jjJ \rightarrow LSJ$  transformation matrix:** If we denote the basis states of the (standard) product basis of the (subshell) states in  $jjJ$ - and  $LSJ$ -coupling by the  $|\text{standard} - jjJ\rangle$  and  $|\text{standard} - LSJ\rangle$ , the transformation of a relativistic CSF can be formally written as

$$\begin{aligned}
 |\gamma_r^{(R)} J\rangle &= \sum_{\text{standard}-jjJ} |\text{standard} - jjJ\rangle \langle \text{standard} - jjJ | \gamma_r^{(R)} J \rangle \\
 &= \sum_{\substack{\text{standard}-jjJ \\ \text{standard}-LSJ \\ \text{atomic}-LSJ}} |\gamma_t^{(NR)}(L_t, S_t) JP\rangle \langle \gamma_t^{(NR)}(L_t, S_t) JP | \text{standard} - LSJ \rangle \langle \text{standard} - LSJ | \text{standard} - jjJ \rangle \langle \text{standard} - jjJ | \gamma_r^{(R)} J \rangle \\
 &= \sum_{Q_p, Q_r} |\gamma_t^{(NR)}(L_t, S_t) J\rangle \times [\text{LSJ} - \text{recoupling}] \times [\text{product of LSJ} - jjJ \text{ matrix elements}] \times [jjJ - \text{recoupling}]
 \end{aligned}$$

Here, a summation over  $\text{standard} - jjJ$ ,  $\text{standard} - LSJ$ , ... implies a summation over all quantum numbers that are not bound due to given values at the left-hand-side of the expression. In the last line, in particular, the summation runs over the (product basis) quantum numbers  $Q_p = [(\alpha_1 \nu_1 L_1 S_1), (\alpha_2 \nu_2 L_2 S_2), \dots]$  as well as over the (re-coupling) quantum numbers as defined by the [LSJ-recoupling] and [jjJ-recoupling] recoupling coefficients,  $Q_r = [J_1, J_2, J_{12}, J_3, J_{123}, \dots, T_1, T_2, T_{12}, T_3, T_{123}, \dots]$  with  $\delta_{J_1, T_1} \delta_{J_2, T_2} \delta_{J_{12}, T_{12}} \dots$ .

- **Re-coupling from the  $LSJ$ -product to the  $(L_t, S_t) J$ -coupled basis:**

$$\begin{aligned}
 \langle \gamma_t^{(NR)}(L_t S_t) J | \text{standard} - LS \rangle &\equiv [\text{LSJ} - \text{recoupling}] \\
 &= \langle ( \cdot ( \cdot ( (L_1, L_2) L_{12}, L_3 ) L_{123}, \dots ) L_t, ( \cdot ( (S_1, S_2) S_{12}, S_3 ) S_{123}, \dots ) S_t ) J | ( \cdot ( ( (L_1, S_1) T_1, (L_2, S_2) T_2 ) T_{12}, (L_3, S_3) T_3 ) T_{123}, \dots ) J \rangle
 \end{aligned}$$

- **Re-coupling from the  $jjJ$ -product to the  $jXJ$ -coupled basis:**

$$\begin{aligned}
 \langle \text{standard} - jj | \gamma_r^{(R)} J \rangle &\equiv [jjJ - \text{recoupling}] \\
 &= \left\langle \left( \cdot \left( \left( (\bar{J}_1, \bar{J}_1) J_1, (\bar{J}_2, \bar{J}_2) J_2 \right) J_{12}, (\bar{J}_3, \bar{J}_3) J_3 \right), J_{123}, \dots \right) J \left| \cdot \left( \left( \left( (\bar{J}_1, \bar{J}_1) \bar{X}_1, \bar{J}_2 \bar{X}_2, \bar{J}_2 \right) \bar{X}_2, \bar{J}_3 \right) \bar{X}_3, \dots \right) J \right\rangle
 \end{aligned}$$

- In practice, the summation over such large sets of quantum numbers is often replaced by a (simpler or even single) summation over the allowed and predefined sets of quantum numbers.

### 3.7.b. Re-coupling coefficients

#### Zero open shell:

➤ CSF with *no* open shell are the same in the  $jjJ$ - and  $LSJ$ -coupling scheme,

$$|\gamma_r^{(\text{R, closed})} \mathbb{J} = 0^+\rangle \equiv |\gamma_r^{(\text{NR, closed})} \mathbb{J} = 0^+\rangle$$

and can be simply replaced with their (expansion) coefficients in the expansion of the wave function.

#### Single open shell:

➤ Re-coupling from the  $LSJ$ -product to the  $(L_t, S_t)J$ -coupled basis:

$$[\text{LSJ} - \text{recoupling}] = \langle (L_1, S_1) J | (L_1, S_1) J \rangle = 1.$$

➤ Re-coupling from the  $jjJ$ -product to the  $jXJ$ -coupled basis:

$$[\text{jjJ} - \text{recoupling}] = \left\langle (\bar{J}_1, J_1) J | (\bar{J}_1, J_1) J \right\rangle = 1.$$

#### Two open shells:

➤ Re-coupling from the  $LSJ$ -product to the  $(L_t, S_t)J$ -coupled basis:

$$\begin{aligned} [\text{LSJ} - \text{recoupling}] &= \langle ((L_1, L_2) L_{12}, (S_1, S_2) S_{12}) J | ((L_1, S_1) T_1, (L_2, S_2) T_2) J \rangle \\ &= \sqrt{[T_1, T_2, L_{12}, S_{12}]} \begin{Bmatrix} L_1 & S_1 & T_1 \\ L_2 & S_2 & T_2 \\ L_{12} & S_{12} & J \end{Bmatrix}, \end{aligned} \quad Q_r = [T_1, T_2].$$

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➤ Re-coupling from the  $jjJ$ -product to the  $jXJ$ -coupled basis:

$$\begin{aligned}
 [\text{jjJ} - \text{recoupling}] &= \left\langle \left( (\bar{J}_1, \bar{J}_1^+) J_1, (\bar{J}_2, \bar{J}_2^+) J_2 \right) J \mid \left( ((\bar{J}_1, \bar{J}_1^+) \bar{X}_1, \bar{J}_2), \bar{X}_2, \bar{J}_2^+ \right) J \right\rangle \\
 &= (-1)^{\bar{J}_2 + \bar{J}_2^+ + J_1 + J} \delta(\bar{J}_1, \bar{J}_1^+, J_1) \sqrt{[J_2, \bar{X}_2]} \begin{Bmatrix} J_1 & \bar{J}_2 & \bar{X}_2 \\ \bar{J}_2^+ & J & J_2 \end{Bmatrix}, \quad Q_r = [J_1, J_2].
 \end{aligned}$$

➤ The summation over  $Q_r$  also implies all further quantum numbers that are needed in order to distinguish different subshell states in  $jjJ$  and to  $LSJ$  product basis, respectively. In JAC, this summation is performed by (1) applying the standard coupling  $L \oplus S = T_{\text{low}}, \dots, T_{\text{up}}$  and (2) by a summation over the corresponding subshell states from a internally (defined) list. An explicit summation over  $J_1, J_2$  can be omitted because of orthogonality, and which gives rise to the Kronecker factors  $\delta_{T_1, J_1}, \delta_{T_2, J_2}$ .

#### 3.7.c. In JAC implemented $jjJ - LSJ$ transformation

##### Transformation, notations & application:

➤ Representation  $c_s^{(\text{NR})}(\alpha)$  of atomic states.

$$|\psi_\alpha\rangle = \sum_t c_t^{(\text{NR})}(\alpha) \left| \gamma_t^{(\text{NR})} LSJ \right\rangle = \sum_r c_r^{(\text{R})}(\alpha) \left| \gamma_r^{(\text{R})} J \right\rangle$$

➤ Unique label:  $(\dots ((n_1 \ell_1)^{N_1} 2S_1+1 L_1; (n_2 \ell_2)^{N_2} 2S_2+1 L_2) 2S_{12}+1 L_{12}, (n_3 \ell_3)^{N_3} 2S_3+1 L_3 \dots) 2S+1 L_J$

Such a label is assigned to each selected level, for instance, by printing

$$[(n_1 \ell_1)^{N_1} 2S_1+1 L_1] 2S_{12}+1 L_{12} : [(n_2 \ell_2)^{N_2} 2S_2+1 L_2] 2S_{12}+1 L_{12} : [(n_3 \ell_3)^{N_3} 2S_3+1 L_3] 2S_{123}+1 L_{123} : \dots$$

➤ In JAC, a  $jjJ - LSJ$  transformation is performed for all selected levels if `jjLS = LSjjSettings(true, ...)` is set in `AsfSettings()`.

➤ In JAC, the  $LSJ$ -coupled expansion of selected levels is printed by default after the diagonalization of the (many-electron) Hamiltonian matrix.

➤ Apart from the  $LSJ$  assignment to atomic levels, JAC can rather easily support also the full transformation of the wave functions, if required for (nonrelativistic) computations. This is achieved by setting the cut-off parameter to *zero* in `LSjjSettings(...)`.

- A proper  $LSJ$  spectroscopic notation has been found useful in order to explain the excitation and decay properties of individual levels, at least qualitatively.
- The  $jjJ - LSJ$  transformation of the ( $jjJ$ -coupled) CSF is supported for all **shell structures with (up to) two open shells, including open  $s$ -,  $p$ -,  $d$ -shells**. Further work is required to fully implement also transformation of open  $f$ -shells.

### Implementation of $jjJ - LSJ$ transformation:

- Internally, a  $jj - LS$  transformation of selected levels from a given multiplet is carried out by the following steps:
  - 1) Generation of a non-relativistic shell list for the given multiplet, based on a proper relativistic subshell list. The program terminates with an error message if the relativistic subshells of the given multiplet are ordered in *non-standard order*, for example, ...,  $2p_{3/2}$ ,  $2p_{1/2}$ , ... or similar.
  - 2) Extract all non-relativistic configurations that contribute to the given multiplet and generate a corresponding set of nonrelativistic CSF states; define an instance of a nonrelativistic basis.
  - 3) Expand (in turn) each relativistic CSF into the non-relativistic basis and add their contributions to the (nonrelativistic) eigenvectors. This expansion is made independently of CSF with zero-, one-, two-, ... open (nonrelativistic) shells.
  - 4) Check and print the  $LSJ$  expansion of all selected levels in a neat format with respect to the (nonrelativistic) configurations as well as CSF.

```

"""
'struct LS_jj_qn' ... defines a struct for the generalized quantum numbers (qn) of a LS_jj matrix element.
+ NN      ::Int64    ... occupation w = NN
+ QQ      ::Int64    ... subshell total quasispin 2*Q = nu
+ LL      ::Int64    ... subshell total angular momentum 2*L
+ SS      ::Int64    ... subshell total angular momentum 2*S
+ JJ      ::Int64    ... subshell total angular momentum 2*J
+ Nm      ::Int64    ... subshell Nm
+ Qm      ::Int64    ... subshell quantum number 2*Qm = nu^-
+ Jm      ::Int64    ... subshell quantum number 2*Jm
+ Qp      ::Int64    ... subshell quantum number 2*Qp = nu^+
+ Jp      ::Int64    ... subshell quantum number 2*Jp
"""

```

## 3.8. Atomic interaction amplitudes

### Concept of interaction amplitudes:

- **Hamiltonian matrix:** The (scalar) Hamiltonian matrix elements in section 3.4 below describe the interaction energy between two ASF  $|\alpha \mathbb{J} M\rangle$  and  $|\alpha' \mathbb{J}' M'\rangle$ , and these Hamiltonian matrix elements are perhaps the most simplest (many-electron) **interaction amplitudes in atomic structure theory**:  $\langle \psi_\alpha | \mathbb{H} | \psi_{\alpha'} \rangle = \langle \alpha \mathbb{J} M || \mathbb{H} || \alpha' \mathbb{J}' M' \rangle = \sum_{rs} c_r^*(\alpha) H_{rs} c_s(\alpha') \delta_{JJ'} \delta_{MM'} \delta_{PP'}$ .
- In relativistic atomic structure theory, in particular,  $E_\alpha = \langle \alpha \mathbb{J} || \mathbb{H}^{(\text{DCB})} || \alpha \mathbb{J} \rangle \approx \langle \alpha \mathbb{J} || \mathbb{H}^{(\text{DC})} || \alpha \mathbb{J} \rangle$  is the total energy of a given level  $|\alpha \mathbb{J}\rangle$ .
- **Many-electron interaction amplitudes:** Using time-dependent perturbation theory and Fermi's Golden rule, most atomic level and transition properties can be expressed rather **similarly as matrix elements of a (so-called) transition or interaction operator  $\mathbb{T}$** :

$$\langle \psi_\alpha | \mathbb{T} | \psi_\beta \rangle = \sum_{rs} c_r^*(\alpha) T_{rs} c_s(\beta).$$

- **Reduced matrix elements:** More often than not, the transition operator  $\mathbb{T}$  can be expressed in terms of **spherical tensor operators of rank  $K$** , and, perhaps, with magnetic projections  $Q = -K, -K+1, \dots, K$ , respectively. These spherical tensor operators transform under rotations like the spherical harmonics  $Y_{KQ}(\vartheta, \varphi)$ :  $\mathbb{T} = \sum_{KQ} a(K, Q) \mathbb{T}_Q^{(K)}$ . – Instead of the complete matrix elements, that include the magnetic projections  $M$  of the given atomic or configuration states as well as of the tensor operators, one often makes then use of the reduced matrix elements by applying the **Wigner-Eckart theorem**:

$$T_{Q,rs}^{(K)} \equiv \langle \alpha_r \mathbb{J}_r M_r | \mathbb{T}_Q^{(K)} | \alpha_s \mathbb{J}_s M_s \rangle = \langle J_r M_r, KQ | J_s M_s \rangle \langle \alpha_r \mathbb{J}_r || \mathbb{T}^{(K)} || \alpha_s \mathbb{J}_s \rangle.$$

- Although these reduced transition amplitudes are frequently utilized in the literature, for instance in order to discuss and analyze various atomic processes, **not many codes make explicit use of these amplitudes as the natural building blocks** for describing the level structure and properties of (open-shell) atoms and ions. – More often than not, a **prior decomposition of these transition amplitudes** into various types of one- and two-particle (reduced) matrix elements or even directly into radial integrals is made well before any implementation or coding is considered. **This prior decomposition has seriously hampered the modelling of complex processes** and is still the reason, why most atomic structure codes support the computation of just a few properties, cross sections and rates.
- Indeed, the different techniques, which are employed in the decomposition of the many-electron amplitudes, as well as the large number of definitions (and notations) of angular- or radial-type integrals in the literature has hampered not only the comparison of different codes

during the last decades but made it difficult also to (re-) use these entities for other properties than those considered by the original program developers.

- In JAC, we attempt to overcome this situation and to make consequent use of the many-electron amplitudes  $\langle \alpha_r \mathbb{J}_r M_r \mid \mathbb{T}_Q^{(K)} \mid \alpha_s \mathbb{J}_s M_s \rangle$  and/or  $\langle \alpha_r \mathbb{J}_r \parallel \mathbb{T}^{(K)} \parallel \alpha_s \mathbb{J}_s \rangle$  as the central building blocks for describing the interaction among the electrons as well as with external particles and fields.
- The concept of these interaction amplitudes has helped to simplify and maintain the code. — Moreover, it (will hopefully) enable us to exploit the many-electron interaction and transition amplitudes also in second- and higher-order processes, once an appropriate (intermediate) basis  $\{\psi_\nu(PJM) \equiv |\alpha_\nu \mathbb{J}_\nu M_\nu\rangle, \nu = 1, \dots, n_\nu\}$  has been constructed for modelling some given process.

### Reduced matrix elements of spherical tensor operators for two CSF:

- For two CSF of well-defined symmetry, the **reduced matrix elements of spherical tensor operators** take always the form

$$\langle \gamma_r P_r J_r \parallel \mathbb{T}^{(K)} \parallel \gamma_s P_s J_s \rangle = \sum_t u(a_t b_t; K) X(a_t b_t; K), \quad \mathbb{T}^{(K)} = \sum_j \mathbb{T}^{(K)}(j; \dots) \quad \dots \text{one part. operators}$$

$$\langle \gamma_r P_r J_r \parallel \mathbb{T}^{(K)} \parallel \gamma_s P_s J_s \rangle = \sum_t v(L_t; a_t b_t c_t d_t; K) X^{L_t}(a_t b_t c_t d_t; K), \quad \mathbb{T}^{(K)} = \sum_{i < j} \mathbb{T}^{(K)}(i, j; \dots) \quad \dots \text{symmetric two part. operators,}$$

and where  $X(a_t b_t; K)$  and  $X^{L_t}(a_t b_t c_t d_t; K)$  are **one- and two-particle (effective) interaction strengths**, cf. section 3.1. These interaction strengths are specific to the particular transition or interaction operator under consideration, and **special care has to be taken that the same phase convention applies, if different interaction amplitudes are to be combined with each other.**

- In practice, there are two steps in the computation of the transition amplitudes above for any pair of CSF, which need to be performed separately: The (pure) angular coefficients are usually calculated by means of some proper program, for instance ANCO (Gaigalas *et al.*, 2001), for either a pair or a whole set of CSF, and all the non-vanishing coefficients are then returned together to the calling routine. For these coefficients, the associated one- or two-particle interaction strengths are evaluated and summed up to form the requested matrix element  $T_{rs}$  of the transition matrix.
- To decrease the computational effort, these (one- and two-particle) interaction strength are often also *stored* and *re-utilized* in various applications, a feature which has not yet been considered in JAC.
- A similar decomposition of the many-electron matrix elements (amplitudes) as for the Hamiltonian matrix elements into a sum of *angular coefficient*  $\times$  *interaction strength* can be made for every physically relevant (interaction) amplitude, independent of the particular rank of

### 3. Many-electron atomic interactions, state functions, density operators and statistical tensors

the corresponding interaction operator. Since most atomic processes can be traced back to just a (very) few of such interaction or transition amplitudes, we made use of them as the *building blocks in the design and implementation of the JAC tools*.

#### Electron-electron interaction:

- The interaction among the electrons is described by the scalar operator that occurs in the Dirac-Coulomb-Breit Hamiltonian; cf. section 3.1

$$\mathbb{V}^{(e-e)} = \mathbb{V}^{(\text{Coulomb})} + \mathbb{V}^{(\text{Breit})} = \sum_{i < j} \left( \frac{1}{r_{ij}} + b_{ij} \right).$$

- The Dirac-Coulomb-Breit Hamiltonian gives rise to the (reduced) interaction amplitudes

$$\langle \gamma_r P_r J_r \parallel \mathbb{V}^{(e-e)} \parallel \gamma_s P_s J_s \rangle = \langle \gamma_r P_r J_r \parallel \mathbb{V}^{(\text{Coulomb})} \parallel \gamma_s P_s J_s \rangle + \langle \gamma_r P_r J_r \parallel \mathbb{V}^{(\text{Breit})} \parallel \gamma_s P_s J_s \rangle.$$

For scalar operators, the full and reduced matrix elements coincide with each other and need not to be distinguished, since the Clebsch-Gordan coefficient in the Wigner-Eckert theorem then simply evaluates to  $\langle J_r M_r, 00 \mid J_s M_s \rangle = \delta_{J_r J_s} \delta_{M_r M_s}$  for  $K = Q = 0$ .

- Since the Coulomb and Breit interaction operators are both scalar operators and since they contribute *additively* to the total electron-electron interaction, the same (pure) angular coefficients  $v(L_t; a_t b_t c_t d_t; K = 0)$  occur in the decomposition of the amplitudes above, although different angular and parity selection rules apply for the various interaction strengths themselves, cf. section 3.1. In the computation of these amplitudes, the Breit interaction can therefore be added quite easily to the (usual) instantaneous Coulomb repulsion, if this appears appropriate for some particular process:

$$X^{(L)}(abcd) = X^{(L, \text{Coulomb})}(abcd) + X^{(L, \text{Breit})}(abcd).$$

- The electron-electron interaction amplitudes above are indeed utilized (and calculated) in many components of the JAC tools.

#### Electron-photon interaction:

- The (relativistic) interaction of an electron with the radiation field is central for describing all photo-excitation, ionization, emission and capture processes.
- The electron-photon interaction is described by the one-particle operator:  $\mathcal{R}_\lambda(\mathbf{k}) = \sum_i \boldsymbol{\alpha}_i \cdot \mathcal{A}_{\lambda,i}(\mathbf{k})$ ; cf. chapter 4 below.



### Electron-nucleus (hyperfine) interactions:

- For a nuclear spin  $I > 0$ , the **hyperfine interaction** describes the interplay of each atomic electron with the electric and magnetic (multipole) fields of the nucleus, apart from the dominant (electric-monopole) field of the nucleus due to its nuclear charge  $Ze$ .
- The two dominant contributions to the hyperfine interaction arise from the nuclear magnetic-dipole field  $\mathbf{A} = \frac{\boldsymbol{\mu} \times \mathbf{r}}{r^3}$  and the electric-quadrupole field  $\Phi(r) = \sum_{ij} \frac{r_i r_j}{2r^5} Q_{ij}$ , while all higher magnetic and electric multipole fields, that may in principle occur for nuclei with spin  $I > 1$ , are typically negligible.

## 3.9. Atomic density operators

### Motivation

- **Density matrix formalism:** In this formalism, the state of a physical system is characterized by means of statistical operators  $\rho$  which describe a single system or an ensemble of equally prepared collision systems in either a pure quantum state or in a mixture of different states with any given degree of coherence.
- The great benefit of using the (atomic) density matrix theory is that it enables one to “accompany” such an ensemble through one or several steps of excitation, decay and/or collision process, and without loss of quantum-mechanical information.
- **Transition operators  $R$  of the density-matrix elements** If one starts from a well-specified initial state of the system, as described by the operator  $\rho_i$ , all information about the atomic interactions, that govern the process, simply resides in the (so-called) *transition operators*  $R$ , and the final state operator just follows from the well-known relation:  $\rho_f = R \rho_i R^+$ .
- **Building blocks:** The electron-photon, electron-electron interaction and electron-nucleus interaction amplitudes from above are the building blocks to form the (scattering or) transition matrix for all atomic processes of interest and, hence, the density matrices of the system at the various steps of some particular excitation or decay process.
- For scattering states with a single electron in the continuum, these interaction amplitudes can be readily calculated by means of the JAC tools.
- Apart from the efficient computation of these building blocks, one also requires a simple book-keeping of these amplitudes in order to set-up the (atomic) density matrices at the various steps. To facilitate this handling and the book-keeping of these amplitudes, they are typically provided in JAC by some (amplitude) array, and together with all quantum numbers that are needed for their unique classification.
- The prior computation of the many-electron amplitudes enables one to start directly from the amplitudes (for one or several elementary processes) and to combine them into an appropriate form, in order to support the prediction of different atomic properties and parameters.
- **Set-up of the density matrix:** In practice, the computation of all interaction amplitudes is performed within three steps: (i) the generation of all necessary bound-state wave functions, (ii) the evaluation of the transition amplitudes between bound states and scattering states in the continuum, and (iii) the set-up and handling of the density matrices, i.e. by applying the formulas from chapter 7.

### Density operator of an atom in level $|\alpha\mathbb{J}\rangle$ :

- **Mixed states of a level  $|\alpha\mathbb{J}\rangle$  :** For an atom in a well-defined and isolated level  $(\alpha\mathbb{J})$ , a general mixed state with regard to its magnetic subspace (i.e. its projections  $M$ ) can always be written as

$$\hat{\rho} = \sum_{MM'} c_{MM'} |\alpha\mathbb{J}M\rangle \langle\alpha\mathbb{J}M'|.$$

- **Statistical tensor:** Often, it is most convenient to represent the intermediate state of the ions in terms of its so-called statistical tensors

$$\rho_{kq}(\alpha\mathbb{J}) = \sum_{MM'} (-1)^{J-M'} \langle JM, J(-M') | kq \rangle \langle\alpha\mathbb{J}M | \rho | \alpha\mathbb{J}M' \rangle.$$

and which are non-zero only for  $0 \leq k \leq 2J$  and  $-k \leq q \leq k$ . These statistical tensors transform like the spherical harmonics.

- Although both, the (reduced, mixed) density matrix and the statistical tensors of some level  $(\alpha\mathbb{J})$  are mathematically equivalent, the latter form can be transformed more easily and analogue to the spherical harmonics of rank  $k$  under a rotation of the coordinates.
- **(Initial) polarization state of an atom:** The polarization state of the initial atom can be characterized by the statistical tensors  $\rho_{kq}^{\mathbf{n}_t}(\alpha_i\mathbb{J}_i)$ , if defined with regard to the axis of the (target) polarization  $\mathbf{n}_t = (\vartheta_t, \varphi_t)$ . While the odd-rank tensors  $k = 1, 3, \dots$  are known to characterize the orientation of the target atom, the even-rank tensors  $k = 2, 4, \dots$  describe its alignment. The maximal rank of these tensors is limited by the condition  $k \leq 2J$ .

## 3.10. Parity- and time-violating atomic interactions

### 3.10.a. Interactions beyond the standard model

#### Standard electro-weak model:

- The success of the standard electro-weak model of elementary particles is indeed extraordinary. It has been tested by many physical processes that cover more than ten orders of magnitude in momentum transfer. The standard model also correctly predicted the existence of new particles, such as the neutral  $Z$  boson.

### 3. Many-electron atomic interactions, state functions, density operators and statistical tensors

- The standard model is sometimes understood also as a low-energy manifestation of a more complete theory that, perhaps, unifies the four forces. Many well-motivated extensions to the standard model have been proposed, such as supersymmetric, technicolour and left-right symmetric models, and they often predict physical phenomena quite different from those of the standard model.
- **CP symmetry:** This symmetry refers to the combined symmetry of charge conjugation  $C$  and parity  $P$ .
- **Violation of CP symmetry:** The violation of CP symmetry was first discovered in 1964 in decays of the neutral K mesons, and this violation is taken into account by the standard electro-weak model by just a **single complex phase in the quark-mixing matrix (the so-called Kobayashi-Maskawa mechanism)**.
- Various extensions of the standard model, such as supersymmetry, predict considerably larger EDM of the particles. **Any observation of these EDM would therefore lead unambiguously to a hint of new physics.**
- In addition, the standard model of particle physics is sometimes said to violate time-reversal ( $T$ ) invariance, but again only through just a single phase in the Cabibbo-Kobayashi-Maskawa matrix, that mixes the quark flavors.
- For a rather long time, K mesons remained the only system in which CP-violation were observed explicitly. In 2001, however, the BaBar and Belle collaborations detected CP violation also for neutral B mesons, although this is still consistent with the predictions of the standard model.
- **CPT symmetry:** The (mathematical) structure of most gauge theories suggest that the combined symmetry of charge conjugation  $C$ , parity change  $P$  and time-reversal violation  $T$  has to be conserved.
- **If CPT is conserved, as suggested by gauge theories, then CP-violation must be strictly accompanied by T (time-reversal) violation.** Until the present, however, there has been no (undisputed) direct observation of T violation and its detection is of fundamental interest by itself. A direct detection of T-violation may shed further light also on the origin of CP-violation.
- **Permanent electric dipole moment (EDM)** **The measurement of a per of neutrons, atoms, or molecule would show explicit evidence of T-violation.**
- **New physics:** theories in physics beyond the standard model predict new sources of T-violation and these sources may lead to EDM that are many orders of magnitude larger than those from the standard model. This may allow EDM measurements by means of current experiments. Already the parameter space of quite popular extensions of the standard model, such as such as supersymmetry, multi-Higgs models, and left-right symmetric models, are strongly restricted by current measurements.
- **Schiff theorem:** If an external electric field acts upon a neutral atom, that consists out of non-relativistic point-like charged particles with some EDM and that only interact via electrostatic forces with each other, the field is screened exactly at each particle. This screening

arises due to the polarization of the atomic electrons by the external field. Therefore, an external electric field cannot induce an atomic EDM. As shown by Schiff, however, this **shielding is incomplete if magnetic or finite-size effects are taken into account**, and this may give rise in principle to an atomic EDM.

### High-precision measurements on fundamental symmetries of atoms:

- **Atomic high-precision experiments:** High-precision measurements on the (violation of fundamental) symmetries in atoms are suitable to test the standard model of elementary particles and to search for new physics beyond it. Such precision experiments in atomic physics complement measurements in high-energy physics.
- Indeed, atomic experiments have played an important role in the verification of the standard model for more than 30 years ago already. While the first evidence for neutral-weak currents, i.e. the existence of the neutral Z boson, was discovered in neutrino scattering, a parity violation was first established in atomic experiments, and before it was later observed also in high-energy electron scattering.
- Today, **atomic physics plays a major role in the search for *new* physics beyond the standard model**.
- **PNC measurements for cesium:** For example, the cesium measurement on parity non-conservation in atomic physics have been found in excellent agreement with the standard model.
- **Atomic PNC measurements:** Instead of more and more accurate PNC measurements and computations, it appears useful today to perform such measurements for different isotopes of the same atom, since the ratio of the PNC signals for two different isotopes is rather insensitive to the details of the electron structure. For such measurements, the largest uncertainty to the PNC amplitudes arise from the unknown neutron distribution of the isotopes (Roberts *et al.*, 2014).

### 3.10.b. Parity-violating (P-odd, T-even) interactions

#### Motivation:

- **Parity non-conservation (PNC) effects:** These effects arise in atoms largely due to the exchange of  $Z^0$ -bosons between the atomic electrons and the nucleus.
- The exchange of  $Z^0$  bosons is formally described by the Weinberg-Salam theory and leads to (observable) phenomena, such as circular dichroism and optical rotation, and which have been investigated in order to detect P-violation in cesium, thallium, bismut and plumbum.
- There is also another contribution to atomic parity violation due to the exchange of  $Z^0$  bosons between each pair of electrons; however, this effect is negligibly small for heavy atoms and well suppressed by a factor  $10^{-3} \dots 10^{-4}$ . This contribution is usually neglected in most PNC studies.
- PNC effects in atoms can be caused also by nuclear moments, i.e. by parity-violating electromagnetic form factors that occur in addition to the usual nuclear magnetic-dipole and electric-quadrupole moments (form factors).
- **Rare-earth atoms:** Several PNC experiments in rare-earth atoms have been suggested for close-lying levels of opposite parity in order to enhance the PNC effects. Moreover, rare-earth often have various stable isotopes and, hence, the dependence on atomic theory can be removed by taking ratios of the measured PNC contributions for different isotopes.
- **Enhancement of PNC signals:** Several physical factors can contribute to an enhancement (or suppression) of the measured parity-violating signal for selected atomic transitions: (i) The PNC amplitude is expected to scale slightly faster than  $Z^3$  with the nuclear charge of the ions and make heavy elements even more favorable; and (ii) the existence of close-lying levels of different parity. For the alkaline-earth elements, there are often levels with a splitting  $\sim 10 \text{ cm}^{-1}$  which need to be compared to the well-studied  $6s - 7s$  PNC transition in Cs with a splitting  $\sim 104 \text{ cm}^{-1}$  (Roberts *et al.*, 2014).

#### Nuclear-spin independent P-odd electron-nucleus interactions:

- **Nuclear-spin independent Hamiltonian:** For non-relativistic nucleons, the nuclear-spin independent Hamiltonian for the P-odd electron-nucleus interaction is given by the effective single-electron operator

$$\mathbb{H}^{(\text{weak}-\text{charge})} = -\frac{G}{\sqrt{2}} \gamma_5 [Z C_{1p} \rho_p(r) + N C_{1n} \rho_n(r)] ,$$

where  $Z$  and  $N$  are the number of protons and neutrons, and where the proton and neutron densities are normalized to unity,  $\int dr \rho_p(r) = 1 = \int dr \rho_n(r)$ .

- If the proton and neutron densities coincide,  $\rho_p = \rho_n \equiv \rho$ , the nuclear-spin independent Hamiltonian simplifies and can be described in terms of a nuclear weak-charge  $Q_W$ ,

$$\mathbb{H}^{(\text{weak-charge})} = -\frac{G}{2\sqrt{2}} \gamma_5 Q_W \rho.$$

- **Nuclear weak charge  $Q_W$ :** This charge is very close to the neutron number and is in lowest order in the electro-weak interaction given by

$$Q_W = -N + Z(1 - 4\sin^2 \theta_W) \approx -N.$$

This value for  $Q_W$  is modified however by radiative corrections.

- The nuclear-spin independent electron-neutron interaction is a scalar that mixes levels (states) with the same total electron angular momentum but different parity.
- For heavy atoms, the nuclear-spin independent interaction due to the nuclear weak-charge gives the dominant contribution to parity violation, when compared to other mechanisms.

### Nuclear-spin dependent P-odd electron-nucleus interactions:

- **P-odd interactions that depend on the nuclear spin  $I$ :** There are different interactions that depend on the nuclear spin  $I$  due to: i) the neutral weak current; ii) hyperfine-induced neutral currents; or iii) the nuclear anapole moment.
- **Spin-dependent electron-nucleus interactions:** The dominating contribution in heavy atoms comes from the nuclear anapole moment ( $\kappa_a$ ), while smaller contributions still arise from the  $Z_0$  exchange due to spin-dependent electron-nucleus weak interactions ( $\kappa_z$ ) as well as the interaction of the weak charge ( $\kappa_Q$ ) with the magnetic hyperfine interaction (Roberts *et al.*, 2014).
- **Interaction with the nuclear anapole:** The anapole moment gives typically the dominant nuclear-spin dependent P-odd electron-nucleus interaction. The nuclear anapole moment arises from P-odd interactions inside the nucleus. These interactions manifest themselves in atoms through a slightly modified electro-magnetic interaction of the nucleus with atomic electrons.
- The anapole moment  $\kappa_a$  increases with the atomic mass number,  $\kappa_a \propto A^{2/3}$  and, thus, dominates in heavy atoms, when compared to other nuclear-spin dependent mechanisms listed above.
- **Hamiltonian for the electron-anapole (-moment) interaction:** For an external nucleon with orbital angular momentum  $\ell$ , the Hamiltonian is given by (Dzuba *et al.*, 2009)

$$\mathbb{H}^{(\text{anapole})} = \frac{G}{\sqrt{2}} \kappa_a \frac{K}{I(I+1)} \boldsymbol{\alpha} \cdot \mathbf{I} \rho(r), \quad K = (I + 1/2) (-1)^{I+1/2-\ell}.$$

### Anapole moment:

- **Anapole moment:** The anapole moment arises from P-odd, T-even nuclear moments (i.e. parity-violating nuclear forces) and usually requires less accurate atomic structure calculations. This may simplify the search for favorable conditions and atomic systems.
- The notion of the anapole moment was introduced by Zel'dovich just after the discovery of parity violation.
- However, an **nuclear anapole moment was unambiguously detected only 30 years later**. In 1997, a group at Boulder measured a nuclear anapole moment in  $^{133}\text{Cs}$  to an accuracy of 14 % in an atomic experiment. This is considered as the first observation of an electromagnetic moment that violates fundamental discrete symmetries (Ginges and Flambaum, 2003).



- A P-odd, T-even anapole moment of the nucleus arises due to the presence of a parity violating weak interaction between nucleons.
- The anapole moment is directed along the nuclear spin  $\mathbf{I}$ :  $\langle \mathbf{a} \rangle = -\pi \langle r^2 \mathbf{j} \rangle = a \mathbf{I}/I$ . Since the current vector  $\mathbf{j}$  changes its sign under reflection of coordinates, but not the spin  $\mathbf{I}$ , **an anapole moment need to violate parity but not time-reversal symmetry.**

### 3.10.c. Time-reversal violating (P-odd, T-odd) interactions

#### Motivation:

- **Time-reversal symmetry:** Many physical systems and processes are asymmetric under time-reversal. In classical mechanics, for example, a velocity reverses its direction under time reversal,  $v \rightarrow -v$ , while an acceleration does not. This time-reversal asymmetry is often associated with some energy dissipation and, hence, with the second law of thermodynamics. Without dissipation, the laws of mechanics are usually considered to be invariant with regard to time-reversal.
- In nuclei, P-odd and T-odd nuclear moments can arise due to an intrinsic EDM of the nucleons, or due to P-odd and T-odd nuclear forces. The P-odd and T-odd nuclear forces induce larger nuclear moments than a single nucleon EDM.
- Various P-odd and T-odd interactions can be considered, especially for heavy atoms and ions. These interaction include the **tensor-pseudotensor electron-nucleon (e-N) interaction, the scalar-pseudoscalar (e-N) interaction, the nuclear Schiff moment and the interaction of the electron EDM with the internal nuclear magnetic field of the atom.**
- An atomic EDM can be induced if the nucleus possesses P-odd and T-odd nuclear moments.

#### P-odd & T-odd Hamiltonians due to electron-nucleus interactions:

- The expressions for  $\mathbb{H}^{(\text{scalar-pseudoscalar})}$  and  $\mathbb{H}^{(\text{weak charge})}$  as well as  $\mathbb{H}^{(\text{I-independent})}$  are quite similar.
- These electron-nucleus interactions mixes atomic states of opposite parity and induces a static electric-dipole moments in atoms.

#### P-odd & T-odd nuclear moments:

- The operators  $\mathbb{H}^{(\text{scalar-pseudoscalar})}$  and  $\mathbb{H}^{(\text{tensor})}$  have both electronic and nuclear components. While these Hamiltonians are overall scalar, the electronic and nuclear operators can be of any (equal) rank. To obtain any non-zero matrix element, therefore, the triangle rule for coupling of angular momenta imposes restrictions on the angular momenta of the electron and nuclear states.

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- **Nuclear moments that violate parity and time-reversal invariance:** electric dipole, magnetic quadrupole, electric octupole. For an nuclear electric-dipole moment, the Hamiltonian must be of the form  $\mathbb{H}^{(\text{Schiff moment})}$  in order to mix electron states of opposite parity.
- **Nuclear electric-dipole (Schiff) moments:** If  $\mathbf{S}$  is nuclear Schiff moment, the associated P-odd and T-odd Hamiltonian is given by

$$\mathbb{H}^{(\text{Schiff-moment})} = -e \phi^{(\text{Schiff-moment})} = -4\pi e \mathbf{S} \cdot \nabla \delta(r).$$

Here  $\phi^{(\text{Schiff-moment})}$  is the electrostatic potential of the nucleus that corresponds to a P-odd and T-odd charge distribution.

- **Nuclear magnetic-quadrupole moments (MQM):** A general expression for the nuclear MQM can be constructed in terms of the total angular momentum of the system  $I$ ,

$$M_{ij} = -e \frac{3}{2} \frac{M}{I(2I-1)} \left[ I_i I_j + I_j I_i - \frac{2}{3} I(I+1) \delta_{ij} \right].$$

The quantity  $M$  is conventionally referred to as the nuclear MQM and is defined as the maximum projection of  $M_{ij}$  upon the nuclear axis,  $M = M_{zz}$ . The magnetic quadrupole moment violates parity and time-reversal invariance.

#### 3.10.d. Time-reversal violating atomic electric-dipole moments

##### Motivation:

- A non-zero electric-dipole moment (EDM) of atoms require P-odd and T-odd interactions, either among the nucleons or the electrons and the nucleus.
- Therefore, the **null measurements of EDM in atoms place severe restrictions upon new sources of CP-violation** that often arise in models beyond the standard model, such as supersymmetry.
- The nuclear EDM exceeds the EDM of single nucleons by one to two orders of magnitude. P-odd and T-odd nuclear forces generate all P-odd and T-odd nuclear moments, such as the Schiff and MQM moments. These nuclear moments can be 10-to-100 times larger than those generated by the presence of a nucleon EDM.
- In the standard model, an atomic EDM is suppressed by many orders when compared to predictions from (so-called) *new* theories. Therefore, the detection of an EDM would be **unambiguous evidence of new physics**, even if it contributes only weakly to a very small signal in atomic parity violation.

- An atomic EDM can arise from various P-odd and T-odd mechanisms (Ginges and Flambaum, 2004): i) an intrinsic EDM of an electron; (ii) a P-odd and T-odd electron-nucleon interaction; (iii) an intrinsic EDM of an external nucleon; (iv) a P-odd and T-odd nucleon-nucleon interaction.
- Until the present, no permanent EDMs in neutrons, atoms or molecules have been detected.
- The best limits on the electron electric dipole moment is presently derived from measurements of atomic EDM.

### Physics of the atomic EDM:

- An atomic EDM is associated with an atomic level  $|\alpha J\rangle$  and arises due to the admixture of levels (states)  $|\nu J_\nu\rangle$  with opposite-parity wave functions. This atomic EDM has the form

$$d^{(\text{atomic EDM})} = 2 \sum_{\nu} \frac{\langle \alpha J | \mathbb{D} | \nu J_\nu \rangle \langle \nu J_\nu | \mathbb{H}^{(\text{PT})} | \alpha J \rangle}{E_\alpha - E_\nu} = d^{(\text{atomic EDM})} \frac{\mathbf{F}}{F},$$

where  $\mathbb{D}$  is the electric dipole operator,  $\mathbb{H}^{(\text{PT})}$  the P-odd and T-odd operator that mixes  $|\alpha J\rangle$  with the set of wave functions  $|\alpha_\nu J_\nu\rangle$ , and  $F$  is the total angular momentum of the atom corresponding to the state  $|\alpha J\rangle$ .

- A typical EDM experiment is performed in parallel electric and magnetic fields and can be described by the Hamiltonian

$$\mathbb{H}^{(\text{em field})} = -\mu \cdot \mathbf{B} - \mathbf{d} \cdot \mathbf{E}.$$

- A linear Stark shift is measured by observing the change in frequency when the electric field is reversed, since it is sensitive to the P-odd and T-odd term  $\mathbf{E} \cdot \mathbf{B}$ .
- Interaction Hamiltonian of the electron EDM with an internal atomic electric field  $\mathbf{E}^{\text{int}}$ : The Hamiltonian for interaction of the electron EDM with internal atomic electric field  $\mathbf{E}^{\text{int}}$  can be written as (Dzuba and Flambaum, 2009)

$$\mathbb{H}^{\text{eEDM-electric}} = -d_e \sum_{i=1}^N (\gamma_o - 1)_i \boldsymbol{\Sigma}_i \cdot \mathbf{E}^{\text{int}}, \quad \boldsymbol{\Sigma}_i = \begin{pmatrix} \boldsymbol{\sigma}_i & 0 \\ 0 & \boldsymbol{\sigma}_i \end{pmatrix}.$$

- Atomic EDM: An atomic EDM can be caused by the interaction of the electron EDM (if it exists) with either the internal electric field of the atom or the scalar-pseudoscalar electron-nucleon (T-odd, P-odd) interaction; an EDM of atoms in their ground state is then given by

$$\mathbf{d}_{\text{atom}} = 2 \sum_{\nu} \frac{\langle 0 | \mathbb{D} | \nu \rangle \langle \nu | \mathbb{H}^{\text{eEDM-...}} | 0 \rangle}{E_o - E_\nu}$$

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where  $|0\rangle$  is the atomic ground state,  $\mathbb{D} = -e \sum_i \mathbf{r}_i$  is the electric-dipole operator and  $\mathbb{H}^{\text{eEDM}}$  the (P-odd, T-odd) interaction operator. Here, the summation need formally to be performed over the complete many-electron spectrum of the intermediate levels  $\nu$  with total energies  $E_\nu$ .

- Measurements of EDMs in paramagnetic atoms, i.e. for a total electron angular momentum  $J \neq 0$ , are most sensitive to leptonic sources of P- and T-violation, and especially the electron EDM.
- In contrast, measurements of EDM in diamagnetic systems with zero total electron angular momentum are most sensitive to P-odd and T-odd mechanisms in the hadronic sector.

### 3.11. Elements from atomic spectroscopy

#### Line profiles in atomic spectroscopy:

- **Experimental photon energy distribution:** In order to fit the observed spectra to Fano's formula, the experimental photon energy distribution (window function) need usually to be taken into account. This experimental window function can often be represented by Gaussian with a full-width-at-half maximum (FWHM) according to the experimental energy spread. Therefore, the Fano profile has to be convoluted with a Gaussian. Analogue convolutions are known from the Voigt line profiles in the emission spectroscopy of hot gases which also includes the convolution of a Lorentzian with a Gaussian line curve.
- **Convolution of a Lorentzian and a Gaussian line profile:** The Voigt profile arises from the convolution of a Lorentzian and a Gaussian line profile with the Lorentzian and Gaussian FWHM  $\Delta_L$  and  $\Delta_G$

$$L(E) = \frac{2A}{\pi} \frac{\Delta_L}{4(E - E_r)^2 + \Delta_L^2}, \quad G(E) = \frac{2}{\Delta_G} \sqrt{\frac{\ln 2}{\pi}} \exp \left[ -\frac{4(\ln 2) E^2}{\Delta_G^2} \right], \quad \int dE L(E) = A, \quad \int dE G(E) = 1$$

$$V(E) = \int_{-\infty}^{\infty} dE' L(E') G(E' - E) = A \frac{4\sqrt{\ln 2}}{\pi^{3/2} \Delta_G} \int_{-\infty}^{\infty} dE' \frac{\Delta_L}{4(E' - E_r)^2 + \Delta_L^2} \exp \left[ -\frac{4(\ln 2) (E - E')^2}{\Delta_G^2} \right].$$

- **Convolution of a Lorentzian and a Gaussian line profile:** The convolution of a Lorentzian profile with a Gaussian function generally results in a **Faddeeva function that can be utilized for a fast and accurate evaluation of the convolution, e.g., in peak fitting routines** (Schippers 2018). With the following definitions, the Voigt profile can be expressed in terms of the Faddeeva function as

$$t = \frac{2\sqrt{\ln 2}(E' - E)}{\Delta_G}, \quad x = \frac{2\sqrt{\ln 2}(E_r - E)}{\Delta_G}, \quad y = \frac{\Delta_L \sqrt{\ln 2}}{\Delta_G}, \quad z = x + iy$$

$$V(E) = A \frac{2\sqrt{\ln 2}}{\Delta_G \sqrt{\pi}} \frac{1}{\pi} \int_{-\infty}^{\infty} dt \frac{y e^{-t^2}}{(t - x)^2 + y^2} dt = A \frac{2\sqrt{\ln 2}}{\Delta_G \sqrt{\pi}} \Re[w(z)]$$

- **Faddeeva function  $w(z)$ :** The Faddeeva function is a scaled complex error function which, for  $\Im z = y > 0$ , is defined as (Schippers 2018)

$$w(z) = e^{-z^2} \operatorname{erfc}(-iz) = \frac{i}{\pi} \int_{-\infty}^{\infty} dt \frac{e^{-t^2}}{z - t}, \quad \Re[w(z)] = \frac{1}{\pi} \int_{-\infty}^{\infty} dt \frac{y e^{-t^2}}{(t - x)^2 + y^2}, \quad \Im[w(z)] = \frac{-1}{\pi} \int_{-\infty}^{\infty} dt \frac{(t - x) e^{-t^2}}{(t - x)^2 + y^2}.$$

$\Re[w(x + iy)]$  and  $\Im[w(x + iy)]$  can be displayed as functions of the scaled energy  $x$  for different ratios  $y$  of Lorentzian and Gaussian widths.

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- **Convolution of a Fano and a Gaussian line profile:** One can re-define the Fano line profile slightly differently from above with a positive amplitude  $a > 0$  and with the same definition of  $t$ ,  $x$  and  $y$  as above. This definition of the Fano profile is consistent with the theoretical treatment of the atomic photoabsorption cross section. For the convolution of a Fano and a Gaussian line profile, this gives rise to (Schippers 2018):

$$F(E) = \frac{2a}{q^2 \Delta_L \pi} \left[ \frac{(q + \epsilon)^2}{1 + \epsilon^2} - 1 \right], \quad \lim_{q \rightarrow \infty} F(E) \rightarrow L(E) .$$

$$C(E) = \int_{-\infty}^{\infty} dE' F(E') G(E' - E) = \frac{2a}{q^2 \Delta_L \pi} \left[ \frac{1}{\sqrt{\pi}} \int_{-\infty}^{\infty} dt \frac{[qy + (t - x)]^2 e^{-t^2}}{(t - x)^2 + y^2} dt - 1 \right], = \frac{a}{q^2} \frac{2\sqrt{\ln 2}}{\Delta_G \sqrt{\pi}} \left\{ (q^2 - 1) \Re(w) - 2q \Im(w) \right\} .$$

## 4. Atomic interactions with the radiation field

### 4.1. Wave equations & optical fields

#### 4.1.a. Homogeneous wave equation

##### General solutions:

- **Monochromatic plane-wave light:** Monochromatic plane-wave light is uniquely determined by either the electric ( $\mathbf{E}$ ) or magnetic ( $\mathbf{B}$ ) field. Without restriction, we here consider a plane-wave along  $\mathbf{k} \parallel \mathbf{e}_z$ .
- **Waves that propagate along the wave vector  $\pm \mathbf{k}$ :** We consider the homogenous wave equation (for  $\psi = \{\mathbf{E}, \mathbf{B}, \Phi, \mathbf{A}, \dots\}$ ) with the known solution (proof by substitution and by means of the chain rule), whose wave fronts propagate along  $\pm \mathbf{k}$

$$\square \psi(\mathbf{r}, t) = 0 \qquad \psi(\mathbf{r}, t) = \psi_+(\mathbf{k} \cdot \mathbf{r} + \omega t) + \psi_- \underbrace{(\mathbf{k} \cdot \mathbf{r} - \omega t)}_{\text{phase}}$$

for arbitrary functions  $\psi_{\pm}(\mathbf{k} \cdot \mathbf{r} \pm \omega t)$ , and if they fulfill the dispersion relation  $\omega = (\pm) c k$ .

- **(Complex) monochromatic plane waves:** If we restrict ourselves to periodic (harmonic) solutions, two obvious solutions are

$$\psi(\mathbf{r}, t) = A_- e^{i(\mathbf{k} \cdot \mathbf{r} - \omega t)} \qquad \text{or} \qquad \psi(\mathbf{r}, t) = A_+ e^{i(\mathbf{k} \cdot \mathbf{r} + \omega t)}$$

$$\psi = \text{const.} \quad \text{for} \quad t = t_o \qquad \Longleftrightarrow \qquad \mathbf{k} \cdot \mathbf{r} = \text{const.}$$

Here, the vector  $\mathbf{k}$  is called the wave or propagation vector and  $\omega$  the frequency of the plane waves.

### 4.1.b. Plane-wave radiation

#### Plane-wave solutions:

➤ **Plane waves:** Plane waves always propagate into the direction of the wave vector and can be written as

$$\mathbf{E}(\mathbf{r}, t) = \mathbf{E}_o e^{i(\mathbf{k} \cdot \mathbf{r} - \omega t)}, \quad \mathbf{B}(\mathbf{r}, t) = \mathbf{B}_o e^{i(\tilde{\mathbf{k}} \cdot \mathbf{r} - \tilde{\omega} t)}$$

However, since the em field must satisfy not only the homogenous wave equation  $\square \psi = 0$  but also Maxwell's equations for the coupling of the  $\mathbf{E}$ - und  $\mathbf{B}$  fields

$$\text{rot } \mathbf{E} = -\frac{\partial \mathbf{B}}{\partial t} \iff i(\mathbf{k} \times \mathbf{E}_o) e^{i(\mathbf{k} \cdot \mathbf{r} - \omega t)} = i\tilde{\omega} \mathbf{B}_o e^{i(\tilde{\mathbf{k}} \cdot \mathbf{r} - \tilde{\omega} t)}, \quad \omega = \tilde{\omega}; \quad \mathbf{k} = \tilde{\mathbf{k}}; \quad \mathbf{k} \times \mathbf{E}_o = \omega \mathbf{B}_o.$$

Of course, this Maxwell equation need to be fulfilled for all times and places.

➤ If we use of the three other Maxwell equations, we furthermore obtain:

$$\text{div } \mathbf{E} = 0 \implies \mathbf{k} \cdot \mathbf{E}_o = 0$$

$$\text{div } \mathbf{B} = 0 \implies \mathbf{k} \cdot \mathbf{B}_o = 0$$

$$\text{rot } \mathbf{B} = \frac{1}{c^2} \frac{\partial \mathbf{E}}{\partial t} \implies \mathbf{k} \times \mathbf{B}_o = -\frac{\omega}{c^2} \mathbf{E}_o \implies (\mathbf{k} \times \mathbf{B}_o)^2 = k^2 B_o^2 = \frac{\omega^2}{c^4} E_o^2 \implies E_o^2 = c^2 B_o^2.$$

Therefore, the **three vectors**  $\mathbf{E}_o$ ,  $\mathbf{B}_o$ ,  $\mathbf{k}$  form (up to any even permutation) an **orthogonal and right-handed system**, i.e. the vectors  $\mathbf{E}$  and  $\mathbf{B}$  are perpendicular on  $\mathbf{k}$  for plane waves.

➤ **Vector potential:**

$$\mathbf{A}(\mathbf{r}, t) = \mathbf{A}_\pm e^{i(\mathbf{k} \cdot \mathbf{r} \pm \omega t)} = \sum_{\lambda} A_{o,\lambda} \mathbf{e}_\lambda e^{i(\mathbf{k} \cdot \mathbf{r} \pm \omega t)}$$

➤ **Equation of a plane:** For plane waves, all spatial points  $\mathbf{r}$  with  $\mathbf{k} \cdot \mathbf{r} = \text{const.}$  have the same value at any time  $t = t_o$ .

➤ **Wave crest:** The wave crest of a plane wave is  $A_\pm e^{i(\mathbf{k} \cdot \mathbf{r} \pm \omega t)}$  and propagates into the direction  $\pm \mathbf{k}$ .



### Conserved properties of plane waves with $\mathbf{k} \parallel \mathbf{e}_z$ :

➤ (Linear) momentum:

$$p_k = -i \partial_k, \quad p_{x,y} \mathbf{A}(\mathbf{r}, t) = 0, \quad p_z \mathbf{A}(\mathbf{r}, t) = k_z.$$

➤ Orbital angular momentum:

$$\ell_z = -i (x \partial_y - y \partial_x) = -i \frac{\partial}{\partial \varphi}, \quad \ell_z \mathbf{A} = 0.$$

➤ Spin (of photons):

$$s_z = -i \begin{pmatrix} 0 & 1 & 0 \\ -1 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \quad s_z \mathbf{A} = \lambda \mathbf{A}.$$

➤ Total angular momentum:

$$j_z = \ell_z + s_z, \quad j_z \mathbf{A} = \lambda \mathbf{A}.$$

➤ For plane waves, the (four) good quantum numbers refer to  $\mathbf{p} = (p_x, p_y, p_z)$  and  $\lambda_c$ .

### 4.1.c. Polarization of plane waves in classical electrodynamics

#### Polarization of plane-waves beams:

➤ Apart from the (well-defined) frequency of plane-wave photons, each photon has generally a spin angular momentum  $\hbar$  that can be (partly) aligned in parallel or antiparallel to the direction of propagation. A full alignment of all photon spins gives then rise to a circularly-polarized light beam.

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- A beam with a circularly-polarized planar wavefront has an azimuthal component of the Poynting vector that is proportional to the radial intensity gradient, even if it has no orbital angular momentum. This gives rise to a finite value, if integrated over the total cross-section of the beam.
- **Polarization of the electric field:** Polarization is an additional property of (em) waves that describes the orientation of their oscillations; by convention, the **polarization**  $-1 \leq p \leq 1$  **refers to the electric field**. The polarization of light is associated with the spin-angular momentum (density) of the em wave.
- The **polarisation is said to be linear if  $p = 0$ , and is left- or right-circular for  $p = -1$  and  $p = +1$** , respectively. For all other values of  $-1 < p < 1$ , the wave is said to be **elliptically** polarized. Sometimes, one also refers to the **chirality or handedness of the waves** in order to denote the left- or right-circular polarization of the light beam.
- **Vector potential of a circularly-polarized plane wave:**

$$\mathbf{A}(\mathbf{r}, t) = \mathbf{e}_\lambda e^{i(kz - \omega t)}; \quad \mathbf{e}_\lambda = (\mathbf{e}_x + i\lambda \mathbf{e}_y) \quad \begin{cases} \lambda = +1 & \text{left circular} \\ \lambda = -1 & \text{right circular} \end{cases}$$

- **Spin angular momentum (SAM):** The circular polarization of light is usually associated with a well-defined SAM.
- **Complex spherical unit vectors  $\mathbf{e}_\pm$ :**

$$\mathbf{e}_\pm = \frac{1}{\sqrt{2}} (\mathbf{e}_x \pm i\mathbf{e}_y) \quad \Longleftrightarrow \quad \begin{aligned} \mathbf{e}_x &= \frac{1}{\sqrt{2}} (\mathbf{e}_+ + \mathbf{e}_-) \\ \mathbf{e}_y &= \frac{-i}{\sqrt{2}} (\mathbf{e}_+ - \mathbf{e}_-) \end{aligned}$$

- **Complex electric field  $\mathbf{E}^{(c)} = (E_x^{(c)}, E_y^{(c)}, 0)$ :** Although a complex notation of em fields is often very convinient, the transverse, real electric field is given by

$$\left. \begin{aligned} E_x^{(c)} &= |E_x^{(c)}| e^{i\gamma} \\ E_y^{(c)} &= |E_y^{(c)}| e^{i(\gamma+\delta)} \end{aligned} \right\} \quad \mathbf{E}^{(c)} = \Re[(E_x^{(c)} \mathbf{e}_x + E_y^{(c)} \mathbf{e}_y) e^{i(kz - \omega t)}] = E_x \mathbf{e}_x + E_y \mathbf{e}_y$$

$$E_x = |E_x^{(c)}| \cos(kz - \omega t + \gamma), \quad E_y = |E_y^{(c)}| \cos(kz - \omega t + \gamma + \underbrace{\delta}_{!!}) \quad \dots \text{real.}$$

- **Relative phase  $\delta$  to distinguish different polarization:** This relative phase  $\delta$  refers to the complex field amplitudes  $E_x^{(c)}$  and  $E_y^{(c)}$

- **Linear polarization along the angle  $\alpha$ :** For  $\delta = n\pi$  ( $n \in \mathbb{N}$ )

$$\mathbf{E}^{(c)} = (|E_x^{(c)}| \mathbf{e}_x \pm |E_y^{(c)}| \mathbf{e}_y) \cos(kz - \omega t + \gamma) = |\mathbf{E}^{(c)}| \mathbf{e}_\alpha \cos(kz - \omega t + \gamma)$$

$$|\mathbf{E}^{(c)}| = \sqrt{|E_x^{(c)}|^2 + |E_y^{(c)}|^2}, \quad \tan \alpha = \pm \frac{|E_y^{(c)}|}{|E_x^{(c)}|} \quad \dots \begin{cases} (+) & n = \text{gerade} \\ (-) & n = \text{ungerade} \end{cases}$$

$$\mathbf{e}_\alpha = \cos \alpha \mathbf{e}_x + \sin \alpha \mathbf{e}_y$$

The angle  $\alpha$  denotes the (polarization-) direction of  $\mathbf{E}$  with regard to the  $x$ -axis.

- **General monochromatic plane wave:** Any plane wave can always be written as superposition of two components

## 4.2. Representation and parametrization of photons in atomic theory

### 4.2.a. Stokes parametrization and density matrix of a photon

#### Stokes parameters:

- **Stokes parameters:** These parameters provide a simple and alternative characterization for the polarization of plane waves and are given in terms of three real parameters that describe the polarization state of (transversal) electromagnetic radiation.
- **Stokes parameters  $P_1, P_2, P_3$ :** The first two Stokes parameters quantify the *relative* asymmetries between intensities  $I_\chi$  of light that is linearly polarized under different angles  $\chi$  with regard to the reaction plane:  $P_1 = (I_o - I_{90})/(I_o + I_{90})$  and  $P_2 = (I_{45} - I_{135})/(I_{45} + I_{135})$ . The parameter  $P_3$  reflects the degree of circular polarization of the emitted photons.
- The Stokes parameters were defined by George Gabriel Stokes in 1852 as a mathematically convenient alternative to the more common description of incoherent or partly polarized radiation in terms of its total intensity  $I$ , the degree of polarization  $P$ , and the orientation of the (so-called) polarization ellipse.
- **Degree of polarization:** The polarization state of the photon is said to be *pure* if they obey the Stokes parameters obey the restriction  $\sum_{i=1}^3 P_i^2 = 1$ .

Table 4.1.: Stokes parametrization of photons for selected polarization of photons with  $\mathbf{k} \parallel \mathbf{e}_z$ .

Selected photon polarization	$P_1$	$P_2$	$P_3$
unpolarized	0	0	0
polarized along $x$ -axis	1	0	0
polarized along $y$ -axis	-1	0	0
right-circularly polarized	0	0	1
left-circularly polarized	0	0	-1

- **Degree of linear polarization:** The parameter  $P_\ell = \sqrt{P_1^2 + P_2^2}$  is called the **degree of linear polarization**.
- **Stokes parameters  $P_3 = P_c$ :** Photons with  $\lambda = +1$  are called **right-circularly polarized photons**, and photons with  $\lambda = -1$  are called **left-circularly polarized photons**.
- Note that the notation in optics is often opposite to that used in quantum electrodynamics for historical reasons.

### Alternative representation of polarized photons:

- Instead of the Stokes parameters  $P_1$  and  $P_2$ , it is often convenient for the analysis of different experimental situations to use the degree of linear polarization  $P_\ell$  and the direction of the principal axis of the polarization ellipse  $\varphi_0$ .
- **Polarization ellipse:** If, in the chosen coordinate frame, the  $x$ -axis is directed along the principal axis of a polarized photon beam ( $\varphi_0 = 0^\circ$ ), the Stokes parameter  $P_1$  takes the largest value  $|P_1| = P_\ell$ , while the Stokes parameter  $P_2$  vanishes. This easily shows that the Stokes parameters  $P_1$  and  $P_2$  depend on the choice of the  $x$ -axis even if the  $z$ -axis is fixed along the beam, while the degree of linear polarization  $P_\ell$  remains invariant.
- **Polarization ellipse:** It is often convenient to introduce the angle  $\varphi_0$  by the relations

$$\cos 2\varphi_0 = \frac{P_1}{P_\ell}, \quad \sin 2\varphi_0 = \frac{P_2}{P_\ell}$$

- **Probability for measuring a linearly polarized photon along  $\varphi$  in the  $xy$ -plane:** For a density matrix with given  $P_\ell$ ,  $\varphi_0$  Balashov *et al.* (2001) shows the probability distribution

$$W(\varphi) = \frac{1}{2}(1 + P_\ell \cos 2(\varphi - \varphi_0))$$

for measuring linearly-polarized photon along  $\varphi$  in the  $xy$ -plane.

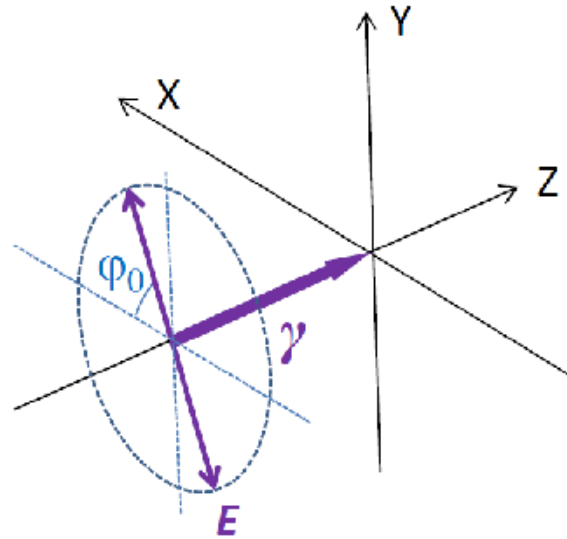


Figure 4.1.: Laboratory coordinate system.

- **Angle of the polarization vector:** The angle  $\varphi_0$  indicates the *principal axis* of the *polarization ellipse* of the photon beam. It describes the direction of the axis (with regard to the  $x$ -axis) which the intensity of the transmitted photon beam is the largest, while it is lowest intensity perpendicular to this direction.

### Photon density matrix:

- **Photon density matrix in helicity representation:** Most naturally, the spin state of an incident photon is expressed within its **helicity representation**

$$\rho_\gamma = \sum_{\lambda\lambda'} c_{\lambda\lambda'} |\mathbf{k}\lambda\rangle \langle \mathbf{k}\lambda'|,$$

where the helicity  $\lambda$  is the spin projection of the photon upon the direction of its propagation, i.e. along the momentum or wave vector  $\mathbf{k}$ .

- Since for a photon (with intrinsic spin  $s \equiv 1$ ) the helicity just takes the values  $\lambda = \pm 1$ , only three real parameters are required to describe the spin state of the photons. These parameters in the photon density matrix are closely connected to the Stokes parameters of light (Rose, 1961)

$$(c_{\lambda\lambda'}) = \langle \mathbf{k}\lambda | \rho_\gamma | \mathbf{k}\lambda' \rangle = \frac{1}{2} \begin{pmatrix} 1 + P_3 & P_1 - iP_2 \\ P_1 + iP_2 & 1 - P_3 \end{pmatrix}.$$

Note that this definition differs from those in Balashov *et al.* (2005) by a factor  $-1$  for the  $P_1$  and  $P_2$  parameters because of the different convention of the spherical unit vectors  $\mathbf{e}_\pm$ .

### 4.2.b. Pure polarization states of photons

#### Linearly-polarized plane-wave photons with $\mathbf{k} \parallel \mathbf{e}_z$ :

- In this section, we consider **pure polarization states of an incoming photon**, i.e. light with a photon density matrix  $P_1^2 + P_2^2 + P_3^2 = 1$ , and which is thus equivalent to a density-matrix with  $\rho^2 = \rho$ .
- With this restriction in mind, **any polarized state of a transverse photon** can be described by the linear combination

$$|e\rangle = g_+ \mathbf{e}_+ + g_- \mathbf{e}_-,$$

where  $\mathbf{e}_\lambda$  ( $\lambda = \pm$ ) are the **unit polarization vectors within the circular basis**, and with  $|g_+|^2 + |g_-|^2 = 1$ .

- The state of a linearly-polarized photon along some given direction  $\varphi_0$  in the  $x - y$  plane can be written in the form

$$|\mathbf{e}_{\varphi_0}\rangle = -\frac{1}{\sqrt{2}} \left( e^{-i\varphi_0} \mathbf{e}_+ - e^{i\varphi_0} \mathbf{e}_- \right).$$

- Pure linearly-polarized light with  $P_\ell = 1$ :  $|g_+| = |g_-|$ .

### Elliptically-polarized plane-wave photons with $\mathbf{k} \parallel \mathbf{e}_z$ :

- Photon density matrix in terms of the  $g_\pm$  coefficients: With these notations, an arbitrarily polarized pure state of a (transversal) photon can be expressed in terms of the  $g_\pm$  coefficients by either a  $2 \times 2$  matrix or by means of the corresponding Stokes parameters:

$$\langle \mathbf{k}\lambda | \rho | \mathbf{k}\lambda' \rangle = \begin{pmatrix} |g_+|^2 & g_+ g_-^* \\ g_+^* g_- & |g_-|^2 \end{pmatrix}, \quad P_1 = -2 \Re(g_+ g_-^*), \quad P_2 = 2 \Im(g_+ g_-^*), \quad P_3 = |g_+|^2 - |g_-|^2.$$

- Polarization ellipse: The linear polarization can be characterized also by  $P_\ell$  and  $\varphi_0$  with

$$P_\ell = 2 |g_+ g_-^*| = 2 |g_+| |g_-| \cos 2\varphi_0 = -\frac{\Re(g_+ g_-^*)}{|g_+ g_-^*|}, \quad \sin 2\varphi_0 = -\frac{\Im(g_+ g_-^*)}{|g_+ g_-^*|}.$$

These simple relations can be used to also find the  $g_\pm$ -coefficients for given parameters  $P_\ell$  and  $\varphi_0$ .

- Fully-polarized light with  $P_\ell \neq 1$ :

$$\text{for } |g_+| > |g_-|: \quad |g_+| = \frac{1}{2} \left( \sqrt{1+P_\ell} + \sqrt{1-P_\ell} \right), \quad |g_-| = \frac{1}{2} \left( \sqrt{1+P_\ell} - \sqrt{1-P_\ell} \right)$$

$$\text{for } |g_+| < |g_-|: \quad |g_+| = \frac{1}{2} \left( \sqrt{1+P_\ell} - \sqrt{1-P_\ell} \right), \quad |g_-| = \frac{1}{2} \left( \sqrt{1+P_\ell} + \sqrt{1-P_\ell} \right)$$

- The relative phase of the coefficients:  $g_+ = |g_+| e^{i\xi_+}$ ;  $g_- = |g_-| e^{i\xi_-}$  can be found from the equation:  $\xi_- - \xi_+ = 2\varphi_0 + \pi$ .
- Polarization vectors  $\boldsymbol{\epsilon}_1, \boldsymbol{\epsilon}_2$  of elliptical-polarized plane-wave radiation: An alternative representation of the polarization vector is given by:

$$\boldsymbol{\epsilon}_j = \left( \mathbf{e}_x \cos\left(\frac{\xi_j}{2}\right) + i \mathbf{e}_y \sin\left(\frac{\xi_j}{2}\right) \exp\left(\frac{i\eta_j}{2}\right) \right), \quad \boldsymbol{\epsilon}_j^* = \left( \mathbf{e}_x \cos\left(\frac{\xi_j}{2}\right) - i \mathbf{e}_y \sin\left(\frac{\xi_j}{2}\right) \exp\left(\frac{-i\eta_j}{2}\right) \right), \quad j = 1, 2.$$

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- **Two elliptically-polarized beams with the same angular momentum but opposite helicity:** The easiest way to describe such a superposition is  $\xi_1 = \xi_2 = \pi/2$ ,  $\eta_1 = \eta_2 = 0$ , and this leads to

$$\epsilon_1 = \epsilon_2 \equiv \epsilon, \quad \epsilon_1^* = \epsilon_2^* \equiv \epsilon^*, \quad \epsilon \cdot \epsilon^* = 1, \quad \epsilon \cdot \epsilon = \epsilon^* \cdot \epsilon^* = 0.$$

- **Two elliptically-polarized beams with the same helicity and opposite angular momentum:** These beams are described by  $\xi_1 = -\xi_2 = \pi/2$ ,  $\eta_1 = \eta_2 = 0$ , and the two polarization vectors then fulfill

$$\epsilon_1 = \epsilon_2^* \equiv \epsilon, \quad \epsilon_1^* = \epsilon_2 \equiv \epsilon^*, \quad \epsilon \cdot \epsilon^* = 1, \quad \epsilon \cdot \epsilon = \epsilon^* \cdot \epsilon^* = 0.$$

#### Beams of fully-polarized light with $\mathbf{k} \parallel \mathbf{e}_z$ :

- **Sequences of completely polarized light pulses:** In experiments, one often considers a sequence of (fully polarized) photon pulses, which propagate however into different directions.
- **Hierarchy of angles in defining the em pulses:** For such a sequence of pulses, all angles are defined with respect to a fixed laboratory frame  $\Sigma = (x, y, z)$ , and where we assume all coordinate systems to be properly ‘right-handed’.
- In the (primed) coordinates  $\Sigma'$  as associated with some given pulse, the polarization coefficients are assumed to be:

$$g_+ = 1, \quad g_- = 0 \quad (\text{for ‘helicity : } +1' \text{)}$$

$$g_+ = 0, \quad g_- = 1 \quad (\text{for ‘helicity : } -1' \text{)}$$

$$g_+ = -\frac{1}{\sqrt{2}} e^{-i\phi_0}, \quad g_- = \frac{1}{\sqrt{2}} e^{i\phi_0} \quad (\text{for ‘linear’ and given } \phi_0).$$

- **First pulse:** The wave vector of the first pulse always defines the  $z$ -axis and its linear polarization ( $\phi_0$ ) the  $x$ -axis of the fixed frame  $\Sigma$ , i.e. the first pulse defines the  $x - z$  scattering plane. For a circularly polarized pulse, the  $x$ -axis is defined indirectly due to the definition of subsequent pulses (see below); it need not to be defined explicitly if there is only a single, circularly-polarized pulse.



➤ Parameters of the first pulse:

$$\phi_1 = 0; \quad \theta_1 = 0; \quad (\phi_1 = 0 \quad \text{for circularly polarized light})$$

## 4.3. Multipole decomposition of the radiation field

### 4.3.a. Elements from the theory of multipole transitions

#### Transition amplitudes:

- Johnson (2008) provide a systematic decomposition of the transition amplitudes into multipole components.
- Transition amplitudes due to a transverse-gauge vector potential  $\mathbf{A}^{(\text{transverse})}(\mathbf{r}, \omega)$ : For a single-electron atom, the transition amplitude from state  $a = |n_a \kappa_a m_a\rangle \rightarrow b = |n_b \kappa_b m_b\rangle$  due to the interaction with a plane-wave radiation field with the transverse-gauge vector potential  $\mathbf{A}^{(\text{transverse})}(\mathbf{r}, \omega) = \boldsymbol{\epsilon} e^{i\mathbf{k} \cdot \mathbf{r}}$ , with  $\boldsymbol{\epsilon} \cdot \mathbf{e}_z = 0$  is

$$T_{ba} = \int d^3\mathbf{r} \, \psi_b^* \boldsymbol{\alpha} \cdot \mathbf{A}^{(\text{transverse})}(\mathbf{r}, \omega) \psi_a.$$

#### Multipole potentials:

- Plane-wave multipole potentials  $\mathbf{a}_{LM}^{(p)}(\mathbf{r})$ : In the multipole expansion of the plane-wave vector potential, the multipole potentials satisfy both, the Helmholtz equation and transversality condition

$$\nabla^2 \mathbf{a}_{LM}^{(p)} + k^2 \mathbf{a}_{LM}^{(p)} = 0, \quad \nabla \cdot \mathbf{a}_{LM}^{(p)} = 0 \quad p = 0, 1.$$

- The multipole potentials  $\mathbf{a}_{LM}^{(p)}$  with  $p = 0$  are the magnetic multipole potentials and with  $p = 1$  the electric multipole potentials.

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- **Electron-multipole interaction operator:** The interaction operator  $\boldsymbol{\alpha} \cdot \mathbf{a}_{LM}^{(p)}$  is an irreducible tensor operator of rank  $L$  and parity  $(-1)^{L+1-p}$ .
- The **multipole expansion of the (transverse-gauge) plane-wave vector potential** still contains all information about the polarization and propagation direction of the photons.
- **Representation of the multipole potentials:** Using the well-known identities of the spherical Bessel functions, the multipole potentials  $\mathbf{a}_{LM}^{(p)}$  can be expressed as (Johnson, 2007)

$$\mathbf{a}_{LM}^{(0)}(\mathbf{r}) = j_L(kr) \mathbf{Y}_{LM}^{(0)}(\mathbf{r}), \quad \mathbf{a}_{LM}^{(1)}(\mathbf{r}) = \left[ j_L'(kr) + \frac{j_L(kr)}{kr} \right] \mathbf{Y}_{LM}^{(1)}(\mathbf{r}) + \sqrt{L(L+1)} \frac{j_L(kr)}{kr} \mathbf{Y}_{LM}^{(-1)}(\mathbf{r}).$$

#### 4.3.b. Single-electron (reduced) multipole-transition matrix elements

##### Johnson's single-electron reduced multipole-transition matrix elements:

- **Transverse (velocity) gauge:** Johnson (2007, section 6.3) derives the **reduced multipole-transition matrix elements** for  $\mathbb{M} = (L, p)$  and  $q = \omega/c$  as:

$$\begin{aligned} \langle a \| \mathbb{T}^{(\mathbb{M}, \text{magnetic})}(\omega) \| b \rangle &= \langle -\kappa_a \| \mathbb{C}^{(L)} \| \kappa_b \rangle \int_0^\infty dr \frac{\kappa_a + \kappa_b}{L+1} j_L(qr) [P_a Q_b + Q_a P_b] \\ \langle a \| \mathbb{T}^{(\mathbb{M}, \text{electric})}(\omega) \| b \rangle &= \langle \kappa_a \| \mathbb{C}^{(L)} \| \kappa_b \rangle \int_0^\infty dr \left\{ -\frac{\kappa_a - \kappa_b}{L+1} \left[ j_L'(qr) + \frac{j_L(qr)}{qr} \right] [P_a Q_b + Q_a P_b] + L \frac{j_L(qr)}{qr} [P_a Q_b - Q_a P_b] \right\}. \end{aligned}$$

- **Length gauge:** While the reduced multipole-transition matrix elements are the same in length gauge, Johnson (2007, section 6.3) derives the reduced multipole-transition matrix elements in length gauge as:

$$\begin{aligned} \langle a \| \mathbb{T}^{(\mathbb{M}, \text{electric})}(\omega) \| b \rangle &= \langle \kappa_a \| \mathbb{C}^{(L)} \| \kappa_b \rangle \int_0^\infty dr \left\{ j_L(qr) [P_a P_b + Q_a Q_b] + j_{L+1}(qr) \left[ \frac{\kappa_a - \kappa_b}{L+1} [P_a Q_b + Q_a P_b] + [P_a Q_b - Q_a P_b] \right] \right\}. \end{aligned}$$

### Johnson's frequency-dependent multipole-moment operators:

- **Multipole-moment matrix elements:** Following Johnson (2007), the **reduced multipole-moment matrix elements** can be calculate for the reduced multipole-transition matrix elements

$$\langle a \parallel \mathfrak{q}^{(\mathbb{M})}(\omega) \parallel b \rangle = \frac{(2L+1)!!}{q^L} \langle a \parallel \mathfrak{t}^{(\mathbb{M})}(\omega) \parallel b \rangle.$$

### 4.3.c. Many-electron (reduced) multipole emission and absorption amplitudes

#### JAC's (standard) multipole amplitudes for photon absorption and emission:

- **Multipole (absorption) amplitude  $|\alpha_i \mathbb{J}_i\rangle + \hbar\omega(\mathbb{M}, \lambda) \rightarrow |\alpha_f \mathbb{J}_f\rangle$ :** For the absorption of a photon with energy  $\omega$ , well-defined multipolarity (angular momentum)  $L$  and helicity  $\lambda = \pm 1$ , we always use the notation

$$\langle \alpha_f \mathbb{J}_f \parallel \mathfrak{O}^{(\mathbb{M}, \text{absorption})}(\omega) \parallel \alpha_i \mathbb{J}_i \rangle.$$

- **Photon emission:** from an atom or ion  $|\alpha_i \mathbb{J}_i\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + \hbar\omega(\mathbb{M})$ .
- Care has to be taken about the many-electron multipole amplitudes which appear in many expressions. While we need formally to distinguish between the absorption  $a_L^\lambda$  and emission operators  $a_L^{\lambda+}$  operators, **all (one- and many-electron) multipole-transition matrix elements are always evaluated in absorption.**
- For any photon emission process, this would then require to 'interchange' the sequence *final-state* — *operator* — *initial-state* explicitly. This is a (very) high price however in such a large software project, since such an interchange affects the intuitive description of many processes in terms of initial-intermediate-final levels as well as still possibly other subsequent (final) states/levels.

## 4.4. Electromagnetic light pulses

### 4.4.a. High-intensity pulses

#### Matter in intense fields:

- The exposure of atoms and matter to high-intensity laser fields is known to give rise to **high-harmonic generation, above-threshold ionization and well as laser-induced dissociation and plasma formation**, and which are all fundamentally different from processes in the weak-field regime.
- These phenomena are often outside of what is described by (standard) atomic-structure methods since the external field, that acts upon the electrons in typical strong-field experiments, is then comparable to the Coulomb potential of the nucleus.
- **Pump-probe experiments:** In these experiments, two short pulses are typically required. While the pump pulse starts the reaction, the probe pulse investigates the state of the system after a defined time delay. These measurements is repeated several times with increasing time delays in order to finally obtain a **molecular movie**.
- If the level structure of the atoms is not affected by the (moderately intense) laser pulses, the time-dependent density matrix may allow to model the (time-dependent) level occupation and coherences of an atoms and, eventually, all the photoexcitation and ionization processes.
- **Two-color absorption in strong laser fields:** The simultaneous absorption of a single high-frequency photon in a strong optical laser field generally leads to equally spaced sideband peaks in the photoelectron spectra due to the (additional) absorption or emission of several laser photons. This (XUV + IR) two-color absorption has been first observed in high-order harmonic generation by overlapping high harmonics with the fundamental laser frequency. With a femtosecond XUV pulse from a FEL, the two-color absorption in strong laser fields can be explored without the contributions from neighboured harmonics (Gramajo *et al.*, 2017).
- **Two-color absorption in strong laser fields:** Two different interaction regimes can be distinguished for the interaction of atoms with a combined and overlapping XUV and IR pulses owing to their relative pulse duration: (i) If the XUV pulse duration  $\tau^{(\text{XUV})} \gtrsim \tau^{(\text{IR})}$  is longer than the IR pulse, the intensity of the sideband peaks in the photoelectron spectrum directly reflect the intensity of the IR laser; (ii) for  $\tau^{(\text{XUV})} \ll \tau^{(\text{IR})}$ , in constrast, the photoelectron spectrum acts as as *streak camera* and reflects the shape and duration of an IR laser pulse. Moreover, variations of the polarization state of the XUV and/or IR laser fields gives generally rise to **dichroic effects in the photoelectron spectrum** (Gramajo *et al.*, 2017).

### 4.4.b. Pulse shapes and optical cycles

#### Pulse shaping in physics and science:

- **Pulse shaping is frequently used in physics, science and technology.** Pulse shaping generally refers to the the process of changing the waveform of a (transmitted) pulse, for instance, in order to make a transmitted signal better suited for a particular communication channel. In telecommunication, for example, pulse shaping is essential in order to ensure that the signal fit to a given frequency band.
- In femtosecond physics and chemistry, pulse shaping describes a technology for generating nearly arbitrary pulses with user defined, ultrafast optical waveforms, and by controlling its phase, amplitude, and polarization. Here, the most widely applied technique is **Fourier transform pulse shaping**.

#### Frequently applied pulse shapes for time-dependent atomic computations:

- In pump- and/or pump-probe experiments, all (symmetric) light pulses are typically defined in terms of a **central frequency  $\omega$  and a (delay) time  $T_d$**  for the arrival of the pulse center with regard to same reference time  $t_0 = 0$ . Moreover, the shape function is always normalized to  $\max[f_s(t)] = \max[f(t)/f_o] = 1$  and, hence, the **intensity of the em field must be captured properly by the constant  $f_o$** .
- In JAC, the following (symmetric) shapes of the em pulses have been pre-defined: "sin<sup>2</sup>", "sin<sup>2</sup>: plateau", "sin<sup>2</sup>: cycles" and "gaussian" facilitate the description of experimentally specified pulses. Further pulses shapes might be added if the need arises from the side of experiment.
- It is important to ensure that the pulse-shape function  $f_s(t)$  (envelope) just contains an **integer number of optical cycles**. Therefore, one often starts with an  $\sin^2$  envelope. For example, an envelope function  $\sin^2 10T$  pulse refers to a pulse with 10 optical cycles.

### 4.4.c. Maximum pulse intensity

#### Vector potential:

- **Amplitude of the vector potential:** This amplitude need to be derived from the (maximum) intensity of a laser pulses that is typically given in  $\text{W}/\text{cm}^2$ . In general, this relation between intensity and amplitude is non-trivial since the relationship between intensity and the electric field depends also slightly on the polarization of the pulse. A simple relation only occurs for a linearly-polarized laser beam.

#### 4. Atomic interactions with the radiation field

- **Average intensity of a pulse:** If only the average intensity is given, the conversion factor depends also on the particular pulse-shape function. In practice, however, experimentalists rarely know the intensity accurately and, hence, a simplified normalization procedure can normally be used.
- **Vector potential of a single pulse:** The interaction of the electrons with the electric field of the incident pulse of frequency  $\omega$ , duration  $T$ , carrier envelope phase  $\phi$  and with linear polarization along  $\mathbf{e}_p$  is often described in terms of the vector potential

$$\mathbf{A}(t) = A_o \sin^2\left(\frac{\pi t}{T}\right) \sin(\omega t + \phi) \mathbf{e}_p, \quad A_o = \frac{1}{\omega} \sqrt{\frac{I}{I_o}},$$

and where  $I_o = 3.5 \cdot 10^{16} \text{ W/cm}^2$ .

#### Intensity and pulse geometry:

- The atomic unit of electric field is  $m_e^2 e^5 / \hbar^4 = 5.1422 \cdot 10^9 \text{ V/cm}$ . From this we have:  $E_0 [a.u.] = 5.34 \cdot 10^{-9} \sqrt{I [W/cm^2]}$ .

#### 4.4.d. Bichromatic laser fields

#### Bichromatic laser fields:

- **Bichromatic laser fields:** Of some special interests in light-atom interaction studies are fields with two components of commensurate frequencies; for instance, the fundamental component  $\omega$  and one of its low harmonics  $2\omega$  or  $3\omega$ . Moreover, these two components are in general out of phase by some angle  $\phi$ .
- **Linearly-polarized bichromatic fields:** For such a field, the electric field can be written in a quite general form as

$$E(t) = F(t) \left[ \cos(\omega t + \phi_1^{(\text{cep})}) + \eta \cos(2\omega t + \phi_2^{(\text{cep})}) \right]$$

with frequencies  $\omega$  of the fundamental radiation and  $2\omega$  due to the second harmonic. Here, the same envelope function  $F(t)$  is applied for both, the fundamental and the second harmonic, while the ratio of the corresponding amplitudes is specified by the real parameter  $\eta$  ( $\eta > 0$ ). The carrier-envelope phases (CEP) are here denoted by  $\phi_1^{(\text{cep})}$  and  $\phi_2^{(\text{cep})}$ , respectively.

## 5. Atomic amplitudes

### 5.1. In JAC implemented amplitudes

#### 5.1.a. Dipole amplitudes (MultipoleMoment)

##### Amplitude, notations & application:

- Formal quantum notation:  $\langle \alpha_f \mathbb{J}_f \parallel \mathbb{D} \parallel \alpha_i \mathbb{J}_i \rangle = \langle \alpha_f \mathbb{J}_f \parallel \sum_i \mathbf{r}_i \parallel \alpha_i \mathbb{J}_i \rangle$  for  $P_f \neq P_i$
- Note that the full dipole operator is often defined by:  $\mathbb{D}' = -|e| \mathbb{D}$ .
- Using JAC: Call `MultipoleMoment.dipoleAmplitude(level_f::Level, level_i::Level)` or `(level_f::Level, level_i::Level; display=true)`, if the value of the amplitude needs to be printed to screen.

#### 5.1.b. Electro-magnetic multipole transition amplitudes (MultipoleMoment, Radiative)

##### Amplitude, notations & application:

- Formal quantum notation:

$$\langle \alpha_f \mathbb{J}_f \parallel \mathbb{O}^{(\mathbb{M}, \text{emission})}(\omega) \parallel \alpha_i \mathbb{J}_i \rangle = \left\langle \alpha_i \mathbb{J}_i \left\| \sum_{k=1}^N \boldsymbol{\alpha}_k a_{k,L}^p \right\| \alpha_f \mathbb{J}_f \right\rangle^* \equiv \langle \alpha_i \mathbb{J}_i \parallel \mathbb{O}^{(\mathbb{M}, \text{absorption})}(\omega) \parallel \alpha_f \mathbb{J}_f \rangle^*$$

$$\langle \alpha_f \mathbb{J}_f \parallel \mathbb{T}^{(\mathbb{M}, \text{absorption})}(\omega) \parallel \alpha_i \mathbb{J}_i \rangle$$

## 5. Atomic amplitudes

which retains the **right order of the initial and final levels with regard to absorption or emission** and which can be obtained from `Jac.Radiative.amplitude()`; cf. section 4.3.c.

- Using JAC: Call `Radiative.amplitude("absorption", Mp::Multipole, gauge::EmGauge, omega::Float64, level_f::Level, level_i::Level)` or `("absorption", Mp::Multipole, gauge::EmGauge, omega::Float64, level_f::Level, level_i::Level; display=true)`, if the value of the amplitude needs to be printed to screen.
- Using JAC: Call `MultipoleMoment.transitionAmplitude(Mp::Multipole, gauge::EmGauge, omega::Float64, level_f::Level, level_i::Level)` or `(Mp::Multipole, gauge::EmGauge, omega::Float64, level_f::Level, level_i::Level; display=true)`, if the value of the amplitude needs to be printed to screen.

### Further remarks:

- A multipole  $\mathbb{M} \equiv (L, p) \equiv (L, \text{electric/magnetic}) = E1, M1, E2, \dots$  hereby contains all information about its multipolarity (angular momentum)  $L$  and type *magnetic* ( $p = 0$ ) or *electric* ( $p = 1$ ), respectively.
- In the module `Radiative`, the multipole-transition amplitudes can be calculated in one of the gauge = {Coulomb, Babushkin, Magnetic} as appropriate for electric- and magnetic-multipole transitions.
- **Johnson's multipole-transition operators:** Following Johnson (2007), the electron-photon interaction with the multipole fields ( $\mathbb{M}$ ) can be expressed in terms of a dimensionless **multipole-transition operator**  $\mathbb{T}_M^{(\mathbb{M})}(\omega) = \sum_j \mathbb{t}_M^{(\mathbb{M})}(\mathbf{r}_j; \omega)$ , an irreducible tensor operator, whose reduced (single-electron) matrix elements are given in Section 4.3.b. The amplitudes to these operators always describe the absorption of a photon with a given multipolarity.
- Johnson's (2007) multipole-transition amplitudes can be calculated in one of the gauge = {Velocity, Length, Magnetic} as appropriate for electric- and magnetic-multipole transitions.

### 5.1.c. Electro-magnetic multipole-moment amplitudes (MultipoleMoment)

#### Amplitude, notations & application:

- **Formal quantum notation:**

$$\langle \alpha_f \mathbb{J}_f \parallel \mathbb{Q}^{(\mathbb{M})}(\omega) \parallel \alpha_i \mathbb{J}_i \rangle$$



- Using JAC: Call `MultipoleMoment.amplitude(Mp::Multipole, gauge::EmGauge, omega::Float64, level_f::Level, level_i::Level)` or `(Mp::Multipole, gauge::EmGauge, omega::Float64, level_f::Level, level_i::Level; display=true)`, if the value of the amplitude needs to be printed to screen.

### Further remarks:

- **Johnson's multipole-moment operator:** Johnson (2007) also defines a **frequency-dependent (irreducible tensor) multipole-moment operator**

$$\mathbb{Q}_M^{(\mathbb{M})}(\omega) = \sum_j \mathbb{Q}_M^{(\mathbb{M})}(\mathbf{r}_j; \omega) \equiv \frac{(2L+1)!!}{(\omega/c)^L} \sum_j \mathbb{t}_M^{(\mathbb{M})}(\mathbf{r}_j; \omega)$$

- These amplitudes are calculated by a call to (Johnson's) multipole-transition amplitudes from above [cf. section 5.1.b] but by taking the frequency-dependence into account.
- Electric multipole-moment amplitudes can be calculated in one of the gauge = {**Velocity**, **Length**, **Magnetic**} as appropriate for electric- and magnetic-multipole transitions.

### 5.1.d. Momentum transfer amplitudes (FormFactor)

#### Amplitude, notations & application:

- **Formal quantum notation:**  $\langle \alpha_f \mathbb{J}_f \parallel \sum_j \exp(i \mathbf{q} \cdot \mathbf{r}_j) \parallel \alpha_i \mathbb{J}_i \rangle = \text{selection rules}$
- Using JAC: Call `FormFactor.amplitude(level_f::Level, level_i::Level, ...)` or `(level_f::Level, level_i::Level, ...; display=true)`, if the value of the amplitude needs to be printed to screen.

### Further remarks:

- The (single-electron) **momentum transfer amplitude** help describe inelastic scattering processs and is closely related also to the (so-called) **generalized oscillator strength**.

## 5. Atomic amplitudes

- The (one-electron) momentum-transfer operator  $\mathbb{T}^{(1)}(\mathbf{q}) = \sum_{j=1}^N \exp(i \mathbf{q} \cdot \mathbf{r}_j)$  can be utilized in order to express the standard form factor for an atom in a given sublevel  $|\alpha \mathbb{J} M\rangle$  in terms of an (many-electron) expectation value; cf. 6.1.d.

## 5.2. In JAC partly-implemented amplitudes

### 5.2.a. Parity non-conservation amplitudes (ParityNonConservation)

#### Amplitude, notations & application:

- **Formal quantum notation:**  $\langle \alpha_f \mathbb{J}_f || \mathbb{H}^{(\text{weak-charge})} || \alpha_i \mathbb{J}_i \rangle$
- **Using JAC:** Call `ParityNonConservation.weakChargeAmplitude(level_f::Level, level_i::Level, model::Nuclear.Model)` or `(level_f::Level, level_i::Level, model::Nuclear.Model; display=true)`, if the value of the amplitude needs to be printed to screen.

#### Further remarks:

- Parity non-conservation (PNC) effects arises in atoms mainly because of the exchange of  $Z^0$ -bosons between atomic electrons and the nucleus. Various observables due to this parity non-conserving interactions can be described in terms of a **nuclear-spin independent Hamiltonian**  $\mathbb{H}^{(\text{weak-charge})}$ , an effective single-electron operator, if non-relativistic nucleons can be assumed, cf. section 3.10.b.
- Especially, if the **proton and neutron densities coincide**,  $\rho_p = \rho_n = \rho$ , the general nuclear-spin independent Hamiltonian simplifies and can be described in terms of a single nuclear- weak charge  $Q_W$ ,

$$\mathbb{H}^{(\text{weak-charge})} = -\frac{G}{2\sqrt{2}} Q_W \sum_j \gamma_{5,j} \rho(r_j),$$

where  $G = 2.22 \cdot 10^{-14}$  a.u. is the Fermi constant,  $\rho(r)$  is the normalized nuclear density with  $\int dV \rho(\mathbf{r}) = 1$ , and  $Q_W \approx -N + Z (1 - 4 \sin^2 \theta_W)$  is the nuclear weak charge with  $\theta_W$  being the Weinberg angle.

### 5.2.b. Schiff-moment amplitudes (ParityNonConservation)

#### Amplitude, notations & application:

- **Formal quantum notation:**  $\langle \alpha_f \mathbb{J}_f \parallel \mathbb{H}^{(\text{Schiff-moment})} \parallel \alpha_i \mathbb{J}_i \rangle = \langle \alpha_f \mathbb{J}_f \parallel \mathbb{H}^{(\text{Schiff-moment})} \parallel \alpha_i \mathbb{J}_i \rangle \delta_{J_f, J_i} \quad P_f \neq P_i$
- Using JAC: Call `ParityNonConservation.schiffMomentAmplitude(level_f::Level, level_i::Level, model::Nuclear.Model)` or `(level_f::Level, level_i::Level, model::Nuclear.Model; display=true)`, if the value of the amplitude needs to be printed to screen.

#### Further remarks:

- **Nuclear electric-dipole (Schiff) moments:** If  $\mathbf{S}$  is the P-odd and T-odd nuclear Schiff moment, the associated P-odd and T-odd Hamiltonian is given by (Dzuba *et al.*, 2009)

$$\mathbb{H}^{(\text{Schiff-moment})} = -e \phi^{(\text{Schiff-moment})} = -4\pi e \mathbf{S} \cdot \nabla \delta(r) = -\frac{3 S^{(\text{SM})} (\mathbf{I}/I) \cdot \mathbf{r}}{R} \rho(r)$$

where  $S^{(\text{SM})} (\mathbf{I}/I)$  is the Schiff-moment vector, i.e. the product of the coupling constant  $S^{(\text{SM})}$  and the nuclear spin  $\mathbf{I}$ , and  $R = \int_0^\infty dr r^4 \rho(r)$ . Moreover,  $\phi^{(\text{Schiff-moment})}$  is the electrostatic potential of the nucleus corresponding to a P-odd and T-odd charge distribution; cf. section 3.10.c.

## 5.3. Further amplitudes, not yet considered in JAC

### 5.3.a. Anapole-moment amplitudes (ParityNonConservation)

#### Amplitude, notations & application:

➤ Formal quantum notation:  $\langle \alpha_f \mathbb{J}_f \parallel \mathbb{H}^{(\text{anapole-moment})} \parallel \alpha_i \mathbb{J}_i \rangle$

#### Further remarks:

- A **P-odd, T-even anapole moment** of the nucleus was first introduced by Zel'dovich and arises due to the presence of a parity-violating weak interaction between nucleons, cf. section 3.10.b. The anapole moment is directed along the nuclear spin  $\mathbf{I}$ , gives typically the **dominant nuclear-spin dependent electron-nucleus interaction** and manifests itself in terms of a slightly modified electro-magnetic interaction between the nucleus and the atomic electrons.
- **Hamiltonian for the interaction of the electrons with the nuclear anapole moment:** This Hamiltonian has the form (Dzuba *et al.*, 2009)

$$\mathbb{H}^{(\text{anapole-moment})} = \frac{G}{\sqrt{2}} (I + 1/2) (-1)^{I+1/2-\ell_N} \kappa_{\text{AM}} \sum_j \frac{\mathbf{I} \cdot \boldsymbol{\alpha}_j}{I(I+1)} \rho(r_j),$$

where  $G = 2.22 \cdot 10^{-14}$  a.u. is the Fermi constant,  $I$  is the nuclear spin,  $\ell_N$  the orbital angular momentum of the outermost nucleon and  $\kappa_{\text{AM}}$  is a dimensionless coupling constants for ...

### 5.3.b. Scalar-pseudo-scalar amplitudes

#### Amplitude, notations & application:

➤ Formal quantum notation:  $\langle \alpha_f \mathbb{J}_f \parallel \mathbb{H}^{(\text{scalar-pseudo-scalar})} \parallel \alpha_i \mathbb{J}_i \rangle$

### Further remarks:

- **Hamiltonian for the P-odd, T-odd scalar-pseudo-scalar interaction:** In coordinate representation, this Hamiltonian can be written as (Dzuba *et al.*, 2009)

$$\mathbb{H}^{(\text{scalar-pseudo-scalar})} = -\frac{G}{\sqrt{2}} \frac{1}{2m_p c} C_s \sum_j \left[ \gamma_0 \frac{d\rho(r)}{dr} \right]_j$$

where  $G$  is the Fermi constant,  $m_p$  is the nucleon mass,  $C_s$  is a (isotope and) nuclear-state dependent constant, and where  $\left[ \gamma_0 \frac{d\rho(r)}{dr} \right]$  is a scalar term that arises from the scalar product of the  $\nabla$ -vector with an particularly averaged spin-vector of the nucleus.

### 5.3.c. Tensor-pseudo-tensor amplitudes

#### Amplitude, notations & application:

- **Formal quantum notation:**  $\langle \alpha_f \mathbb{J}_f \parallel \mathbb{H}^{(\text{tensor-pseudo-tensor})} \parallel \alpha_i \mathbb{J}_i \rangle$

### Further remarks:

- **Hamiltonian for the P-odd, T-odd tensor-pseudo-tensor interaction:** In coordinate representation, this Hamiltonian can be written as (Dzuba *et al.*, 2009)

$$\mathbb{H}^{(\text{tensor-pseudo-tensor})} = i \sqrt{2} G C_t \sum_j [\gamma \rho(r)]_j$$

where  $G$  is the Fermi constant,  $C_t$  is a (isotope and) nuclear-state dependent constant, and where  $[\gamma \rho(r)]$  is a scalar term that arises from the scalar product of the  $\gamma$ -vector with an averaged spin-vector of the nucleus.

- **Total P-odd, T-odd electron-nucleus interaction Hamiltonian:**  $\mathbb{H}^{(\text{P-odd, T-odd electron-nucleus})} = \mathbb{H}^{(\text{tensor-pseudo-tensor})} + \mathbb{H}^{(\text{tensor-pseudo-tensor})}$

### 5.3.d. Nuclear magnetic-quadrupole-moment amplitudes due to internal $B$ -field

#### Amplitude, notations & application:

➤ **Formal quantum notation:**  $\langle \alpha_f \mathbb{J}_f \parallel \mathbb{H}^{(\text{nuclear-MQM})} \parallel \alpha_i \mathbb{J}_i \rangle$

#### Further remarks:

- This interaction (Hamiltonian) arise from the interaction of the electron EDM  $d_e$  with the electro-magnetic fields of the nucleus as seen by the electron. If the treatment is restricted to second order, only the interaction of the electron EDM with the magnetic field  $\mathbf{B}$ , which is created by the nuclear magnetic-dipole moment, need to be taken into account.
- **Nuclear-MQM Hamiltonian:** In coordinate representation, this Hamiltonian can be written as (Dzuba *et al.*, 2009)

$$\mathbb{H}^{(\text{nuclear-MQM})} = -i d_e \boldsymbol{\gamma} \cdot \mathbf{B} = -i d_e \sum_j \boldsymbol{\gamma}_j \cdot \left[ \nabla_j \times \frac{\mathbf{M} \times \mathbf{r}_j}{r_{>,j}^3} + \dots \right]$$

## 5.4. Composed many-electron amplitudes, not yet considered in JAC

### 5.4.a. Parity-violating (non-diagonal, second-order) amplitudes

#### Amplitude, notations & application:

- **Parity violation:** The (effective) parity-violating interaction is caused by the exchange of  $Z_o$  bosons between the atomic electrons and the nucleus. The parity-violating interaction typically results in non-vanishing off-diagonal electric-dipole matrix element between two atomic states of the same parity; cf. section 3.10.b.

### 5.4.b. Charge-parity-violating (diagonal, second-order) amplitudes

#### Amplitude, notations & application:

- **Charge-parity violation:** An intrinsic electron electric dipole moment  $d_e$  leads to CP-violation in atoms and also induces an **atomic electric dipole moment**. cf. section 3.10.d.

### 5.4.c. Electric-dipole moment enhancement factor

#### Amplitude, notations & application:

- An intrinsic electric-dipole moment (EDM) of the electron can enhanced the EDM of the atom. In heavy atoms, in particular, an induced atomic EDM can be strongly enhanced, when compared to the electron EDM itself.
- A typical goal is to find the atomic EDM  $d_{\text{atom}} = |\mathbf{d}_{\text{atom}}| \equiv d_{\text{atom}} |\mathbf{F}/F|$ , where  $\mathbf{F} = \mathbf{I} + \mathbf{J}$  is the total angular momentum of an atomic (hyperfine) level.
- The value of the atomic EDM, compared to the electron EDM, is expressed by means of an enhancement factor  $K = d_{\text{atom}}/d_e$ , and which increases with nuclear charge  $Z$  faster than  $Z^3$ .





## 6. Atomic properties

### 6.1. In JAC implemented level properties

#### 6.1.a. Transition probabilities for a single multiplet (Einstein)

##### Properties, notations & application:

- **Photon emission:** from an atom or ion  $A^* \longrightarrow A^{(*)} + \hbar\omega$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + \hbar\omega$
- Using JAC: Perform an `Atomic.Computation(..., properties=[EinsteinX, ...], configs=[...], einsteinSettings=Einstein.Settings(...), ...)` or call directly functions from the module `Einstein`.
- In the JAC program, the transition probabilities, oscillator strength (in emission) and radiative lifetimes are calculated and tabulated by default for all selected transitions from a given single multiplet (i.e. list of configurations). Apart from the multipoles, the user can also specify an overall shift in the photon energies as well as a minimum and maximum transition energy, for which transitions are taken into account.
- In the `Einstein` module, the emission (transition) amplitudes  $\langle \alpha_f \mathbb{J}_f || \mathbb{O}^{(\mathbb{M}, \text{emission})} || \alpha_i \mathbb{J}_i \rangle$  or absorption amplitudes  $\langle \alpha_f \mathbb{J}_f || \mathbb{O}^{(\mathbb{M}, \text{absorption})} || \alpha_i \mathbb{J}_i \rangle$  can be obtained directly by call of the function `Jac.Einstein.amplitude()`.

##### Further remarks:

- Einstein  $A$  and  $B$  coefficients and oscillator strength for levels  $|\alpha \mathbb{J}\rangle$  from a single multiplet are frequently needed for various applications and in different contexts.

## 6. Atomic properties

- Although these coefficients are *no* (single) level property, the **Einstein** module still supports the computation of these coefficients for **levels from a single multiplet**, and which is generated by a single CSF basis.
- Although the **Einstein** module does not allow to include relaxation effects, this module helps obtain either a quick estimate of the Einstein coefficients between levels from just a few configurations or to deal with cascade computations.
- **JAC's standard multipole amplitudes for photon emission**: Although **all (one- and many-electron) matrix elements are always evaluated in absorption** within the JAC toolbox, we keep the intuitive description of the matrix elements as *final-state – operator – initial-state*, cf. section 4.3.c.

### 6.1.b. Hyperfine parameters and hyperfine representations (Hfs)

#### Properties, notations & application:

- **Hyperfine splitting** of an atomic level into hyperfine (sub-) levels:  $|\alpha\mathbb{J}\rangle \longrightarrow |\alpha(J)\mathbb{F}\rangle, \quad F = |I - J|, \dots, I + J - 1, I + J$ .
- **Formal quantum notation**:  $|\alpha\mathbb{J}\rangle \longrightarrow |\alpha\mathbb{F}\rangle \equiv |(I, \alpha'\mathbb{J})\mathbb{F}\rangle \equiv |\alpha(I\mathbb{J})\mathbb{F}\rangle \equiv |\alpha(IJP)\mathbb{F}\rangle$  or  $|\alpha\mathbb{J}M_J\rangle \longrightarrow |\alpha\mathbb{F}M\rangle \equiv |(I, \alpha'\mathbb{J})\mathbb{F}M\rangle \equiv |\alpha(I\mathbb{J})\mathbb{F}M\rangle \equiv |\alpha(IJP)\mathbb{F}M\rangle$ .
- Using JAC: Perform an `Atomic.Computation(.., properties=[HFS, ..], configs=[..], hfsSettings=Hfs.Settings(..), ..)` or call directly functions from the module `Hfs`. The Boolean values `calcT1` and `calcT2` in `Hfs.Settings` determine which parts of the hyperfine Hamiltonian are taken into account into the computations.
- In the JAC program, the hyperfine *A* and *B* coefficients as well as the (electric) hyperfine amplitudes  $\langle \alpha\mathbb{J} || \mathbb{T}^{(1)} || \alpha\mathbb{J} \rangle, \langle \alpha\mathbb{J} || \mathbb{T}^{(2)} || \alpha\mathbb{J} \rangle$  are calculated and tabulated by default for all selected levels together with the energy shifts  $\Delta E_F$  of the hyperfine levels with regard to the (electronic) level energy  $E_{\alpha\mathbb{J}}$ .
- In JAC, the hyperfine amplitudes  $\langle \beta_r\mathbb{J}_r || \mathbb{T}^{(1)} || \beta_s\mathbb{J}_s \rangle$  and  $\langle \beta_r\mathbb{J}_r || \mathbb{T}^{(2)} || \beta_s\mathbb{J}_s \rangle$  can be obtained from the function `Jac.Hfs.amplitude()`.
- In JAC, moreover, an explicit representation of a hyperfine multiplet is obtained by diagonalizing the Hamiltonian  $\mathbb{H} = \mathbb{H}^{(\text{DCB})} + \mathbb{H}^{(\text{hfs})}$  within the atomic hyperfine *IJF*-coupled basis if `calcIJFexpansion=true` is set in `Hfs.Settings`.
- In JAC, all the (hyperfine) level energies  $E_{\alpha\mathbb{F}}$  and the energies relative to the lowest (hyperfine) level are tabulated if `printDeltaEF=true` in `Hfs.Settings`.

- In JAC, all (diagonal and nondiagonal) hyperfine amplitudes are calculated and tabulated explicitly if `calcNondiagonal=true` in `Hfs.Settings`.

### Further remarks:

- For a nucleus with *non-zero* spin  $I > 0$ , a hyperfine splitting of (all) atomic levels generally occurs since each atomic electron also interacts with the electric and magnetic (multipole) fields of the nucleus. This is sometimes described also as the interaction of the magnetic moments of the electrons and nucleus.
- In atomic physics, this interaction is better known as ‘**hyperfine interaction**’ whose two dominant contributions arise from the nuclear magnetic-dipole field  $\mathbf{A} = \frac{\boldsymbol{\mu} \times \mathbf{r}}{r^3}$  and the electric-quadrupole field  $\Phi(r) = \sum_{ij} \frac{r_i r_j}{2r^5} Q_{ij}$ , respectively. In these expressions,  $\boldsymbol{\mu}$  is the nuclear magnetic moment operator and  $Q_{ij}$ ,  $i, j = 1..3$  are the Cartesian components of the nuclear quadrupole operator.
- Apart from the hyperfine parameters  $A(\alpha \mathbb{J})$  and  $B(\alpha \mathbb{J})$ , we also provide in JAC a **representation of the atomic hyperfine levels** that can be utilized, for instance, to compute **hyperfine quenched transitions probabilities**.

### Hyperfine interaction Hamiltonian:

- **Relativistic hyperfine interaction Hamiltonian:** For many-electron systems, this Hamiltonian can be written as

$$H^{(\text{hfs})} = \sum_K \mathbb{W}^{(K)} \cdot \mathbb{T}^{(K)}$$

where  $\mathbb{W}^{(K)}$  and  $\mathbb{T}^{(K)}$  represent the **spherical tensor operators of rank  $K$  that occur in the nucleonic and electronic sectors**, respectively.

- **Reduced nuclear matrix elements:** These matrix elements are determined geometrically by

$$\langle I \| \mathbb{W}^{(1)} \| I \rangle = \mu_I \sqrt{\frac{I+1}{I}}, \quad \langle I \| \mathbb{W}^{(2)} \| I \rangle = \frac{Q}{2} \sqrt{\frac{(I+1)(2I+3)}{I(2I-1)}},$$

while, in contrast, the corresponding **electronic amplitudes require detailed atomic structure calculations**.

- Values of nuclear magnetic dipole and electric quadrupole moments can be found in the compilation by Stone (2005).

## 6. Atomic properties

- **Electronic tensor operators:** For an  $N$ -electron atom or ion, these tensor operators are given by (Andersson and Jönsson, 2008)

$$\mathbb{T}^{(1)} = \sum_j^N \mathbb{t}^{(1)}(j) = \sum_j^N -i\sqrt{2}\alpha \frac{(\boldsymbol{\alpha}_j \mathbb{C}^{(1)}(j))^{(1)}}{r_j^2}$$

interaction with the magnetic – dipole field

$$\mathbb{T}^{(2)} = \sum_j^N \mathbb{t}^{(2)}(j) = \sum_j^N -\frac{\mathbb{C}^{(2)}(j)}{r_j^3}$$

interaction with the electric – quadrupole field,

and where  $\alpha$  is the fine-structure constant,  $\boldsymbol{\alpha}_j$  the vector of Dirac matrices of the  $j$ -th electron, and the  $\mathbb{C}^{(K)}$  are (normalized) spherical tensors as before.

### Atomic (hyperfine) $IJF$ -coupled basis:

- This (geometrically fixed) basis is obtained by the standard coupling of the nuclear states  $|IM_I\rangle$  and the ASF  $|\alpha \mathbb{J}M_J\rangle$

$$|\alpha(I, J) \mathbb{F}M\rangle \equiv |(I, \alpha' \mathbb{J}) \mathbb{F}M\rangle = \sum_{M_I M_J} |IM_I\rangle |\alpha' \mathbb{J}M_J\rangle \langle IM_I, JM_J | FM\rangle.$$

- The atomic (hyperfine)  $IJF$ -coupled basis, also called the  $IJF$ -coupled ASF basis, is applied internally to represent all hyperfine levels, see below.

### Atomic hyperfine amplitudes and levels:

- **Atomic hyperfine levels:** For the combined system ‘nucleus+electrons’, the atomic hyperfine states (magnetic sub-levels) can be formed as linear combination

$$|\alpha \mathbb{F}M\rangle = \sum_{r=1} \tilde{c}_r(\alpha) |(I, \beta_r \mathbb{J}_r) \mathbb{F}M\rangle$$

of hyperfine (basis) states  $|(I, \beta_r \mathbb{J}_r) \mathbb{F}M\rangle$  of the same total angular momentum  $F$  and (the same electronic) parity  $P$ , and where  $\{\tilde{c}_r(\alpha)\}$  denotes the representation of these states in the **atomic hyperfine-coupled basis** (Johnson, 2010).

- Analogue as for the ASF  $|\alpha\mathbb{J}\rangle$  in the standard MCDF ansatz, the representation  $\{\tilde{c}_r(\alpha)\}$  of an atomic hyperfine state is obtained by diagonalizing the Hamiltonian  $\mathbb{H} = \mathbb{H}^{(\text{DFB})} + \mathbb{H}^{(\text{hfs})}$  of the combined system ‘nucleus+electrons’ in either the *IJF*-coupled CSF or ASF basis.
- In the *IJF*-coupled ASF basis, we can make use of  $\mathbb{H}^{(\text{DFB})} |(I, \beta_r \mathbb{J}_r) \mathbb{F} M\rangle = E(\beta_r \mathbb{J}_r) |(I, \beta_r \mathbb{J}_r) \mathbb{F} M\rangle \equiv E_{\beta_r \mathbb{J}_r} |(I, \beta_r \mathbb{J}_r) \mathbb{F} M\rangle$ .
- For the hyperfine part of the Hamiltonian, the (reduced) matrix elements can be written after some standard angular momentum algebra as

$$\begin{aligned} \langle (I, \beta_r \mathbb{J}_r) \mathbb{F} M | H^{(\text{hfs})} | (I, \beta_s \mathbb{J}_s) \mathbb{F}' M' \rangle &= \delta_{MM'} \langle (I, \beta_r \mathbb{J}_r) \mathbb{F} | H^{(\text{hfs})} | (I, \beta_s \mathbb{J}_s) \mathbb{F}' \rangle \\ &= \delta_{MM'} \delta_{\mathbb{F}\mathbb{F}'} (-1)^{I+J_r+F} \sum_K \begin{Bmatrix} I & J_r & F \\ J_s & I & K \end{Bmatrix} \langle \beta_r \mathbb{J}_r || \mathbb{T}^{(K)} || \beta_s \mathbb{J}_s \rangle \langle I || \mathbb{W}^{(K)} || I \rangle, \end{aligned}$$

with  $\delta_{\mathbb{F}\mathbb{F}'} = \delta_{PP'} \delta_{FF'}$ , and if all nuclear excitations are ignored right from the beginning.

- **Hyperfine amplitudes:** These amplitudes refer to the (reduced) electronic matrix elements  $\langle \beta_r \mathbb{J}_r || \mathbb{T}^{(K)} || \beta_s \mathbb{J}_s \rangle$ . In the JAC program, we compute the **hyperfine amplitudes for both, the magnetic-dipole and the electric-quadrupole operators** of the electron-nucleus interaction.
- These hyperfine amplitudes are utilized to set-up and diagonalize the Hamiltonian matrix within the **atomic *IJF*-coupled basis** for  $\mathbb{H} = \mathbb{H}^{(\text{DCB})} + \mathbb{H}^{(\text{hfs})}$  in order to determine the representation  $\{\tilde{c}_r(\alpha)\}$  of the atomic hyperfine states.

## Hyperfine energies and interaction constants:

- **Hyperfine interaction constants:** Usually the hyperfine splitting is considered independently for each atomic level  $|\alpha\mathbb{J}\rangle$  and without the need to specify all the hyperfine level energies  $E_{\alpha\mathbb{F}}$  explicitly. Instead, the energy splitting of an atomic level  $|\alpha\mathbb{J}\rangle$  into hyperfine levels  $|\alpha\mathbb{F}\rangle$  can be expressed conveniently in terms of the (hyperfine interaction) constants

$$A(\alpha \mathbb{J}) = \frac{\mu_I}{I} \frac{1}{\sqrt{J(J+1)}} \langle \alpha \mathbb{J} || \mathbb{T}^{(1)} || \alpha \mathbb{J} \rangle, \quad B(\alpha \mathbb{J}) = 2Q \sqrt{\frac{J(2J-1)}{(J+1)(2J+3)}} \langle \alpha \mathbb{J} || \mathbb{T}^{(2)} || \alpha \mathbb{J} \rangle$$

- **Hyperfine energy splitting:** With these constants, the hyperfine energy shifts with regard to the electronic level energy  $E(\alpha\mathbb{J})$  is given by

$$\Delta E_F = \frac{A(\alpha \mathbb{J}) C}{2} + B(\alpha \mathbb{J}) \frac{3/4 C(C+1) - I(I+1) J(J+1)}{2I(2I-1) J(2J-1)}, \quad C = F(F+1) - J(J+1) - I(I+1).$$

### 6.1.c. Isotope-shift parameters (IsotopeShift)

#### Properties, notations & application:

- **Isotope shift** of an atomic level for two isotopes  $A, A'$  with nuclear masses  $M < M'$ :  $E(\alpha\mathbb{J}; A) \longrightarrow E(\alpha\mathbb{J}; A')$ .
- **Formal quantum notation:**  $\Delta E(\alpha\mathbb{J}; A, A') / h = \frac{M' - M}{M M'} K^{(\text{MS})} + F \delta \langle r^2 \rangle$ .
- Using JAC: Perform an `Atomic.Computation(.., properties=[Isotope, ..], configs=[..], isotopeSettings=IsotopeShift.Settings(..), ..)` with `calcNMS=true`, `calcNMS=true` or `calcF=true` and/or call directly functions from the module `IsotopeShift` in order to include the various contribution of the recoil Hamiltonian.
- In JAC, we calculate and tabulate by default the (relativistic) mass-shift parameters  $K^{(\text{MS})}$ ,  $K^{(\text{NMS})}$ ,  $K^{(\text{SMS})}$  and the field-shift parameter  $F$  for all selected levels.

#### Further remarks

- Owing to the conservation of the total momentum, the (total non-relativistic) mass shift is a sum of two parts, the (so-called) **normal mass shift (NMS)** and the **specific mass shift (SMS)**:

$$\mathbb{H}^{(\text{MS}, \text{nr})} = \mathbb{H}^{(\text{NMS}, \text{nr})} + \mathbb{H}^{(\text{SMS}, \text{nr})} = \frac{1}{2M} \sum_j \mathbf{p}_j^2 + \frac{1}{2M} \sum_{i \neq j} \mathbf{p}_i \cdot \mathbf{p}_j$$

While the normal mass shift operator is obviously a one-particle operator, the specific mass shift is a two-particle operator. The sum of these two parts is known also as the **total recoil operator within the nonrelativistic theory**.

## Relativistic recoil (Hamilton) operator:

- **Recoil Hamiltonian:** Within the lowest-order relativistic approximation ( $\sim v^2/c^2$ ) and in first order of  $m/M$  for a nucleus with mass  $M$ , the recoil corrections are given by means of the (recoil) Hamiltonian (Tupitsyn *et al.*, 2003):

$$\mathbb{H}^{(\text{recoil})}(M) = \frac{1}{2M} \sum_{ij} \left[ \mathbf{p}_i \cdot \mathbf{p}_j - \frac{\alpha Z}{r_i} \left( \boldsymbol{\alpha}_i + \frac{(\boldsymbol{\alpha}_i \cdot \mathbf{r}_i) \mathbf{r}_i}{r_i^2} \right) \cdot \mathbf{p}_j \right]$$

- The expectation value of  $\mathbb{H}^{(\text{recoil})}(M)$  with regard to the relativistic wave function from  $\mathbb{H}^{(\text{DCB})}$  gives the **recoil correction to the energy of the atomic energy  $|\alpha\mathbb{J}\rangle$**  in first order of  $m/M$ .
- **Decomposition of the relativistic recoil Hamiltonian:** Similar as in the non-relativistic theory, the (relativistic) recoil operator above can be written as (Tupitsyn *et al.*, 2003; Gaidamauskas *et al.*, 2011)

$$\begin{aligned} \mathbb{H}^{(\text{recoil})}(M) &= \mathbb{H}^{(\text{NMS, relativistic})} + \mathbb{H}^{(\text{SMS, relativistic})} \\ &= \frac{1}{2M} \sum_i \left[ \mathbf{p}_i^2 - \frac{\alpha Z}{r_i} \left( \boldsymbol{\alpha}_i + \frac{(\boldsymbol{\alpha}_i \cdot \mathbf{r}_i) \mathbf{r}_i}{r_i^2} \right) \cdot \mathbf{p}_i \right] + \frac{1}{2M} \sum_{i \neq j} \left[ \mathbf{p}_i \cdot \mathbf{p}_j - \frac{\alpha Z}{r_i} \left( \boldsymbol{\alpha}_i + \frac{(\boldsymbol{\alpha}_i \cdot \mathbf{r}_i) \mathbf{r}_i}{r_i^2} \right) \cdot \mathbf{p}_j \right] \\ &= \frac{1}{2M} \sum_i \left[ \mathbf{p}_i^2 - \frac{\alpha Z}{r_i} \boldsymbol{\alpha}_i \cdot \mathbf{p}_i - \frac{\alpha Z}{r_i} (\boldsymbol{\alpha}_i \cdot \mathbb{C}_i^{(1)}) \mathbb{C}_i^{(1)} \cdot \mathbf{p}_i \right] + \frac{1}{2M} \sum_{i \neq j} \left[ \mathbf{p}_i \cdot \mathbf{p}_j - \frac{\alpha Z}{r_i} \boldsymbol{\alpha}_i \cdot \mathbf{p}_j - \frac{\alpha Z}{r_i} (\boldsymbol{\alpha}_i \cdot \mathbb{C}_i^{(1)}) \mathbb{C}_i^{(1)} \cdot \mathbf{p}_j \right] \\ &= \mathbb{H}^{(\text{NMS})} + \mathbb{H}^{(\text{SMS, A})} + \mathbb{H}^{(\text{SMS, B})} + \mathbb{H}^{(\text{SMS, C})} \end{aligned}$$

- The relativistic correction to the recoil effect is typically strongly overestimated if the nonrelativistic recoil operator is applied just together with relativistic wave functions.
- The full relativistic theory of the nuclear recoil effect can be formulated only within the framework of quantum electrodynamics. Such a theory was first formulated by Shabaev (1985), where the complete  $\alpha Z$ -dependence of the recoil corrections to an energy level  $|\alpha\mathbb{J}\rangle$  was derived in first order of  $m/M$ .

### Reduced one-electron matrix elements of the recoil (Hamilton) operator:

- **Recoil Hamiltonian:** While the normal mass-shift Hamiltonian is a one-particle operator, the specific mass-shift is a (symmetric) two-particle operator

$$\mathbb{H}^{(\text{NMS})} = \sum_j \mathbb{h}^{(\text{NMS})}(\mathbf{r}_j), \quad \mathbb{H}^{(\text{SMS}, k)} = \sum_{i \neq j} \mathbb{h}^{(\text{SMS}, k)}(\mathbf{r}_i, \mathbf{r}_j) = \sum_{i \neq j} g(r_i, k) g(r_j, k) \left( \mathbb{T}_i^{(1)} \cdot \mathbb{T}_j^{(1)} \right), \quad k = A, B, C$$

The particular tensorial structure of the two-particle specific mass-shift Hamiltonian ensures that the reduced (one-electron) matrix elements can be expressed in terms of the **first-rank effective interaction strength**  $X^{(1; \text{SMS}, k)}$  that are specified below for  $k = A, B, C$ .

- **Reduced one-electron matrix element for the normal mass-shift:** With the definition of the angular coefficients in GRASP, this one-electron matrix element is given with  $P' = \frac{\partial P}{\partial r}$  (Gaidamauskas *et al.*, 2011)

$$\begin{aligned} \langle a \| \mathbb{h}^{(\text{NMS})} \| b \rangle &= \delta_{\kappa_a, \kappa_b} \frac{1}{2M} \int_0^\infty dr \left( P'_a P'_b + Q'_a Q'_b + \frac{\ell_b (\ell_b + 1) P_a P_b + (2j_b - 1) 2j_b Q_a Q_b}{r^2} \right. \\ &\quad \left. - (2\alpha Z) \frac{Q_a P'_b + Q_b P'_a}{r} - (\alpha Z) \left( \frac{\kappa_b - 1}{r^2} \right) (Q_a P_b + Q_b P_a) \right) \end{aligned}$$

### Parametrization of the isotope shift:

- **Mass shift of two isotopes:** Within lowest order of  $m/M$ , the **(isotope) mass shift** for an atomic level  $|\alpha \mathbb{J}\rangle$  is determined by the difference of the expectation values of  $\mathbb{H}^{(\text{recoil})}(M)$  for two different isotopes:

$$\Delta E^{(\text{MS})}(M, M') = \langle \alpha \mathbb{J} | \mathbb{H}^{(\text{recoil})}(M) - \mathbb{H}^{(\text{recoil})}(M') | \alpha \mathbb{J} \rangle = \frac{M' - M}{M M'} K^{(\text{MS})}$$

$$\frac{K^{(\text{MS})}}{M} = \langle \alpha \mathbb{J} | \mathbb{H}^{(\text{recoil})}(M) | \alpha \mathbb{J} \rangle = \langle \alpha \mathbb{J} \| \mathbb{H}^{(\text{recoil})}(M) \| \alpha \mathbb{J} \rangle,$$

and where the Wigner-Eckardt theorem is used here in a special form that is consistent with GRASP and that is utilized for the implementation of corresponding angular coefficients.



- The conversion factor between the mass-shift parameters  $K^{(\text{MS})}$  in the frequently applied units is:

$$K^{(\text{MS})}/[\text{GHz u}] = 3609.4824 K^{(\text{MS})}/[\text{a.u.}].$$

- Field shift of two isotopes: is caused by their different nuclear charge distributions and can be parametrized by:

$$\Delta E^{(\text{FS})}(M, M') = - \left\langle \alpha \mathbb{J} \left| \sum_{j=1}^N (\mathbb{V}^{(\text{nuc})}(r_j; R') - \mathbb{V}^{(\text{nuc})}(r_j; R)) \right| \alpha \mathbb{J} \right\rangle = \left\langle \alpha \mathbb{J} \left\| \sum_{j=1}^N \delta \mathbb{V}^{(\text{nuc})}(r_j; R', R) \right\| \right\rangle = -F \delta \langle r^2 \rangle$$

$$F = \frac{\left\langle \alpha \mathbb{J} \left| \sum_j (\mathbb{V}^{(\text{nuc})}(r_j; R') - \mathbb{V}^{(\text{nuc})}(r_j; R)) \right| \alpha \mathbb{J} \right\rangle}{R' - R}$$

where  $\delta \langle r^2 \rangle = R' - R$  is the difference of the *root-mean-square* (rms) charge radii of the two isotopes with masses  $M', M$ .

#### 6.1.d. Atomic form factors (FormFactor)

##### Properties, notations & application:

- Form factor  $F(q; \alpha \mathbb{J})$  of an atom in level  $|\alpha \mathbb{J}\rangle$  with an (assumed) spherical-symmetric charge distribution:.
- Formal quantum notation:  $F^{(\text{standard})}(q; \alpha \mathbb{J}), F^{(\text{modified})}(q; \alpha \mathbb{J})$ .
- Using JAC: Perform an `Atomic.Computation(..., properties=[FormF, ...], configs=[...], formSettings=FormFactor.Settings(...), ...)` or call directly functions from the module `FormFactor`.
- In JAC, the standard and modified form factors are calculated and tabulated by default for all selected levels in `FormFactor.Settings` and for all specified  $q$ -values [in a.u.]; the default values of the momentum transfer are  $q = 0.1, 1.0, 10.0$  a.u.
- Atomic form factor: In atomic physics, the form factors are generally taken as a fast approximation to the (scattering) amplitude in forward direction at  $\vartheta \simeq 0$  if an incident quantum wave is scattered by an isolated atom with spherical-symmetric charge distribution.
- In practice, of course, the definition of the atomic form factor depends on both, the particular scattering (elastic or inelastic) as well as the type of incident radiation, such as photons, electrons or neutrons. As usual, we here restrict ourselves to the atomic form factor for the elastic (x-ray) scattering of photons.

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- **Momentum transfer:** Most easily, the atomic form factor can be written as function of the momentum transfer  $q = |\mathbf{q}|$ , that is closely related also to the scattering angle:  $\hbar q = 2k \sin(\vartheta/2) = \frac{2}{\lambda} \sin(\vartheta/2)$ .
- **Momentum-transfer variable:** Instead of the momentum transfer  $q$ , some tabulations of x-ray scattering amplitude make use also of the (so-called) momentum-transfer variable  $x = [\sin(\vartheta/2)] / \lambda [\text{\AA}]$  with  $\lambda [\text{\AA}] = 12.398520/E [\text{keV}]$ .
- **Application of form factors:** The elastic scattering of photons by atoms, ions or molecules has been found important in many fields of physics, such as crystallography, plasma physics and astrophysics. For high-photon energies well above the K-shell threshold energy of the atom or ion, both the standard form factor (FF) or the modified form factor (MFF) approximations are widely applied and were found moderately successful in estimating the scattering cross sections. However, **both form factor approximations typically fail for small photon energies.**

### Definition of atomic form factors:

- **Standard atomic form factor:** For an atom in level  $|\alpha \mathbb{J}\rangle$  with a spherically-symmetric charge distribution  $\rho(r)$ , the standard atomic form factor is given by

$$F^{(\text{standard})}(q; \alpha \mathbb{J}) = 4\pi \int dr r^2 \rho(r) \frac{\sin(qr)}{qr} = \langle \alpha \mathbb{J} M = 0 | \mathbb{T}^{(1)}(\mathbf{q}) | \alpha \mathbb{J} M = 0 \rangle.$$

It is equal to the Fourier transform of the charge density  $\rho(r)$ , if the charge density is just normalized to the number of bound electrons:  $4\pi \int dr r^2 \rho(r) = N$ .

- As seen from above, this standard form factor can be expressed also as (many-electron) expectation value of the single-electron) momentum-transfer operator  $\mathbb{T}^{(1)}(\mathbf{q}) = \sum_{j=1}^N \exp(i \mathbf{q} \cdot \mathbf{r}_j)$  for a given sublevel  $|\alpha \mathbb{J} M = 0\rangle$ .
- **Modified atomic form factor:** For an atom or ion with a spherically-symmetric charge distribution, a modified atomic form factor can be defined via the charge distributions of the individual (sub-) shells of the element

$$F^{(\text{modified})}(q; \alpha \mathbb{J}) = 4\pi \sum_{i=1}^N \int dr r^2 \rho_i(r) \frac{\sin(qr)}{qr} \frac{mc^2}{\varepsilon'_i - V(r)},$$

and where  $\rho_i(r)$  is the charge density of the  $i$ -th electron,  $\varepsilon'_i = \varepsilon_i + mc^2$  its energy (including the rest mass), and  $V(r)$  the (full atomic) potential of the  $i$ -th electron as seen at position  $r$ . Here, the summation is taken over all electrons, and the charge distribution must fulfill again the constraint that  $4\pi \sum_i \int dr r^2 \rho_i(r) = N$  is equal to the number of bound electrons.

- **Generalized atomic form factor:** A generalized form factor can be defined for any pair of states; it is often calculated for the ground state  $|\alpha_o \mathbb{J}_o M_o\rangle$  and some excited state  $|\alpha_e \mathbb{J}_e M_e\rangle$  by:

$$F^{(\text{generalized})}(q; \alpha_e \mathbb{J}_e M_e; \alpha_o \mathbb{J}_o M_o) = \langle \alpha_e \mathbb{J}_e M_e | \mathbb{T}^{(1)}(\mathbf{q}) | \alpha_o \mathbb{J}_o M_o \rangle .$$

This generalized atomic form factor help describe also inelastic scattering processs and is closely related to the **generalized oscillator strength**, cf. Section 7.3.m.

### 6.1.e. Energy shifts in plasma environments (PlasmaShift)

#### Properties, notations & application:

- **Plasma shift** of an atomic level:  $E(\alpha \mathbb{J}) \longrightarrow E(\alpha \mathbb{J}; \text{plasma model and parameters})$  .
- **Formal quantum notation:**  $\Delta E^{(\text{plasma})}(\alpha \mathbb{J}) = E(\alpha \mathbb{J}; \text{plasma model and parameters}) - E(\alpha \mathbb{J}; \text{unperturbed})$  .
- **Spectral line shifts:** The interaction between the plasma and the bound electrons of an ion embedded into the plasma alters the ionic structure and transition properties. Therefore, **a plasma environment generally results in spectral line shifts, line broadening and changes in the line shapes**, when compared with the free ion.
- For realistic plasma conditions, the number of bound states of the ions often becomes finite and the electrons less tightly bound by the nucleus, if the plasma coupling parameter increases.
- Using JAC: Perform an `Atomic.Computation(..., properties=[Plasma, ...], configs=[...], plasmaSettings=PlasmaShift.Settings(...), ...)` or call directly functions from the module `PlasmaShift`.
- In JAC, the plasma (energy) shifts are calculated and tabulated by default for a Debye-Hückel plasma model with screening parameter  $\lambda = 0.25$ . Hereby, the Debye-Hückel interaction is incorporated only in first-order perturbation theory within the CI matrix.

#### High-density plasma:

- Relativistic corrections to the plasma screening may become necessary for a detailed analysis of atomic spectra, especially if lines of multiply-charged ions are used in some spectral analysis.

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- **Collisional-radiative model:** The effects of the plasma-environment can be further taken into account also by kinetic *or* collisional-radiative models, and where the energy shifts are often introduced by semi-empirical formulas, such as the formula by Stewart and Pyatt.
- **Application of simple plasma models:** Apart from plasma diagnostics, a **proper treatment of a (plasma) screening potential** has been found useful also in other areas of physics, such as nuclear and elementary particle physics, solid state physics, or even the design of nanostructures.
- **Strongly-coupled plasma:** For a **strongly-coupled dense plasma**, the (averaged) electrostatic energy between neighbouring particles in the plasma exceeds their thermal energy. The coupling parameter  $\Gamma = q^2/4\pi\epsilon_0 R_0 kT > 1$  is used in order to characterize such a strongly-coupled plasma. This **coupling parameter  $\Gamma$  simply describes the ratio of the electrostatic and the thermal energy of the plasma particles, and where  $R_o = (3/4 \pi n_e)^{1/3}$ .**
- **Laser-induced plasma with  $\Gamma \gg 1$ :** Strongly-coupled plasma obeys generally classical statistics and has been observed for a large number of experimental and astrophysical conditions. In the laboratory, for instance, laser-induced plasma with  $\Gamma \gg 1$  has been investigated in high-compression fusion experiments. Such a laser-produced plasma provides a very typical example of a high-density, strongly-coupled plasma with electronic densities well beyond  $10^{24} \text{ cm}^{-3}$ .
- Theoretical (plasma) models for hot and dense plasma are necessary to simulate its properties, such as the equation of state, radiative transfer coefficients or the conduction coefficients in plasma. These properties are relevant for inertial confinement fusion and astrophysics.
- **Average-atom (AA) model:** The AA model has been utilized to describe the electronic structure of hot and dense plasmas. This model was developed in order to describe **bound electrons in terms of the (atomic) shell model, and to deal with the free electrons within the Thomas-Fermi statistical model.**
- **Plasma models:** Plasma screening effects have been investigated extensively by employing various analytical models, such as the ion sphere, Debye-Hückel, Stewart-Pyatt as well as Ecker-Kröll models.

### Frequently applied plasma models in atomic structure calculations:

- **Debye-Hückel model:** This model has been commonly used in order to incorporate plasma effects into atomic structure calculations. However, the **validity of this model remains rather questionable** because it is valid only when the *correlation* time of the ion is much longer than the lifetime of excited atomic states. Therefore, **this perturbative approach is at best valid for weakly-coupled plasmas and should not be applied for modeling of high-density plasmas.**
- **Coulomb interaction in the Debye-Hückel model:** In this model, the (pairwise) Coulomb interaction among the charged particles is modified

(screened) for an ion with nuclear charge  $Z$  to

$$V^{(\text{DH})}(r, \lambda) = -\sum_i^N \frac{Ze^{-\lambda r_i}}{r_i} + \sum_{i>j}^N \frac{e^{-\lambda r_{ij}}}{r_{ij}} = \sum_i^N V^{(\text{e-n, DH})}(r_i, \lambda) + \sum_{i>j}^N V^{(\text{e-e, DH})}(r_{ij}, \lambda).$$

Here,  $N$  is the number of bound electrons in the ion,  $r_i$  is the distance of the  $i$ -th electron from the nucleus, and  $r_{ij}$  the distance between the electrons  $i$  and  $j$ .

- **Plasma screening parameter  $\lambda$ :** This parameter is the inverse of the Debye shielding length for a certain plasma environment and can be readily expressed in terms of the electron density  $n_e$  and the temperature  $T_e$  of the plasma by:

$$\lambda = \left[ \frac{4\pi n_e}{k T_e} \right]^{\frac{1}{2}}.$$

- **Debye-Hückel model:** To obtain the modified two-particle integrals due the plasma screening, we can write

$$V^{(\text{e-e, DH})}(r_{ij}, \lambda) = -\lambda \sum_{l=0}^{\infty} (2l+1) j_l(i\lambda r_{<}) h_l^1(i\lambda r_{>}) P_l(\cos\theta)$$

in terms of the *larger* ( $r_{>} = \max(r_i, r_j)$ ) and *smaller* radii ( $r_{<} = \min(r_i, r_j)$ ) of the one-particle radii  $r_i$  and  $r_j$ , respectively, and where  $j_l$  denotes a spherical Bessel function and  $h_l^1$  a Hankel function of the first kind.

- **Debye-Hückel model:** While the Debye-Hückel screening of the electron-nucleus interaction usually destabilizes the binding of the electron, the screening of the electron-electron repulsion counteracts this trend.
- Saha and Fritzsche (2006) demonstrated for beryllium-like ions that the Debye-Hückel screening (if applicable) should be incorporated into both, the electron-nucleus and electron-electron interaction in order to obtain reliable results for the plasma shifts.
- **Ion-sphere model:** In this model, the ion is typically enclosed in a spherically symmetric cell that contains the exact number of electrons in order to ensure neutrality. In the ion-sphere model, one needs to define an electron-density distribution that either obeys self-consistency equations or need to make some simple hypothesis, such as an uniform density within the cell.

## Plasma in oscillating electric fields:

- **Plasma in oscillating electric fields:** The collective movement of plasma particles, such as the electron plasma waves (Langmuir oscillations) or ion plasma waves, is mainly determined by the plasma creation and/or heating device but lead also to quite significant fields (Peyrusse 1977).

## Free-electron distributions in plasma:

- In a plasma, the free-electron (radial) density  $n_e(r)$  can follow different distributions:
- **Maxwell-Boltzmann distribution:** The Maxwell-Boltzmann statistics has been often applied to describe the free-electron distribution but is **not well justified for high-density or low-temperature plasma**.
- **Fermi-Dirac distribution:** This distribution is preferable if the free electrons are degenerate

$$n_e(r) = \frac{4}{\sqrt{\pi}} \int_{r_o}^{\infty} dr \frac{\sqrt{r}}{e^{r-r_o - \mu/kT_e} + 1} = 2 \lambda_{th}^{-3} \mathcal{F}_{1/2} \left( -\frac{V(r)}{kT_e} + \frac{\mu}{kT_e}; -\frac{V(r)}{kT_e} \right).$$

Here  $\mu$  is the chemical potential of the plasma,  $\mathcal{F}_{1/2}(x, y)$  the incomplete Fermi-Dirac integral and  $\lambda_{th} = \sqrt{2\pi/kT_e}$  the thermal de-Broglie thermal wavelength.

- **Uniform electron-gas distribution:** This distribution considers a sphere of radius  $R_o = (3/4\pi n_e)^{1/3}$  around the ion with  $N_b$  bound electrons, in which the ionic charge is completely neutralized by the additional  $N_f = Z - N_b$  free electrons. Although all electrons inside of the sphere will strongly interact with the embedded ion, the free electrons are here assumed to be distributed uniformly. With this assumption in mind, the electrostatic potential for the bound electrons with radial coordinate  $r$  is given by

$$V^{(IS)}(\mathbf{r}; \mathbf{R}_0) = \begin{cases} -\frac{Z}{r} + \frac{Z-N_b}{2R_0} \left[ 3 - \left( \frac{r}{R_0} \right)^2 \right] & \text{for } \mathbf{r} \leq \mathbf{R}_0 \\ 0 & \text{for } \mathbf{r} > \mathbf{R}_0. \end{cases}$$

In this potential above, the first term describes the interaction of the bound electron with the nucleus, while the second **(repulsive) term arises from the plasma background and always causes a shift of all the level energies towards the continuum**.

- **Collision-induced localization:** The plasma environment typically leads to a broadening of the momentum distribution of free electrons and partly also to its localization. Indeed, random collisions of the continuum electron with particles from the plasma environment may

results in a loss of coherence of the outgoing, photoionized electrons when leaving the parent ions. Such a collision-induced localization resembles the Anderson localization of randomly scattered electrons in solids as well as the dynamical localization of Rydberg electrons in plasma.

- **Collision-induced localization:** The collision-induced energy and momentum broadening of stationary states can reduce also the lifetime of the bound states. At a given plasma temperature, the momentum broadening increases with the plasma density or, in other words, the electron localization is more pronounced in a dense plasma.

### Gordon-Kim theory:

- In the Gordon-Kim theory (1972), the total energy of the system consists out of the direct Coulomb potential energy, the exchange Coulomb potential energy, the kinetic energy as well as the correlation energy. The direct Coulomb potential energy is calculated by integrating the Coulomb potential for the (approximated) electron density, while the other contributions to the total energy were calculated by using a temperature-dependent local density functional (LDA) approach.

### 6.1.f. Level-dependent fluorescence and Auger yields (DecayYield)

#### Properties, notations & application:

➤ Fluorescence *versus* Auger decay of an (excited) atomic level:  $A^{q+*} \longrightarrow \begin{bmatrix} A^{q+(*)} + \hbar\omega \\ A^{(q+1)+(*)} + e_a^- \\ \dots \end{bmatrix}.$

The corresponding level-specific or configuration-averaged fluorescence and Auger yields just describe the fraction of atoms in a particular level or configuration that decay either by fluorescence *or* Auger electron emission.

➤ Formal quantum notation:  $|\alpha\mathbb{J}\rangle \longrightarrow \begin{bmatrix} \{ |\alpha_r\mathbb{J}_r\rangle + \hbar\omega(\{\mathbb{M}\}) \} \\ \{ |\alpha_a\mathbb{J}_a\rangle + |\varepsilon\kappa\rangle \} \\ \dots \end{bmatrix}.$

- Using JAC: Perform an `Atomic.Computation(.., properties=[Yields, ..], configs=[..], yieldSettings=DecayYield.Settings(..), ..)` or call directly functions from the module `DecayYield`.
- In JAC, we calculate and tabulate by default the fluorescence yields  $\omega_r(\alpha\mathbb{J})$  and Auger yields  $\omega_a(\alpha\mathbb{J})$  for all selected levels.
- In JAC, the computation of all fluorescence and Auger yields is always traced back to single-step cascade computations. Therefore, the same (cascade) approaches `{ AverageSCA, SCA, ... }` can be applied as in cascade computations, cf. section 8.1. The particular approach, that is to be applied in the given computation above, need to be specified however as `String` in the `DecayYield.Settings("AverageSCA", ...)`.
- In JAC, the explicit settings of the underlying `Cascade.Computation` can be overwritten in the function `Cascade.determineSteps()`; no attempt has been made so far to support such control features directly by means of the `DecayYield.Settings(...)`.
- Fluorescence yields  $\omega_r(\alpha\mathbb{J})$ : This yield typically represents the probability of an inner-shell (core) hole to be filled under photon emission, in contrast and competition with other nonradiative processes.
- Auger yields  $\omega_a(\alpha\mathbb{J})$ : This yield describes the (complementary) probability for an inner-shell hole level to decay by electron emission, hence:  $\omega_r(\alpha\mathbb{J}) + \omega_a(\alpha\mathbb{J}) = 1$ . The Auger yield is sometimes further partitioned into an Auger and Coster-Kronig yield:  $\omega_a(\alpha\mathbb{J}) = \omega^{(\text{Auger})}(\alpha\mathbb{J}) + \omega^{(\text{Coster-Kronig})}(\alpha\mathbb{J})$ .



- If not explicitly overwritten within the code [cf. the module `Cascade`], only the electric-dipole (E1) transitions and the **Coulomb interaction**  $\mathbb{V}(\text{Coulomb})$  are taken into account in the evaluation of all Auger and radiative transition amplitudes and rates.
- More often than not, the fluorescence and Auger yields are computed and analyzed for either  $K$ - or  $L$ -shell (holes) states, and for which usually no distinction need to be made for the different fine-structure levels ( $\alpha\mathbb{J}$ ) that arise only due to the couplings of the valence shells.
- **Typically, it is completely sufficient to include only electric-dipole decay channels**, although other multipoles could be incorporated explicitly in the computations, as far this is supported by the underlying cascade approach, or invoked explicitly within the code.

## 6.2. In JAC partly-implemented level properties

### 6.2.a. Approximate many-electron Green function for atomic levels (GreenFunction)

#### Properties & notations:

- **Green function representation for the levels of one or several (given) configurations:** Such a representation refers to just one or several lists of many-electron levels more multiplets  $\left\{ E_{\nu_i}, |\psi_{\nu_i}(\beta_i \mathbb{J}_i)\rangle, \nu_i = 1, \dots, \nu_i^{(\max)} \right\}$  of total symmetry  $\mathbb{J}_i$  and with a maximum number of  $n = 0, \dots, n_{\max}$  electrons within the continuum. Therefore, these levels formally belong to both, the bound-state spectrum of the atom ( $E_{\nu_i} < 0$ ) as well as to the continuum spectrum (scattering states) with 1, 2, ... electrons in the continuum. All these levels are constructed from a single set of one-electron orbitals, and which is properly chosen to represent the bound-state levels of the given configurations. **Each Green function representation typically consists of several (Green function) channels and is classified by the (maximum) number of free electrons together with a pre-specified (de-) excitation scheme in order to generate from the given configurations all nonrelativistic configurations as to be considered in the many-electron (CSF) basis of the Green function representation.**
- **Formal quantum notation:**  $[\{E_{\nu_1}, |\psi_{\nu_1}(\beta_1 \mathbb{J}_1)\rangle\}, \{E_{\nu_2}, |\psi_{\nu_2}(\beta_2 \mathbb{J}_2)\rangle\}, \dots] \longrightarrow \mathbb{G}(E) = \int \sum_{\nu_1 \nu_2 \dots} \frac{|\psi_{\nu_i}(\beta_i \mathbb{J}_i)\rangle \langle \psi_{\nu_i}(\beta_i \mathbb{J}_i)|}{E - E_{\nu_i}}$
- Using JAC: Perform an `Atomic.Computation(..., properties=[Green, ...], configs=[...], greenSettings=GreenFunction.Settings(...), ...)` or call directly functions from the module `GreenFunction`.
- In JAC, we provide an approximate representation of the many-electron Green functions for one or several given electron configurations, and which is suitable for numerical computations.

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- In JAC, the approximation of the requested Green function representation is determined by an abstract `approach::AbstractGreen Approach`. Further details can be specified for the requested Green function by the maximum principle quantum number  $n^{(\max)}$  and a list of orbital angular momenta  $[\ell_1, \ell_2, \dots]$  that are allowed within the given (de-) excitation scheme in order to generate the many-electron (CSF) basis. These parameters can be specified in `GreenFunction.Settings(...)`.
- **Classification of the many-electron Green functions:** In JAC, the major classification of Green function representations is based on the number of free electrons in the continuum. The many-electron basis must hereby include both, bound-type CSF as well as CSF with up to the maximum number of free electrons, and constructed in a systematic manner. *At present, we restrict ourselves to Green function representations with at least  $N - 1$  bound electrons and, hence, just a single electron in the continuum.*
- **Green function channel:** A channel refers to the set of atomic levels  $\{E_{\nu_i}, |\psi_{\nu_i}(\beta_i \mathbb{J}_i)\rangle\}$  with same total symmetry  $\mathbb{J}_i$ .
- **Green function approaches:** Several approaches are distinguished in JAC in order to determine how the set of many-electron levels  $\{E_{\nu_i}, |\psi_{\nu_i}(\beta_i \mathbb{J}_i)\rangle\}$  are to be generated and how much of the electron-electron interaction is taken into account to represent *together* the bound and continuum spectra of the atom or ion:
  - (a) **Diagonal CSF basis without configuration interaction (SingleCSFwithoutCI):** This is a fast though very rough approximation, in which each CSF with total symmetry  $\mathbb{J}_i$  also just represents a single level  $|\psi_{\nu_i}\rangle$  of the spectrum. This approximation fully omits all configuration interactions within the CSF basis.
  - (b) **Configuration interaction only between bound-state orbitals (CoreSpaceCI):** This approach diagonalizes (in turn) the Hamiltonian matrices of symmetry  $(\beta_i \mathbb{J}_i)$  by including the electron-electron interaction only between bound-state orbitals, while their interaction with and among the free electrons are neglected.
  - (c) **Damped configuration-interaction approach (DampedCI):** This approach again diagonalizes the Hamiltonian matrices of symmetry  $(\beta_i \mathbb{J}_i)$  by *including the electron-electron interaction for all pairs of electrons but damped by some exponential factor  $e^{-\tau r}$* . The use of such a *damping function* in the Slater integrals (and, possibly the Breit integrals – though this has not been realized so far) ensures that the bound-bound, bound-free and free-free interactions are treated equally without that a *continuation* of the continuum orbitals for  $r \rightarrow \infty$  need to be considered in the numerical evaluation of the radial integrals.
- **(De-) Excitation schemes of electrons for the representation of approximate Green functions:** Since many-electron Green functions are very complex entities with many ‘infinities’ already within their formal definition, further control about the generation of such Green function approximations is obtained by specifying a proper (de-) excitation scheme of the electrons with regard to the given configurations:
  - (i) **(De-) excitation of a single electron from a given set of (nonrelativistic) configurations (DeExciteSingleElectron):** Cf. section 3.3
  - (\*) Further (de-) excitation schemes will be implemented as the needs arises in selected applications.

- **Approximate many-electron Green functions play an essential role in studying all second- and higher-order perturbation processes as well as for various dynamic processes in atomic systems.** For many-electron atoms and ions, however, the construction of proper Green function approximations is still rare in the literature because of their complexity and the required computational resources. More often than not, Green functions have instead been used in the past in order to describe either single-electron or quasi-single-electron atoms.
- **‘Infinities’ in Green function representations:** Already a single-electron Green function generally implies three *infinities* due to the electron’s spatial degree of freedom, and which are often treated numerically by restricting the principal quantum number  $n^{(\max)}$  and the angular momentum quantum numbers  $j^{(\max)}$ ,  $m_j^{(\max)}$ . An  $N$ -fold multiple of such infinities arise for  $N$ -electron Green functions and make their approximation a real challenge. In JAC, approximation many-electron Green functions are restricted by:
  - orbitals with a principal quantum number  $n^{(\max)}$  which applies for to all electrons;
  - a specified list of orbital angular momenta  $[\ell_1, \ell_2, \dots]$ ;
  - the coupling of the  $N$  orbitals to total symmetries  $(\beta_1 \mathbb{J}_1), (\beta_2 \mathbb{J}_2), \dots$ .
 Each of these symmetries represent a many-electron continuum, in which the total Hamiltonian is diagonal.
- Apart from the wave function, which describe the atom or ion in some particular quantum states, one often also needs a summation over a (quasi-) complete many-electron spectrum, especially if parts of the atomic interaction are treated as a perturbation.
- Although such a summation over a complete spectrum appears rather straightforward to be carried out explicitly, the large number of levels in a typical Green function representation and the need of *continuum-continuum* matrix elements may hamper such an approach.

### Green functions for scattering processes:

- **Green functions for the study of scattering processes:** Formally, most scattering processes can be analyzed by simply solving the Schrödinger equation for some given many-electron Hamiltonian  $\mathbb{H}$ , the collision energy  $E$  and by including the proper scattering boundary conditions

$$\mathbb{H} |\Psi^+\rangle = E |\Psi^+\rangle.$$

- **Lippman-Schwinger representation:** A formal solution to the Schrödinger equation above is given by the Lippman-Schwinger equation which provides an integral form of the scattering wave function  $|\Psi^+\rangle$  in terms of the initial state  $|\Psi_o\rangle$  and the many-electron Green function  $\mathbb{G}^+(E)$  (cf. Tong *et al.*, 2009)

$$|\Psi^+\rangle = |\Psi_o\rangle + \frac{1}{E - \mathbb{H} + i\eta} \mathbb{V} |\Psi_o\rangle + \dots = |\Psi_o\rangle + \mathbb{G}^+(E) \mathbb{V} |\Psi_o\rangle + \dots$$

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Here,  $\eta > 0$  is a small positive infinitesimal which ensures that the boundary conditions of an outgoing wave are properly taken into account, and where a (radiative) loss rate can be incorporated into the representation of  $|\Psi^+\rangle$  by means of a complex energy  $E$ .

- **Green function with optical potential:** The Greens function of the scattering system with total scattering energy  $E$  can be formally written as

$$\mathbb{G}^+(E) = \sum_{\nu} \frac{|\psi_{\nu}\rangle \langle \psi_{\nu}|}{E - E_{\nu} + i\eta},$$

and where the summation runs over all (generalized) eigenfunctions of the Hamiltonian  $\mathbb{H}$ . This Green function can be utilized also for scattering and dynamical processes if a proper (sub-) set of many-electron levels  $\{E_{\nu}, |\psi_{\nu}\rangle\}$  is taken into account.

- **Optical potential:** If a proper optical potential is added to the (inner-) atomic potential, the outgoing wave is formally absorbed due to the occurrence of complex single-electron energies. Tong *et al.* (2009) suggest an energy-dependent optical potential of the form

$$V^{(\text{optical})}(r; \varepsilon) = \begin{cases} 0 & \text{for } r < r_c \\ V_o \left( \frac{r - r_c}{r_{\max} - r_c} \right) & \text{for } r \geq r_c \end{cases}, \quad \frac{\varepsilon^{1/2}}{r - r_c} < V_o < \varepsilon^{1/2} (r_{\max} - r_c),$$

and where the optical potential becomes *nonzero* at  $r_c$ . However, little experience exists so far how the parameters  $r_c$ ,  $r_{\max}$ ,  $\varepsilon$  need to be chosen in order to obtain a realistic representation of the Green function, and how useful such an approximation is in different physical contexts.

- **Diagonalization of the Hamiltonian matrix with an optical potential:** For an optical potential, the **Hamiltonian matrix is no longer hermitian but symmetric complex**, and it generally includes both diagonal and off-diagonal matrix elements due to the electron-electron interaction. From the diagonalization of the Hamiltonian matrix in a CSF basis with well-defined total symmetry  $\mathbb{J}_i$ , one then obtains the eigenenergies  $\{E_{\nu}\}$  and (atomic) eigenvectors  $\{E_{\nu_i}, |\psi_{\nu_i}\rangle\}$ , but which need to be calculated just once.

### 6.2.b. Scattering amplitudes and scattering factors (FormFactor)

#### Properties, notations & application:

- Scattering amplitude  $S(q; \alpha \mathbb{J})$  of an atomic level with supposed spherical-symmetric charge distribution:.
- Formal quantum notation: .

#### Anomalous scattering factors:

- Coherent Rayleigh scattering factor:
- Using the modified form factor for describing the scattering in forward-scattering angle  $\vartheta \approx 0$ , the real  $g'(\omega)$  and imaginary  $g''(\omega)$  anomalous scattering factors can be approximately defined in terms of the photoionization cross sections and oscillator strength (Zhou *et al.*, 1990). This formally requires a summation over the complete spectrum of the atom, although the major contributions arise from just a few resonances.
- Anomalous scattering factor: The real  $g'(\omega)$  and imaginary  $g''(\omega)$  anomalous scattering factors are frequently employed in order to express the coherent (Rayleigh) scattering factor  $G(\omega)$  for incident photons with energy  $\hbar\omega$ .
- Coherent Rayleigh scattering factor: For incident photons with energies  $\hbar\omega$ , the coherent (Rayleigh) scattering factor in forward direction has been written as (Zhou *et al.*, 1990)

$$G(\omega) = F(0) + g'(\omega) + i g''(\omega),$$

where  $F(0) = N$  is the atomic form factor. Therefore, the (Rayleigh) scattering amplitude and scattering factor can be expressed also in terms of the standard (atomic) form factor (FF) or the modified form factor (MFF).

#### Incoherent scattering function $S(\mathbf{q})$ :

- Incoherent scattering function:
- Incoherent scattering function: This scattering function can be expressed in terms of the generalized atomic form factor

$$S(\mathbf{q}; \alpha_o \mathbb{J}_o M_o) = \sum_{e \neq o} |F^{(\text{generalized})}(q; \alpha_e \mathbb{J}_e M_e; \alpha_o \mathbb{J}_o M_o)|^2 = \langle \alpha_o \mathbb{J}_o | \mathbb{T}^{(2)}(\mathbf{q}) | \alpha_o \mathbb{J}_o \rangle - |F^{(\text{standard})}(q; \alpha_o \mathbb{J}_o)|^2.$$

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but where the summation over  $e$  includes all discrete levels as well as an integration over the continuum. In the second equivalence, the two-electron momentum-transfer operator  $\mathbb{T}^{(2)}(\mathbf{q}) = \sum_{m,n=1}^N \exp[i \mathbf{q} \cdot (\mathbf{r}_m - \mathbf{r}_n)]$  occurs, but then just requires to compute the expectation value of this operator with regard for the ground state  $|\alpha_o \mathbb{J}_o M_o\rangle$ .

### 6.2.c. Lande $g_J$ factors and Zeeman splitting of fine-structure levels (LandeZeeman)

#### Properties, notations & application (not yet fully implemented):

- Zeeman splitting of an atomic level into Zeeman (sub-) levels:  $|\alpha \mathbb{J}\rangle \longrightarrow |\alpha \mathbb{J} M\rangle$ ,  $M = -J, \dots, J-1, +J$ .
- Formal quantum notation:  $|\alpha \mathbb{J} M\rangle$ .
- Using JAC: Perform an `Atomic.Computation(.., properties=[LandeJ, ..], configs=[..], zeemanSettings=LandeZeeman.Settings(..), ..)` with `calcLandeJ=true` or call directly functions from the module `LandeZeeman`.
- In JAC, we calculate and tabulate by default the Lande  $g_J \equiv g(\alpha \mathbb{J})$  factors for the selected fine-structure levels.
- In the JAC program, the Zeeman amplitudes  $\langle \beta_r \mathbb{J}_r \| \mathbb{N}^{(1)} \| \beta_s \mathbb{J}_s \rangle$  and  $\langle \beta_r \mathbb{J}_r \| \Delta \mathbb{N}^{(1)} \| \beta_s \mathbb{J}_s \rangle$  can be obtained from the function `Jac.LandeZeeman.amplitude()`.

#### Further remarks:

- The Zeeman effect describes the (level) splitting of an atomic level  $|\alpha \mathbb{J}\rangle$  into its sub-levels  $|\alpha \mathbb{J} M\rangle$  in the presence of a static magnetic field.
- The various components  $|\alpha \mathbb{J} M\rangle$  of a Zeeman-split line usually have different intensities, and with some of them possibly even forbidden within the dipole approximation.
- The energy splitting of the Zeeman sub-levels depends of course on the magnetic field strength. — Therefore, the Zeeman effect has been applied for measuring magnetic field strengths, e.g. in astrophysics at the surface of the Sun or in laboratory plasmas.

### Zeeman Hamiltonian $\mathbb{H}^{(\text{mag})}$ :

- If the  $z$ -axis is chosen along the magnetic field  $\mathbf{B} \parallel \mathbf{e}_z$  and if we neglect diamagnetic contributions, the **magnetic (Zeeman) Hamiltonian** can be written in terms of spherical tensors (Andersson and Jönsson, 2008)

$$\mathbb{H}^{(\text{mag})} = \left( \mathbb{N}_0^{(1)} + \Delta \mathbb{N}_0^{(1, \text{QED})} \right) B$$

$$\mathbb{N}^{(1)} = \sum_j^N \mathfrak{n}^{(1)}(j) = \sum_j^N -i \frac{\sqrt{2}}{2\alpha} r_j (\boldsymbol{\alpha}_j \mathbb{C}^{(1)}(j))^{(1)}$$

$$\Delta \mathbb{N}^{(1)} = \sum_j^N \Delta \mathfrak{n}^{(1)}(j) = \sum_j^N \frac{g_s - 2}{2} \beta_j \boldsymbol{\Sigma}_j, \quad \boldsymbol{\Sigma} = \begin{pmatrix} \boldsymbol{\sigma} & 0 \\ 0 & \boldsymbol{\sigma} \end{pmatrix},$$

and where the second term in the magnetic Hamiltonian above is (so-called) Schwinger's QED correction. Here,  $g_s = 2.00232$  is the  $g$ -factor of the electron, including some leading QED corrections.

### Lande factor $g_J \equiv g(\alpha\mathbb{J})$ and Zeeman splitting of an electronic level $|\alpha\mathbb{J}\rangle$ :

- **Zeeman splitting of an atomic level  $|\alpha\mathbb{J}\rangle$  :** If the energy splitting due to the magnetic field is small, when compared with the fine-structure splitting, it can be expressed in first-order perturbation theory by

$$\begin{aligned}\Delta E^{(\text{mag})}(\alpha\mathbb{J}M) &= E(\alpha\mathbb{J}M) - E(\alpha\mathbb{J}, M=0) \\ &= \left\langle \alpha\mathbb{J}M \left| \mathbb{N}_0^{(1)} + \Delta \mathbb{N}_0^{(1, \text{QED})} \right| \alpha\mathbb{J}M \right\rangle = (-1)^{J-M} \begin{pmatrix} J & 1 & J \\ -M & 0 & M \end{pmatrix} \sqrt{2J+1} \left\langle \alpha\mathbb{J} \left\| \mathbb{N}^{(1)} + \Delta \mathbb{N}^{(1, \text{QED})} \right\| \alpha\mathbb{J} \right\rangle B \\ &= \frac{M}{\sqrt{J(J+1)}} \left\langle \alpha\mathbb{J} \left\| \mathbb{N}^{(1)} + \Delta \mathbb{N}^{(1, \text{QED})} \right\| \alpha\mathbb{J} \right\rangle B\end{aligned}$$

and with regard to the unperturbed energy  $E(\alpha\mathbb{J}, M=0) = E(\alpha\mathbb{J}, B=0)$ .

- **Lande  $g_J$  factor of an atomic level  $|\alpha\mathbb{J}\rangle$  :** Usually, the dependence on the  $M$  quantum number is factored out and the energy splitting is expressed for a single level  $|\alpha\mathbb{J}\rangle$  in terms of its Lande factor  $g_J \equiv g(\alpha\mathbb{J})$

$$\Delta E^{(\text{mag})}(\alpha\mathbb{J}M) = g(\alpha\mathbb{J}) M \frac{B}{2}, \quad g(\alpha\mathbb{J}) \equiv g_{\mathbb{J}} = 2 \frac{\left\langle \alpha\mathbb{J} \left\| \mathbb{N}^{(1)} + \Delta \mathbb{N}^{(1, \text{QED})} \right\| \alpha\mathbb{J} \right\rangle}{\sqrt{J(J+1)}}$$

### 6.2.d. Lande $g_F$ factors and Zeeman splitting of hyperfine levels (LandeZeeman)

#### Properties, notations & application (not yet fully implemented):

- **Zeeman splitting** of an atomic hyperfine level into Zeeman (sub-) levels:  $|\alpha(J)\mathbb{F}\rangle \longrightarrow |\alpha(J)\mathbb{F}M\rangle, \quad M = -F, \dots, F-1, +F$ .
- **Formal quantum notation:**  $|\alpha(J)\mathbb{F}M\rangle$ .



- Using JAC: Perform an `Atomic.Computation(.., properties=[LandeJ, ..], configs=[..], zeemanSettings=LandeZeeman.Settings(..), ..)` with `calcLandeF=true` or call directly functions from the module `LandeZeeman`.
- In JAC, we calculate and tabulate by default the Lande  $g_F \equiv g(\alpha\mathbb{F})$  factors for the selected hyperfine-structure levels.

### Further remarks:

- The Lande  $g$ -factor (also known as  $g$  value or dimensionless magnetic moment) is a **dimensionless entity that characterizes the gyromagnetic ratio of an atomic or hyperfine level**. The Lande  $g$ -factor is basically a proportionality constant how efficient a particle, spin or quantum state with given total angular momentum  $s$  contributes to the splitting in an external magnetic field  $B$ .

### Lande factor $g_F \equiv g(\alpha\mathbb{F})$ of an hyperfine level $|\alpha\mathbb{F}\rangle$ :

- **Zeeman splitting of an hyperfine level  $|\alpha\mathbb{F}\rangle$  :** If the energy splitting due to the magnetic field is small, when compared with the hyperfine-splitting, it can be expressed in first-order perturbation theory in terms of the reduced electronic Zeeman amplitudes by

$$\begin{aligned}
 \Delta E^{(\text{mag})}(\alpha\mathbb{F}M) &= E(\alpha\mathbb{F}M) - E(\alpha\mathbb{F}, M=0) \\
 &= \left\langle (I, \alpha\mathbb{J}) FM \left| \mathbb{N}_0^{(1)} + \Delta \mathbb{N}_0^{(1, \text{QED})} \right| (I, \alpha\mathbb{J}) FM \right\rangle \\
 &= M \frac{F(F+1) + J(J+1) - I(I+1)}{2F(F+1)} \frac{\left\langle \alpha\mathbb{J} \left\| \mathbb{N}_0^{(1)} + \Delta \mathbb{N}_0^{(1, \text{QED})} \right\| \alpha\mathbb{J} \right\rangle}{\sqrt{J(J+1)}} \\
 &= M \frac{F(F+1) + J(J+1) - I(I+1)}{2F(F+1)} g(\alpha\mathbb{J}) \frac{B}{2},
 \end{aligned}$$

and expressed with regard to the unperturbed hyperfine energy  $E(\alpha\mathbb{F}, M=0) = E(\alpha\mathbb{F}, B=0)$ .

- **Lande factor  $g_F \equiv g(\alpha\mathbb{F})$  and Zeeman energy splitting of an hyperfine level  $|\alpha\mathbb{F}\rangle$  :**

$$g(\alpha\mathbb{F}) = \frac{F(F+1) + J(J+1) - I(I+1)}{2F(F+1)} g(\alpha\mathbb{J}) \quad \implies \quad \Delta E^{(\text{mag})}(\alpha\mathbb{F}M) = g(\alpha\mathbb{F}) M \frac{B}{2}$$

### 6.2.e. Sensitivity of level energies with regard to variations of $\alpha$ (AlphaVariation)

#### Properties, notations & application (not yet fully implemented):

- **Differential sensitivity parameter  $\Delta q(\delta\alpha; \beta\mathbb{J})$  of an atomic level:** For an atomic level  $|\beta\mathbb{J}\rangle$ , the level energy depends of course on the fine-structure constant  $\alpha$ . Therefore, an energy shift arises for any variation  $\delta\alpha = \alpha - \alpha_o$  of the fine-structure constant and for a non-zero differential sensitivity parameter  $\Delta q(\delta\alpha; \beta\mathbb{J})$ :

$$\Delta E(\delta\alpha; \beta\mathbb{J}) = \Delta q(\delta\alpha; \beta\mathbb{J}) \left[ \left( \frac{\alpha}{\alpha_o} \right)^2 - 1 \right],$$

and where  $\alpha_o$  is the (current) reference value of the fine-structure constant.

- **Formal quantum notation:**  $\Delta E(\delta\alpha; \beta\mathbb{J}), \Delta q(\delta\alpha; \beta\mathbb{J}), K(\beta\mathbb{J})$ .
- **Using JAC:** Perform an `Atomic.Computation(..., properties=[AlphaVar, ...], configs=[...], alphaSettings=AlphaVariation.Settings(...), ...)` or call directly functions from the module `AlphaVariation`.
- In JAC, we calculate and tabulate by default the differential sensitivity parameter  $\Delta q(\delta\alpha; \beta\mathbb{J})$  and the enhancement factor  $K(\beta\mathbb{J})$  for all selected levels.

#### Further remarks:

- **Fine-structure constant  $\alpha$ :** The dimensionless constant  $\alpha$  becomes dynamical, i.e. explicitly time-dependent, in a number of (quantum-field) theories that go beyond the standard model and general relativity.
- **Search for variations of fundamental constants:** This search is closely related to several, so-far still unexplained phenomena of the Universe, such as the nature of dark matter.
- From the individual level-dependent sensitivity parameters  $\Delta q(\delta\alpha; \beta\mathbb{J})$  and enhancement factors  $K(\beta\mathbb{J})$ , the analogue values for any atomic transition can be simply obtained by taking the sum of the corresponding level parameters.

- Experimentally, a variation of  $\alpha$  are currently measured by monitoring for two ions the ratio of (two) clock frequencies with quite different values of  $K$ .

### Enhancement and sensitivity of atomic levels with regard to variations of $\delta\alpha$ :

- **Enhancement of energy shifts due to  $\delta\alpha$  variations:** The sensitivity of an atomic level energy  $E(\beta\mathbb{J}; \alpha)$  on  $\delta\alpha$  can be re-written also in terms of the dimensionless enhancement factor  $K = 2 \Delta q / \Delta E$  as:

$$\frac{E(\beta\mathbb{J}; \alpha) - E(\beta\mathbb{J}; \alpha_o)}{E(\beta\mathbb{J}; \alpha_o)} = K \frac{\alpha - \alpha_o}{\alpha_o} = K \frac{\delta\alpha}{\alpha_o}$$

- The reasons for  $\delta\alpha$  variations can be of different type, such as temporal, spatial, slow drift, or of oscillatory nature, and it can depend on the gravity-potential or others reasons. The theoretical reasoning about such changes is typically outside of atomic theory.

## 6.3. Further properties, not yet considered in JAC

### 6.3.a. Multipole polarizabilities (MultipolePolarizability)

#### Properties & notations:

- General multipole polarizability of an atomic level: .
- Dynamic polarizability of an atomic level: .

#### Multipole polarizabilities:

- Frequency-dependent (multipole) polarizabilities are important in various fields of physics.
- The ac Stark shift of a clock transition, for example, is determined by the frequency-dependent electric-dipole polarizabilities of the clock states. The (so-called) **magic wavelengths** of a clock transition can be determined by finding the frequencies where the ac electric-dipole polarizabilities of the two clock states are the same.
- Dynamic multipole (ac  $2^L$ -pole) polarizability: of a level  $|0\rangle$  is (Porsev *et al.*, 2018)

$$\alpha^{(\mathbb{M})}(\omega, \beta_o \mathbb{J}_o) = \frac{L+1}{L} \frac{2L+1}{[(2L+1)!!]^2} (\alpha \omega)^{2L-2} \sum_{\nu} \frac{(E_{\nu} - E_o) |\langle \alpha_{\nu} \mathbb{J}_{\nu} | \mathbb{T}^{(\mathbb{M})} | \beta_o \mathbb{J}_o \rangle|^2}{(E_{\nu} - E_o)^2 - \omega^2}$$

where  $\mathbb{T}^{(\mathbb{M})}$  is the corresponding multipole operator for E1, M1, ... transitions.

#### Dynamic electric-dipole polarizability:

- The (dynamic) electric-dipole polarizability is defined as (Derivianko *et al.*, 1999)

$$\alpha^{(\text{E1})}(i\omega, \beta_o \mathbb{J}_o) = \frac{2}{3} \sum_{\nu} \frac{E_{\nu} - E_o}{(E_{\nu} - E_o)^2 + \omega^2} |\langle \alpha_{\nu} \mathbb{J}_{\nu} | \mathbf{R} | \beta_o \mathbb{J}_o \rangle|^2$$

and which includes for  $\alpha^{(\text{E1})}(\omega \equiv 0)$  the static electric-dipole polarizability.

- **Dynamic (dipole) polarizability  $\alpha(\omega)$**  : The dynamic polarizability describes the response of an atoms and nanoobjects to an (external) electromagnetic disturbance as long as this field is weaker than the (internal) atomic field:  $E \ll E^{(\text{atomic})} = m^2 e^5 / \hbar^4 \approx 5.14 \cdot 10^9 \text{ V/cm}$ , and if the wavelength of the field is much larger than the size of the object (Astapenko 2013).
- Mathematically, the dynamic polarizability  $\alpha(\omega)$  can be expressed by a second-rank tensor  $(\alpha_{ij})$  that connects the induced dipole moment  $\mathbf{d}$  of the atom with the strength of the external electric field  $\mathbf{E}$  at the frequency  $\omega$

$$d_i = \sum_j \alpha_{ij} E_j(\omega).$$

- For spherically symmetrical systems, this tensor simplifies to a scalar

$$\alpha_{ij}(\omega) = \alpha(\omega) \delta_{ij} \quad \longrightarrow \quad \mathbf{d}(\omega) = \alpha(\omega) \mathbf{E}.$$

- **Static polarizability of level  $|\beta \mathbb{J} M\rangle$**  : In a dc electric field  $\mathbf{E} = E \mathbf{e}_z$ , the static polarizability is defined as (Porsev *et al.* 1999)

$$\Delta E(\beta \mathbb{J} M) = -\frac{1}{2} \alpha(\beta \mathbb{J} M) E^2 = -\frac{1}{2} \left( \alpha^{(0)}(\beta \mathbb{J}) + \frac{3M^2 - J(J+1)}{J(2J-1)} \alpha^{(0)}(\beta \mathbb{J}) \right) E^2$$

$$\alpha(\beta \mathbb{J} M) = -2 \sum_{\nu} \frac{|\langle \beta_{\nu} \mathbb{J}_{\nu} M_{\nu} | \mathbb{D}_z | \beta \mathbb{J} M \rangle|^2}{E_{\beta} - E_{\nu}}$$

$$\alpha^{(0)}(\beta \mathbb{J}) = \frac{-2}{3(2J+1)} \sum_{\nu} \frac{|\langle \beta_{\nu} \mathbb{J}_{\nu} | \mathbb{D} | \beta \mathbb{J} \rangle|^2}{E_{\beta} - E_{\nu}}$$

$$\alpha^{(2)}(\beta \mathbb{J}) = \left( \frac{40 J (2J-1)}{3(2J+3)(2J+1)(J+1)} \right)^{1/2} \sum_{\nu} (-1)^{J+J_{\nu}+1} \begin{Bmatrix} J & 1 & J_{\nu} \\ 1 & J & 2 \end{Bmatrix} \frac{|\langle \beta_{\nu} \mathbb{J}_{\nu} | \mathbb{D} | \beta \mathbb{J} \rangle|^2}{E_{\beta} - E_{\nu}}$$

- **Clausius-Mossotti relation**: For a system embedded into a medium, the dynamic polarizability of atoms also defines the **dielectric permittivity  $\varepsilon(\omega)$  of the medium** with number density  $n_a$  of atoms, and if the medium is assumed to just consists out of a single type of atoms

$$\frac{\varepsilon(\omega) - 1}{\varepsilon(\omega) + 2} = \pi n_a \alpha(\omega)$$

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- **Measurement of the dynamic polarizability:** Experimentally, the dynamic polarizability  $\alpha(\omega) = \alpha(\omega, \beta_o \mathbb{J}_o)$  can be determined by measuring the refraction index of the substance  $n(\omega) = \sqrt{\varepsilon(\omega)}$ .
- **Quadratic Stark effect:** The dynamic polarizability also defines the energy shift  $\Delta E_n$  of an atomic level in an external electric field. In second-order perturbation theory and for a spherically-symmetric state  $|\beta_o \mathbb{J}_o\rangle$ , this level shift can be written for

$$\Delta E_n^{(2)}(\beta_o \mathbb{J}_o) = -\frac{1}{2} \alpha_n(\omega, \beta_o \mathbb{J}_o) E^2$$

- **Linear Stark effect:** If the frequency of the external field coincides with the eigenfrequency of an atom,  $\omega \approx \omega^{(atom)}$ , the energy shift is found to be linear in electric field intensity.
- **Static polarizability  $\alpha(0)$ :** For a static external field, the polarizability  $\alpha(\omega = 0)$  just describes the level shift in a constant electric field that arises in addition to the interatomic interaction potential of the atoms at long distances (van-der-Waals interaction). This static polarizability also describes the **polarization potential of a neutral atom with for slow particles with charge  $e_o$  at long distance** as well as the **scattering cross section for the elastic scattering at such a charged particle**

$$V^{(pol)} = -e_o^2 \frac{\alpha(0)}{2r^4}, \quad \sigma^{(elastic-scattering)} = 2\pi e_o \sqrt{\frac{\alpha(0)}{2E}}.$$

In deriving this cross section, one assumes a semi-classical motion of the incident particle with energy  $E$  along a well-defined trajectory with impact parameter  $b$ , and over which one need to average.

- **Atomic dipole moment:** Because of its spherical symmetry, the dipole moment of a atom is *zero* without an external field (Schiff theorem); therefore, the **induced dipole moment** is a direct measure for the disturbance of an atom by a field and for low-enough field strengths, the response of an atom to electromagnetic field can be simply characterized by its polarizability.
- **Dynamic polarizability  $\alpha(\omega)$ :** The dynamic polarizability of an atom in level  $|\beta_o \mathbb{J}_o\rangle$  can be expressed in terms of the transition frequencies and oscillator strength related to this level

$$\alpha(\omega, \beta_o \mathbb{J}_o) = \frac{e^2}{m} \sum_n \frac{f_{no}}{\omega_{no}^2 - \omega^2 - i\omega \delta_{no}} \quad \Rightarrow \quad \alpha_0 \equiv \alpha(\omega = 0) = \frac{e^2}{m} \sum_n \frac{f_{no}}{\omega_{no}^2}.$$

This expressions shows that the **dynamic polarizability of an atom is generally a complex value with the dimension of a volume**. While the imaginary part of the dynamic polarizability describes damping (constants) of the transition oscillators, i.e. the absorption of radiation, the

real part defines the refraction of an electromagnetic wave in a medium. Moreover, the summation includes both, the discrete as well as the continuous (parts of the) energy spectrum.

- **Static polarizability  $\alpha(0)$ :** This is a real and positive value which may take rather large numerical values, if there are transitions with high oscillator strength and low eigenfrequency in the atomic spectrum as, for instance. for the alkaline-earth atoms.

### Scalar and tensor polarizability:

- The scalar polarizability of a level  $\alpha\mathbb{J}$  in a light field at frequency  $\omega$  can be expressed as the sum over intermediate states that are allowed by the electric-dipole selection rules (Safronova *et al.*, 2018)

$$\alpha_0(\omega) = \frac{2}{3(2J+1)} \sum_k \frac{(E_\nu - E_o) \langle \alpha_\nu \mathbb{J}_\nu \| \mathbb{D} \| \beta_o \mathbb{J}_o \rangle}{(E_\nu - E_o)^2 - \omega^2},$$

and where the frequency  $\omega$  of the incident light is assumed to be off resonance by at least several linewidths with regard to any of the intermediate levels  $\{|\alpha_\nu \mathbb{J}_\nu\rangle\}$ .

- The expression for the tensor polarizability has a similar structure.

### Electric-dipole and magnetic-dipole susceptibilities:

- **Susceptibility:** The susceptibility  $\chi^{(\text{electric, magnetic})}$  is generally a (dimensionless) proportionality constant that characterizes the induced polarization of a medium in response to an applied electric or magnetic field. The greater the susceptibility, the greater the response of the material to an applied field.
- **Electric susceptibility  $\chi_e \equiv \chi^{(\text{electric})}$ :** A proportionality **constant or tensor which relates the induced dielectric polarization density  $\mathbf{P}$  to the applied field** and which is very closely related to the atomic or molecular polarizability  $\alpha$ :

$$\mathbf{P} = \varepsilon_o \chi_e \mathbf{E}, \quad \mathbf{p} = \varepsilon_o \alpha \mathbf{E}_{\text{local}}$$

Here  $\mathbf{p}$  refers to the induced atomic or molecular electric-dipole moment. Analogue expressions also exist for the magnetic susceptibility  $\chi^{(\text{magnetic})}$ , the magnetic field  $\mathbf{B}$  and the magnetization  $\mathbf{M}$  as well as for the magnetic-dipole moment  $\mathbf{m}$ .

### 6.3.b. Dispersion coefficients

#### Remarks & notations:

- **Lennard-Jones model:** This model refers to a mathematically simple potential for describing the interaction between a pair of neutral atoms or molecules. It was proposed first in 1924 by John Lennard-Jones.
- The long-range interactions between atoms and molecules play a prominent role especially in the low-energy and low-temperature collision experiments.
- **Dispersion coefficients  $C_n$ :** The interaction of atoms at large interatomic distances, for instance for cold gases in a cold traps, is usually described in terms of **dispersion (van der Waals) coefficients  $C_n$** .
- **Alkaline-earth dimers:** For two alkaline-earth atoms in their ground state, the long-range potential as function of the internuclear distance  $R$  can be written as (Porsev and Derivianko, 2006)

$$V(R) = -\frac{C_6}{R^6} - \frac{C_8}{R^8} - \frac{C_{10}}{R^{10}} - \dots$$

#### Atomic computation of dispersion coefficients:

- **Atom-wall coefficient,  $C_3$ :** For an atom in a spherically-symmetric ground state  $|\psi_o\rangle = |\beta_o \mathbb{J}_o M_o\rangle$ , the dispersion coefficient  $C_3$  of the Lennard-Jones interaction between an atom and a perfectly conducting wall is given by (Derivianko *et al.*, 1999)

$$C_3 = \frac{1}{4\pi} \int_0^\infty d\omega \alpha^{(E1)}(i\omega; \beta_o \mathbb{J}_o M_o) = \frac{1}{12} \langle \beta_o \mathbb{J}_o M_o | \mathbf{R} \cdot \mathbf{R} | \beta_o \mathbb{J}_o M_o \rangle, \quad \mathbf{R} = \sum_{i=1}^N \mathbf{r}_i.$$

- **Van-der-Waals coefficient,  $C_6$ :** For an atom in a spherically-symmetric ground state  $|\psi_o\rangle \equiv |\beta_o \mathbb{J}_o M_o\rangle$ , the dispersion coefficient  $C_6$  of the Lennard-Jones interaction can be expressed as integral over the dynamic electric-dipole polarizability (Derivianko *et al.*, 1999)

$$C_6 = \frac{3}{\pi} \int_0^\infty d\omega [\alpha^{(E1)}(i\omega; \beta_o \mathbb{J}_o M_o)]^2.$$



- Since the van-der-Waals coefficients  $C_6$  contains the dipole matrix element to the fourth power, these dipole amplitudes must usually be calculated with high accuracy.
- **Long-range interaction of two atoms:** For two atoms  $a$  and  $b$  in spherical-symmetric ground states,  $|\psi_a\rangle = |\beta_a \mathbb{J}_a M_a\rangle$  and  $|\psi_b\rangle = |\beta_b \mathbb{J}_b M_b\rangle$ , the long-range interaction can be written as (Porsev and Derivianko, 2006)

$$V(R) = - \sum_{n=3} \frac{C_{2n}^{ab}}{R^{2n}}, \quad C_{2n}^{ab} = \frac{(2n-2)!}{2\pi} \sum_{k=1}^{n-2} \frac{1}{(2k)!(2k')!} \int_0^\infty d\omega \alpha_k^{(\text{E1})}(i\omega; \psi_a) \alpha_{k'}^{(\text{E1})}(i\omega; \psi_b), \quad k' = n - k - 1,$$

and where  $\alpha_k^{(\text{E1})}(i\omega; \psi_c)$  is the  $2^k$ -pole dynamic (electric-dipole) polarizability of each atom  $c$  in its given state  $|\alpha_c \mathbb{J}_c M_c\rangle$ .

- **Dispersion coefficients  $C_n^{(AB)}$  of mixtures:** For mixed interactions of atomic pairs  $(AB)$ , a simple rule is often applied:

$$C_n^{(AB)} = \sqrt{C_n^{(A)} C_n^{(B)}}.$$

- **Van-der-Waals coefficients  $C_{12}$  and  $C_{14}$ :** The high-order dispersion coefficients, i.e. the van der Waals coefficients  $C_{12}$  and  $C_{14}$  are usually not so easily available for most systems of interest. They can be approximated from the low-order dispersion coefficients  $C_{12}$  and  $C_6, C_8, C_{10}$  by using certain dimensionless ratios of such coefficients as discussed for instance by Douketis *et al.* (1982).

### 6.3.c. Stark shifts and ionization rates in static electric fields

#### Remarks & notations:

- **LoSurdo-Stark effect:** refers to the shifting and broadening of atomic and molecular states under the influence of an external electric field.
- (Dynamic) Stark shift:** The dynamic Stark shift is caused by the interaction of an atom with the electro-magnetic field.
- Such a dynamic (time-dependent) Stark shift can be induced by a laser pulses but also by a single-photon wave packet.
- **Keldish adiabacity parameter  $\gamma_K$ :** Different strong-field ionization mechanisms can be distinguished for different laser frequencies  $\omega_L$ , (electric) field strength  $E_o = |\mathbf{E}_o|$  as well as atomic ionization potentials  $I_p$ . They are often distinguished by the **Keldish adiabacity parameter**,  $\gamma_K = \omega \sqrt{2I_p}/E_o$ .
- **Multiphoton regime,  $\gamma_K \gg 1$ :** This regime is realized at high frequency and low or moderate field strength and gives rise to a **a power law for the ionization rate**,  $\Gamma \propto E_o^n$ .

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- **Tunnel regime,  $\gamma_K \ll 1$ :** This regime occurs at low laser frequency and high-field strength and gives rise to an **exponential dependence of  $\Gamma$  on  $E_o$** . In the tunnel regime, the ionization is often considered as a quasi-static process due to the low frequency, and the **total ionization yield is obtained as integral over the time-dependent field strength**:

$$A = 1 - \exp \left[ 1 - \int_{-\infty}^{\infty} dt \Gamma(E_o(t)) \right].$$

- In the tunnel regime, atomic structure computation can provide reliable predictions for both, the Stark shifts in an (quasi-static) electric field as well as for the ionization rates.
- **Ammosov-Delone-Krainov (ADK) formula:** This formula for the ionization rate has often been used for many-electron atoms and has been extended also towards molecules.
- **Quasi-static regime:** For  $\gamma_K \ll 1$ , one often further distinguishes between the **(tunnel ionization** and the **above-barrier ionization** due to the **(semi-classical) field strength  $E^{(\text{above-barrier})} = I_p^2/4$** . From a simplified classical viewpoint, the electron can leave the atom without tunneling for  $E_o > E^{(\text{above-barrier})}$  due to the suppression of the (atomic) Coulomb potential. In practice, however, this *barrier* is not sharp since an outgoing (scattered) electron is still partly reflected even for a classical suppressed barrier.
- **(Quasi-) Static ionization rates:** Within the tunneling regime, a variety of semiempirical expressions exist for the static ionization rates  $\Gamma(E_o)$ . Many of them are based on Landau's seminal formula for the ionization rate of the hydrogen atom in its ground state, such as the ADK rule.

### 6.3.d. Dressed atomic (Floquet) states and quasi-energies in slowly varying laser fields

#### Remarks & notations:

- **Dressed atomic states:** In a laser pulse of moderately varying intensity, the dressed states follow adiabatically the temporal evolution of the intensity during the passage of the pulse. The slow variation of the intensity then results into fairly slow changes in the (coefficients of the) superposition of the dressed states when compared with the optical period of the laser.
- **Quasi-stationary wave function of dressed states:** For sufficiently slow variations of the laser intensity, each dressed state can be written in the **Floquet form, i.e. by means of a space-independent complex exponential *multiplied by a time-harmonic function that oscillates with the same (optical) period as the incident field.***

- **Floquet equations:** If the incident field is expanded in terms of a Fourier series, the time-dependent Schrödinger equation can be decomposed and re-written as a system of time-independent equations, the (so-called) **Floquet equations**. For multiphoton ionization, this system should be solved together with the radiation (Siegert) boundary conditions.
- The Floquet ansatz generally enables one to adopt the (time-independent) **dressed-state approach to pulses with a sufficiently slowly varying field intensity**, and where the time-dependence is captured by slow changes in the coefficients of the superposition.
- **The (complex) eigenvalues of the Floquet equations provide both, the quasienergies of the dressed states and the resonance poles of the scattering amplitudes for laser-assisted scattering.**
- The Floquet approximation is typically adequate for infrared laser pulses as well as for atoms in either the ground or some Rydberg level. It has been found inappropriate for UV laser frequencies and for low-lying excited levels because of avoided crossings between the dressed levels and large non-ponderomotive ac Stark shifts at such high frequencies.
- The computation of the widths of quasi-energies often requires to go beyond first-order perturbation theory.
- **Quasi-dressed states of atoms in a constant dc-field of constant intensity and frequency:** The dressed states of atoms in a field of constant intensity and frequency refer to particular quasi-stationary solutions of the time-dependent Schrödinger equation.
- In atomic physics and if the laser pulses are neither ultra-intense nor ultra-short, these pulses can be modeled as a rapidly oscillating electric field but with a well-defined and slowly varying envelope (field amplitude). For such quasi-constant dc-fields, the behaviour of atoms can be described in terms of **superpositions of dressed bound states that decay under the emission of photoelectrons**.

### Neutral atoms in direct-current (dc) fields:

- **Complex-eigenvalue Schrödinger equation:** The system "neutral atom + dc field" defines a rather complicated many-electron problem. Because of the (additional) time-dependent field, the discrete spectrum of the atom changes into a resonance spectrum whose states are obtained as solutions of a complex-eigenvalue Schrödinger equation

$$(H - z_o)\psi = 0, \quad H = H^{(\text{atom})} + \sum_i \mathbf{F} \cdot \mathbf{r}_i, \quad z_o = E_o + \Delta(F) - \frac{i}{2} \Gamma(F).$$

- **Solution of the complex eigenvalue equation:** This equation can be solved in terms of square-integrable function spaces, if the transformation  $\rho = r e^{i\theta}$  is applied.

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- This coordinate rotation for  $\rho$  has been applied to the Hamiltonian operators for the computation of the LoSurdo-Stark ground-state resonance of hydrogen atoms, and the solution was achieved by **direct diagonalization of the non-Hermitian Hamiltonian on a large square-integrable basis set**. However, such a direct diagonalization of the the non-Hermitian Hamiltonian matrix is quite unrealistic for the decaying states of many-electron atoms and molecules.

### 6.3.e. Fano profiles of continuum-embedded resonances

#### Remarks & notations:

- In physics, a **Fano resonance generally refers to a resonant scattering process (phenomenon)**, and which gives rise to an asymmetric line-shape due to the interference between a constant background and a resonant scattering amplitude. Fano-type resonances can be found across many areas of physics and engineering because resonant scattering is a general wave phenomenon.
- In atomic physics, the background amplitude often arise from the direct (ionization) amplitude to the continuum, while the resonant amplitude arise from the (indirect) coupling of the atom to a discrete, localized state that is embedded into the continuum.

### 6.3.f. Light shifts of atomic levels

#### Remarks & notations:

- **Light shifts for clock levels:** The frequency of an atomic clock can be affected by (i) the thermal bath of blackbody radiation (BBR), (ii) the electric-quadrupole coupling of the ions to the trapping field and with other residual electro-magnetic fields, (iii) by ambient magnetic fields as well as (iv) through various Stark shifts (Yudin *et al.*, 2014).
- In highly-charged ions, the gross energy intervals nominally grow  $\propto \tilde{Z}^2 \equiv (Z - N)^2$  as  $\tilde{Z}$  increases. This leads to a substantial suppression of the ac-Stark and BBR shifts that are  $\propto 1/\tilde{Z}^4$ . Moreover, the **quadrupolar moments of the bound-state electronic density scale with  $1/\tilde{Z}^2$  and, this also reduces the quadrupolar shifts.**

### 6.3.g. Frequency-dependent ac Stark shifts of atomic levels

#### Perturbative treatment of ac-Stark shifts $E_n(\omega)$ :

- **Atoms and ions in intense laser field:** A perturbative description of the ac Stark shift of atoms in fields with intensity  $I > 10^{13} \text{ W/cm}^2$  has been a subject of continuing concern.
- Although, perturbation theory is often considered inadequate for the analysis of multiphoton ionization experiments with atoms and ions, little is known about the breakdown of perturbation theory (Pan *et al.*, 1988). Nevertheless, a perturbative treatment of the many-electron atom in intense fields is still of interest for various scenarios and intensities.

### 6.3.h. Black-body radiation shifts

#### Remarks & notations:

- From radiation studies it is known that object emit radiation whose wavelength depends on its temperature.
- **Black-body radiation shifts:** More generally, the interaction with an environment at non-zero absolute temperature causes a level shift in the atom.

### 6.3.i. Hyperpolarizability

#### Remarks & notations:

- The hyperpolarizability of atoms refers to a nonlinear-optical property of matter and is closely related to the second-order electric susceptibility (per unit volume).
- Any accurate calculation of static hyperpolarizabilities is still a great challenge for atomic theory, and this even applies to simple atoms, such as lithium with only three electrons. As reported in the literature, both the choice of the basis set (grid) as well as the electronic correlations can have a large effects upon the (accuracy of the) computed hyperpolarizabilities

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- **Hyperpolarizability of closed-shell atoms:** For a neutral, closed-shell atom in its ground state  $|\beta^1S_0\rangle$ , the energy (level) shift in the presence of a homogeneous static electric field of strength  $F$  is given by (Kassimi and Thakkar, 1994)

$$E(F; \beta\mathbb{J}) - E_o = -\alpha \frac{F^2}{2!} - \gamma_2 \frac{F^4}{4!} - \gamma_4 \frac{F^6}{6!} - \gamma_6 \frac{F^8}{8!} - \dots,$$

where  $\alpha$  is the (static) dipole polarizability and  $\gamma_k$ ,  $k = 2, 4, \dots$  the hyperpolarizability parameters of the atom.

- **Induced dipole moment  $\mu$ :** In the same field static field  $F$ , the dipole moment is then given by

$$\mu(F; \beta\mathbb{J}) = \alpha F + \gamma_2 \frac{F^3}{3!} + \gamma_4 \frac{F^5}{5!} + \dots$$

- **Finite-field procedure:** In the computational procedure, the polarizabilities and hyperpolarizabilities are obtained from energy and/or dipole moment calculations of the atom in the presence of an electric fields of various strengths, and by using truncations of the given expansion above.

## 6.4. Other topics closely related to atomic properties

### 6.4.a. Laser cooling, precision spectroscopy and quantum control

#### Quantum logic spectroscopy:

- **Quantum logic spectroscopy:** Quantum logic is typically implemented by the laser-induced coupling of the internal and external degrees of freedom of two ions, namely the **spectroscopy** and the **logic** ion. By making use of appropriate laser pulses, the information about the internal state of the spectroscopic ion can then be transferred from the spectroscopic to the logic ion.
- Such a quantum-logical transfer scheme can be realized with two ions in a linear Paul trap.
- **Quantum-logic techniques:** Many limitations of precision spectroscopy can be overcome if simple quantum-logic techniques are used and **if the spectroscopic and logic ion is controlled together**. In this technique, especially, the readily-controllable logic ion provides sympathetic cooling, internal state preparation and detection of the ion with an interesting spectroscopic transition.
- Simply speaking, the logic ion can be viewed as both, a remote control and a sensitive quantum sensor for the internal and external state of the spectroscopy ion.
- **Quantum logic can also be used to laser-cool molecules to near their rotational and vibrational ground state** by avoiding spontaneous emission of photons from the molecule.

#### Laser cooling and precision spectroscopy:

- **Laser-cooling techniques:** The recent development of laser-cooling techniques made cold and ultra-cold atomic samples available for optical spectroscopy. This development nowadays allowed for long interrogation times and significantly reduced Doppler shifts and gave eventually rise to an unprecedented uncertainty in atomic clocks, exceeding one part in  $10^{-15}$ .
- In 1995, a controlled-NOT (CNOT) gate between a trapped atomic ion and its motional mode in the trap was first experimentally realized. Since then, the **coherent control over the internal and external degrees of freedom in an atomic system** has been demonstrated.
- In quantum information processing with trapped ions, precision spectroscopy has led to new applications, such as teleportation, error-correction as well as the implementation of the Deutsch-Jozsa and semi-classical Fourier transformation algorithms.
- **Internal vs. motional degrees of freedom:** Laser spectroscopy helped develop methods to simultaneously **control both, the internal electronic and motional degrees of freedom of individual atoms, molecules and singly-charged ions in traps with high accuracy**.

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- Most laser spectroscopy with clouds of trapped ions have explored a rather small class of atoms and atomic ions, including hydrogen, the alkali and alkaline-earth atoms and ions as well as few other species (Schmöger *et al.*, 2015). These atoms and ions have electronic transitions that are suited for laser cooling and that help localize the atoms for precision spectroscopy.
- **Coulomb crystals:** A signature for trapped and well-localized ions is the formation of Coulomb crystals, i.e. **spatially ordered ensembles of mutually repelling ions that are confined in a common external trapping potential**.
- Coulomb crystals form when the thermal kinetic energy becomes (much) smaller than the ion-ion electrostatic energy.
- **Sympathetic cooling:** If atoms and ions do not have a direct cooling transition, **sympathetic cooling by use of co-trapped ions has been found a successful strategy in the past for singly and doubly charged atoms and molecules**.

### 6.4.b. Atomic clocks

#### Search and application of atomic clocks:

- An atomic clock is typically based on an electron transition of either microwave, optical or ultraviolet frequency, and which serves as frequency standard for its time-keeping element.
- Atomic clocks are the most accurate time and frequency standards; they are frequently applied as primary standards for international time distribution services, to control the wave frequency of television broadcasts as well as the global navigation satellite systems such as GPS.
- While earlier atomic clocks were based on masers at room temperature, modern clocks are operated near to absolute zero temperature by **probing an atomic fountain in a microwave-filled cavity**. The NIST-F1 atomic clock, for example, one of the national primary time and frequency standards of the United States.
- Recent developments in the **set-up and control of optical atomic clocks has improved their precision by a factor of 1000** in less than two decades.
- One particular goal in developing atomic clocks is to establish a **clock with a high sensitivity to the variation of the fundamental fine-structure constant  $\alpha$** .
- **Applications of atomic clocks:** Various applications of atomic clocks are now in reach or have already been realized by the improved precision and the high stability of these clocks. These applications include: (1) the study of many-body physics and quantum simulations; (2) relativistic geodesy; (3) very long baseline interferometry; (4) searches for the variation of the fundamental constants; (5) searches for dark matter candidates; (6) tests of the Lorentz invariance and several others.



- These and further ideas and applications, such as the use of atomic clocks for gravitational wave detection, will require even more precise clocks in the future.
- The accuracy and stability of optical atomic clocks has been advanced significantly over the past ten years. The systematic uncertainty of the Sr optical lattice clock has been reduced to  $2.1 \times 10^{-18}$  in (so-called) fractional frequency units. A similar small (systematic) uncertainty of  $3.2 \times 10^{-18}$  was reported also for a single-trapped ion atomic clock, based on an electric-octupole transition in  $\text{Yb}^+$ .
- **Application of atomic clocks:** Atomic clocks with even improved accuracy will find their application in many fields of science and technology, such as in relativistic geodesy, (very) long baseline interferometry, gravitational wave detection tests of the Lorentz invariance, the search for variations of the fundamental constants or the re-definition of the second. Further improvement of clock precision is needed for several of these applications.

### Types of optical atomic clocks:

- There are currently **two types of optical atomic clocks under development:** (i) based on neutral atoms in optical lattices or (ii) based on a **single trapped ion**. For both types, a similar uncertainty have been reached, namely:  $2.1 \times 10^{-18}$  for a Sr neutral atom clock (Nicholson *et al.*, 2015) and  $3.2 \times 10^{-18}$  for a  $\text{Yb}^+$  trapped ion clock (Huntemann *et al.*, 2016), that have been operated on a particular electric-octupole (E3) transition.
- **Proposals for a Yb atomic clock:** The  $4f^{14}6s^2\ ^1S_0 - 4f^{14}6s6p\ ^3P_0$  transition in neutral Yb can be induced by hyperfine mixing in isotopes with nuclear spin  $I \neq 0$  and has served already as some frequency standard. However, neutral Yb has still another and very interesting  $4f^{14}6s6p\ ^3P_0 - 4f^{13}6s^25d\ (J=2)$  electric-quadrupole (E2) transition at an easily accessible wavelength of 1695 nm, which might be suitable as well for the development of another frequency standard in this atom (Safronova *et al.*, 2018).
- **$^{229\text{m}}\text{Th}$  clock:** Thorium has one exceptional nuclear state that is known for the last 40 years and that has a very low energy of presumably below 10 eV. The currently assumed excitation energy of this isomeric state is  $7.8 \pm 0.5$  eV, or correspondingly  $159 \pm 11$  nm or  $\sim 1900$  THz. **This isomeric state conceptually allows for a direct laser excitation of the nucleus by using solid-state laser technology.** This state has therefore been **proposed for the development of a nuclear clock of extremely high stability** owing to the expected high resilience against external influences and its radiative lifetime in the range of minutes to hours (von der Wense *et al.*, 2017).

### 6.4.c. Atomic partition functions

#### Motivation:

- The atomic partition function  $U(T)$  includes a summation over all bound levels of the atom and is defined as

$$U(T) = \sum_j g_j \exp\left(-\frac{E_i}{kT}\right).$$

### 6.4.d. Atom-atom and atom-ion interaction potentials

#### Motivation:

- **Interatomic forces and potentials:** Accurate interatomic potentials are of fundamental importance for understanding and modelling the static and dynamic properties of gases, liquids or even solids.
- In particular, a good **knowledge of the interatomic forces in (closed-shell) molecules has been found a key for understanding the behavior of various systems**, such as the geometry and stability of molecular solids, the properties of liquids and their phase transitions as well as collisions between molecules in gases or in molecular beams.
- While the theory of interatomic forces has been developed over the last four decades, fast and reasonably reliable predictions are still often not easily available, especially if a whole range of intermolecular separations need to be considered.
- While the scattering of charged particles has already a rather long history, collisions of ions and neutral particles at ultracold temperatures has emerged only recently as an exciting (new) field in AMO physics.
- **Ion-atom interactions at cold temperatures:** At cold temperatures of just a few Kelvin or even lower, ion-atom interactions are often affected by (i) rapid energy variations in the potentials as induced by long-range polarization contributions, (ii) the generally large number of contributing partial waves and (iii) a rather sensitive dependence of the interatomic interactions upon the details of the short-range potential. Further difficulties may arise from (iv) other weak interactions, such as hyperfine interactions and other.
- The standard numerical methods for dealing with intermolecular forces are often not only inefficient in addressing these difficulties but they sometimes even miss important physics, such as extremely narrow resonances.

- **Intermolecular interaction:** The (inter-molecular) interaction energies are often partitioned into the SCF parts of the individual atoms and the correlation energy of the molecule as a whole. For atoms, the SCF computation can be carried out quite easily, while the correlation energies need to be obtained semi-empirically by means of standard long range multipolar expansion (Douketis *et al.*, 1982).
- **Ion transport coefficients in a gaseous environment:** These coefficients can often be derived from the ion-neutral interaction potentials and such computations are now routine for atomic ion-atom systems.
- The (attractive) interatomic interaction is **dominated by correlation or dispersion effects at large internuclear distances  $R$** , where perturbation theory gives rise to a series expansion  $R^{-2n}$ ,  $n \geq 3$ . However, this perturbation series is known to be asymptotically divergent and is best truncated at its smallest term, following the work of Dalgarno and Lewis.
- Because of these difficulties, many accurate interatomic potentials have been determined experimentally by combining data from beam scattering experiments, virial and transport coefficient measurements as well as liquid and solid state properties. In practise, however, the **interatomic potentials are still best probed in the intermediate regions**.

### Gordon-Kim (1972) theory for atom-atom interactions:

- Gordon and Kim (1972) developed a model to calculate the interactions between closed-shell atoms, ions and molecules. In this model, the **electron density of the system is taken as the sum of the two separate densities**, and the non-Coulombic part of the interaction potential is calculated from this density by using energy expressions from the free-electron gas.
- **Assumptions of the Gordon-Kim theory:** Three basic assumptions were made by Gordon and Kim (1972):
  - i) No rearrangement or distortion of the separate atomic densities occurs when the two atoms are brought together and, hence, **the total electron density is just the sum of the two (spherically symmetric) atomic densities**. This assumption neglects of course the rearrangement of the electron density and therefore restricts the theory to systems with a strong chemical bond, including covalent chemical bonds and to interatomic distances that are smaller than about half of the equilibrium separation.
  - ii) The interatomic interaction can be evaluated from the combined atomic density, including the Coulomb interactions between all charges.
  - iii) Gordon and Kim applied originally Hartree-Fock wave-functions in order to obtain the electron densities of the separate atoms, although other, and perhaps more accurate, wave functions should work as well.
- **Total interaction of the two (atomic) charge distributions:** The total energy includes the **contributions from the kinetic, exchange and correlation energies as well as the (direct) Coulomb interactions energy**. Below, we briefly summarize how these contributions can be calculated by just using atomic wave functions.

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- Apart from the Coulomb interaction of the two (atomic) charge distributions, the kinetic, exchange and correlation energy contributions to the total interaction energy are obtained from the electron density and the free-electron gas approximation.
- The regions around the nuclei (in which the density varies rather rapidly) do overall not contribute much to the interaction energy, even if they strongly affect the total energy. These contributions to the total energy cancel out if the atomic energies are subtracted. Only the outer regions of the atoms, in which the atomic densities overlap, contribute significantly to the interaction energy. For these reasons, the uniform electron gas model should be reasonably accurate (Gordon and Kim, 1972).
- The rather simple Gordon-Kim potential have been used, for example, in the treatment of ionic compounds such as alkali-halide and alkaline-earth-dihalide molecules in order to predict (surprisingly) accurate molecular properties (Kim and Gordon, 1974). In this work, Gordon and Kim made the following two basic assumptions in addition to those assumptions already made for their pair potentials: (a) the systems are made of free ions; (b) the interactions of ions are pairwise additive, i.e. the total interaction is the sum of all the pair interactions, neglecting the many-body interactions.

### Perturbative methods:

- Different regions of the molecular potential curves have been investigated by different theoretical methods. Perturbation theory has been applied, in particular, to describe the long-range attraction, and where the first- and second-order perturbation energies are often approximated by just the first few terms in a multipole-series expansion of the Coulomb interaction between the molecules.
- **Failure of perturbation theory:** However, the perturbation method fails to describe the potential at short and intermediate distances since it: (i) does not include exchange interactions between the electrons, (ii) diverges for strong interatomic interactions, and (iii) because a multipole expansion generally fails when the electron distributions overlap each other.
- At intermediate distances, i.e. in the region of the minimum of the intermolecular potential, only SCF computations typically predict the potentials reasonably well if the molecule does not form strong chemical bonds.

## 6.4.e. Dispersive interactions in liquid and solids

### Dispersive interactions in molecular-dynamics and Monte-Carlo simulations:

- **Dispersive interactions:** Dispersive forces and interactions in media often refer to contributions to the intermolecular interactions that arise from the polarization of one the atoms in the (fluctuating) instantaneous multipole field of other atoms. Therefore, dispersive interactions exist between each pair, triple, etc. of atoms or ions.

- **Application of dispersive interactions:** Apart from inert-gas or crystallized solids of various organic molecules, the dispersive interactions are small, though not negligible. Reasonable estimates of the magnitude of the dispersion interaction are therefore necessary for many solid-state studies, such as the cohesion of matter, the behaviour of ionic crystals, for molecular liquids, the adsorption on surfaces or in porous media (Pellenq and Nicholson, 1998).
- In Monte-Carlo and molecular-dynamics simulations, for example, a reasonable accurate potential energy function is required and need to be based on quantum-mechanical estimates, even if the dynamics of the particles is described classically.
- For dense-phase and solid-state systems, full quantum mechanical calculations of the interaction energy are generally not feasible and, hence, require to describe the interactions by means of inter-molecular or inter-ionic potentials; cf. section 6.4.d.
- Expressions for the dispersion coefficients can be derived from perturbation theory; however, the theory of inter-molecular interaction also shows clearly that the dispersive interactions in media are generally not pair-additive. Apart from the two-particle interactions, the three-body terms are known to play a significant role in the thermodynamics of condensed phases.

### Two-body dispersive interactions:

- **Long-range dispersive interactions:** If we consider two interacting neutral atoms, time-dependent perturbation theory gives rise to a long-range dispersive interaction due to charge fluctuations of the interacting species.
- **Multipole expansion:** For two spherically symmetric atoms or molecules  $A$  and  $B$ , that are separated by a distance  $R$ , a multipole expansion is given by (Pellenq and Nicholson, 1998)

$$V^{(\text{dispersion, AB})}(R) = - \left[ \frac{C_6}{R_6} + \frac{C_8}{R_8} + \frac{C_{10}}{R_{10}} + \dots \right].$$

- In this expansion, the dispersion coefficient  $C_6$  describes the interaction between two instantaneous dipoles,  $C_8$  the interaction between a quadrupole and a dipole and  $C_{10}$  both, the interaction between an octopole and a dipole as well as between two quadrupoles.

### Three-body dispersive interactions:

- In the multipole expansion, the terms with triplets of species  $A$ ,  $B$  and  $C$  are obtained from perturbation theory at the third and fourth orders. For such a triplet, the total dispersion energy is however non-additive:

$$V^{(\text{dispersion})} = V^{(\text{dispersion, AB})} + V^{(\text{dispersion, BC})} + V^{(\text{dispersion, CA})} + V^{(\text{dispersion, ABC})} + \dots$$

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where  $V^{(\text{dispersion, ABC})}$  represents a sum of several terms describing the three-body interaction.

### 6.4.f. Transport coefficients for ion mobility and diffusion in gases

#### Transport coefficients in gases:

- **Ion-mobility spectrometry (IMS):** The IMS is an analytical technique that is used to separate and identify ionized molecules in the gas phase, based on their mobility within a carrier buffer gas.
- IMS has been heavily employed for military or security purposes, such as detecting drugs and explosives, but also for the analysis of both small and large biomolecules.
- **Ion mobility  $K$ :** The ion mobility is defined as the proportionality factor between the drift velocity  $v_d$  of the ions and the applied electric field  $E$ :  $v_d = K E$ . These ion mobilities are usually reported as **reduced (ion) mobilities for a standard gas density  $n_0$** , i.e. for standard temperature  $T = 273$  K and pressure  $p_o = 1013$  hPa.
- Ion mobilities also help in analysis of collision broadening in ion cyclotron resonance, the study of ion-molecule reactions as a function of electric field strength in flow-drift tubes, and especially the use of mobility data to test or determine ion-neutral interaction potentials.
- Since the reduced ion mobility just refers to the standard density, it is still temperature dependent  $K_o = K \frac{n}{n_o} = K \frac{T_o}{T} \frac{p}{p_o}$ .
- **Kinetic theory of ion mobility in neutral gases:** Viehland and Mason (1978) present the first rigorous kinetic theory of ion mobility in neutral gases which is valid for electric fields of arbitrary strength and without restriction to the ion-neutral mass ratio or interaction potential. The results have proved useful in a number of applications: the calculation of low-field ion mobilities in high-temperature gases, if measurements as function of electric field strength were made before in low-temperature gases.

### 6.4.g. Polarizability and optical absorbance of nanoparticles

#### Optical properties of nanoparticles:

- Noble-metal nanoparticles with a size, that is comparable to the wavelength of the incident light, show characteristic colors owing to the strong absorption and scattering of light in the visible region.
- **Localized surface plasmon resonance (LSPR).** The LSPR are often explained as collective oscillation of free electrons in the nanoparticle that are induced by the electro-magnetic waves. Therefore, the **color of the absorbed light can be tuned by the size, shape, materials as well as the surrounding environment of the nanoparticles**, and with applications in chemistry, physics and biomedical fields.
- Nanoparticles with a size below 10 nm have an extremely high surface-to-volume ratios and, thus, substantially different optical properties when compared to large nanoparticles or the bulk material. These optical properties also depend on the atomistic structure.

#### Methods for calculating the polarizability of nanoparticles:

- The polarizability of a discretized medium is in general proportional to its volume.
- **Polarizability of the i-th atom  $\alpha_i$ :** In a given nanoparticle, the polarizability of the i-th atom  $\alpha_i$  can be approximated by multiplying the polarizability of the free atom  $\alpha^{(\text{free})}$  with the relative volume (ratio), i.e. by dividing the effective volume  $V^{(\text{eff})}$  by the volume of the free atom  $V^{(\text{free})}$ ,

$$\alpha_i = \frac{V^{(\text{eff})}}{V^{(\text{free})}} \alpha^{(\text{free})} = \frac{V^{(\text{eff})}}{V^{(\text{free})}} \frac{e^2}{m} \sum_n \frac{f_n}{\omega_n^2 - \omega^2 - i\omega\delta}$$

where  $\omega_n$  is the n-th absorption frequency,  $f_n$  its oscillator strength and where the frequency shift  $\delta$  depends on the coordination factor within the nanoparticle.

- **Optical absorbance of nanoparticles:** The optical absorbance of the nanoparticle is obtained from the total polarizability by

$$\sigma^{(\text{absorption})} = \frac{4\pi\omega}{c} \Im \{ \alpha^{(\text{total})}(\omega) \} = \frac{4\pi\omega}{c} \Im \left\{ \sum_i^N \alpha_i^{(\text{SCF})}(\omega) \right\}$$

### 6.4.h. Laser-produced plasma

#### Properties & applications of laser-produced plasma (LPP):

- Laser-produced plasmas (LPP) have attracted recent interest because of their use in developing laboratory ion sources and pulsed light sources at short wavelengths.
- In LPP, hot electrons can subsequently ionize the much cooler atoms, forming an x-ray-emitting, high-temperature plasma spark at the surface of the solid (Murnane *et al.*, 1991).
- The emitted x-rays from such a laser-produced plasma are generally incoherent but have a high brightness because of the small size, short lifetime and the high temperature of the radiating plasma.
- LPP have been proposed as source of both, x-rays and multiply-charged ions. Whereas x-ray sources aim for a high radiation loss from the plasma due to the generation of highly stripped atoms, such higher charge states of ions are mainly produced by a high plasma temperature and a low radiation level (loss).
- To determine the properties of a LPP, such as temperature, density or its radiation emission, one needs a proper model for the coupling of the laser light with the heavy element (solid-state) target.
- LPP will find applications in extreme ultraviolet (EUV) lithography, EUV metrology as well as in the modification of surfaces.
- LPP sources have been developed for x-rays in the water window, though with limited conversion efficiency and/or brightness so far. These plasmas can be used for novel microscope designs as well as for cell tomography.
- **Scaling laws for LPP:** Scaling laws predict that brighter and more efficient x-ray sources will be obtained if (more) intense laser pulses are employed. These sources can then be used for time-resolved x-ray scattering studies and for the development of x-ray lasers.
- **Scaling laws for LPP:** Simple arguments can be used to derive scaling laws predict the response of the LPP if the laser parameters are varied. If any hydrodynamic expansion is neglected during the interaction with the laser pulse, we can equate the incident heating flux to the cooling flux that occurs due to classical thermal conduction.
- Scaling laws can be used also to predict the radiation output from a LPP. For a short pulse of about 100 fs, the heat can penetrate for about 50 nm, and the hot radiating region can be considered as optically thin (Murnane *et al.*, 1991). Therefore, the Planck mean-free path  $X$  of the radiation is much longer than the characteristic length  $L$ , and the plasma radiates as a diluted black body with power  $P_r = A \sigma T^4 L / \lambda$ , and where  $\sigma$  is the Stefan-Boltzmann constant and  $A$  the radiating area.



### Ultrafast x-ray pulses from laser-produced plasma:

- A high-temperature plasma is created when an intense laser pulse is focused onto the surface of a solid (Murnane *et al.*, 1991). From such a plasma, an **ultrafast pulse of x-ray radiation is emitted when the laser pulse length is less than about a picosecond**.
- X-ray pulses with a duration of a picosecond or less have been observed, more than an order of magnitude shorter than those produced by any other x-ray source.
- A direct comparison of fast x-ray pulses from LPP with other x-ray sources is difficult; in general, however, LPP yield x-rays with much higher peak power but lower average power than synchrotrons and other high-average power sources such as rotating anode tubes.

### Ionization models for LPP:

- **Collisional-radiative (CR) ionization model:** A CR model can be applied to LPP with temperatures above a few tens of eV. At these temperatures, the average ionization stage  $A^{q+}$  of a given element is dominated by the electron plasma temperature due to electron and photon-impact excitation and ionization processes.
- Indeed, **two main types of excitation and ionization processes occur in LPP**: collisional and radiative processes, and where the corresponding de-excitation and recombination refer to the same types of processes. The ionization-recombination equations for both processes are:

$$A^{(q-1)+} + e^- \longleftrightarrow A^{q+} + e_1^- + e_2^-, \quad A^{(q-1)+} + \hbar\omega \longleftrightarrow A^{q+} + e_p^-.$$

- **Ionization-recombination equation:** In these ionization-recombination equations, one often only considers collisions with (free) electrons since these collisions are much more efficient than those with heavier particles. Obviously, these equations also neglect autoionization and its (inverse) dielectronic recombination.

### Deflagration models for LPP:

- **Deflagration from the dictionary:** Deflagration refers to an explosion in which the speed of burning is lower than the speed of sound in the surroundings or, in other words, a deflagration is a fire in which a flame travels rapidly, but at subsonic speed, through a gas.
- The radiative deflagration (model) describes the rapid interaction of the incident laser radiation with a very thin layer near to the target surface. In this layer, the electron density reaches a cut-off  $n^{(\text{cut-off})} = 10^{21} / \lambda^2 [\mu\text{m}]$ . This rapid interaction also leads to a hot plasma that expands into the vacuum and which *vice versa*, results in a shock wave into the solid owing to momentum conservation.

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- The properties of the hot plasma in the (thin) interaction layer can then be determined mainly from the power density in the focal zone and the wavelength of the incident radiation.

### 6.4.i. Plasma diagnostics

#### Interpretation of line spectra:

- More often than not, any detailed interpretation of the line spectra from laboratory plasmas and astrophysical sources has been found a great challenge since it requires one to go often beyond a local thermodynamical equilibrium (LTE) approach. If no equilibrium is reached, the spatially-dependent photon emission and absorption need to be taken into account in the solution of the rate equations throughout the emitting medium.
- **Laboratory *versus* astrophysical plasma:** Laboratory plasmas can differ from astrophysical sources in various ways:
  - Laboratory plasmas can be short-lived so that they do not reach statistical equilibrium.
  - Even if often much larger densities occur in laboratory plasma, when compared to typical astrophysical sources, still not all lines are optically thick and, in particular, the re-bound and free-free transitions (continuum radiation) are usually optically thin.
  - In laboratory plasma, most atomic levels are often collisionally rather than radiatively dominated.
  - The (large) density of laboratory plasma can dramatically affect the line profiles, and this should be taken into account in the radiative-transfer calculations.
- However, in order to determine the relevant physical parameters that describe both, the astrophysical and laboratory plasmas, i.e. its electron temperature, density distribution, ion and element abundances, etc., we need to compare the observed data to a theoretical model spectra.
- In general, any collision between particles cause energy exchange and changes in the state distribution changes.
- **Non-local thermodynamic equilibrium (NLTE):** NLTE exists in a wide variety of astrophysical and laboratory-created plasmas. Examples of NLTE astronomical plasmas are the stellar corona, interstellar nebulae and some other low-density ionized plasmas. In the laboratory, NLTE exists in laser-produced plasmas, tokamaks and Z-pinch based experiments.
- X-ray satellites, such as CHANDRA and XMM-NEWTON have provided large amounts of high-resolution spectra from astronomical objects, and with many of them in NLTE.

### Diagnostics of astrophysical plasma:

- Hot plasmas are present in the universe in a variety of astrophysical objects, from stellar coronae to the birth of stars and up to the intergalactic medium in clusters of galaxies. The high-quality data, that have been obtained during the last few decades from space observatories, like EINSTEIN, ROSAT, ASCA, and SAX, enabled one to study in great detail the physical condition of several astrophysical plasmas.
- **Solar corona plasma:** The solar corona plasma is low-density plasma and is believed to be dominated by the spontaneous decay and radiative recombination. For these two processes, the corresponding decay rates are much higher than the collisional decay and three-body recombination and, hence, the plasma is not in a local thermodynamic equilibrium and cannot be described by Saha or Boltzmann equations.
- **Static stellar atmospheres:** In the standard models of static stellar atmospheres, the atomic (or ionic) populations, the radiation field, the temperature and density are calculated **self-consistently by means of rate equations and radiative transfer equations but owing to certain constraints, such as statistical equilibrium, radiative equilibrium or a hydrostatic equilibrium.**

### Diagnostics of laboratory plasma:

- For ITER, various concepts have been developed to reduce the peak heat at the plasma-facing components of the tokamak. Indeed, much of the current experimental and theoretical work in fusion research focus on the question of how the heat from the plasma need to be transferred to the walls of the main and divertor chambers.
- **Heat transfer process:** To analyse this heat transfer and to enroll the relative importance of the primary atomic processes, simple quasi-analytic models and refined rate coefficients and cross sections have been utilized. Important heat transfer processes refer to the radiation losses through bremsstrahlung, impurity radiation losses from the plasma edge, the charge exchange and radiation losses from hydrogen in outer layer as well as to radiation losses due to impurities in the divertor plasma.
- In large fusion devices such as ITER, most of the heating power strikes the divertor plate and results in a high peak-heat loads of all plasma-facing components.
- The ITER divertor has been designed in order to enhance the charge exchange processes, hydrogen and impurity line radiation, ionization, or elastic collisions between the recycling gas and the plasma in the diverted plasma and, hence, to spread out the heat and momentum transfer over the device.

### Diagnostics of tungsten plasma:

- Atomic physics has been crucial for understanding the plasma-energy balance and for diagnostic development. In particular, tungsten is nowadays very frequently used on present-day tokamaks in preparation for ITER.
- Tungsten has become of high importance since it will be a major constituent of ITER plasmas. Tungsten will be applied as plasma-facing component that is able to withstand high heat loads and with a lower tritium retention than other possible materials.
- **ITER Core Imaging x-ray Spectrometer (CIXS):** The CIXS has been designed to measure the ion temperature and the motion of ITERs plasma core, based especially on the x-ray emission of neon-like  $W^{64+}$  ions (Beiersdorfer *et al.*, 2015). The emission from tungsten will be measured by extreme ultraviolet (EUV) and optical spectrometers in order to determine its concentration in the plasma and to assess power loss and the sputtering rate of tungsten.

### Diagnostics of lanthanide ions:

- The radiative properties of the lanthanide ions have numerous applications in solid-state laser materials, photonics and the lighting industry. Radiative transitions of lanthanide ions were observed in doped crystals and in solution with small shifts relative to those of the free ions.
- Suzuki *et al.* (2018) have systematically observed extreme ultraviolet (EUV) spectra from the highly-charged ions of nine lanthanide elements with nuclear charge  $Z = 60..70$  in optically thin plasmas as produced in the Large Helical Device (LHD). The wavelengths of the main peaks in the quasi-continuum features were found to agree well previously measured singlet transitions of Pd-like ions as well as with the discrete spectral lines from Cu-like and Ag-like ions.
- The bright emission spectra from highly-charged ions of lanthanide elements ( $Z = 57 - 71$ ) are known to appear in the extreme ultraviolet (EUV) or soft x-ray wavelength range due to  $n = 4 - 4$  transitions of lanthanide ions in the wavelength range of 5-12 nm. These emission spectra are typically strongly affected by relativistic effects and multi-electron correlations. Likely, these emission spectra will play an important role in application to the next generation light sources for EUV lithography.
- **Observationos with the Hubble Space Telescope:** In astrophysics, observationos with the Hubble Space Telescope of spectra of chemically peculiar stars have proven the presence of lanthanide ions up to doubly-charged stages. More recently, the shortly delayed observation of gravitational waves and the emission of electromagnetic waves during a neutron-star merger have renewed the interest for radiative properties of higher-charged lanthanide ions.

### 6.4.j. Radial distribution functions for plasma and liquid models

#### General remarks:

- **Distribution functions for plasma:** The radial distribution function for a plasma is often required and need to be calculated. In these computations, quantum effects have to be taken into account for small inter-particle separations, either by using a modified path-integral technique or, equivalently, by a **summation over states, together with the assumption that the charged particles interact via a shielded Coulomb potential**.
- General expressions for the radial distribution function in an electron-ion plasma outside of an thermal equilibrium have been obtained in the literature in first-order of the plasma parameters.
- **Distribution functions for fluids:** Radial distribution function and static structure factors have been computed at several places in the literature. The radial distribution function can be utilized to characterize a fluid, for instance in terms of the pronounced maxima at some inter-particle distance, and followed by successive minima and maxima with reduced amplitudes.

#### Quantal hypernetted-chain approximation:

- **Hypernetted-chain equation:** In statistical mechanics, the **hypernetted-chain equation** arises as a closure relation in order to solve the **Ornstein-Zernike equation, i.e. a relation between the direct and total correlation functions**. The hypernetted-chain equation is commonly applied in fluid theory in order to obtain, e.g. expressions for the **radial distribution function**.
- The hypernetted-chain approximation is an integral-equation method from statistical physics that arises from the theory of liquids.
- A quantal version of the hypernetted-chain equation (QHNC) was derived in the literature by using the density-functional method, and extended to treat a liquid metal and a plasma as an ion-electron mixture.

### 6.4.k. Average-atom model for warm-dense matter

#### Warm-dense matter:

- **Warm-dense matter (WDM):** This term generally refers to some state of matter under plasma conditions, and which can range from condensed matter to **weakly-coupled plasmas with typical temperatures from a few to a few hundred eV as well as with densities from a few hundredths to about a hundred times of the solid density**.

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- In modeling WDM, great challenges arise from the partial ionization of the plasma, the degeneracy of the electron states, the bound-state level shifts, the pressure ionization in plasma as well as the strong coupling of the ion-ion pairs that must be taken into account self-consistently.
- **Equation of state:** A good understanding of the properties and behaviour of WDM, such as the equation of state, its radiation opacity, or transport properties, has been found important, for modeling for example astrophysical objects or inertial confinement fusion experiments (Hou *et al.*, 2015).
- **Electronic structure of impurities:** Studies on the electronic structure of impurities within an electron gas have been performed in the literature as function of their density and temperature. For these investigations, one need to know the **dependence of screening effects on the plasma parameters**.
- **Friedel oscillations:** At low temperatures and metallic densities, especially, the electron density often exhibits Friedel oscillations as known also from solid-state physics.
- Hou *et al.* (2015) combined the average-atom model with the **hypernetted chain approximation** (to a so-called AAHNC approximation) in order to describe the electronic and ionic structure in the WDM regime.

### Average-atom model:

- **Average-atom (AA) model:** This model divides the plasma into separate **neutral Wigner-Seitz cells with a nucleus of charge  $Z$  and  $Z$  free electrons**. Various versions of the AA model have been implemented in plasma physics for describing warm- and hot-dense plasma.
- In the AA model, each ion or atom is described within a finite (ion) sphere and by using a central-field approximation for the bound electrons. For heavy atoms and ions, a (radial) Dirac equation can be applied but where the potential  $V(r)$  need again to be calculated self-consistently.
- **Electron density in the AA model:** For an isolated atom or ion, the electron density is calculated in the AA model in order to include the level broadening of the energy levels within a plasma environment. This density is applied also for studying the temperature and density effects upon the electron distributions within a statistical way.
- **Radial potential in the AA model:** (Hou *et al.*, 2015)

$$V(r) = -\frac{Z}{r} + \int d^3r' \frac{\rho_b(r')}{|\mathbf{r} - \mathbf{r}'|} + V_{xc}[\rho_b(r) + \rho_{e,0}] - V_{xc}(\rho_{e,0}) - \frac{\rho_{e,0}}{\beta} \int d^3r' C_{ee}(|\mathbf{r} - \mathbf{r}'|) h_{ie}(r') - \frac{\rho_{i,0}}{\beta} \int d^3r' C_{ie}(|\mathbf{r} - \mathbf{r}'|) h_{ii}(r')$$

where the first four terms constitute the contributions from a single ion: electron-nucleus interaction, the electro-static repulsion with the other bound electrons as well as the exchange and correlation potential, which are often evaluated by using some local density approximation. The last two terms in this expression represent the interaction of the ion with the surrounding free electrons and with the other ions. Here,  $\rho_{e,0} = \rho(r_b)$  is the uniform electron density.

### Spectral emission from plasma in a non-local thermodynamic equilibrium (NLTE):

- In a non-local thermodynamic equilibrium (NLTE), in which the levels of each configuration are in LTE among each other, the population  $N_a$  of a given configuration follows as solution of a system of rate equations

$$\frac{\partial N_a}{\partial t} = -N_a \sum_b R_{ab} + \sum_b R_{ba} N_b,$$

where  $R_{ab}$  is a global transition rate that connects the configurations  $a$  and  $b$ .

- Many collisional-radiative models, that are applied for a realistic modelling of NLTE line emission, often account for individual levels explicitly for spectroscopic reasons, while the levels of other configurations are lumped together. These **superlevels** have no spectroscopic meaning but are included to describe the population kinetics.

#### 6.4.1. Equation-of-state relations for astro physics and condensed matter

##### Equation-of-state:

- **Equation-of-state relations:** More often than not, the equation-of-state refers to the pressure-density or pressure-temperature relation for matter under (more or less) extreme conditions.
- When the atomic density is rather low, we can deal with isolated atoms that are in equilibrium with the surrounding free electrons, while the statistical or Thomas-Fermi-Dirac model has been found useful for a very high density.

##### Liberman's self-consistent model:

- Liberman (1979) describes a model for condensed matter in which the ions surrounding a particular atom are replaced by a positive charge distribution which is constant outside of a sphere containing the atom and zero inside. This models enables one to separate the quantities that pertain to the atom from those of the electron gas (in which it is imbedded) and hence to derive the desired equation-of-state data.

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- The model assumes that there are sufficient electrons in the system in order to give overall electrical neutrality together with the (required) electrical neutrality inside the sphere.
- Instead of a polyhedral cell, which are surrounded by other identical cells, a **nearly equivalent spherical cell is assumed in Liberman's model, and surrounded by a uniform electron gas**. The density of this electron gas is the same as the mean density of ionic charges. Therefore, the electron gas outside will not alter the charge distribution inside the cell, if its density is correctly chosen, nor will the atom affect the electron gas.

### 6.4.m. Radiation damage of DNA by electron impact

#### Radiation damage of DNA:

- **Radiation damage:** This term often refers to the damage of biological material on rather short time scales, i.e. reactions and processes that occur within ns or even less following the interaction of a high energy quantum with a living cell. These initial events may ultimately lead to the collapse of such cells. In practice, they result in the death of the individual cells within hours or days but may affect the matter also on much longer time scales.
- **Low-energy electrons (LEE) that interact with the DNA:** The absorption of high-energetic photons by biological matter often leads to low-energy electrons (LEE) that interact with the DNA. These LEE can cause specific resonant processes and may lead eventually either to single- or double-strand breaks in DNA materials, to the damage of its molecular components and possibly even to biological apoptosis (Bacarelli *et al.*, 2011).
- Several theoretical and computational approaches have been developed for studying the molecular processes that occur in the various steps of the energy deposition by LEE.
- In contrast, an alteration of the genetic expression of DNA may cause the onset of some diseases.

#### Electron-molecule dynamics:

- To understand the electron-impact of bio-molecules, a quite sizeable computational machinery need to be developed. The multiple inelastic scattering of these electrons will result in a re-distribution in the energy content and may lead to a fragmentation as well as the excitation of various vibrational modes.



## 6.4.n. Exotic atoms and ions

### Muonic atoms and ions:

- **Muonic atom:** When a muon is captured by the nucleus, it generally forms a hydrogen-like muonic ion, although it is typically still surrounded by other atomic electrons.
- Measurements of the transition energies in muonic atoms help determine nuclear parameters, such as charge radii, quadrupole moments and magnetic hyperfine constants. One of the most precise measurements of the nuclear root-mean-square radius by means of muonic atoms refer to the radius of  $^{208}\text{Pb}$  at the 0.2 % level.
- **Muonic hydrogen:** In muonic hydrogen, the electron is replaced by a muon  $\mu$  with a 200 times larger mass. This large mass results orbits that are 200 times closer to the nucleus, when compared to the electron in regular hydrogen.
- Because of the finite extent of the proton, muons in their  $s$ -state have therefore some enhanced probability to be inside the proton, and where it sees also a slightly reduced electric charge of the proton. Therefore, the muon is less bound as larger the proton is.



## 7. Atomic processes

### 7.1. In JAC implemented processes

#### 7.1.a. Photoemission. Transition probabilities (PhotoEmission)

##### Process, notations & application:

- **Photoemission:** from an atom or ion  $A^* \longrightarrow A^{(*)} + \hbar\omega$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + \hbar\omega$
- **JAC's standard multipole amplitudes for photon emission:** Although all (one- and many-electron) electron-photon interaction matrix elements are always evaluated in *absorption* within JAC, we wish and need to retain the intuitive description of matrix elements as typically applied in quantum mechanics:  $\langle \text{final} - \text{state} | \text{operator} | \text{initial} - \text{state} \rangle$ . In JAC, we therefore introduce an explicit and independent notation for a *standard (reduced) emission as well as absorption multipole matrix element*

$$\langle \alpha_f \mathbb{J}_f || \mathbb{O}^{(\mathbb{M}, \text{emission})} || \alpha_i \mathbb{J}_i \rangle = \left\langle \alpha_i \mathbb{J}_i \left\| \sum_{k=1}^N \alpha_k a_{k,L}^p \right\| \alpha_f \mathbb{J}_f \right\rangle^* \equiv \langle \alpha_i \mathbb{J}_i || \mathbb{O}^{(\mathbb{M}, \text{absorption})} || \alpha_f \mathbb{J}_f \rangle^*$$

which retains the right order and can be obtained from `Jac.PhotoEmission.amplitude()`; cf. section 4.3.c. A multipole  $\mathbb{M} \equiv (L, p) = \text{E1, M1, E2, ...}$  hereby contains all information about its multipolarity (angular momentum)  $L$  and type *magnetic* ( $p = 0$ ) or *electric* ( $p = 1$ ).

- Using JAC: Perform an `Atomic.Computation(.., process=Jac.Radiative, processSettings=PhotoEmission.Settings(..), ..)` or call directly functions from the module `PhotoEmission`.

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- In JAC, the transition probabilities and radiative lifetimes are tabulated by default for all selected transitions  $i \rightarrow f$ .
- In JAC, the anisotropy (structure) parameters  $f_k(\alpha_i \mathbb{J}_i, \alpha_f \mathbb{J}_f)$  are calculated and tabulated if the flag `calcAnisotropy = true` is set in `PhotoEmission.Settings`.

### Transition probabilities and oscillator strengths:

- With the notation of the (standard reduced) emission multipole matrix element from above, the transition probability and oscillator strengths are given by

$$W_{i \rightarrow f} = \frac{8\pi \alpha \omega}{2J_i + 1} \sum_{\mathbb{M}} |\langle \alpha_f \mathbb{J}_f \| \mathbb{O}^{(\mathbb{M}, \text{emission})} \| \alpha_i \mathbb{J}_i \rangle|^2, \quad f_{i \rightarrow f} = \frac{2\omega}{3(2J_i + 1)} \sum_{\mathbb{M}} |\langle \alpha_f \mathbb{J}_f \| \mathbb{O}^{(\mathbb{M}, \text{emission})} \| \alpha_i \mathbb{J}_i \rangle|^2,$$

and where the sum over  $\mathbb{M} = \text{E1, M1, E2, ...}$  runs over all requested multipoles in a given computation.

### Angular distribution and anisotropy parameters of the fluorescence radiation:

- The fluorescence (radiation) from an excited atom is characterized not only by its frequency  $\hbar\omega$  but also by the **angular distribution and polarization of the emitted radiation**. Both of these properties of the emitted photons depend on the sublevel occupation of the excited atom or ion and, hence, on its reduced statistical tensors  $\mathcal{A}_{kq}(\alpha_i \mathbb{J}_i)$  which are **associated with the prior excitation process**.
- **Angular distribution for the photoexcitation of initially unpolarized but aligned atoms:** For initially unpolarized but aligned atoms, the angular distribution of the (characteristic) fluorescence radiation is given in perturbation theory by:

$$W(\vartheta) = \frac{W_o}{4\pi} \left( 1 + \sum_{k=2,4,\dots} f_k(\alpha_i \mathbb{J}_i, \alpha_f \mathbb{J}_f) \mathcal{A}_{k0}(\alpha_i \mathbb{J}_i) P_k(\cos\vartheta) \right),$$

where  $W_o$  is the total decay rate and  $\vartheta$  the angle of the photons with regard to the alignment axis (quantization axis).

- **Anisotropy (structure) parameters  $f_k(\alpha_i \mathbb{J}_i, \alpha_f \mathbb{J}_f)$ :** In the angular distribution above, the **anisotropy (or structure) parameters  $f_k(\alpha_i \mathbb{J}_i, \alpha_f \mathbb{J}_f)$  are independent of the particular excitation process of the level  $|\alpha_i \mathbb{J}_i\rangle$**  and merely reflects the electronic structure of

the ion in the two levels involved in the fluorescence

$$f_k(\alpha_i \mathbb{J}_i, \alpha_f \mathbb{J}_f) = \left[ \sum_{\mathbb{M}} |\langle \alpha_i \mathbb{J}_i || \mathbb{O}^{(\mathbb{M}, \text{emission})} || \alpha_f \mathbb{J}_f \rangle|^2 \right]^{-1} \frac{\sqrt{2J_i + 1}}{2} \sum_{\mathbb{M} \mathbb{M}'} i^{L' + p' - L - p} (-1)^{J_f + J_i + 1 + k} [L, L']^{1/2} \langle L1, L' - 1 | k0 \rangle \\ \times \left( 1 + (-1)^{L + p + L' + p' - k} \right) \begin{Bmatrix} L & L' & k \\ J_i & J_i & J_f \end{Bmatrix} \langle \alpha_f \mathbb{J}_f || \mathbb{O}^{(\mathbb{M}, \text{emission})} || \alpha_i \mathbb{J}_i \rangle \langle \alpha_f \mathbb{J}_f || \mathbb{O}^{(\mathbb{M}', \text{emission})} || \alpha_i \mathbb{J}_i \rangle^*,$$

- Of course, there is usually one multipole term that dominates the radiative decay (fluorescence) for any given pair of initial and final-bound states.

### Semi-empirical calculations of transition probabilities:

- Semi-empirical calculations of transition probabilities are often performed by using Cowan's code and, especially, the components RCN/RCN2/RCG/RCE. In a first step, Hartree-Fock calculations of radial integrals are carried out, including relativistic corrections (HFR) to generate all required atomic orbitals and average energies of configurations. The diagonalization of the Hamiltonian matrix then leads to approximate energy values and eigenvectors of levels in intermediate coupling.
- When, in addition, also experimental energies are available, RCE is run to perform an iterative least-squares fit with the aim to minimize the differences between the calculated and experimental energies, and where the radial integrals are applied as fitting parameters. The mean error of such a fit is defined by  $\Delta E = \sqrt{\sum_i (E_i^{(\text{exp})} - E_i^{(\text{calc})})^2 / (N_i - N_p)}$  where  $N_i$  is the number of experimentally known energies and  $N_p$  the number of free parameters.

### Unresolved transition arrays (UTA):

- **Unresolved transition arrays:** The UTA approach has become the method of choice in order to calculate the radiative properties of high- $Z$  elements. Within this approach, a group of lines (i.e. a so-called array) which belongs to the same pair of electron configurations are treated globally by a Gaussian distributions whose width and position are given by analytical formulae.
- The UTA approach is computationally efficient and accurate for plasma conditions where the line broadening is such, that individual line profiles merge. This approach has been further extended to a group of arrays in the so-called super-transition array method (STA) where set of arrays are treated globally.

## 7. Atomic processes

- **Resolved transition arrays (RTA):** The RTA method was proposed to overcome the shortcoming of the UTA method that just a few transitions often dominate the spectral emission distribution of the array. This RTA method improves on the UTA approach by resolving the transition array into individual lines by using random lines but without performing an explicit atomic structure calculation.
- **UTA method:** For a group of lines with known transition energies  $E_i$  and corresponding the oscillator strengths  $g_i$ , the UTA is obtained by calculating the first two moments for the array of  $N$  transitions:

$$\mu^{(1)} = \frac{\sum_i^N g_i E_i}{\sum_i g_i}, \quad \mu^{(2)} = \frac{\sum_i^N g_i E_i^2}{\sum_i g_i},$$

- These two moments can be utilized together with the variance  $\sigma^2 = \mu^{(2)} - (\mu^{(1)})^2$  in order to define the **spectral distribution of the UTA**

$$f(\mathcal{E}) = \frac{1}{\sigma \sqrt{\pi}} \exp \left[ -\frac{(\mathcal{E} - \mu^{(1)})^2}{2\sigma^2} \right].$$

- By avoiding the explicit calculation of the properties of each line, that belongs to a given array, the UTA method enables one to account for a large number of transitions for a rather low computational cost.
- Overall, however, the use of such UTA appears quite tedious, and the application of the formulas derived by Bauche-Arnould and coworkers (1985, 1988) might lose its importance as more powerful computers become available.

### Two-electron one-photon transitions (TEOP):

- Two-electron one-photon transitions were postulated almost 100 years ago and are well-known from the optical region.
- In the early 1970s, Wölfl *et al.* reported results for nickel and iron, for which high-energetic satellites arose due to two-electron one-photon transition. In these experiments, two vacancies in the K-shell were simultaneously filled by two L-shell electrons under the emission of a single photon.
- Since the interaction with the electro-magnetic field is governed by an one-electron operator, the contribution of two-electron one-photon (TEOP) transitions can usually be ignored in comparison to other, E1 allowed one-electron transitions.

### Fractional occupation number (FON) approach:

- **Fractional occupation number (FON) approach:** The FON approach involves the prescription of a single, fictitious configuration with FONs for each  $(n, \kappa)$  subshell in order to represent an entire set of physical configurations within a given model. This fictitious configuration is used to first generate a single set of bound wave functions. Second, integer occupation numbers are then combined with the (averaged) wave functions in order to describe all the physical configurations that are enumerated in a particular model (Sampson *et al.*, 2009).

### 7.1.b. Photoexcitation or photoabsorption (PhotoExcitation)

#### Process, notations & application:

- **Photoexcitation** of an atom or ion:  $A + \hbar\omega \longrightarrow A^*$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + \hbar\omega(\mathbf{k}, \lambda) \longrightarrow |\alpha_f \mathbb{J}_f\rangle$
- Using JAC: Perform an `Atomic.Computation(.., process=Jac.PhotoExc, processSettings=PhotoExcitation.Settings(..), ..)` or call directly functions from the module `PhotoExcitation`.
- In JAC, the photoexcitation cross sections are calculated by default for resonant and completely linearly-polarized plane-wave radiation.
- In JAC, the photoexcitation cross sections are calculated also for plane-wave radiation with given Stokes parameters  $(P_1, P_2, P_3)$  if the flag `calcStokes = true` are set in `PhotoExcitation.Settings` and if the Stokes parameters are given explicitly by `PhotoExcitation.Settings.stokes`.
- In JAC, the statistical tensors  $\rho_{kq}(\alpha_f \mathbb{J}_f)$  and the alignment parameters  $\mathcal{A}_{kq}(\alpha_f \mathbb{J}_f)$  are calculated and tabulated for unpolarized atoms and plane-wave photons with given Stokes parameters, if the flag `calcTensors = true` is set in `PhotoExcitation.Settings` and if the Stokes parameters are given explicitly by `PhotoExcitation.Settings.stokes`.

### Photoexcitation cross sections:

- **Total photoabsorption cross section:** For **initially unpolarized atoms** and incident photons with given (photon) density matrix  $\langle \mathbf{k}\lambda | \rho_\gamma | \mathbf{k}\lambda' \rangle$ , the total photoabsorption cross section can be most easily defined in the **limit of a zero linewidth** as

$$\begin{aligned} \sigma(\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f; \omega) &= \frac{4\pi^2\alpha}{\omega(2J_i+1)} \sum_{M_i M_f, \lambda \lambda'} \langle \mathbf{k}\lambda | \rho_\gamma | \mathbf{k}\lambda' \rangle \langle \alpha_f \mathbb{J}_f M_f | \mathbb{T}^{(\text{absorption})}(\mathbf{k}, \lambda) | \alpha_i \mathbb{J}_i M_i \rangle \\ &\quad \times \langle \alpha_f \mathbb{J}_f M_f | \mathbb{T}^{(\text{absorption})}(\mathbf{k}, \lambda') | \alpha_i \mathbb{J}_i M_i \rangle^* \delta(\omega + E_i - E_f), \end{aligned}$$

and where  $\mathbb{T}^{(\text{absorption})}(\mathbf{k}, \lambda)$  formally describes the electron-photon interaction operator for absorbing a photon  $\hbar\omega(\mathbf{k}, \lambda)$ .

- In order to allow for a **finite natural width  $\Gamma_f$  of the excited level**, we need to replace the  $\delta(\omega + E_i - E_f)$  function in the above expression by a (Lorentzian) spectral distribution

$$\begin{aligned} \sigma(\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f; \omega) &= \frac{4\pi^2\alpha}{\omega(2J_i+1)} \sum_{M_i M_f, \lambda \lambda'} \langle \mathbf{k}\lambda | \rho_\gamma | \mathbf{k}\lambda' \rangle \langle \alpha_f \mathbb{J}_f M_f | \mathbb{T}^{(\text{absorption})}(\mathbf{k}, \lambda) | \alpha_i \mathbb{J}_i M_i \rangle \\ &\quad \times \langle \alpha_f \mathbb{J}_f M_f | \mathbb{T}^{(\text{absorption})}(\mathbf{k}, \lambda') | \alpha_i \mathbb{J}_i M_i \rangle^* \frac{1}{\pi} \frac{\Gamma_f/2}{(\omega + E_i - E_f)^2 + \Gamma_f^2/4}. \end{aligned}$$

- **Photoexcitation cross sections for initially unpolarized atoms:** For incident plane-wave photons with (photon) density matrix  $(c_{\lambda\lambda'}) = \frac{1}{2} \begin{pmatrix} 1 + P_3 & P_1 - iP_2 \\ P_1 + iP_2 & 1 - P_3 \end{pmatrix}$ , which are resonant to the given transition  $|\alpha_i \mathbb{J}_i\rangle \rightarrow |\alpha_f \mathbb{J}_f\rangle$ , the total photoexcitation cross section is

$$\sigma(\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f; \omega) = \frac{2\pi^3\alpha}{\omega(2J_i+1)} \sum_{\mathbb{M}, \lambda} c_{\lambda\lambda} \left| \langle \alpha_f \mathbb{J}_f | \mathbb{O}^{(\mathbb{M}, \text{absorption})}(\lambda) | \alpha_i \mathbb{J}_i \rangle \right|^2 \times \begin{cases} \delta(E_i + \omega - E_f) \\ \frac{1}{\pi} \frac{\Gamma_f/2}{(E_i + \omega - E_f)^2 + \Gamma_f^2/4} \end{cases}.$$

- **Photoexcitation cross sections for initially unpolarized atoms and linear-polarized, circularly-polarized or unpolarized incident photons:** For resonant and unpolarized or completely linearly-polarized or circularly-polarized plane-wave photons, especially, the total photoexcita-



tion cross section is given by

$$\sigma(\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f; \omega) = \frac{8\pi^3 \alpha}{\omega (2J_i + 1)} \sum_{\mathbb{M}} |\langle \alpha_f \mathbb{J}_f || \mathbb{O}^{(\mathbb{M}, \text{absorption})}(\lambda) || \alpha_i \mathbb{J}_i \rangle|^2 \times \begin{cases} \delta(E_i + \omega - E_f) \\ \frac{1}{\pi} \frac{\Gamma_f/2}{(E_i + \omega - E_f)^2 + \Gamma_f^2/4} \end{cases}$$

### Final-level statistical tensors and alignment parameters:

- **Statistical tensors:** Instead of the (final-level) density matrix, it is often more convenient to describe the population of the magnetic sublevels in terms of (so-called) statistical tensors of the atom or ion

$$\rho_{kq}(\alpha_f \mathbb{J}_f) = \sum_{M_f M'_f} (-1)^{J_f - M'_f} \langle J_f M_f, J_f (-M'_f) | kq \rangle \langle \alpha_f \mathbb{J}_f M_f | \rho | \alpha_f \mathbb{J}_f M'_f \rangle.$$

and which are **non-zero only for**  $0 \leq k \leq 2J_f$  and  $-k \leq q \leq k$ .

- **Reduced statistical tensors or alignment parameters:** The reduced statistical tensors of the excited levels or so-called **alignment parameters** are defined by

$$\mathcal{A}_{kq}(\alpha_f \mathbb{J}_f) = \frac{\rho_{kq}(\alpha_f \mathbb{J}_f)}{\rho_{00}(\alpha_f \mathbb{J}_f)}.$$

- **Statistical tensors for the photoexcitation of initially unpolarized atoms and incident plane-wave radiation:** For incident plane-wave photons with (photon) density matrix  $(c_{\lambda\lambda'}) = \frac{1}{2} \begin{pmatrix} 1 + P_3 & P_1 - iP_2 \\ P_1 + iP_2 & 1 - P_3 \end{pmatrix}$ , which are resonant to the given transition  $|\alpha_i \mathbb{J}_i\rangle \rightarrow |\alpha_f \mathbb{J}_f\rangle$ , the reduced statistical tensors can be written as

$$\rho_{kq}(\alpha_f \mathbb{J}_f) = \dots$$

- **Statistical tensors for the photoexcitation of initially unpolarized atoms by incident plane-wave radiation with well-defined helicity  $\lambda$ :** For

incident plane-wave photons with helicity  $\lambda$ , the reduced statistical tensors can be written as (Surzhykov *et al.*, 2015)

$$\begin{aligned} \rho_{kq}(\alpha_f \mathbb{J}_f) &= \delta_{q0} C \sum_{\mathbb{M}\mathbb{M}'} i^{L+L'} [L, L]^{1/2} (i\lambda)^p (-i\lambda)^{p'} (-1)^{L'+J_f+J_i+k} \langle L\lambda, L' - \lambda | kq \rangle \begin{Bmatrix} L & L' & k \\ J_f & J_f & J_i \end{Bmatrix} \\ &\times \langle \alpha_f \mathbb{J}_f \parallel \mathbb{O}^{(\mathbb{M}, \text{absorption})} \parallel \alpha_i \mathbb{J}_i \rangle \langle \alpha_f \mathbb{J}_f \parallel \mathbb{O}^{(\mathbb{M}', \text{absorption})} \parallel \alpha_i \mathbb{J}_i \rangle^* . \end{aligned}$$

### Photoabsorption and line-shape function:

- **Absorption coefficient  $\alpha(\omega)$ :** For a beam of light that propagates in  $x$ -direction, the absorption coefficient is phenomenologically defined in terms of changes of the spectral intensity  $I(\omega)$ , cf. Hilborn (2002)

$$\frac{1}{I(\omega)} \frac{dI(\omega)}{dx} = -\alpha(\omega) .$$

The spectral intensity  $I(\omega)$  is spectroscopically a time-averaged quantity over, at least, several optical cycles.

- **Line-shape function  $g(\omega)$ :** In most practical applications, the **frequency dependence of the absorption and emission processes (coefficient) is important but is often approximated simply by means of a normalized line-shape function  $g(\omega)$  with  $\int_{-\infty}^{\infty} d\omega g(\omega) = 1$** , and where negative frequencies have no special physical meaning. Here, the use of  $-\infty$  as the lower limit in the normalization integral greatly simplifies the normalization calculation. — For atoms in motion or in collision, the line-shape function  $g(\omega)$  describes the appropriate ensemble-averaged line shape. The line-shape function  $g(\omega)$  has the dimensions of 1/angular frequency and, hence, the explicit functional form of  $g(\omega)$  and  $\tilde{g}(\lambda)$  is not equivalent.
- **Absorption cross section  $\sigma^{(\text{absorption})}(\omega)$ :** If we assume  $\alpha(\omega)$  to be proportional to the number  $N_o$  of atoms per unit volume in the lower level  $|\beta_o \mathbb{J}_o\rangle$ , which the beam intercepts, the absorption cross section is given by:

$$\alpha(\omega) = N_o \sigma^{(\text{absorption})}(\omega) .$$

- **Relation between absorption cross section and Einstein's A and B coefficients:** (Hilborn 2002)

$$\begin{aligned}\sigma^{(\text{absorption})}(\omega) &= \frac{g_f}{4 g_i} \lambda_{fi}^2 g(\omega) A_{fi}, & \sigma_o^{(\text{absorption})}(\omega) &= \frac{g_f}{4 g_i} \lambda_{fi}^2 A_{fi} \\ \sigma^{(\text{absorption})}(\omega) &= \frac{\hbar \omega}{c} B_{12}^\omega g(\omega), & \sigma_o^{(\text{absorption})} &= \frac{\hbar \omega_{21}}{c} B_{12}^\omega.\end{aligned}$$

- **Relation between line strength and Einstein A coefficient for a spontaneous  $i \rightarrow f$  emission:**

$$S_{if} = \frac{3 \epsilon_o \hbar c^3}{2 \omega_{if}^3} g_i A_{if}.$$

- **Weighted oscillator strength:** From the principle of detailed balance and in order to avoid confusion about the absorption or emission viewpoint, one often reports the weighted oscillator strength  $g f = g_1 f_{12}^{(\text{absorption})} \equiv -g_2 f_{21}^{(\text{emission})}$ .
- **Rabi frequency for on-resonance excitations:** At resonance  $\omega = \omega_{fi}$  of the incident (linearly-polarized) light field and for two non-degenerate levels  $(i, f)$ , the Rabi frequency can be expressed in terms of the electric field amplitude of the (and the transition dipole moment  $\mu_{fi}$  by

$$\Omega_R = \frac{\mu_{fi} E}{\hbar}, \quad \mu_{fi} = \langle \alpha_f \mathbb{J}_f M = 0 | \mathbb{O}^{(\mathbb{M}, \text{absorption})} | \alpha_i \mathbb{J}_i M = 0 \rangle$$

For degenerate upper and lower levels, there are generally several Rabi frequencies for each combination  $\langle \alpha_f \mathbb{J}_f M_f | \mathbb{O}^{(\mathbb{M}, \text{absorption})} | \alpha_i \mathbb{J}_i M_i \rangle$ , and this gives usually rise to a rather complicated dynamical behaviour (Hilborn 2002).

## Photoexcitation by twisted light:

- **Photoexcitation by twisted light:**

$$\hbar \omega^{(\text{twist})} + A \longrightarrow A^*$$

- **Photoexcitation of macroscopic targets:** For macroscopic targets, the sublevel-dependent photoexcitation cross sections  $\sigma_{m_f}^{(\text{twist})} = f(\vartheta_k)$  as well as the cross section ratio with regard to incoming plane-wave radiation  $\sigma^{(\text{twist})}/\sigma^{(\text{plane})}$  have been calculated.

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- **Finite, localized targets:** For finite, localized targets, the sublevel-dependent photoexcitation cross sections as well as alignment and polarization parameters all depend on the impact parameter of the target with regard to the beam axis,

$$\sigma_{m_f}^{(\text{twist})} = f(\vartheta_k, \dots, \mathbf{b}, m), \quad \frac{\sigma^{(\text{twist})}}{\sigma_{(\text{plane})}} = f(\mathbf{b}, \text{target size}), \quad \mathcal{A}_2, P_1 = f(\mathbf{b}, \text{target size}).$$

- The photoexcitation of atoms by a Bessel beam of twisted light has been investigated for low- $Z$  hydrogen-like ions. Markgraf-Scholz *et al.* (2014) have calculated especially the *partial* (subshell) cross sections and the alignment of the excited states, if the beam interacts with macroscopic hydrogen target.

### 7.1.c. Atomic photoionization (PhotoIonization)

#### Process, notations & application:

- **Photoionization** of an atom or ion:  $A + \hbar\omega \longrightarrow A^{+*} + e_p^-$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + \hbar\omega(\mathbb{M}) \longrightarrow |\alpha_f \mathbb{J}_f\rangle + e_p^-(\mathbf{p}, m_s) \quad \text{or} \quad |\alpha_i \mathbb{J}_i\rangle + \hbar\omega(\mathbb{M}) \longrightarrow |\alpha_f \mathbb{J}_f\rangle + e_p^-(\varepsilon \kappa)$
- Using JAC: Perform an `Atomic.Computation(.., process=Jac.Photo, processSettings=PhotoIonization.Settings(..), ..)` or call directly functions from the module `PhotoIonization`.
- In JAC, the total photoionization cross sections are calculated and tabulated by default, while the partial cross sections are printed only if the flag `calcPartialCs = true` is set in `PhotoIonization.Settings`.
- In JAC, the statistical tensors of the photoion are calculated and tabulated for the photoionization of unpolarized atoms by plane-wave photons with given Stokes parameters, if the flag `calcTensors` is set in `PhotoIonization.Settings` and if the Stokes parameters are given explicitly in `Ionization.Settings.stokes`.
- Using JAC: Perform an `Atomic.Computation(.., process=Jac.PhotoInPlasma, processSettings = PlasmaShift.PhotoSettings(..), ..)` or call directly functions from the modules `PhotoIonization` and `PlasmaShift`. Until the present, however, only a Debye-Hückel plasma model is supported in JAC, and where the Debye-Hückel screening of the electron-nucleus and electron-photon interaction is incorporated only in first-order perturbation theory into the CI matrix.
- In JAC, the reduced matrix elements above are the main *building blocks* to represent and discuss most photo-ionization properties as they arise from the interaction of atoms and ions with the radiation field.
- **Quantization axis:** While the magnetic quantum numbers of the ion and photo electron usually refer to the adopted *quantization axis* of the overall system (*z*-axis), the **helicity  $\lambda$  of the photon is *always* defined with regard to the wave vector  $\mathbf{k}$  of the photon**, i.e. with regard to the direction of the incident photons.
- **Electron-electron correlations:** **Interesting correlation effects with regard to exchange, interchannel coupling, relaxation and polarization** occurs often in the vicinity of the photoionization threshold.
- **Dichroism:** A dichroism in the photoelectron emission is conventionally related to a change in the polarization of the incident light, and while the polarization state of the target is kept fixed.

## Photoionization amplitude:

- **Emission of an electron with well-defined asymptotic momentum  $\mathbf{p}$  and spin projection  $m_s$ :** If the atom undergoes the transition  $|\alpha_i \mathbb{J}_i M_i\rangle \rightarrow |\alpha_f \mathbb{J}_f M_f\rangle$  during the photoionization process, the transition amplitude for the absorption of a photon with wave vector  $\mathbf{k}$  and helicity  $\lambda = \pm 1$  is given by:

$$\mathcal{M}(M_f, m_s, \lambda, M_i) = \left( \frac{4\pi^2 \alpha}{\omega} \right)^{1/2} \left\langle \alpha_f \mathbb{J}_f M_f, \mathbf{p} m_s \left| \sum_{i=1}^N \boldsymbol{\alpha}_i \mathbf{u}_\lambda e^{i\mathbf{k} \cdot \mathbf{r}_i} \right| \alpha_i \mathbb{J}_i M_i \right\rangle.$$

- The population of the magnetic sublevels  $|\alpha_f \mathbb{J}_f M_f\rangle$  of the photoion depends (of course) on the spin projection  $\lambda$  of the incident photons as well as the magnetic quantum numbers  $M_f$  of the final state of the photoion, while one need to average over the initial sublevels and the spin projections of the emitted electron.
- **(Reduced many-electron) photoionization amplitude:** For a partial-wave decomposition of the outgoing photoelectron and with a multipole expansion of the electron-photon interaction, the reduced photoionization amplitude describes the contribution of individual multipole components of the radiation field to the formation of a scattering state  $|(\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J}_t\rangle$  with well-defined total symmetry and free-electron energy  $\varepsilon$

$$\langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J}_t | \mathbb{O}^{(\mathbb{M}, \text{photoionization})} | \alpha_i \mathbb{J}_i \rangle = i^{-\ell} \exp(i \Delta_\kappa) \langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J}_t | \mathbb{O}^{(\mathbb{M}, \text{absorption})} | \alpha_i \mathbb{J}_i \rangle$$

Here, the photoionization amplitude also accounts for the proper phase (or the asymptotic boundary condition) of the outgoing electron.

- In JAC, the reduced photoionization amplitudes above are the main *building blocks* to represent and discuss most photoionization properties as they arise from the interaction of atoms and ions with the radiation field.

## Partial and total cross sections:

- **Partial photoionization cross sections for the photoionization of an initially unpolarized atom:** For the photoionization of an initially unpolarized atom from level  $|\alpha_i \mathbb{J}_i\rangle \rightarrow |\alpha_f \mathbb{J}_f M_f\rangle$  by unpolarized plane-wave photons with wave vector  $\mathbf{k} \parallel \mathbf{e}_z$  (quantization axis), the

partial cross sections can be expressed as:

$$\begin{aligned}
\sigma(\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f M_f) &= \frac{1}{2(2J_i + 1)} \sum_{M_i \lambda m_s} \int d\Omega_p |\mathcal{M}(M_f, m_s, \lambda, M_i)|^2 \\
&= \frac{4\pi^2 \alpha}{2\omega (2J_i + 1)} \sum_{M_i \lambda} \sum_{\kappa m} \sum_{JJ'M} \langle J_f M_f jm | JM \rangle \langle J_f M_f jm | J'M \rangle \left\langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J} M \left| \sum_{i=1}^N \boldsymbol{\alpha}_i \mathbf{u}_\lambda e^{i\mathbf{k} \cdot \mathbf{r}_i} \right| \alpha_i \mathbb{J}_i M_i \right\rangle \\
&\quad \times \left\langle (\alpha_f J_f, \varepsilon \kappa) \mathbb{J}' M \left| \sum_{i=1}^N \boldsymbol{\alpha}_i \mathbf{u}_\lambda e^{i\mathbf{k} \cdot \mathbf{r}_i} \right| \alpha_i J_i M_i \right\rangle^* .
\end{aligned}$$

Here, the integration over all emission angles  $\Omega_p = (\vartheta_p, \varphi_p)$  of the photoelectron results into a *incoherent* summation over the partial waves of the many-electron scattering states with different  $\kappa$ 's of the electron.

- **Partial photoionization cross sections for initially unpolarized atoms by unpolarized plane-wave photons:** For unpolarized atoms and unpolarized incident plane-wave photons, the partial cross sections from above can be simplified and expressed as:

$$\begin{aligned}
\sigma(\alpha_i \mathbb{J}_i \rightarrow \alpha_f J_f M_f) &= \frac{8\pi^3 \alpha}{2\omega (2J_i + 1)} \sum_{t\lambda} \sum_{\kappa} \sum_{JJ'} \sum_{\mathbb{M}\mathbb{M}'} i^{L-L'} (i\lambda)^p (-i\lambda)^{p'} (-1)^{L+L'} \sqrt{(2L+1)(2L'+1)(2J+1)(2J'+1)} \\
&\quad \times \langle L'\lambda J_f M_f | t, \lambda + M_f \rangle \langle L\lambda J_f M_f | t, \lambda + M_f \rangle \begin{Bmatrix} j & J' & J_f \\ J & J_i & L \\ J_f & L' & t \end{Bmatrix} \\
&\quad \times \left\langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J} \left\| \mathbb{O}^{(\mathbb{M}, \text{photoionization})} \right\| \alpha_i \mathbb{J}_i \right\rangle \left\langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J}' \left\| \mathbb{O}^{(\mathbb{M}', \text{photoionization})} \right\| \alpha_i \mathbb{J}_i \right\rangle^* .
\end{aligned}$$

- **Total photoionization cross sections for initially unpolarized atoms by unpolarized plane-wave photons:** For the photoionization of an initially unpolarized atom from level  $|\alpha_i \mathbb{J}_i\rangle \rightarrow |\alpha_f \mathbb{J}_f\rangle$  by unpolarized plane-wave photons with wave vector  $\mathbf{k} \parallel \mathbf{e}_z$ , the total cross sections

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can be expressed as:

$$\sigma(\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f) = \sum_{M_f} \sigma(\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f M_f) = \frac{4\pi^2 \alpha \omega}{2(2J_i + 1)} \sum_{\kappa} \sum_J \sum_{\mathbb{M}} |\langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J} \parallel \mathcal{O}^{(\mathbb{M}, \text{photoionization})} \parallel \alpha_i \mathbb{J}_i \rangle|^2.$$

### Final-state density matrix 'photo-ion + photo-electron':

- The final-state density matrix of the overall system 'photo-ion + photo-electron' is not so useful in general and is needed only for entanglement studies and related topics. Here, we just summarize some formulas but further work is needed to find the proper expression for different properties of the incident light.
- **Final-state density matrix of the 'photo-ion + photo-electron':** After the absorption of the photon, we have a free electron with asymptotic linear momentum  $\mathbf{p}$  and with spin projection  $m_s$ , while the photoion is left in a fine level  $|\alpha_f \mathbb{J}_f\rangle$ . Using a basis with well-defined (angular) momenta  $J_i$  and  $J_f$  of the initial as well as the residual ion and the asymptotic linear momentum  $\mathbf{p}$  and spin projection  $m_s$  of the ejected electron, the final-state density matrix can be written as

$$\begin{aligned} & \langle \alpha_f J_f M_f, \mathbf{p} m_s \mid \rho_f \mid \alpha_f J_f M'_f, \mathbf{p} m'_s \rangle \\ &= \sum_{M_i M'_i \lambda \lambda'} \langle \alpha_f \mathbb{J}_f M_f, \mathbf{p} m_s \mid \hat{\mathcal{R}} \mid \alpha_i \mathbb{J}_i M_i, \mathbf{k} \lambda \rangle \langle \alpha_i \mathbb{J}_i M_i, \mathbf{k} \lambda \mid \rho_i \mid \alpha_i \mathbb{J}_i M'_i, \mathbf{k} \lambda' \rangle \langle \alpha_i \mathbb{J}_i M'_i, \mathbf{k} \lambda' \mid \hat{\mathcal{R}}^\dagger \mid \alpha_f J_f M'_f, \mathbf{p} m'_s \rangle \\ &= \sum_{M_i M'_i \lambda \lambda'} c_{M_i M'_i} c_{\lambda \lambda'} \langle \alpha_f \mathbb{J}_f M_f, \mathbf{p} m_s \mid \boldsymbol{\alpha} \cdot \mathbf{u}_\lambda e^{i\mathbf{k} \cdot \mathbf{r}} \mid \alpha_0 J_0 M_0 \rangle \langle \alpha_f J_f M'_f, \mathbf{p} m'_s \mid \boldsymbol{\alpha} \cdot \mathbf{u}_{\lambda'} e^{i\mathbf{k} \cdot \mathbf{r}} \mid \alpha_i \mathbb{J}_i M'_i \rangle^*, \end{aligned}$$

where  $\langle \alpha_f \mathbb{J}_f M_f, \mathbf{p} m_s \mid \boldsymbol{\alpha} \cdot \mathbf{u}_\lambda e^{i\mathbf{k} \cdot \mathbf{r}} \mid \alpha_i \mathbb{J}_i M_i \rangle$  represents the photoionization amplitude from above.

### Angular distribution of photoelectrons:

- **Measurement operator for angular distributions:** If the detector is not sensitive to the spin state of the photoelectron and the residual photoion, a proper measurement (projection) operator is obtained if we sum over the electron spin projections  $m_s$  as well as the magnetic quantum numbers  $M_f$  of the ion. In the density matrix formalism, this projection operator for measuring the angular distribution of the



emitted photoelectrons therefore is:

$$P = \sum_{M_f m_s} |\alpha_f \mathbb{J}_f M_f, \mathbf{p} m_s\rangle \langle \alpha_f \mathbb{J}_f M_f, \mathbf{p} m_s|,$$

- With this measurement operator, the well-known angular distribution of the photoelectrons is just the normal trace of the final-state density matrix:

$$W(\vartheta) = \text{Tr}(P \rho_f) = \sum_{M_f m_s} \langle \alpha_f \mathbb{J}_f M_f, \mathbf{p} m_s | \rho_f | \alpha_f \mathbb{J}_f M_f, \mathbf{p} m_s \rangle.$$

### Reduced statistical tensors of the photoion:

- **Statistical tensors of the photoion in level  $|\alpha_f \mathbb{J}_f\rangle$  after the photoionization of unpolarized atoms by unpolarized plane-wave photons:**  
If the incident light propagates along the  $x$ -axis (quantization axis, because of its typical use in experimental discussions) and if the photoelectron remains unobserved, the statistical tensors are given by Kämpfer *et al.* (2016)

$$\begin{aligned} \rho_{kq}(\alpha_f \mathbb{J}_f) &= \frac{\pi}{2J_i + 1} \delta_{q0} \sum_{\mathbb{M} \mathbb{M}'} \sum_{\kappa} \sum_{JJ'} \sum_{\lambda=\pm 1} i^{L+p-L'-p'} \lambda^{p+p'} [L, L', J, J']^{1/2} (-1)^{J+J'+J_f+J_i+j+1} \langle L\lambda, L' - \lambda | kq \rangle \begin{Bmatrix} J_f & j & J' \\ J & k & J_f \end{Bmatrix} \begin{Bmatrix} J' & J_i & L' \\ L & k & J \end{Bmatrix} \\ &\quad \times \langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J} \parallel \mathbb{O}^{(\mathbb{M}, \text{photoionization})} \parallel \alpha_i \mathbb{J}_i \rangle \langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J}' \parallel \mathbb{O}^{(\mathbb{M}', \text{photoionization})} \parallel \alpha_i \mathbb{J}_i \rangle^*. \end{aligned}$$

- **Statistical tensors of the photoion in level  $|\alpha_f \mathbb{J}_f\rangle$  after the photoionization of unpolarized atoms by plane-wave photons with given density matrix  $(c_{\lambda\lambda'})$ :** If the incident plane-wave light with well-defined (photon) density matrix  $(c_{\lambda\lambda'})$ , i.e. with the given Stokes parameters

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$(P_1, P_2, P_3)$ , propagates along the  $z$ -axis, the statistical tensors are given more generally by (Sharma *et al.*, 2010)

$$\begin{aligned} \rho_{kq}(\alpha_f \mathbb{J}_f) &= \frac{\pi}{2J_i + 1} \sum_{\mathbb{M}\mathbb{M}'} \sum_{\kappa J J'} \sum_{\lambda \lambda'} [\delta_{\lambda \lambda'} (1 + \lambda P_3) + (1 - \delta_{\lambda \lambda'}) (P_1 - i\lambda P_2)] i^{L-L'+p-p'} \lambda^p (\lambda')^{p'} [L, L', J, J']^{1/2} \\ &\times (-1)^{(J+J'+J_f+J_i+j+1)} \langle L\lambda L' - \lambda' | kq \rangle \left\{ \begin{matrix} J_f & j & J' \\ J & k & J_f \end{matrix} \right\} : \left\{ \begin{matrix} J' & J_i & L' \\ L & k & J \end{matrix} \right\} \\ &\times \left\langle (\alpha_f J_f, \varepsilon \kappa) J \left\| \mathbb{O}^{(\mathbb{M}, \text{photoionization})} \right\| \alpha_i J_i \right\rangle \left\langle (\alpha_f J_f, \varepsilon \kappa) J' \left\| \mathbb{O}^{(\mathbb{M}', \text{photoionization})} \right\| \alpha_i J_i \right\rangle^*, \end{aligned}$$

- In JAC, the statistical tensors of the photoion are calculated and tabulated for the photoionization of unpolarized atoms by plane-wave photons with given Stokes parameters, if the proper information is given by the `PhotoIonization.Settings`.

### Photoeionization by twisted light:

- Photoionization by twisted light:

$$A + \hbar\omega^{(\text{twist})} \longrightarrow A^+ + e_p^-$$

### 7.1.d. Radiative recombination (PhotoRecombination)

#### Process, notations & application:

- **Radiative recombination** or photorecombination of an atom or ion:  $A^{q+} + e^- \longrightarrow A^{(q-1)+} + \hbar\omega$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + |\varepsilon \kappa\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + \hbar\omega(\mathbb{M})$  or  $|\alpha_i \mathbb{J}_i\rangle + |\mathbf{p} m_s\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + \hbar\omega(\mathbb{M})$
- Using JAC: Perform an `Atomic.Computation(..., process=Jac.Rec, processSettings=PhotoRecombination.Settings(...), ...)` or call directly functions from the module `PhotoRecombination`.
- In JAC, the total radiative recombination cross sections are calculated and tabulated by default for all selected recombination lines.
- In JAC, the (reduced) photorecombination amplitudes are the building blocks to obtain most of the properties below.
- In JAC, the anisotropy parameters  $\beta_\nu^{(\text{REC})}$ ,  $\nu = 1, \dots, 4$  are calculated and tabulated for all selected photorecombination lines if the flag `calcAnisotropy = true` is set in `PhotoRecombination.Settings`.
- In JAC, the statistical tensors for the recombination of unpolarized ions with unpolarized plane-wave electrons are calculated and tabulated if the flag `calcTensors = true` is set in `PhotoRecombination.Settings`.
- The photorecombination of ions is known also as **radiative recombination (RR)** or **radiative electron capture (REC)** in the literature.
- The radiative recombination (RR) can be viewed as the time-reversed photoionization process and occurs frequently in stellar and laboratory plasmas as well as in ion-trap and storage-ring experiments.
- If a highly-charged ion circulates in a storage ring, it may capture electrons from rest-gas atoms in the ring and will then be misbent by subsequent steering magnets. Therefore, the **radiative recombination or electron capture typically results in a loss of the ion from the ring**, and where the lifetime is usually proportional to the inverse of the product of the capture cross section and the target pressure.
- **Magnetic sublevel population  $|\alpha_f \mathbb{J}_f M_f\rangle$  of the recombined ion:** The radiative recombination amplitude from above can be used also to describe the (magnetic) sublevel population of the recombined ion. In the framework of density matrix theory, this sublevel population of the ionic (or atomic) states is described most conveniently in terms of the *statistical tensors*  $\rho_{kq}(\alpha_f \mathbb{J}_f)$  of the recombined ion in its final level  $|\alpha_f \mathbb{J}_f\rangle$ .
- *Quantization axis:* Similar as in the photoionization process, the particular expression for the statistical tensors depend on the choice of the *quantization axis* of the overall system ( $z$ -axis) as well as the geometry of the experiments.

### Photorecombination amplitude:

- **Photorecombination amplitude:** In first-order perturbation theory, most radiative recombination properties can be traced back to the evaluation of the (reduced photon emission) matrix element

$$\langle \alpha_f \mathbb{J}_f \parallel \mathbb{O}^{(\mathbb{M}, \text{recombination})} \parallel (\alpha_i \mathbb{J}_i, \varepsilon \kappa) \mathbb{J}_t \rangle = \langle (\alpha_i \mathbb{J}_i, \varepsilon \kappa) \mathbb{J}_t \parallel \mathbb{O}^{(\mathbb{M}, \text{photoionization})} \parallel \alpha_f \mathbb{J}_f \rangle^*$$

which describes the capture of a free electron with energy  $\varepsilon$  and angular momentum  $\kappa$  under the (simultaneous) emission of a photon with multipolarity  $\mathbb{M} = (L, p)$ .

### Final-state density matrix ‘recombined ion + photon’:

- **Final-state density matrix ‘recombined ion + photon’:** For an unpolarized ion with nuclear spin  $I = 0$  and for unpolarized plane-wave electrons, this density matrix can be expressed as:

$$\begin{aligned} & \langle \alpha_f \mathbb{J}_f M_f, \mathbf{k} \lambda \mid \rho_f \mid \alpha_f \mathbb{J}_f M'_f, \mathbf{k} \lambda' \rangle \\ &= \frac{8\pi}{2(2J_i + 1)} \sum_{\mathbb{M} \mathbb{M}'} \sum_{J J'} \sum_{\kappa \kappa'} \sum_{\nu \mu_1 \mu_2 s s'} D_{\mu_1 \mu_2}^\nu(\mathbf{k} \rightarrow \mathbf{e}_z) i^{L'-L} (i\lambda')^{\pi'} (-i\lambda)^\pi (-1)^{1/2+J_i+J'+L+s+s'-M'_{J,f}} [\ell, \ell', L, L', j, j', J, J', s]^{1/2} \\ & \quad \times [s'] \begin{Bmatrix} j' & 1/2 & \ell' \\ \ell & s & j \end{Bmatrix} \begin{Bmatrix} j' & J_i & J' \\ J & s & j \end{Bmatrix} \begin{Bmatrix} J & L & J_f \\ \nu & s' & L' \end{Bmatrix} \begin{Bmatrix} J & J' & s \\ J_f & s & L' \end{Bmatrix} \langle \ell' 0 \ell 0 \mid s 0 \rangle \langle J_f M_f s' - M'_f \mid \nu \mu_1 \rangle \\ & \quad \times \langle s 0 s' M'_f \mid J_f M'_f \rangle \langle L' \lambda' L - \lambda \mid \nu \mu_2 \rangle \langle \alpha_f \mathbb{J}_f \parallel \mathbb{O}^{(\mathbb{M}, \text{recombination})} \parallel (\alpha_i \mathbb{J}_i, \varepsilon \kappa) \mathbb{J} \rangle \langle \alpha_f \mathbb{J}_f \parallel \mathbb{O}^{(\mathbb{M}', \text{recombination})} \parallel (\alpha_i \mathbb{J}_i, \varepsilon \kappa') \mathbb{J} \rangle^*. \end{aligned}$$

and if a nuclear spin  $I = 0$  is assumed.

- In JAC, no attempt has been made so far to implement the density matrix of the overall system ‘recombined ion + photon’ because of its little relevance for practical computations. However, this density matrix could be implemented quite easily by using the photorecombination amplitudes from above.

### Total recombination cross sections:

- **Total recombination cross section:** If the density of states and the principle of detailed balance between the capture and the (photo-) ionization of an electron is taken into account, the total cross section reads as

$$\sigma(\alpha_i \mathbb{J}_i, \alpha_f \mathbb{J}_f) = \frac{8\alpha^3 \pi^3 \omega}{(2J_i + 1) \beta^2 \gamma^2} \left[ \sum_{\mathbb{M} J \kappa} \left| \langle \alpha_f \mathbb{J}_f \parallel \mathbb{O}^{(\mathbb{M}, \text{recombination})} \parallel (\alpha_i \mathbb{J}_i, \varepsilon \kappa) \mathbb{J} \rangle \right|^2 \right],$$

where  $\beta = v/c$  and  $\gamma = \sqrt{1 - \beta^2}$  arise from the Lorentz transformation in going from the projectile into the laboratory frame.

### Angular distribution of emitted photons:

- **Angular distribution of emitted photons:** For initially unpolarized ions and unpolarized incident plane-wave electrons, the angular distribution of the emitted (RR) photons can be readily obtained from the final-state density matrix by taking the trace  $\text{Tr} (P_{\mathbf{k}} \rho_f)$  with regard to the measurement (projection operator) of plane-wave photons:

$$W(\theta) = \frac{\sigma}{4\pi} \left( 1 + \sum_{\nu=1} P_{\nu}(\cos \theta) \beta_{\nu}^{(\text{REC})}(\alpha_i \mathbb{J}_i, \alpha_f \mathbb{J}_f) \right).$$

- **Anisotropy parameters of the emitted photons:**

$$\begin{aligned} \beta_{\nu}^{(\text{REC})} = & -\frac{1}{2} \sum_{JJ' \kappa \kappa'} \sum_{\mathbb{M} \mathbb{M}'} i^{L+p-L'-p'} (-1)^{J_i-1/2-J_f} [L, L', \ell, \ell', j, j', J, J']^{1/2} \langle \ell 0, \ell' 0 \mid \nu 0 \rangle \langle L 1, L' - 1 \mid \nu 0 \rangle \left( 1 + (-1)^{L+p+L'+p'-\nu} \right) \\ & \times \left\{ \begin{matrix} J & J' & \nu \\ L' & L & J_f \end{matrix} \right\} \left\{ \begin{matrix} J & J' & \nu \\ j' & j & J_i \end{matrix} \right\} \left\{ \begin{matrix} j & j' & \nu \\ \ell' & \ell & 1/2 \end{matrix} \right\} \langle \alpha_f \mathbb{J}_f \parallel \mathbb{O}^{(\mathbb{M}, \text{recombination})} \parallel (\alpha_i \mathbb{J}_i, \varepsilon \kappa) \mathbb{J} \rangle^* \\ & \times \left\langle \alpha_f \mathbb{J}_f \parallel \mathbb{O}^{(\mathbb{M}', \text{recombination})} \parallel (\alpha_i \mathbb{J}_i, \varepsilon \kappa') \mathbb{J}' \right\rangle \left[ \sum_{J \kappa \mathbb{M}} \left| \langle \alpha_f \mathbb{J}_f \parallel \mathbb{O}^{(\mathbb{M}, \text{recombination})} \parallel (\alpha_i \mathbb{J}_i, \varepsilon \kappa) \mathbb{J} \rangle \right|^2 \right]^{-1}. \end{aligned}$$

### (Reduced) Statistical tensors of the recombined ion in the level $|\alpha_f \mathbb{J}_f\rangle$ :

- **Statistical tensors  $\rho_{k0}(\alpha_f \mathbb{J}_f)$  of the final ion following the recombination of unpolarized atoms with unpolarized plane-wave electrons:** If the quantization axis is adopted along the incoming electron momentum  $\mathbf{p} \parallel \mathbf{e}_z$ , the statistical tensors of the recombined ion in the level  $|\alpha_f \mathbb{J}_f\rangle$  are given by

$$\rho_{k0}(\alpha_f \mathbb{J}_f) = \frac{32\pi^3}{2J_i + 1} \sum_{\mathbb{M}} \sum_{JJ' \kappa \kappa'} [\ell, \ell', j, j', J, J']^{1/2} (-1)^{J_i + L - J_f + J - J' - 1/2} \langle \ell 0, \ell' 0 | k 0 \rangle \begin{Bmatrix} j & j' & k \\ \ell' & \ell & 1/2 \end{Bmatrix} \begin{Bmatrix} j & j' & k \\ J' & J & J_i \end{Bmatrix} \begin{Bmatrix} J & J' & k \\ J_f & J_f & L \end{Bmatrix}$$

$$\langle \alpha_f \mathbb{J}_f \parallel \mathbb{O}^{(\mathbb{M}, \text{recombination})} \parallel (\alpha_i \mathbb{J}_i, \varepsilon \kappa) \mathbb{J} \rangle^* \langle \alpha_f \mathbb{J}_f \parallel \mathbb{O}^{(\mathbb{M}, \text{recombination})} \parallel (\alpha_i \mathbb{J}_i, \varepsilon \kappa') \mathbb{J}' \rangle$$

Here, the summation over  $(\mathbb{M} J \kappa)$  includes all allowed combinations of the multipoles of the radiation field as well as the many-electron continua  $|(\alpha_i \mathbb{J}_i, \varepsilon \kappa) \mathbb{J}\rangle$  of the initial ion.

- **Reduced statistical tensors:** Instead of using the statistical tensors  $\rho_{k0}(\alpha_f \mathbb{J}_f)$ , it is often more convenient to describe the sublevel population of the residual ions in terms of the (so-called) *reduced* tensors:

$$\mathcal{A}_{k0}(\alpha_f \mathbb{J}_f) = \frac{\rho_{k0}(\alpha_f \mathbb{J}_f)}{\rho_{00}(\alpha_f \mathbb{J}_f)},$$

which are directly related to the population cross sections  $\sigma(\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f M_f)$  for the electron capture into an ionic sublevel  $|\alpha_f \mathbb{J}_f M_f\rangle$ .

### Photorecombination of ions with twisted electrons:

- **Angular distribution and the linear polarization of photons:** Expressions have been derived for the angular distribution and the linear polarization of the emitted photons due to the capture of twisted electrons into the ground state of (hydrogen-like) ions. These expressions showed that the angular and polarization distributions are sensitive to both, the transverse momentum and the topological charge of the incident electron beam.
- **Photorecombination with twisted electrons:** For twisted electrons with well-defined values of the linear momentum  $p_z$  and with topological charge  $m = 0, \pm 1, \pm 2$ , i.e. with well-defined projection of the OAM upon the propagation direction, the initial-state wave function is

given by:

$$\psi_i^{(\text{tw})}(\mathbf{r}) \equiv \psi_{\kappa m p_z}^{(\text{tw})}(\mathbf{r}) = \int \frac{d^2 \mathbf{p}_\perp}{(2\pi)^2} a_{\kappa m}(\mathbf{p}_\perp) e^{i(\mathbf{p}_\perp \cdot \mathbf{r}_\perp + p_z z)},$$

where  $a_{\kappa m}(\mathbf{p}_\perp)$  refer to the Fourier coefficients that include also the transverse momentum profile of the beam.

➤ Final bound-state wave functions:

$$\psi_f(\mathbf{r} - \mathbf{b}_\perp) = e^{-i\mathbf{p} \cdot \mathbf{b}_\perp} \psi_f(\mathbf{r}).$$

### 7.1.e. Auger and autoionization processes (AutoIonization)

#### Process, notations & application:

- Auger emission (autoionization) of an atom or ion:  $A^{q+*} \longrightarrow A^{(q+1)+(*)} + e_a^-$
- Formal quantum notation:  $|\alpha_i \mathbb{J}_i\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon \kappa\rangle$
- Using JAC: Perform an `Atomic.Computation(..., process=Jac.Auger, processSettings=AutoIonization.Settings(...), ...)` or call directly functions from the module `Auger`.
- In JAC, the particular kind of the Auger operator  $\mathbb{V}^{(\text{Auger})}$  can be specified by selecting `AutoIonization.Settings.operator = ("Coulomb", "Breit", "Coulomb+Breit", ...)`.
- In JAC, the Auger rates and (Auger) lifetimes are tabulated by default for all selected (Auger) transitions.
- In JAC, the intrinsic anisotropy parameters are calculated and tabulated if the flag `calcAnisotropy = true` is set in `AutoIonization.Settings`.
- Using JAC: Perform an `Atomic.Computation(..., process=Jac.AutoIonizationInPlasma, processSettings=PlasmaShift.AutoIonizationSettings(...), ...)` or call directly functions from the modules `AutoIonization` and `PlasmaShift`. Until the present, however, only a Debye-Hückel plasma model is supported and where the Debye-Hückel screening of the electron-nucleus and electron-photon interaction is incorporated only in first-order perturbation theory within the CI matrix.

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- An autoionization of an (inner-shell) excited bound state can occur (only) if the **initial level is energetically embedded into the *continuum* of the next higher charge state.**
- Formally, an electron emission from such an embedded resonance (level) arise due to its coupling to one or several scattering states, which are degenerate with the initial state and which involve one (or more) *free* electrons.
- While the Coulomb repulsion in the Auger amplitude above is typically sufficient in order to describe the autoionization of light and medium elements, the Breit interaction has been found important for electron emission and capture processes of highly-charged ions.
- The **restriction to the electron-electron interaction in the computation (evaluation) of the Auger amplitudes is quite standard in all presently available Auger codes**, even if the orbital functions of the resonant scattering state  $|\alpha_t \mathbb{J}_t\rangle$  and the final-ionic state  $|\alpha_i \mathbb{J}_i\rangle$  are not quite orthogonal to each other.
- In JAC, we generally support the computation of nonradiative (Auger) rates, relative Auger intensities and intrinsic angular anisotropy parameters. To this end, first the (reduced) Auger amplitudes are calculated for all selected pairs of initial and final levels and, later, these amplitudes are combined into Auger rates, intrinsic anisotropy or spin-polarization parameters.
- Since the number of possible Auger transitions often increases very rapidly, especially if several inner-shell holes and/or valence electrons occur, JAC provides several additional control features for reducing the complexity: Apart from (1) the selection of particular Auger lines, these features include (2) the definition of an **energy threshold (i.e. a minimum energy) of the free electron** as well as (3) a **maximum angular momentum** (negative  $\kappa$  quantum number) for the partial waves of the outgoing electron. A proper use of (3) mainly affects the number of the scattering states, that are taken into account, and can be chosen independently of the formally allowed couplings of the free-electron.
- These control features are in line with the **common experience that partial waves with high angular momenta do either not contribute at all** (i.e. have *zero* Auger amplitudes in most computational models) **or are simply negligible**. In addition, JAC supports the use of an **overall energy shift of all Auger transition energies** which can be specified interactively.
- Often, the number of the possible scattering states  $|(\alpha_f \mathbb{J}_f, \varepsilon \kappa) \alpha_t \mathbb{J}_t\rangle$  of a system increases rapidly since the free electrons may couple in quite different ways to the bound-state electrons from the final (Auger) ions.
- **Reaction plane of the Auger emission:** For a prior photo-absorption or particle impact, this reaction plane is usually defined by the directions of the incident photon (particle) beam and the emitted Auger electron.



## Auger amplitudes:

- **Auger amplitude:** Following Åberg and Howat (1982), the **Auger transition amplitude for the autoionization of an excited bound state**  $|\alpha_i \mathbb{J}_i M_i\rangle$  into the final scattering state  $|\alpha_t \mathbb{J}_t M_t\rangle$ , built from an ion in the final level  $|\alpha_f \mathbb{J}_f\rangle$ , can be expressed most generally as

$$\begin{aligned} \langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J}_t M_t | \mathbb{V}^{(\text{Auger})} | \alpha_i \mathbb{J}_i M_i \rangle &= \langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J}_t | \mathbb{V}^{(\text{Auger})} | \alpha_i \mathbb{J}_i \rangle \delta_{J_i, J_t} \delta_{M_i, M_t} \delta_{P_i, P_t} \\ &= \langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J}_t | \mathbb{H} - E | \alpha_i \mathbb{J}_i \rangle \delta_{J_i, J_t} \delta_{M_i, M_t} \delta_{P_i, P_t}. \end{aligned}$$

This amplitude is independent of the projection  $M_i$  of the initial state owing to the rotational symmetry of free atoms and the scalar character of the Hamiltonian and/or interelectronic interactions.

- In an autoionization process, the **final scattering states**  $|\alpha_t \mathbb{J}_t M_t\rangle \equiv |(\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J}_t M_t\rangle$  arise from the coupling of the final state  $|\alpha_f \mathbb{J}_f M_f\rangle$  of the ion with well-defined symmetry with the partial-wave  $|\varepsilon \kappa m_j\rangle$  of the outgoing electron with energy  $\varepsilon$  and (one-electron) angular momentum  $\kappa = \kappa(j, l)$  as well as magnetic projection  $m_j$ .
- In JAC, as in most other tools for computing autoionization processes, a **common set of orthonormal orbitals** is often supposed for the representation of the initial  $|\alpha_i \mathbb{J}_i\rangle$  and final (ionic) bound levels  $|\alpha_f \mathbb{J}_f\rangle$  in the evaluation of the Auger amplitudes. In this case, the Auger transition operator  $\mathbb{V}^{(\text{Auger})} \equiv \mathbb{H} - E \approx \mathbb{V}^{(\text{e-e})}$  simplifies to the electron-electron interaction operator, while the (reduced) Auger amplitude becomes

$$\langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J}_t | \mathbb{V}^{(\text{Auger})} | \alpha_i \mathbb{J}_i \rangle \simeq \langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J}_t | \mathbb{V}^{(\text{e-e})} | \alpha_i \mathbb{J}_i \rangle \delta_{J_i, J_t} \delta_{M_i, M_t} \delta_{P_i, P_t}.$$

- Within the framework of the Dirac-Coulomb-Breit Hamiltonian, the **interelectronic interaction is the sum of the Coulomb repulsion and Breit interaction**,  $\mathbb{V}^{(\text{e-e})} = \mathbb{V}^{(\text{Coulomb})} + \mathbb{V}^{(\text{Breit})} \simeq \mathbb{V}^{(\text{Coulomb})}$ , although it is often approximated by just the Coulomb part.

### Auger rates:

- **Auger rates:** If the (continuum) interaction between the different Auger channels is neglected, the Auger rates are simply given by

$$A(i \rightarrow f) = 2\pi \sum_{\kappa} \left| \langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J}_i \parallel \mathbb{V}^{(\text{Auger})} \parallel \alpha_i \mathbb{J}_i \rangle \right|^2 ,$$

where the summation runs over all partial waves  $|\varepsilon \kappa\rangle$  that fulfill the triangular rule  $\delta(J_f, j, J_i)$  and where the Auger interaction operator  $\mathbb{V}^{(\text{Auger})}$  still need to be specified due to the (computational) framework and additional approximations that are made in practice.

- In JAC, the Auger rates and (Auger) lifetimes are tabulated by default for all selected (Auger) transitions.

### Angular distribution of Auger electrons:

- **Angular distribution of emitted electrons:** If only the Auger electrons are detected from the autoionization of an (inner-shell) excited ion in the level  $|\alpha_i \mathbb{J}_i\rangle$ , and which is fully characterized by the reduced statistical tensors  $\mathcal{A}_{kq}(\alpha_i \mathbb{J}_i)$ , the angular distribution of the emitted electrons of the Auger transition  $|\alpha_i \mathbb{J}_i\rangle \rightarrow |\alpha_f \mathbb{J}_f\rangle$  can be written in the general form (Balashov *et al.* 2000, equation 3.10)

$$W(\vartheta, \varphi; \alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f) = \frac{A(i \rightarrow f)}{4\pi} \left[ 1 + \sum_{k=2,4,\dots}^{k_{\max}} \sqrt{\frac{4\pi}{2k+1}} \alpha_k(\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f) \sum_{q=-k}^k \mathcal{A}_{kq}(\alpha_i \mathbb{J}_i) Y_{kq}(\vartheta, \varphi) \right] ,$$

and where  $\alpha_k(\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f)$  are the (so-called) **intrinsic anisotropy parameters** of the electron emission

$$\begin{aligned} \alpha_k(\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f) = & \left[ \sum_{\kappa} \langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J}_i \parallel \mathbb{V}^{(\text{Auger})} \parallel \alpha_i \mathbb{J}_i \rangle \right]^{-1} (-1)^{J_i+J_f+k-1/2} \sqrt{2J_i+1} \sum_{\kappa \kappa'} [\ell, \ell', j, j']^{1/2} \langle \ell 0, \ell' 0 \mid k 0 \rangle \\ & \times \left\{ \begin{matrix} J_i & j & J_f \\ j' & J_i & k \end{matrix} \right\} \left\{ \begin{matrix} \ell & j & 1/2 \\ j' & \ell' & k \end{matrix} \right\} \langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J}_i \parallel \mathbb{V}^{(\text{Auger})} \parallel \alpha_i \mathbb{J}_i \rangle \langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa') \mathbb{J}_i \parallel \mathbb{V}^{(\text{Auger})} \parallel \alpha_i \mathbb{J}_i \rangle^* \end{aligned}$$

- Here, the factor  $A(i \rightarrow f)$  is the total probability of the Auger transition  $|\alpha_i \mathbb{J}_i\rangle \rightarrow |\alpha_f \mathbb{J}_f\rangle$ , integrated over the ejection angles. The **anisotropy parameters**  $\alpha_k(\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f)$  contain all information about the dynamics of the decay, while the tensors  $\mathcal{A}_{kq}(\alpha_i \mathbb{J}_i)$  just describe the polarization properties of the initial level.

- **Angular distribution of emitted Auger electrons averaged over several final levels:** If several final levels cannot be resolved by the detector, the total angular distribution is an incoherent sum of the angular distributions of all unresolved final levels with total angular momenta  $J_f$ :

$$W(\vartheta, \varphi; \alpha_i \mathbb{J}_i \rightarrow \{\alpha_f \mathbb{J}_f\}) = \sum_{\alpha_f \mathbb{J}_f} W(\vartheta, \varphi; \alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f) .$$

In this case, the angular distribution can be expressed also by means of some **average anisotropy parameter** (within the electric-dipole approximation)

$$\bar{\beta}_2 = \frac{\sum_{\alpha_f \mathbb{J}_f} \beta_2(i \rightarrow f) A(i \rightarrow f)}{\sum_{\alpha_f \mathbb{J}_f} A(i \rightarrow f)}, \quad \beta_2(i \rightarrow f) = \alpha_2(\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f) A_{20}(\alpha_i \mathbb{J}_i).$$

### Spin polarization of Auger electrons:

- **The spin polarization of Auger electrons is usually described by an polarization vector  $\mathbf{P}$ .** The three components of  $\mathbf{P}$  can be expressed in terms of the statistical tensors of the initial level as well as the intrinsic anisotropy parameters of the considered Auger transition.
- Two coordinate systems,  $S(XYZ)$  and  $S'(X'Y'Z')$  are often employed in order to characterize the spin polarization of Auger electrons; cf. the figure below and Kabachnik *et al.* (2007).

### Two-electron single Auger emission (TESA):

- Non-resonant three-electron Auger transitions have been suggested several decades ago, though clear evidence for such TESA transition were found in Ar only with very low intensity relative to the normal Auger lines.
- The observed TESA transitions usually start with the creation of a double vacancy, i.e. with two vacancies in an inner shell, and which are filled simultaneously by two outer-shell electrons, while the released energy is passed to a third emitted electron.
- Such inner-shell double vacancies are normally at energies above the minimum fourfold ionization energy and can easily be filled also by two separate successive two-electron Auger transitions involving four electrons in total.

### 7.1.f. Dielectronic recombination (Dielectronic)

#### Process, notations & application:

- **Dielectronic recombination (DR)** of an atom or ion:  $A^{q+} + e^- \longrightarrow A^{(q-1)+*} \longrightarrow A^{(q-1)+(*)} + \hbar\omega$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + |\varepsilon \kappa\rangle \longrightarrow |\alpha_d \mathbb{J}_d\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + \hbar\omega(\mathbb{M})$  or  
 $|\alpha_i \mathbb{J}_i\rangle + |\mathbf{p}_i m_i\rangle \longrightarrow |\alpha_d \mathbb{J}_d\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + \hbar\omega(\mathbb{M})$
- Using JAC: Perform an `Atomic.Computation(..., process=JAC.Dierec, processSettings=Dielectronic.Settings(...), ...)` or call directly functions from the module `Dielectronic`.
- In JAC, both the individual  $S(i \rightarrow d \rightarrow f)$  and total resonance strengths  $S(i \rightarrow d)$  are calculated and tabulated by default for all selected pathways.
- In JAC, the capture (Auger) and radiative amplitudes and rates are finally computed by the corresponding modules `Auger` and `PhotoEmission`.
- In most computations, the **dielectronic recombination of a  $N$ -electron target ion in the (initial) level  $|\alpha_i \mathbb{J}_i\rangle$**  is typically handled as a **two-step** process, in which first an electron is captured resonantly from the continuum into a  $[(N+1)\text{-electron}]$  resonance level  $|\alpha_d \mathbb{J}_d\rangle$ , embedded into the continuum of the initial ion and with often two or more electrons in excited (one-electron) orbitals. In a second step, this excited level  $|\alpha_d \mathbb{J}_d\rangle$  then either decays radiatively by the emission of photons to some (final) level  $|\alpha_f \mathbb{J}_f\rangle$ , below of the ionization threshold of the ion, or it returns by autoionization back into the charge stage of the initial target ion.
- Apart from the computation of accurate Auger rates, **helpful calculation of DR spectra critically depends on the simple control (and handling) of all the — radiative and non-radiative — decay branches of the resonantly excited ion.** In JAC, these branches are treated equally within the `Dielectronic` module.
- As seen from the (total) resonance strength below, the **experimentally observed strength of the dielectronic recombination for a given resonance level  $|\alpha_d \mathbb{J}_d\rangle$**  is proportional to the capture rate  $A_a(i \rightarrow d)$  as well as the (total) radiative rate.
- If, in particular, the intermediate level  $|\alpha_d \mathbb{J}_d\rangle$  can autoionize only back to the ground state of the target ion, as for example for the  $3s$  valence excitation of a sodium-like target ion, the total Auger rate is then — up to statistical factors — equal to the capture rate.
- In practice, the **number of the possible scattering states  $|(\alpha_i \mathbb{J}_i, \varepsilon \kappa), \mathbb{J}_d M_d\rangle$  of the recombined ion** often increases very rapidly since, the free electrons may couple in many different ways to the bound-state electrons of the initial ion.

## Strength of dielectronic recombination (DR) resonances:

- **Partial resonance strength:** If the interference between the radiative and dielectronic (nonradiative) capture of the electron is negligible in the field of the target, the integrated DR cross section for an isolated resonance level  $|\alpha_d \mathbb{J}_d\rangle$ , the (so-called) **partial resonance strength**, can be expressed in terms of the Auger and radiative rates of this level by

$$S(i \rightarrow d \rightarrow f) \equiv \int_{-\infty}^{\infty} dE \sigma^{(\text{DR})}(E) = \frac{2\pi^2 \hbar}{k_i^2} \frac{A_a(i \rightarrow d) A_r(d \rightarrow f)}{\Gamma_d}$$

where  $k_i$  denotes the wave number of the incident electron and  $\Gamma_d$  the half-widths of the resonance level  $|\alpha_d \mathbb{J}_d\rangle$ , embedded into the continuum. Since this **resonance strength refers to the area under the DR cross sections**, it is usually given in units of:  $[S] = \text{cm}^2 \text{eV}$ .

- In the resonance strength above,  $A_a(i \rightarrow d)$  is the rate for the **inverse Auger capture** from the initial into the (doubly-excited) resonance level  $|\alpha_d \mathbb{J}_d\rangle$ , while  $A_r(d \rightarrow f)$  refers to the rate for the **radiative stabilization** to the final level  $|\alpha_f \mathbb{J}_f\rangle$  of the  $N + 1$ -electron ion.
- **Total width of the resonance level  $|\alpha_d \mathbb{J}_d\rangle$ :** The total width  $\Gamma_d$  is determined by all possible decay channels of the resonance level  $|\alpha_d \mathbb{J}_d\rangle$  and is given in first-order perturbation theory by

$$\Gamma_d = \hbar \left( \sum_j A_a(d \rightarrow j) + \sum_{f'} A_r(d \rightarrow f') \right),$$

i.e. by taking the sum over all the individual Auger and radiative rates (widths) of the intermediate level  $|\alpha_d \mathbb{J}_d\rangle$ .

- The **Auger and inverse capture rates**  $A_a(i \rightarrow d) = g_i/g_d A_a(d \rightarrow i)$  are proportional to each other owing to the principle of **kinetic balance**, and with the  $g$ 's being the statistical factors of the initial and intermediate resonance levels.
- The use of the individual resonance strength  $S(i \rightarrow d \rightarrow f)$  is appropriate especially, if the energy-dependent DR cross section has a Lorentzian profile

$$\sigma^{(\text{DR})}(E) = \frac{S}{\pi} \frac{\Gamma/2}{(E_r - E)^2 + \Gamma^2/4}$$

around the resonance energy  $E_r = E_d - E_i$  as well as a **natural width less or comparable to the energy spread of the electron beam**. In general, this condition is well fulfilled for most  $\Delta n \geq 1$  resonances (with an excitation of a bound electron from the shell  $n_i \rightarrow n_f$ , for

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which the total widths are small in most cases), but this condition can usually be utilized also for the  $\Delta n = 0$  lines if the shape of the resonances is not to be analyzed in detail.

- **Threshold behaviour:** Owing to the energy of the incident electron [cf. the factor  $1/k_i^2$  above], the **resonance strength increases rapidly towards the threshold of the dielectronic recombination and makes the process sensitive to low kinetic energies of the incoming electrons.**
- **Total resonance strength:** At storage rings, the radiative stabilization of the ions in the (intermediate) resonance level  $|\alpha_d \mathbb{J}_d\rangle$  is often not observed explicitly and, hence, the radiative rate for the individual transition  $d \rightarrow f$  in the resonance strength above has to be replaced by the total radiative rate of the (doubly-excited) resonance level  $|\alpha_d \mathbb{J}_d\rangle$  :

$$A_r(d \rightarrow f) \longrightarrow \sum_{f'} A_r(d \rightarrow f') \quad \Longrightarrow \quad S(i \rightarrow d \rightarrow f) \longrightarrow S(i \rightarrow d) = \sum_{f'} S(i \rightarrow d \rightarrow f')$$

- In JAC, both the individual  $S(i \rightarrow d \rightarrow f)$  and total resonance strengthes  $S(i \rightarrow d)$  are calculated and tabulated by default for all selected pathways.

### 7.1.g. Photoexcitation & fluorescence (PhotoExcitationFluores)

#### Process, notations & application:

- **Photoexcitation of an atom or ion with subsequent fluorescence emission:**  $A + \hbar\omega_i \longrightarrow A^* \longrightarrow A^{(*)} + \hbar\omega_f$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + \hbar\omega(\mathbf{k} || \mathbf{e}_z, \lambda; \{\mathbb{M}\}) \longrightarrow |\alpha_e \mathbb{J}_e\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + \hbar\omega_f(\Omega; \mathbf{P})$
- **Using JAC:** Perform an `Atomic.Computation(..., process=PhotoExcFluor, processSettings=PhotoExcitationFluores.Settings(...), ...)` or call directly functions from the module `PhotoExcitationFluores`.
- In JAC, the reduced density matrix of a fluorescence photon can be calculated at a selected solid angle  $\{\Omega = \vartheta, \varphi\}$  by calling the function `PhotoExcitationFluores.computePhotonDM`, and if the solid angle is specified explicitly.
- The photoexcitation and (subsequent) fluorescence emission from atoms and ions has been investigated in good detail in the literature for both, incident plane-wave radiation as well as for twisted Bessel and Laguerre-Gaussian beams.

- Since the **properties of the photo-excited atoms** can be calculated by the module **PhotoExcitation** and the **properties of the radiative emission** by the module **Radiative**, here we shall focus only upon properties that are specific to the **overall photoexcitation & fluorescence process**, such as the **angular distribution or density matrix of the fluorescence photons**.
- In JAC, the alignment parameters of the (intermediate) excited level  $\mathcal{A}_{kq}(\alpha_e \mathbb{J}_e)$  can be calculated and tabulated by the module **PhotoExcitation**, while the structure functions  $f_k(\alpha_e \mathbb{J}_e, \alpha_f \mathbb{J}_f)$  can be obtained by using the module **PhotoEmission**.
- In JAC, the statistical tensors  $\rho_{kq}(\alpha_e \mathbb{J}_e)$  of the intermediate excited level can be calculated and tabulated by the module **PhotoExcitation**. These tensors are also required for describing the photoexcitation & fluorescence process.

### Density matrix of the fluorescence photon $\hbar\omega_f(\Omega)$ :

- **(Reduced) density matrix of the fluorescence photons for the photoemission from an aligned or oriented atom:** For photons emitted under the angle  $\Omega = (\vartheta, \varphi)$  with regard to the propagation direction of the incident plane-wave radiation  $\mathbf{k} \parallel \mathbf{e}_z$ , the (reduced) density matrix can be readily expressed in terms of the statistical tensors of the excited level  $\rho_{kq}(\alpha_e \mathbb{J}_e)$  as

$$\begin{aligned} \langle \mathbf{k}_0 \lambda \mid \rho_{\omega_f} \mid \mathbf{k}_0 \lambda' \rangle &= 2\pi \sum_{k, qq'} \sum_{\mathbb{M} \mathbb{M}'} D_{-qq'}^k(\varphi, \vartheta, 0) \rho_{kq}(\alpha_e \mathbb{J}_e) i^{L'+p'-L-p} \lambda^p (\lambda')^{p'} \sqrt{2L+1} \sqrt{2L'+1} (-1)^{J_f+J_e+k+q+1} \\ &\times \langle L\lambda, L'(-\lambda') \mid k - q' \rangle \left\{ \begin{matrix} L & L' & k \\ J_e & J_e & J_f \end{matrix} \right\} \langle \alpha_f \mathbb{J}_f \parallel \mathbb{O}^{(\mathbb{M}, \text{emission})}(\omega_f) \parallel \alpha_e \mathbb{J}_e \rangle \langle \alpha_f \mathbb{J}_f \parallel \mathbb{O}^{(\mathbb{M}', \text{emission})}(\omega_f) \parallel \alpha_e \mathbb{J}_e \rangle^* \end{aligned}$$

- In JAC, the reduced density matrix of a fluorescence photon can be calculated at a selected solid angles  $\{\Omega = (\vartheta, \varphi)\}$ , see above.

### Angular distribution and Stokes parameters of the fluorescence photon $\hbar\omega_f(\Omega; \mathbf{P})$ :

- **Angular distribution of the fluorescence light for the photoemission from an aligned or oriented atom:** Most generally, the angular distribution can be obtained from the photon density matrix above by a summation over the two polarization states:

$$W(\vartheta, \varphi) = \sum_{\lambda} \langle \mathbf{k}_0 \lambda \mid \rho_{\omega_f} \mid \mathbf{k}_0 \lambda \rangle .$$

- Using the expressions from above, this **angular distribution can be readily calculated for a set of angles  $\{\Omega_1 = (\vartheta_1, \varphi_1), \Omega_2, \dots\}$** , see above.

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- **Angular distribution of the fluorescence light:** For initially unpolarized atoms and unpolarized photons in the photo-excitation process, in particular, the angular distribution is given by the same formulas as for the photoemission (alone), i.e. in terms of the alignment parameters  $\mathcal{A}_{k0}(\alpha_e \mathbb{J}_e)$  and the structure function of the (intermediate) excited level  $|\alpha_e \mathbb{J}_e\rangle$  :

$$W(\vartheta) = \frac{W_o}{4\pi} \left( 1 + \sum_{k=2,4,..} f_k(\alpha_e \mathbb{J}_e, \alpha_f \mathbb{J}_f) \mathcal{A}_{k0}(\alpha_e \mathbb{J}_e) P_k(\cos\vartheta) \right).$$

### 7.1.h. Photoexcitation & autoionization (PhotoExcitationAutoion)

#### Process, notations & application:

- **Photoexcitation & autoionization** of an atom or ion:  $A + \hbar\omega \longrightarrow A^* \longrightarrow A^{(*)} + e_a^-$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + \hbar\omega(\mathbf{k} || \mathbf{e}_z, \lambda; \{\mathbb{M}\}) \longrightarrow |\alpha_r \mathbb{J}_r\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon\kappa\rangle$
- Using JAC: Perform an `Atomic.Computation(..., process=PhotoExcAuto, processSettings=PhotoExcitationAutoion.Settings(...), ...)` or call directly functions from the module `PhotoExcitationAutoion`.
- In JAC, the partial photoexcitation & autoionization cross sections  $\sigma(i \rightarrow r \rightarrow f)$  are calculated and tabulated by assuming that the direct photoionization cross sections  $\sigma^{(\text{photoionization})}(i \rightarrow f)$  is negligible.
- The photoexcitation and (subsequent) autoionization of atoms and ions has been explored in good detail, either as the dominant (photo-) ionization process within the vicinity of certain (photoionization) resonances or simply as competitive process to atomic photoionization.
- Often, the photoexcitation & autoionization can be considered independently from the *direct* photoionization and then leads to **partial and total photoexcitation & autoionization cross sections**.
- If the direct photoionization is negligible, the photoexcitation & autoionization is the time-reversed process to the dielectronic recombination.



### Fano profiles in photo-ionization cross sections:

- **Fano's parametrization:** For an isolated resonance  $|\alpha_r \mathbb{J}_r\rangle$ , the **total photoionization cross section** can be (formally) parametrized in the **vicinity of this resonance** by the energy-dependent cross section

$$\sigma(\mathcal{E}) = \sigma^{(\text{EA})} \frac{(q + \mathcal{E})^2}{1 + \mathcal{E}^2} + \sigma^{(\text{D})}$$

where  $\sigma^{(\text{D})}$  is the direct photoionization cross section and  $\sigma^{(\text{EA})}$  the cross section due to the photoexcitation & autoionization of the **resonance**. Moreover,  $\mathcal{E} = (E_\omega - E_r)/(\Gamma/2)$  indicates the detuning of the incident photon energy  $E_\omega$  from the resonance energy  $E_r$ , in units of the half-width  $(\Gamma/2)$  of the resonance.

- **Fano  $q$  parameter:** This parameter determines the overall profile of the cross section, ranging from a near-Lorentzian ( $q \rightarrow \infty$ ), to a window-like ( $q \rightarrow 0$ ), and up to a completely asymmetric resonance ( $q \rightarrow \pm 1$ ).
- In JAC, no attempt has been made so far to calculate the Fano profiles explicitly for selected resonances, although this is possible quite readily by combining proper calls to functions from the modules `PhotoIonization`, `PhotoExcitation` as well as `Auger`.

### Total photoionization amplitude and cross sections across the resonance level $|\alpha_r \mathbb{J}_r\rangle$ :

- The photoexcitation and (subsequent) autoionization of atoms and ions obviously contributes to the photoionization and, hence, to its amplitudes and cross sections.
- **Total photoionization amplitudes across the resonance level  $|\alpha_r \mathbb{J}_r\rangle$  :**

$$\begin{aligned} \mathcal{M}(\omega; (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J}_t, \alpha_r \mathbb{J}_r, \alpha_i \mathbb{J}_i) \\ = \langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J}_t \parallel \mathbb{O}^{(\mathbb{M}, \text{photoionization})} \parallel \alpha_i \mathbb{J}_i \rangle + \frac{\langle (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J}_t \parallel \mathbb{O}^{(\text{Auger})} \parallel \alpha_r \mathbb{J}_r \rangle \langle \alpha_r \mathbb{J}_r \parallel \mathbb{O}^{(\mathbb{M}, \text{absorption})} \parallel \alpha_i \mathbb{J}_i \rangle}{E_i + \omega - E_r + i \Gamma_r/2} \end{aligned}$$

In this notation of the **resonant photoionization**, we assume the photoelectron to escape by means of the partial wave  $|\varepsilon \kappa\rangle$  with kinetic energy  $\varepsilon = E_i + \omega - E_f$  and with well-defined angular momentum and parity.

- Apart from the *direct* photoionization amplitude (first term), the total photoionization amplitude also contains the photoexcitation & autoionization amplitudes (second term).

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- From these amplitudes, the **total photoionization cross section (intensity)** is obtained as before by the summation over all the possible scattering states of the final system ‘photoion + electron’:

$$\sigma(\omega; \alpha_i \mathbb{J}_i, \alpha_f \mathbb{J}_f) = \frac{4\pi^2 \alpha \omega}{3(2J_i + 1)} \sum_{\alpha_r \mathbb{J}_r, \kappa \mathbb{J}_t} |\mathcal{M}(\omega; (\alpha_f \mathbb{J}_f, \varepsilon \kappa) \mathbb{J}_t, \alpha_r \mathbb{J}_r, \alpha_i \mathbb{J}_i)|^2,$$

This includes both, the summation over all relevant resonances as well as the summation over the partial waves of the photoelectron and the total angular momenta and parities of the total scattering states  $|\alpha_t \mathbb{J}_t\rangle$ .

- In JAC, no attempt has yet been made to calculate and tabulate the resonant photoionization cross sections or angular anisotropy parameters, although this will be possible since all required many-electron matrix elements are internally quite readily available.

### Angular distribution of resonantly emitted Auger electrons:

- Of course, the angular distribution of the emitted Auger electrons critically depend on the (properties of the) prior excitation process; **in this PhotoExcitationAutoion module, we assume a photoexcitation of the atoms or ions for calculating subsequent properties of Auger electrons. We here also assume that the direct photoionization is negligible near to the considered resonance level  $|\alpha_r \mathbb{J}_r\rangle$ .**
- If the resonantly-excited level  $|\alpha_r \mathbb{J}_r\rangle$  is axially symmetric along the  $z$ -axis,  $\mathcal{A}_{kq}(\alpha_r \mathbb{J}_r) \sim \delta_{q0}$ , the angular distribution only depends on the angle  $\vartheta$  between the symmetry ( $z$ -) axis and the linear momentum of the ejected electron. Then, the angular distribution simplifies to

$$W(\vartheta; \alpha_r \mathbb{J}_r \rightarrow \alpha_f \mathbb{J}_f) = \frac{A_a(r \rightarrow f)}{4\pi} \left[ 1 + \sum_{k=2,4,\dots}^{k_{\max}} \alpha_k(\alpha_r \mathbb{J}_r \rightarrow \alpha_f \mathbb{J}_f) \mathcal{A}_{k0}(\alpha_r \mathbb{J}_r) P_k(\cos \vartheta) \right].$$

- In JAC, no attempt has been made so far to tabulate the angular distribution  $W(\vartheta; \alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f)$  explicitly for selected scenarios; cf. section 7.1.e.
- **Photoexcitation of initially unpolarized atoms within the dipole approximation:** For the photoexcitation of initially unpolarized atoms the Auger angular distribution further simplifies to

$$W(\vartheta; \alpha_r \mathbb{J}_r \rightarrow \alpha_f \mathbb{J}_f) = \frac{A_a(r \rightarrow f)}{4\pi} [1 + \beta_2(r \rightarrow f) P_2(\cos \vartheta)] \equiv \frac{A_a(r \rightarrow f)}{4\pi} [1 + \alpha_2(\alpha_r \mathbb{J}_r \rightarrow \alpha_f \mathbb{J}_f) \mathcal{A}_{20}(\alpha_r \mathbb{J}_r) P_2(\cos \vartheta)].$$

In this case, the overall form of the angular distribution for the photoinduced Auger emission coincides with those for the (direct) photoelectron emission, although the asymmetry parameter  $\beta$  is now a product of two factors:  $\beta_2(r \rightarrow f) = \alpha_2(\alpha_r \mathbb{J}_r \rightarrow \alpha_f \mathbb{J}_f) \mathcal{A}_{20}(\alpha_r \mathbb{J}_r)$ . However, this factorization of the  $\beta_2$  anisotropy parameter is only valid if the initial level is (sufficiently) well isolated.

- The angular distribution of the emitted Auger electrons is sensitive to the phase difference between the Auger amplitudes, in contrast to the total probability  $A(i \rightarrow f)$ . Since the coefficients are symmetric with regard to an interchange  $\kappa \leftrightarrow \kappa'$  in the summation of the angular distribution, this distribution is invariant also with respect to the inversion:  $\vartheta \rightarrow \pi - \vartheta$ ,  $\varphi \rightarrow \varphi + \pi$ , simply because it only contains the spherical harmonics of even rank.

### Spin polarization of resonantly emitted Auger electrons:

- The spin polarization of Auger electrons is usually described by an polarization vector  $\mathbf{P}$ . The three components of  $\mathbf{P}$  can be expressed in terms of the statistical tensors of the initial level as well as the intrinsic anisotropy parameters of the considered Auger transition.
- Two coordinate systems,  $S(XYZ)$  and  $S'(X'Y'Z')$  are often employed in order to characterize the spin polarization of Auger electrons; cf. the figure below and Kabachnik *et al.* (2007).

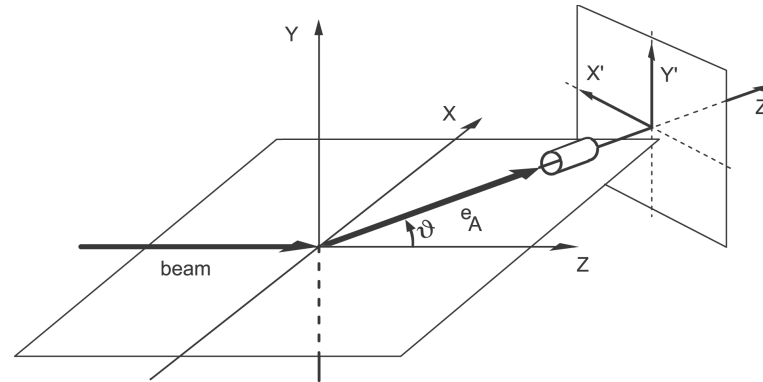


Figure 7.1.: The coordinate systems that are usually utilized for describing the Auger electron spin polarization. From Kabachnik *et al.* (2007).

- (Longitudinal) spin component of the emitted Auger electrons: If the initial level  $|\alpha_i \mathbb{J}_i\rangle$  is produced by photoabsorption of unpolarized target atoms with monochromatic plane-wave photons, the (longitudinal) spin component of the emitted Auger electrons along their propagation direction can be written as:

$$P_{z'} = \frac{\mathcal{A}_{10}(\alpha_i \mathbb{J}_i) \delta_1 \cos \vartheta}{1 + \alpha_2(\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f) \mathcal{A}_{20}(\alpha_i \mathbb{J}_i) P_2(\cos \vartheta)}.$$

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- **Transverse spin component:** The transverse spin component within the  $(x - z)$  reaction plane  $P_{x'}$  and the transverse component perpendicular to the reaction plane  $P_{y'}$  take the form

$$P_{x'} = \frac{\mathcal{A}_{10}(\alpha_i \mathbb{J}_i) \xi_1 \sin \vartheta}{1 + \alpha_2(\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f) \mathcal{A}_{20}(\alpha_i \mathbb{J}_i) P_2(\cos \vartheta)}, \quad P_{y'} = P_y = \frac{\mathcal{A}_{20}(\alpha_i \mathbb{J}_i) \xi_2 \sin 2\vartheta}{1 + \alpha_2(\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f) \mathcal{A}_{20}(\alpha_i \mathbb{J}_i) P_2(\cos \vartheta)},$$

- **For the photoionization by unpolarized or linearly polarized light**, the photo ion is aligned along the beam direction or photon polarization direction, respectively. From these expressions above for the components of the spin-polarization vector, it follows then that the spin of the Auger electron can be oriented only perpendicular to the (reaction) plane that contains the alignment axis and the electron emission direction. This has been called **dynamic polarization** (Klar 1980, Kabachnik 1981).
- **For the photoionization by circularly-polarized light, the photon ion is not only aligned but also oriented.** In this case, the spin components of the Auger electrons in the reaction plane do not vanish. This is an example of an **polarization transfer** process (Klar 1980).

## 7.2. In JAC partly-implemented processes

### 7.2.a. Rayleigh & Compton scattering of light (RayleighCompton)

#### Process, notations & application:

- Rayleigh or Compton scattering of photons at an atom or ion:  $A + \hbar\omega_i \longrightarrow A^{(*)} + \hbar\omega_f$  ... including  $\omega_i = \omega_f$
- Formal quantum notation:  $|\alpha_i \mathbb{J}_i M_i\rangle + \hbar\omega_i(\mathbf{k}_i, \lambda_i; \{\mathbb{M}_i\}) \longrightarrow |\alpha_f \mathbb{J}_f M_f\rangle + \hbar\omega_f(\mathbf{k}_f, \lambda_f; \{\mathbb{M}_f\})$   
 $\omega_i = \omega_f$  ... for elastic Rayleigh scattering
- Using JAC: Perform an `Atomic.Computation(..., process=Compton, processSettings=RayleighCompton.Settings(...), ...)` or call directly functions from the module `RayleighCompton`.

#### Motivation:

- The elastic scattering of light by the bound electrons of an atom or ion is commonly known as Rayleigh scattering and has been found a versatile tool in the past for studying the electronic structure of atoms and their (chemical) environment.
- The inelastic (Compton or Raman) scattering of light on atoms, ions or molecules is associated in contrast with an excitation of the target and has also been explored since the 1920s, both experimentally and theoretically.
- First theoretical investigations on the elastic scattering of photon by bound electrons date back to the mid-1930s. While initially rather simple approximations were applied, for instance based on atomic form factors, a more rigorous quantum-electrodynamical (QED) approach was developed later in terms of the relativistic second-order S-matrix amplitude.
- In JAC, the angular-differential cross sections for the elastic Rayleigh and inelastic Compton scattering are calculated and tabulated for all selected levels of  $(i \rightarrow f)$  line.

#### Second-order scattering amplitude

- Compton scattering amplitude: In second-order perturbation theory, the elastic Rayleigh and inelastic Compton process can be described in terms of the two-photon transition amplitudes for a transition from substates  $|\alpha_i \mathbb{J}_i M_i\rangle \rightarrow |\alpha_f \mathbb{J}_f M_f\rangle$  owing to the absorption and emission

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of photons with well-defined wave vectors  $\mathbf{k}_{i,f}$  and polarization vectors  $\boldsymbol{\epsilon}_{i,f}$ . This amplitude is given in terms of the electron-photon interaction operator by (Jahrsetz *et al.*, 2015)

$$\begin{aligned} \mathcal{M}_{fi}^{(\text{Compton})}(\alpha_f \mathbb{J}_f M_f, \alpha_i \mathbb{J}_i M_i; \mathbf{k}_f, \boldsymbol{\epsilon}_f, \mathbf{k}_i, \boldsymbol{\epsilon}_i) = & \sum_{\alpha_\nu \mathbb{J}_\nu M_\nu} \left[ \frac{\langle \alpha_f \mathbb{J}_f M_f | \hat{R}^\dagger(\mathbf{k}_f, \boldsymbol{\epsilon}_f) | \alpha_\nu \mathbb{J}_\nu M_\nu \rangle \langle \alpha_\nu \mathbb{J}_\nu M_\nu | \hat{R}(\mathbf{k}_i, \boldsymbol{\epsilon}_i) | \alpha_i \mathbb{J}_i M_i \rangle}{E_i + \omega_i - E_\nu} \right. \\ & \left. + \frac{\langle \alpha_f \mathbb{J}_f M_f | \hat{R}(\mathbf{k}_i, \boldsymbol{\epsilon}_i) | \alpha_\nu \mathbb{J}_\nu M_\nu \rangle \langle \alpha_\nu \mathbb{J}_\nu M_\nu | \hat{R}^\dagger(\mathbf{k}_f, \boldsymbol{\epsilon}_f) | \alpha_i \mathbb{J}_i M_i \rangle}{E_i - \omega_f - E_\nu} \right]. \end{aligned}$$

Here, the energies of the photons involved are related by  $\omega_f = \omega_i - E_f + E_i$  to each other and to the energies of the initial and final states of the atoms or ions.

- Owing to the time-ordering in (time-dependent) perturbation theory, there are obviously two contributions to this second-order amplitude for the *absorption-emission* and *emission-absorption* of the two photons involved.
- Although these second-order amplitudes can be readily written down, they are less useful for practical computations as they are based on different representations of the atomic and photon state(s); we therefore first need to perform a multipole expansion of the electron-photon interaction operators and to do all necessary simplifications.
- **Compton (scattering) amplitude for the scattering of photons with well-defined helicity and multipolarity:**

$$\mathcal{M}_{fi}^{(\text{Compton})}(\alpha_f \mathbb{J}_f M_f, \alpha_i \mathbb{J}_i M_i; \mathbb{M}_f, \lambda_f, \mathbb{M}_i, \lambda_i) = \dots \frac{\langle \alpha_f \mathbb{J}_f || \mathbb{O}^{(\mathbb{M}_f, \text{emission})} || \alpha_\nu \mathbb{J}_\nu \rangle \langle \alpha_\nu \mathbb{J}_\nu || \mathbb{O}^{(\mathbb{M}_i, \text{absorption})} || \alpha_i \mathbb{J}_i \rangle}{E_i + \omega_f - E_\nu} + \dots$$

### Angle-differential and total cross sections:

- **Angle-differential cross section for the Compton scattering of unpolarized light by initially unpolarized atoms:** For unpolarized incident plane-wave photons and if the polarization of the scattered photons as well as the magnetic sublevel population of the residual ions both

remain unobserved, the angle-differential cross section is given by

$$\begin{aligned} \frac{d\sigma^{(\text{Compton})}}{d\Omega}(\vartheta; \omega_i) \\ = \frac{1}{2} \frac{\omega_f}{\omega_i} \frac{1}{2J_i + 1} \sum_{\lambda_i \lambda'_i \lambda_f \lambda'_f} \sum_{M_i, M_f} c_{\lambda_i \lambda'_i} \mathcal{M}_{fi}^{(\text{Compton})}(\alpha_f \mathbb{J}_f M_f, \alpha_i \mathbb{J}_i M_i; \mathbb{M}_f, \lambda_f, \mathbb{M}_i, \lambda_i) \mathcal{M}_{fi}^{(\text{Compton})}(\alpha_f \mathbb{J}_f M_f, \alpha_i \mathbb{J}_i M_i; \mathbb{M}_f, \lambda'_f, \mathbb{M}_i, \lambda'_i) \end{aligned}$$

Unfortunately, there is no simple parametrization of this angle-differential cross section  $\frac{d\sigma^{(\text{Compton})}}{d\Omega}(\vartheta; \omega_i)$ .

- The angle-differential Rayleigh cross sections can be obtained from the formulas above by using  $|\alpha_f \mathbb{J}_f\rangle = |\alpha_i \mathbb{J}_i\rangle$  in the two-photon amplitudes and irreducible tensors.
- In JAC, the angle-differential Rayleigh and Compton scattering cross sections are calculated by default for all selected initial levels and ( $i \rightarrow f$ ) lines, if a proper spectrum is specified in `RayleighCompton.Settings.spectrum`.

### 7.2.b. Multi-photon ionization (MultiPhotonIonization)

#### Process, notations & application:

- **Multi-photon ionization:** of an atom or ion  $A + n \hbar \omega \longrightarrow A^{(*)} + e_p^-$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + \hbar \omega_1(\mathbb{M}_1) + \hbar \omega_2(\mathbb{M}_2) + \dots + \hbar \omega_n(\mathbb{M}_n) \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon \kappa\rangle$
- If the photon energy is smaller than the ionization threshold of the atom, it can still be ionized due to its simultaneous interaction with several photons. However, the probability generally decreases rapidly with the number of photons  $N$ , that are required to overcome the ionization threshold.
- Multi-photon ionization is studied mainly by means of intense, pulsed lasers. The ionization probability is for moderate intensities proportional to the  $N$ -th power of the intensity,  $I^N$ , while this power law breaks down at high intensities due to the ac-Stark effect.
- In first strong-field laser experiments, a rather large probability was found for a multiple ionization of atoms, relative to single-ionization probability, and gave rise to various speculation about *collective modes* in the electronic motion of atoms and ions.

### 7.2.c. Multi-photon excitation and decay (MultiPhotonDeExcitation)

#### Process, notations & application:

- Multi-photon excitation of an atom or ion:  $A + n \hbar \omega \longrightarrow A^*$
- Multi-photon decay of an atom or ion:  $A^* \longrightarrow A + n \hbar \omega$
- Formal quantum notation:  $|\alpha_i \mathbb{J}_i\rangle + \hbar \omega_1(\mathbf{k}_1, \lambda_1; \{\mathbb{M}_1\}) + \hbar \omega_2(\mathbf{k}_2, \lambda_2; \{\mathbb{M}_2\}) + \dots + \hbar \omega_n(\mathbf{k}_n, \lambda_n; \{\mathbb{M}_n\}) \longrightarrow |\alpha_f \mathbb{J}_f\rangle$  or  
 $|\alpha_i \mathbb{J}_i\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + \hbar \omega_1(\mathbf{k}_1, \lambda_1; \{\mathbb{M}_1\}) + \hbar \omega_2(\mathbf{k}_2, \lambda_2; \{\mathbb{M}_2\}) + \dots + \hbar \omega_n(\mathbf{k}_n, \lambda_n; \{\mathbb{M}_n\})$
- Using JAC: Perform an `Atomic.Computation(..., process=MultiPhoton, processSettings=MultiPhotonDeExcitation.Settings(...), ...)` or call directly functions from the module `MultiPhotonDeExcitation`.

#### Motivation:

- The two-photon decay of few-electron ions  $|\alpha_i \mathbb{J}_i\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + \hbar \omega_1 + \hbar \omega_2$  has often been considered, both experimentally and theoretically. For this two-photon decay, the total decay rate as well as the spectral (energy) distribution of the emitted photons have been explored.
- In JAC, only the initial and final levels are tabulated by default, together with their symmetry and total transition energy, while the proper flags must be set in `MultiPhotonDeExcitation.Settings` for all further cross section information.

#### Two-photon decay amplitude and total rate:

- Two-photon amplitude for the emission of photons with well-defined helicity and multipolarity:

$$\mathcal{M}_{fi}^{(2-\text{emission})}(\alpha_f \mathbb{J}_f M_f, \alpha_i \mathbb{J}_i M_i; \mathbb{M}_2, \lambda_2, \mathbb{M}_1, \lambda_1; \omega) = \frac{\langle \alpha_f \mathbb{J}_f \| \mathbb{O}^{(\mathbb{M}_2, \text{emission})} \| \alpha_\nu \mathbb{J}_\nu \rangle \langle \alpha_\nu \mathbb{J}_\nu \| \mathbb{O}^{(\mathbb{M}_1, \text{emission})} \| \alpha_i \mathbb{J}_i \rangle}{E_i - \omega - E_\nu}$$



- **Two-photon energy-differential emission rate:** For the transition of the atom from level  $|\alpha_i \mathbb{J}_i\rangle \rightarrow |\alpha_f \mathbb{J}_f\rangle$  can be written as (Goldman and Drake, 1988; Surzhykov 2005)

$$\begin{aligned}
\frac{dW}{d\omega_1}(\omega_1; \alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f) &= \frac{4\pi}{2J_i + 1} \frac{\omega_1 \omega_2}{(2\pi)^3 c^2} \sum_{M_i M_f} \sum_{\lambda_1, \lambda_2} |\mathcal{M}_{fi}(\alpha_f \mathbb{J}_f M_f, \alpha_i \mathbb{J}_i M_i; \mathbb{M}_2, \lambda_2, \mathbb{M}_1, \lambda_1; \omega)|^2 d\Omega_1 d\Omega_2 \\
&= \frac{(4\pi)^5}{2J_i + 1} \frac{\omega_1 \omega_2}{(2\pi)^3 c^2} \sum_{\mathbb{M}_1, \mathbb{M}_2} \sum_{\mathbb{J}_\nu} \left\{ \frac{1}{2J_\nu + 1} \left( |S_{\mathbb{J}_\nu}^{(2-\text{emission})}(\omega_1; \mathbb{M}_2, \mathbb{M}_1; fi)|^2 + |S_{\mathbb{J}_\nu}^{(2-\text{emission})}(\omega_2; \mathbb{M}_1, \mathbb{M}_2; fi)|^2 \right) \right. \\
&\quad \left. + 2 \sum_{\mathbb{J}'_\nu} (-1)^{J_\nu + J'_\nu + L_1 + L_2} \begin{Bmatrix} J_f & J'_\nu & L_1 \\ J_i & J_\nu & L_2 \end{Bmatrix} S_{\mathbb{J}_\nu}^{(2-\text{emission})}(\omega_2; \mathbb{M}_1, \mathbb{M}_2; fi)^* S_{\mathbb{J}'_\nu}^{(2-\text{emission})}(\omega_1; \mathbb{M}_2, \mathbb{M}_1; fi) \right\}
\end{aligned}$$

Here, the summations over  $\mathbb{J}_\nu$  and  $\mathbb{J}'_\nu$  are restricted, and their allowed values can be quite easily read-off from the standard selection rules for the multipole matrix elements.

- **Reduced two-photon (emission) amplitude:**

$$\begin{aligned}
S_{\mathbb{J}_\nu}^{(2-\text{emission})}(\omega; \mathbb{M}_2, \mathbb{M}_1; \alpha_f \mathbb{J}_f, \alpha_i \mathbb{J}_i) &\equiv S_{\mathbb{J}_\nu}^{(2-\text{emission})}(\omega; \mathbb{M}_2, \mathbb{M}_1; fi) \\
&= \sum_{\alpha_\nu} \frac{\langle \alpha_f \mathbb{J}_f \| \mathcal{O}^{(\mathbb{M}_2, \text{emission})} \| \alpha_\nu \mathbb{J}_\nu \rangle \langle \alpha_\nu \mathbb{J}_\nu \| \mathcal{O}^{(\mathbb{M}_1, \text{emission})} \| \alpha_i \mathbb{J}_i \rangle}{E_i - \omega - E_\nu} \\
&= S_{\mathbb{J}_\nu}^{(2-\text{absorption})}(???\omega; \mathbb{M}_2, \mathbb{M}_1; fi ???) = S_{\mathbb{J}_\nu}^{(2-\text{Compton})}(???\omega; \mathbb{M}_2, \mathbb{M}_1; fi ???)
\end{aligned}$$

- In JAC, the two-photon energy-differential emission rate are calculated and tabulated if ...

## 7.2.d. Coulomb excitation (CoulombExcitation)

Process, notations & application:

- Coulomb excitation of an atom or ion by fast, heavy ions:  $A + Z_p \longrightarrow A^* + Z_p$  (projectile remains unaffected)
- Formal quantum notation:  $|\alpha_i \mathbb{J}_i\rangle + Z_p(\beta_p) \longrightarrow |\alpha_f \mathbb{J}_f\rangle, \quad \beta_p = \frac{v_p}{c}$
- Using JAC: Perform an `Atomic.Computation(.., process=CoulExc, processSettings=CoulombExcitation.Settings(..), ..)` or call directly functions from the module `CoulombExcitation`.

Motivation

- In Coulomb excitation processes, one often assumes that a single electron from a (fast-moving) projectile ion with relative velocity  $\beta_p = v_p/c$  (with regard to the target) is excited by a target nucleus at rest.
- Of course, an analogue process occurs also when we consider the excitation of the target that may occur due to a fast-moving (bare) projectile ion, and in which case the role of the projectile and target need simply to be interchanged in all formulas below.
- The Coulomb excitation or ionization of the projectile (ion) will be affected, in addition, also by the target electrons (or, *vice versa*, the projectile electrons) but such rather small effects are often not considered in detail.
- The Coulomb excitation of few-electron ions in fast ion-atom collisions leads to the formation of excited ionic states. In typical ion-atom collision experiments, a *prefared* direction is hereby defined for the overall systems, both at storage rings and the EBIT, and this typically results in an *alignment of the excited ions* along this distinct direction.
- The first step in calculating energy- and angle-differential as well as total Coulomb excitation cross sections usually refers to determining the *impact-parameter dependence of the excitation probability  $P(b)$* . The weighted  $K$ -shell Coulomb excitation probability  $2\pi b P(b)$  is known to have a maximum in the vicinity of the  $K$ -shell radius. *This is quite similar as for the Coulomb ionization probabilities.*
- In JAC, the partial and total Coulomb excitation cross sections are calculated and tabulated by default for all selected pairs of initial and final levels.

## Coulomb excitation amplitude of few-electron projectile ions in position space

- **Coulomb (excitation) amplitude for the excitation of a (single) projectile electron:** If a one-electron projectile moves at a trajectory with impact parameter  $b$ , the amplitude for an excitation from level  $i \rightarrow f$  is given by Eichler & Meyerhof (1995, Eq. 6.2)

$$\mathcal{M}_{fi}^{(\text{Coul})}(b) = i \frac{\gamma Z_t e^2}{\hbar} \int dt \exp\left(\frac{i}{\hbar} (E_f - E_i) t\right) \int d^3r \psi_f^\dagger(\mathbf{r}) \frac{(1 - \beta_p \alpha_z)}{r'(t)} \psi_i(\mathbf{r}),$$

$$r'(t) = \sqrt{(x-b)^2 + y^2 + \gamma^2 (z - vt)^2}, \quad \beta_p = \frac{v_p}{c},$$

where  $\mathbf{r}$  is the coordinate of the projectile electron,  $r'(t)$  its (time-dependent) distance from the target nucleus, and where  $\psi_f(\mathbf{r})$  and  $\psi_i(\mathbf{r})$  denote the final- and initial-state eigenfunctions of the projectile electron with energies  $E_f$  and  $E_i$ .

- **Coulomb (excitation) amplitude for the excitation of a many-electron projectile (ion):** For an Coulomb excitation of a many-electron projectile from the sublevel  $|\alpha_i \mathbb{J}_i M_i\rangle \rightarrow |\alpha_f \mathbb{J}_f M_f\rangle$ , the excitation amplitude for given impact parameter  $b$  can then be written in first-order perturbation theory and natural units ( $\hbar = m_e = c = 1$ ) as:

$$\mathcal{M}_{fi}^{(\text{Coul})}(b; \alpha_f \mathbb{J}_f M_f, \alpha_i \mathbb{J}_i M_i) = i \gamma \alpha Z_t \int dt e^{i(E_f - E_i)t} \left\langle \alpha_f \mathbb{J}_f M_f \left| \sum_{k=1}^N \frac{1 - \beta_p \hat{\alpha}_z(k)}{r'_k(t)} \right| \alpha_i \mathbb{J}_i M_i \right\rangle.$$

- From this transition amplitude, the cross section for an excitation of the projectile from (sub-) level  $|\alpha_i \mathbb{J}_i\rangle \rightarrow |\alpha_f \mathbb{J}_f M_f\rangle$  can be obtained by integrating over all the impact parameters  $b$  and by taking the average over the initial magnetic sublevels  $M_i$ .
- In JAC, we do not use the Coulomb excitation amplitude in position space but make use of a representation of this amplitude in momentum space instead, see below.

## Coulomb excitation amplitude of few-electron projectile ions in momentum space:

- **Coulomb (excitation) amplitude for the excitation of a (single) projectile electron:** To evaluate Coulomb excitation cross sections for one- and many-electron atoms, it is often more convenient to work in momentum space and to express all properties in terms of the

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**Fourier transform of the Coulomb excitation amplitude in position space.** If we make use of the momentum transfer  $\mathbf{q} = (\mathbf{q}_b, q_z)$ , the time-integration in this Coulomb (excitation) amplitude can be carried out explicitly (Eichler & Meyerhof 1995, Eq. 6.4)

$$\mathcal{M}_{fi}^{(\text{Coulex})}(b) = \frac{i \nu_p}{\pi} \int \frac{d^2 q_b}{q_b^2 + (1 - \beta^2) q_0^2} e^{-i \mathbf{q}_0 \cdot \mathbf{b}} \mathcal{K}_{fi}^{(\text{Coulex})}(\mathbf{q}); \quad \mathcal{K}_{fi}^{(\text{Coulex})}(\mathbf{q}) = \langle \psi_f(\mathbf{r}) | (1 - \beta \alpha_z) e^{i \mathbf{q} \cdot \mathbf{r}} | \psi_i(\mathbf{r}) \rangle.$$

This formula removes the complicated time dependence in  $r'(t)$  and replaces it instead by an explicit wave number (momentum transfer) dependence.

- Note, that in momentum space, the (momentum-space) Coulomb excitation amplitude  $\mathcal{K}_{fi}^{(\text{Coulex})}(\mathbf{q})$  does no longer depend on the impact parameter  $b$  but now simply contains an integration over the momentum transfer  $\mathbf{q}$ , and which starts from the **minimum momentum transfer for the excitation** of the initial ion from level  $|i\rangle \rightarrow |f\rangle$ :  $q_0 = (E_f - E_i)/v_p$ .
- In momentum space, indeed, most Coulomb excitation cross sections can be expressed directly in terms of the momentum-space Coulomb excitation amplitude  $\mathcal{K}_{fi}^{(\text{Coulex})}(\mathbf{q})$ .
- **(Momentum-space) Coulomb excitation amplitude for many-electron projectiles:** This amplitude can be further simplified by using the Wigner-Eckert theorem (Surzhykov and Fritzsche, 2008)

$$\begin{aligned}
\mathcal{K}_{fi}^{(\text{Coulex})}(\mathbf{q}; \alpha_f \mathbb{J}_f M_f, \alpha_i \mathbb{J}_i M_i) &= \left\langle \alpha_f \mathbb{J}_f M_f \left| \sum_k^N (1 - \beta \alpha_z(k)) e^{i \mathbf{q} \cdot \mathbf{r}_k} \right| \alpha_i \mathbb{J}_i M_i \right\rangle \\
&= \sum_{LM} i^L Y_{LM}^*(\arccos(q_0/q), 0) \left\langle \alpha_f \mathbb{J}_f M_f \left| \sum_{k=1}^N (1 - \beta_p \hat{\alpha}_z(k)) j_L(q r_k) Y_{LM}(\vartheta_k, \varphi_k) \right| \alpha_i \mathbb{J}_i M_i \right\rangle, \\
&= \sum_t \frac{1}{\sqrt{2J_f + 1}} \langle J_i M_i, t (M_f - M_i) | J_f M_f \rangle \sum_L i^L Y_{L, M_f - M_i}^*(\arccos(q_0/q), 0) \left\langle \alpha_f \mathbb{J}_f \left\| H_{tL}^{(\text{Coulex})}(q) \right\| \alpha_i \mathbb{J}_i \right\rangle \\
&= \sum_t \frac{1}{\sqrt{2J_f + 1}} \langle J_i M_i, t (M_f - M_i) | J_f M_f \rangle \sum_L i^L Y_{L, M_f - M_i}^*(\arccos(q_0/q), 0) \\
&\quad \times (\delta_{tL} \langle \alpha_f \mathbb{J}_f \left\| \mathbb{K}^{(L, \text{Coulex}, jY)} \right\| \alpha_i \mathbb{J}_i \rangle - \beta_p \langle LM, 10 | tM \rangle \langle \alpha_f \mathbb{J}_f \left\| \mathbb{K}^{(t, \text{Coulex}, jT)} \right\| \alpha_i \mathbb{J}_i \rangle)
\end{aligned}$$

- In JAC, these (momentum-space) Coulomb excitation amplitudes  $\mathcal{K}_{fi}^{(\text{Coulex})}(\mathbf{q}; \alpha_f \mathbb{J}_f M_f, \alpha_i \mathbb{J}_i M_i)$  are implemented explicitly and the integration over  $q$  is typically performed by a Gauß-Legendre integration.

### Coulomb excitation operators and interaction strength:

- **Coulomb excitation (interaction) strength:** The (full) Coulomb excitation amplitude in momentum space above combines two kinds of reduced (many-electron) matrix elements due to the two interaction operators:  $\sum_{k=1}^N j_L(q r_k) Y_L(\vartheta_k, \varphi_k)$  of rank  $L$  as well as  $\sum_{k=1}^N j_L(q r_k) \hat{\alpha}(k) \cdot \mathbf{T}_{tL}(\vartheta_k, \varphi_k)$  of rank  $t$ , respectively. These two reduced matrix elements are the actual building blocks that are needed in order to represent and to explore the properties of the target or projectile excitation process.
- **Coulomb excitation (interaction) strength:** The building blocks for the computation of the Coulomb excitation and ionization amplitudes

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are given by the reduced matrix elements

$$\begin{aligned}\langle \alpha_f \mathbb{J}_f \parallel \mathbb{K}^{(L, \text{Coulex}, jY)} \parallel \alpha_i \mathbb{J}_i \rangle &= \left\langle \alpha_f \mathbb{J}_f \parallel \sum_{k=1}^N j_L(q r_k) Y_L(\vartheta_k, \varphi_k) \parallel \alpha_i \mathbb{J}_i \right\rangle \\ \langle \alpha_f \mathbb{J}_f \parallel \mathbb{K}^{(t, \text{Coulex}, jT)} \parallel \alpha_i \mathbb{J}_i \rangle &= \left\langle \alpha_f \mathbb{J}_f \parallel \sum_{k=1}^N j_L(q r_k) \hat{\boldsymbol{\alpha}}(k) \cdot \mathbf{T}_{tL}(\vartheta_k, \varphi_k) \parallel \alpha_i \mathbb{J}_i \right\rangle\end{aligned}$$

### Energie-differential, partial and total Coulomb excitation cross sections:

- **Partial Coulomb excitation cross section:** If the Fourier transformation from the position to momentum coordinates is performed, the integrals over the time  $t$  and the impact parameter  $b$  can be evaluated *analytically*. Then, the partial (excitation) cross section for the Coulomb excitation from (sub-) levels  $|\alpha_i \mathbb{J}_i\rangle \rightarrow |\alpha_f \mathbb{J}_f M_f\rangle$  can be written as (Surzhykov and Fritzsche, 2008)

$$\begin{aligned}\sigma(\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f M_f) &= \frac{2\pi}{2J_i + 1} \sum_{M_i} \int_0^\infty db \, b \left| \mathcal{M}_{fi}^{(\text{Coulex})}(b; \alpha_f \mathbb{J}_f M_f, \alpha_i \mathbb{J}_i M_i) \right|^2 \\ &= 2\pi \left( \frac{8\pi Z_t \alpha}{\beta_p} \right)^2 \frac{1}{2J_i + 1} \sum_{M_i} \int_{q_0}^\infty dq \frac{q}{(q^2 - q_0^2 \beta_p^2)^2} \left| \mathcal{K}_{fi}^{(\text{Coulex})}(\mathbf{q}; \alpha_f \mathbb{J}_f M_f, \alpha_i \mathbb{J}_i M_i) \right|^2.\end{aligned}$$

### Alignment of Coulomb-excited ions:

- **Alignment of Coulomb-excited ions in level  $|\alpha_f \mathbb{J}_f\rangle$ :** For a well-defined ion-atom collision axis, a Coulomb-excited ion in level  $|\alpha_f \mathbb{J}_f\rangle$  is generally aligned; this alignment depends on the partial Coulomb cross sections  $\sigma(\alpha_f \mathbb{J}_f M_f)$  can be described in terms of one (or several)

parameters  $\mathcal{A}_k(\alpha_f \mathbb{J}_f)$ :

$$\mathcal{A}_k(\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f) = \frac{\sqrt{2J_f + 1}}{\sigma(\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f)} \sum_{M_f} (-1)^{J_f - M_f} \langle J_f M_f J_f - M_f | k 0 \rangle \sigma(\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f M_f),$$

### 7.2.e. Photoionization & fluorescence (PhotoIonizationFluor)

#### Process, notations & application:

- **Photoionization** of an atom or ion with subsequent fluorescence emission:  $A + \hbar\omega \longrightarrow A^* + e_p^- \longrightarrow A^{(*)} + e_p^- + \hbar\omega'$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + \hbar\omega \longrightarrow |\alpha_e \mathbb{J}_e\rangle + |\varepsilon\kappa\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon\kappa\rangle + \hbar\omega_f(\mathbb{M})$
- Apart from astrophysical interest, emphasis on the photoionization of inner-shell electrons from multiple and highly charged ions and its subsequent photon emission arises also from the diagnostics of various *laboratory* plasmas.

#### Density operator of the fluorescence photon $\hbar\omega_f(\mathbb{M})$ :

- After the inner-shell photoionization, the photoion appears to be in an excited level  $|\alpha_e \mathbb{J}_e\rangle$  that decays subsequently to some energetically lower level  $|\alpha_f \mathbb{J}_f\rangle$  by the emission of a characteristic photon.
- **Characteristic photon density matrix in the helicity representation:** For this characteristic photon, the density matrix can be expressed in the form  $\langle \mathbf{k}_0 \lambda | \rho_\gamma | \mathbf{k}_0 \lambda' \rangle$ , in which  $\mathbf{k}_0 \equiv (\vartheta_0, \varphi_0)$  denotes the wave vector along the propagation direction of the fluorescence photon and  $\lambda = \pm 1$  its helicity. Note that the helicity representation of this density matrix also describes the photon polarization.

#### Stokes parameter of the fluorescence photon $\hbar\omega_f(\mathbb{M})$ :

- The density matrix of the fluorescence photon  $\hbar\omega_f(\mathbb{M})$  is usually parametrized in terms of the so-called Stokes parameters (Blum 1981; Balashov *et al.*, 2001)

$$\langle \mathbf{k}_0 \lambda | \rho_{\omega_f} | \mathbf{k}_0 \lambda' \rangle \equiv c_{\lambda, \lambda'} = \frac{1}{2} \begin{pmatrix} 1 + P_3 & -P_1 + iP_2 \\ -P_1 - iP_2 & 1 - P_3 \end{pmatrix},$$

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and which are utilized to characterize both the degree of linear ( $P_1$  and  $P_2$ ) and circular ( $P_3$ ) polarization of the light.

- Apart from the population of the excited sublevels  $|\alpha_e \mathbb{J}_e M_e\rangle$  and the angle  $\vartheta$  of the emitted photon, the degree of linear polarization of the characteristic x-ray radiation also depends on the total angular momenta of the excited level  $|\alpha_e \mathbb{J}_e\rangle$  and the final level  $|\alpha_f \mathbb{J}_f\rangle$  of the characteristic transition as well as on its *multipolarity*.

### 7.2.f. Photoionization & autoionization (PhotoIonizationAutoIon)

#### Process, notations & application:

- **Photo ionization** of an atom or ion with subsequent autoionization:  $A + \hbar\omega \longrightarrow A^{+,*} + e_p^- \longrightarrow A^{(*)} + e_p^- + e_a^-$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + \hbar\omega(\mathbb{M}) \longrightarrow |\alpha_r \mathbb{J}_r\rangle + |\varepsilon_p \kappa_p\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon_p \kappa_p\rangle + |\varepsilon_a \kappa_a\rangle$

#### Photo-Auger electron coincidence spectrometry:

- Photo-Auger electron coincidence spectrometry has been found a versatile tool for studying the structure and dynamics of atoms, molecules and solids. In atomic physics, for example, it has been utilized to **realize a so-called complete experiment in the photoionization of atoms as well as for exploring small effects of coherence and post-collision interactions in photo-induced Auger processes**.
- For a full analysis of the photoionization and subsequent autoionization process, both electrons should be detected in coincidence by measuring the energy and angular distributions, i.e. the **photo-Auger electron correlation function**.
- The **photoionization and subsequent autoionization process differs from the direct double photoionization** due to the formation of a (well-defined) intermediate ionic state and, thus, the emission of the two electrons can usually be considered within a two-step model.
- Investigations of the **magnetic circular dichroism in the photoemission from solids** have been found useful for studying the magnetic properties of solids.

#### Photo-Auger electron correlation function:

- **Sequential ionization of polarized targets:** There are three relevant directions that need to be distinguished in the sequential ionization of polarized targets: the direction of the photoelectron  $\mathbf{n}_p = (\vartheta_p, \varphi_p)$ , those of the Auger electron  $\mathbf{n}_a = (\vartheta_a, \varphi_a)$  and the direction of the



target polarization  $\mathbf{n}_t = (\vartheta_t, \varphi_t)$ . Because of these three relevant but independent directions, the (triple-differential) cross sections can be expanded in terms of **tripolar spherical harmonics**

$$\{\mathbb{Y}_{k_0}(\mathbf{n}_t) \otimes \{\mathbb{Y}_{k_1}(\mathbf{n}_p) \otimes \mathbb{Y}_{k_2}(\mathbf{n}_a)\}_k\}_{k'q'} = \langle \dots | \dots \rangle \langle \dots | \dots \rangle Y_{k_0 q_0}(\vartheta_t, \varphi_t) Y_{k_1 q_1}(\vartheta_p, \varphi_p) Y_{k_2 q_2}(\vartheta_a, \varphi_a)$$

### 7.2.g. Dielectronic recombination & fluorescence (DielectronicFluores)

#### Process, notations & application:

- **Dielectronic recombination** of an atom or ion with subsequent fluorescence:

$$A^{q+} + e_s^- \longrightarrow A^{(q-1)+*} \longrightarrow A^{(q-1)+(*)} + \hbar\omega \longrightarrow A^{(q-1)+(*)} + \hbar\omega + \hbar\omega'$$

- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + |\varepsilon_s \kappa_s\rangle \longrightarrow |\alpha_r \mathbb{J}_r\rangle \longrightarrow |\alpha_d \mathbb{J}_d\rangle + \hbar\omega_d(\mathbb{M}_d) \longrightarrow |\alpha_f \mathbb{J}_f\rangle + \hbar\omega_d(\mathbb{M}) + \hbar\omega_f(\mathbb{M}_f)$

### 7.2.h. Electron-impact (de-) excitation (ImpactExcitation)

#### Process, notations & application:

- **Electron-impact excitation** of an atom or ion:  $e_s^- + A \longrightarrow A^* + e_s^{-'}$

- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + |\varepsilon_i \kappa_i\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon_f \kappa_f\rangle$  or  $|\alpha_i \mathbb{J}_i\rangle + |\mathbf{p}_i m_i\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\mathbf{p}_f m_f\rangle$

- Electron-impact excitation is formally an inelastic scattering process in which a free incident electron  $|\varepsilon_i \kappa_i\rangle$  is scattered at the atom and leaves as final electron  $|\varepsilon_f \kappa_f\rangle$ .
- Electron-impact ionization cross sections are required in different field, for example, for calculating level populations and spectral line intensities of non-local-thermodynamic-equilibrium (non-LTE) plasmas.

## Computation of electron-impact processes:

- The electron-impact excitation of atoms is often described in first-order Born approximation, in which one can easily distinguish individual excitation channels. This approximation neglects however the coupling of the *continuum channels*, but which can later be incorporated perturbatively.
- **Distorted-wave Born approximation (DWBA):** We here make use of the DWBA that accounts for the distortion of the continuum orbitals due to (local) potential of the nucleus and all the electrons of the target. The DWBA typically gives better results than the *pure* Born or the Coulomb-Born approximation. The current implementation is similar to the codes by Zhang *et al.*(1989) as well as to the FAC code (Gu 2008).
- For a reliable evaluation of plasma parameters, such as temperature, density or level populations, the excitation cross sections need to be known quite accurately.
- **Computational methods for electron-impact excitation & ionization:** Three methods are commonly used in the literature to calculate electron-impact excitation and ionization cross sections: i) Coulomb-Born (CB) approximation, where the continuum orbital just represent a free electron in a Coulomb potential; ii) distorted-wave (DW) approximation that includes a more realistic potential for the continuum orbitals; iii) close-coupling (CC) approximation. **In the first two approximations, all excitation channels are treated independently and give overall rise to a scattering matrix which is not necessarily unitary.**
- **Electron-impact excitation in plasma diagnostics:** In plasma-diagnostics measurements, the population of the excited states, relevant for plasma diagnostics, is largely driven by the direct electron-impact excitation of the ground and low-lying metastable states, and followed by their radiative decay.

## Collision strength and cross sections:

- **Collision strength:** The name *collision strength* was first suggested by Seaton (1953). For a given (atomic target) transition  $\alpha_i \mathbb{J}_i \rightarrow \alpha_f \mathbb{J}_f$ , the collision strength is related to the cross section  $Q$  by

$$\Omega(\alpha_i \mathbb{J}_i, \alpha_f \mathbb{J}_f) = \frac{4\pi g_i}{\lambda_i^2} Q(\alpha_i \mathbb{J}_i, \alpha_f \mathbb{J}_f),$$

if the continuum orbitals are normalized per unit energy. The statistical weight  $g_i$  of level  $|\alpha_i \mathbb{J}_i\rangle$  ensures the detailed-balance relation  $\Omega(\alpha_i \mathbb{J}_i, \alpha_f \mathbb{J}_f) = \Omega(\alpha_f \mathbb{J}_f, \alpha_i \mathbb{J}_i)$ .

- **Collision strength:** If  $|\varepsilon_i, \kappa_i\rangle$  and  $|\varepsilon_f, \kappa_f\rangle$  denote the partial waves of the incident and scattered electrons, the collision strength is given by

$$\Omega_{if} = \sum_{\kappa_i, \kappa_f} \sum_{J_t} [J_t] \left| \langle (\mathbb{J}_f, \epsilon_f, \kappa_f) \mathbb{J}_t M_t | V^{(e-e)} | (\mathbb{J}_i, \epsilon_i, \kappa_i) \mathbb{J}_t M_t \rangle \right|^2.$$

- In the JAC program, the standard decomposition of the electron-electron interaction matrix elements is utilized to compute the collision strength.
- **Normalization of continuum orbitals:** A continuum orbital is said to be **normalized per unit energy** if it has an asymptotic amplitude  $\sqrt{k/\varepsilon}$  (or,  $\sqrt{2/k}$  in the non-relativistic limit) or, equivalently, if it fullfills, equivalently,

$$\int dr [P_\varepsilon(r) P_{\varepsilon'}(r) + Q_\varepsilon(r) Q_{\varepsilon'}(r)] = \pi \delta(\varepsilon - \varepsilon').$$

- For electron-impact processes, the collision strength  $\Omega(\alpha_i \mathbb{J}_i, \alpha_f \mathbb{J}_f)$  is typically of the order of unity, though still large variation may still occur owing to its dependence on the particular ion as well as the transition and impact energy. Therefore, it is usually not justified to assume  $\Omega = 1$  as sometimes done in astrophysical codes.
- **Simplification for heavy ions with open  $d$ - and  $f$ -shells:** For heavy and ionized atoms with open  $d$ - and  $f$ -shells, the fine-structure of just (mixed) single initial- and final-state configurations may give rise to thousands of (fine-structure) transitions, for which the radial integrals often depend only weakly on the transition energy. For such transition arrays, the total cross section computations can be simplified considerably by either neglecting the weak dependence of the radial integrals or by some simple interpolation.
- Accurate calculations of electron-impact excitation of atom and ions are overall still a challenge for contemporary electronic methods, such as convergent close-coupling, R-matrix close-coupling or time-dependent close-coupling scheme, and even by using the largest (parallel) computers.

### Collisional-excitation rate coefficients:

- The rate coefficient of collisional excitation is given by

$$\alpha^{(\text{impact:excitation})} = \int_{\Delta E}^{\infty} dv f(v) \sigma^{(\text{impact-excitation})}(v)$$

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where  $f(v)$  is velocity distribution of electrons which is typically assumed to have a Maxwellian distribution with electron temperature  $T_e$ ,  $\sigma^{(\text{impact-excitation})}(\alpha_i \mathbb{J}_i, \alpha_f \mathbb{J}_f)$  is the electron-impact excitation cross-section from level  $i \rightarrow f$  at velocity  $v$ , and  $\Delta E = E_f - E_i$  is the excitation energy.

➤ For a given energy of the incident electrons, the collisional-excitation rate coefficient can be expressed also as

$$\alpha^{(\text{impact-excitation})} = \sqrt{\frac{\pi}{E}} \int_{\Delta E}^{\infty} dE E \sigma^{(\text{impact-excitation})}(E) \exp\left(-\frac{E}{T_e}\right).$$

### 7.2.i. Electron-impact excitation & autoionization (ImpactExcitationAutoIon)

#### Process, notations & application:

- **Electron-impact excitation with subsequent autoionization** of an atom or ion:  $A + e_s^- \longrightarrow A^* + e_s^{-'} \longrightarrow A^{+(*)} + e_s^{-'} + e_a^-$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + |\varepsilon_i \kappa_i\rangle \longrightarrow |\alpha_r \mathbb{J}_r\rangle + |\varepsilon_r \kappa_r\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon_r \kappa_r\rangle + |\varepsilon_a \kappa_a\rangle$
- Electron-impact excitation with subsequent autoionization occurs frequently in plasma and has been utilized to study the autoionization of atoms and ions.

### 7.2.j. Radiative-Augur decay (RadiativeAuger)

#### Process, notations & application:

- **Radiative-Augur (autoionization)** of an atom or ion:  $A^{q+*} \longrightarrow A^{(q+1)+,(*)} + (e_a^- + \hbar\omega)$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + \hbar\omega(\mathbb{M}) + |\varepsilon_a \kappa_a\rangle$
- The radiative Auger process results in the simultaneous emission of an electron and photon, and mainly occurs for inner-shell excited atoms and ions. There are no characteristic electron or photon lines associated with this process, since the transition energy is shared between the photon and the electron.

### 7.2.k. Multi-photon double ionization (MultiPhotonDoubleIon)

#### Process, notations & application:

- **Multi-photon double ionization:** of an atom or ion:  $A + n \hbar \omega \longrightarrow A^{(*)} + e_{p_1}^- + e_{p_1}^-$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + \hbar \omega_1 (\mathbb{M}_1) + \hbar \omega_2 (\mathbb{M}_2) + \dots + \hbar \omega_n (\mathbb{M}_n) \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon_1 \kappa_1\rangle + |\varepsilon_2 \kappa_2\rangle$
- The energy of several photons together can also lead to the emission of two electrons: This emission can occur either sequential, direct (non-sequential) or via some given resonances of the atom.
- In this section, we only consider the direct (non-sequential) multi-photon double ionization in which the energy of all photons is continuously shared by the two emitted electrons:  $n \hbar \omega = E_f - E_i + \varepsilon_1 + \varepsilon_2$ .

### 7.2.l. Internal conversion (InternalConversion)

#### Process, notations & application:

- **Internal conversion** of an atom or ion:  $A^{q+} + \text{nucleus}^* \longrightarrow A^{(q+1)+*} + e_c^-$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + \text{nucleus}^* (\{\mathbb{M}\}, \mathcal{E}) \longrightarrow |\alpha_f \mathbb{J}_f\rangle + e_c^-(\mathbf{p}, m_s)$  or  
 $|\alpha_i \mathbb{J}_i\rangle + \text{nucleus}^* (\{\mathbb{M}\}, \mathcal{E}) \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon \kappa\rangle, \quad \varepsilon = \mathcal{E} - (E_i - E_f)$
- Here, the excited nucleus decays under the release of the energy  $\mathcal{E}$  by (several) multipoles  $\{\mathbb{M}\}$  and , leading to an emitted electron with well-defined kinetic energy  $\varepsilon$ .
- **Internal conversion:** formally refers to the decay an excited nucleus to its ground or some lower-lying level, and where the excitation energy is given to a (bound) electron, leading to an ionization of the atom or ion. In general, the internal (energy) conversion competes with the gamma-ray emission as well as the formation of a electron-positron pair for nuclear excitation energies above 1.022 MeV.
- Often, the continuum wave of the outgoing electron has been calculated in the (final-state) atomic potential with a vacancy in the inner shell. Sometimes, however, this hole has been disregarded, leading to differences in the internal conversion coefficients as large as  $\sim 15\%$  for low kinetic energies  $\varepsilon \sim 1$  keV and  $L = 5$  transitions.

### Nuclear decay by internal conversion:

- The internal conversion is more than a simple photoionization process, in which an emitted  $\gamma$  photon leads to the emission of electrons, but refers to an **alternative de-excitation of excited nuclei**.
- **Electron conversion coefficient:** In the internal conversion process, an atomic electron is ejected from one of the atomic shells. The electron conversion coefficient is defined as the probability ratio for emitting an atomic electrons from shell  $x$  to the emission of a  $\gamma$ -ray,  $a_x = P_x / P_\gamma$ .
- The kinetic energy of the emitted electron  $\varepsilon = \Delta E_{\text{nuc}} - E_b$  can be deduced from the nuclear transition energy, the binding energy of the atomic electron and the recoil energy of the emitting atom, which is typically very small. Transitions involving conversion electrons are only possible if  $\varepsilon > 0$ .
- The internal conversion process (ICP) has been found a versatile tool for studying nuclear structure. In particular, the measurement and analysis of **conversion electron spectra reveal possible transitions between nuclear levels**. From the comparison of these spectra with calculated internal conversion coefficients, it is **often possible to assign a unique multipolarity to the nuclear gamma radiation** and, hence, a total angular momentum and parity to the excited nuclear states.
- The internal conversion coefficients can provide detailed information about the atomic nucleus if theoretical and experimental values are compared for transitions of different multipolarity and mixing ratios. A detailed knowledge of these coefficients is needed, for instance, for deriving absolute transition rates and for the normalization of decay schemes, for Mössbauer spectroscopy, for nuclear reaction computations as well as for the calculation of the decay heat of the fuel cells of nuclear reactors.
- **The internal conversion process is largely independent of the nuclear structure, although a non-zero probability requires a finite nuclear size and a finite probability of the electron to be found inside of the nucleus**. For highly suppressed transitions, this finite-size effect can become quite significant. For these transitions, Pauli developed formulae to correct the theoretical conversion coefficient for both, the electric and magnetic multipolarities.

### General theory:

- The internal conversion can be described as QED process by the retarded interaction of charges.
- In particular, the internal conversion is a second-order quantum-electrodynamical process in which a virtual photon is exchanged between the nucleus (proton) and an electron.

- **Internal conversion amplitude:** The (single-electron) matrix element of the conversion transition can be written in terms of the Hamiltonian (Listengarten 1961; Band *et al.*, 2002)

$$\langle f | \mathbb{H}^{(\text{conversion})} | i \rangle = -e \int d^3\mathbf{r} \psi_f^*(\mathbf{r}) [\Phi(\mathbf{r}) + \boldsymbol{\alpha} \cdot \mathbf{A}(\mathbf{r})] \psi_i(\mathbf{r}),$$

where  $\Phi(\mathbf{r})$  and  $\mathbf{A}(\mathbf{r})$  are the time-independent parts of the scalar and vector potentials of the electro-magnetic field that arise due to the **nuclear transition charges and currents**, and by including the retardation of these interactions.

- The retarded potentials in the internal conversion amplitude above are

$$\Phi(\mathbf{r}, t) = e^{-i\omega t} \Phi(\mathbf{r}), \quad \Phi(\mathbf{r}) = \int d^3\mathbf{R} \frac{\rho(\mathbf{R})}{|\mathbf{r} - \mathbf{R}|} e^{ik|\mathbf{R} - \mathbf{r}|} = \sum_{LQ} \Phi_{LQ}(\mathbf{r})$$

$$\mathbf{A}(\mathbf{r}, t) = e^{-i\omega t} \mathbf{A}(\mathbf{r}), \quad \mathbf{A}(\mathbf{r}) = \frac{1}{c} \int d^3\mathbf{R} \frac{\mathbf{J}(\mathbf{R})}{|\mathbf{R} - \mathbf{r}|} e^{ik|\mathbf{R} - \mathbf{r}|} = \sum_{LQ} \mathbf{A}_{LQ}(\mathbf{r})$$

with the wavenumber  $k = \omega/c = \mathcal{E}/\hbar c$ , and where  $\mathbf{r} = (r, \vartheta, \varphi)$  and  $\mathbf{R}$  are the coordinates of the electron and nucleus (proton), respectively.

- Here, we shall not go into further details for the representation of these potentials but just note that these interaction potentials can be written as a sum over multipole contributions as well as in terms of the nuclear charge density  $\rho(\mathbf{R})$  and current density  $\mathbf{J}(\mathbf{r})$  of the bound electrons, respectively.

### Internal conversion coefficients:

- **Internal conversion coefficient:** The internal conversion coefficient is defined as the ratio of the number of electrons  $N_e$  ejected from the atomic shell  $(n\ell)$  to the number of gamma quanta  $N_\gamma$  leaving the atom during the same time:

$$\alpha^{(\text{conversion})}(n\ell) = \frac{N_e}{N_\gamma}, \quad \alpha^{(\text{conversion})}(\alpha_i \mathbb{J}_i) = \sum_f \alpha^{(\text{conversion})}(n\ell)$$

The total internal conversion coefficient of a given (initial) level  $\alpha^{(\text{conversion})}(n\ell)$  of the nucleus is the sum of partial conversion coefficient associated with all possible ionization channels (lines) of the atom or ion.

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- Internal conversion coefficients have been tabulated in a larger number of tables and compilations; in these tabulations, the screening of the nuclear electric field by the atomic electrons was often treated in the framework of either some statistical Thomas-Fermi-Dirac or Dirac-Fock-Slater models.
- For different  $x_\alpha$  pre-factors in the Dirac-Fock-Slater (type) potentials, the differences in the calculated conversion coefficients are typically small  $\sim 1\%$  for  $K$ -shell electrons but may increase to  $\sim 70\%$  for outer electrons or even larger for low kinetic energies of the conversion electron ( $\leq 1$  keV). The question of whether the electronic hole should be taken into account into the atomic potential has been explored previously but seem to have a minor role only.
- A new internal conversion coefficient database BRICC has been developed that includes various tabulations of internal electron conversion (IECC), internal electron-positron pair conversion coefficients (IPCC) as well as the *electronic* factors  $\Omega(E0)$ .
- Theoretical conversion coefficients  $\alpha^{(\text{internal-conversion})}$  are associated with two kinds of uncertainties: (i) due to the physics model that is used in the computations and (ii) due to the spline interpolation that is frequently used in order to generate the tabulations. The physics model is typical based on assumptions about the atomic nucleus and the electron density.

### 7.2.m. Electron capture with nuclear decay

#### Process, notations & application:

- **Electron capture** of an atom or ion by nuclear decay:  $A(Z+1) + e^- \longrightarrow A(Z) + \nu_e$
- **Formal quantum notation:**  $E_\nu = Q^+ - E_f - E_b > 0$
- Here, the nucleus decays by capturing an atomic electron (often from the K-shell), and under the emission of an electron neutrino. Here,  $Q^+$  is the energy difference due to the rest masses of the parent and daughter nucleus,  $E_f$  the energy of the final nuclear state of the daughter nucleus and  $E_b$  the binding energy of the captured electron. The release energy  $E_\nu$  will be shared between the emitted neutrino and, possible, some bremsstrahlungs photon or the shake-up/shake-off of a valence electron.
- For allowed nuclear transitions, nearly all vacancies occur in the  $ns$  shells of the atoms or ions, i.e. in the K,  $L_1$ ,  $M_1$ , ... shells.



## 7.3. Further processes, not yet considered in JAC

### 7.3.a. Coulomb ionization (CoulombIonization)

#### Process, notations & application:

- Coulomb ionization by fast, heavy ions:  $A^{(q+1)+} + Z_p \longrightarrow A^{(q+1)+(*)} + e^- + Z'_p$
- Formal quantum notation:  $|\alpha_i \mathbb{J}_i\rangle + Z_p(\dots) \longrightarrow |\alpha_f \mathbb{J}_f\rangle + e^-$
- The Coulomb ionization of projectile ions by target atoms leads to a different charge state of the ions and often determines their lifetime in storage rings.
- If one of the projectile electrons is ionized, the ion is lost from the ring since it will be misbent by subsequent steering magnets.

#### Perturbative calculation of ionization cross sections:

- Semi-relativistic treatment of the Coulomb ionization: In their semi-relativistic treatment, Anholt and Becker (1987) distinguish three contributions to the ionization cross sections due to Coulomb ionization, transverse ionization, and the ionization due to the spin-flip of the ionized electron. These three contributions have different asymptotic behaviour in ultra-relativistic collisions. — In this approach, the Coulomb ionization cross section is the same as in the plane-wave Born approximation for nonrelativistic projectiles.
- The Coulomb part is dominant at nonrelativistic energies, while the transverse part mainly arises from the magnetic interactions at large impact parameters.
- The cross section for the (projectile) ionization of a  $1s$  electron increases as  $\ln \gamma$  due to the transverse interaction for high projectile energies and if the target screening is neglected. This transverse interaction is however screened and reduced if the charge of the projectile is smaller than those of the target,  $Z_p < Z_t$ .
- For fast collisions and for impact parameters  $b$  comparable with, or larger than the ionic  $K$ -shell radius, the perturbation of the target atom by the projectile can be treated in first-order time-dependent perturbation theory, even for high- $Z$  projectiles.
- Perturbation theory breaks generally down for high- $Z$  projectiles (or high- $Z$  targets). For these high- $Z$  projectiles, an approximate scaling rule for the ionization probability at high relativistic energies ( $\gamma > 5$ ) and fixed impact parameters has been established by numerical calculations. For  $b = 0$ , for example, the ionization probability  $P(b = 0) = 1.8 \times 10^{-4} Z_p^2$  is independent of the target charge within about  $\pm 10\%$  up to the heaviest elements.

### 7.3.b. Bremsstrahlung (BremsStrahlung)

#### Process, notations & application:

- Bremsstrahlung emitted by an electron in the field of an atom or ions  $A^{q+} + e_i^- \longrightarrow A^{(q)+*} + e_f^- + \hbar\omega$
- Formal quantum notation:  $|\alpha_i \mathbb{J}_i\rangle + |\varepsilon_i \kappa_i\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon_f \kappa_f\rangle + \hbar\omega(\mathbb{M})$
- Bremsstrahlung arises generally if a charged particle is decelerated or deflected by another charged particle, often an electron or nucleus, and if parts of its kinetic energy is converted into radiation.
- More generally speaking, bremsstrahlung or braking radiation refers to any radiation that arises due to the deceleration of charged particle, including synchrotron radiation from relativistic particles or cyclotron radiation from non-relativistic particles.
- In atomic physics, bremsstrahlung often refers to the radiation from electrons that are slowed down in matter.
- Atomic bremsstrahlung: This term sometimes refers also to the polarized radiation that arise from the Coulomb field of the incident charged particle.
- Bremsstrahlung has a continuous spectrum and a peak intensity that shifts toward higher frequencies with an increasing energy transfer from the decelerated particles.

### 7.3.c. Inverse bremsstrahlung

#### Process, notations & application:

- Inverse bremsstrahlung of a charged particle in the field of an atom or ions  $\hbar\omega + A^{q+} + e_i^- \longrightarrow A^{(q)+} + e_f^-$
- Formal quantum notation:  $\hbar\omega(\mathbb{M}) + |\alpha_i \mathbb{J}_i\rangle + |\varepsilon_i \kappa_i\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon_f \kappa_f\rangle$
- In the inverse bremsstrahlung process, a free electron gains kinetic energy due to the absorption of a photon.

### 7.3.d. Radiative recombination & fluorescence

#### Process, notations & application:

- Radiative recombination & fluorescence of an ion:  $A^{q+} + e_s^- \longrightarrow A^{(q-1)+,*} + \hbar\omega_r \longrightarrow A^{(q-1)+} + \hbar\omega_r + \hbar\omega_f$
- Formal quantum notation:  $|\alpha_i \mathbb{J}_i\rangle + |\varepsilon_i \kappa_i\rangle \longrightarrow |\alpha_r \mathbb{J}_r\rangle + \hbar\omega_r(\mathbb{M}_r) \longrightarrow |\alpha_f \mathbb{J}_f\rangle + \hbar\omega_r(\mathbb{M}_r) + \hbar\omega_f(\mathbb{M}_f)$
- In the photorecombination & fluorescence process, also known radiative recombination (RR) or radiative electron capture (REC), an electron is captured by the ion into an excited state that subsequently decay under fluorescence emission.

### 7.3.e. Resonant two-color (two-photon, single-electron) ionization

#### Process, notations & application:

- Resonant two-color (two-photon, single-electron) ionization of an ion:  $A + \hbar\omega + \hbar\omega' \longrightarrow A^* + \hbar\omega' \longrightarrow A^{+*} + e_p^-$
- Formal quantum notation:  $|\alpha_i \mathbb{J}_i\rangle + \hbar\omega(\mathbb{M}) + \hbar\omega'(\mathbb{M}') \longrightarrow |\alpha_r \mathbb{J}_r\rangle + \hbar\omega'(\mathbb{M}') \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon \kappa\rangle$
- In the resonant two-color photoionization, an atom or ion is resonantly excited by one photon and subsequently ionized by a second photon of the same or some different frequency.

### 7.3.f. Interference of multi-photon ionization channels (MultiPhotonInterference)

#### Process, notations & application:

- Interference of *non-resonant* one-photon ( $2\omega$ ) and *resonant* two-photon ( $\omega$ ) ionization of an atom or ion:

$$A + \left[ \begin{array}{c} \hbar(2\omega) \\ \hbar\omega + \hbar\omega \end{array} \right] \longrightarrow A^{+(*)} + e_p^-$$

## 7. Atomic processes

- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + \left[ \begin{array}{c} \hbar(2\omega) \\ \hbar\omega + \hbar\omega \end{array} \right] \longrightarrow |\alpha_i \mathbb{J}_i\rangle + |\varepsilon\kappa\rangle$
- The two-pathway quantum interference has been explored in order to better understand the **quantum control of atoms and ions**.
- Such quantum control studies were initially stimulated by the small fraction of the second harmonic that typically arise at XFEL and that cannot so easily be filtered out, though it may strongly influence the experimental from two-photon ionization experiments. Despite of the rather small intensity of the second harmonic, the ionization by photons with frequency  $2\omega$  (second harmonic, first-order process) can readily compete with, or even dominate, the two-photon ionization by photons with frequency  $\omega$ .

### Interference of non-res. one-photon ( $2\omega$ ) and resonant two-photon ( $\omega$ ) ionization:

- **Experimental signals:** For a linearly-polarized bichromatic beam, the interference due to different quantum paths manifests itself in an asymmetry of the photoelectron angular distributions (PAD) with respect to the plane that is perpendicular to the electric field of the incident radiation; cf. Grum-Grzhimailo *et al.*(2015).
- **Photoelectron angular distribution (PAD):** In the dipole approximation and for an isotropic target, this angular distribution must be axially symmetric with regard to the polarization direction

$$\frac{dW}{d\Omega} = \frac{W_o}{4\pi} \left[ 1 + \sum_{k=1}^K \beta_k P_k(\cos \vartheta) \right]$$

where  $\vartheta$  is the angle of the photoelectron with regard to the polarization direction (**E**-field) and  $\beta_k$  the corresponding anisotropy parameters, which can expressed in terms of the photoionization amplitudes.

- For the interference of one-photon and two-photon ionization paths, all terms with  $k = 1, 2, 3, 4$  must be taken into account in the PAD above. The odd polynomials arises from the photoelectron partial waves with opposite parities due to the ininteraction of the atom with the fundamental and second harmonic of the radiation.
- The quantum interference of absorption amplitudes with an even and odd number of photons, such as  $(\omega + 2\omega)$  photoionization, does not modify the total yield but only affect the angle-resolved observations. The interference between an even or odd number of photons lead to photoelectron partial waves with opposite parities.

### 7.3.g. Two-color multi-photon interference ionization (TwoColorInterferenceIon)

#### Process, notations & application:

- Two-color multi-photon interference ionization of an ion:  $A + m \hbar \omega + n \hbar \omega' \longrightarrow A^{+*} + e_p^-$
- Formal quantum notation:  $|\alpha_i \mathbb{J}_i\rangle + m \hbar \omega (\mathbb{M}) + n \hbar \omega' (\mathbb{M}') \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon \kappa\rangle$
- In the two-color multi-photon interference ionization, an atom or ion emits an electron whose angular distribution is affected by the interference of different ionization pathes.

### 7.3.h. Double photoionization (DoublePhotoIonization)

#### Process, notations & application:

- Simultaneous double photoionization of an atom or ion:  $A^{q+*} + \hbar \omega \longrightarrow A^{(q+2)+(*)} + (e_{p_1}^- + e_{p_2}^-)$
- Formal quantum notation:  $|\alpha_i \mathbb{J}_i\rangle + \hbar \omega (\mathbb{M}) \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon_1 \kappa_1\rangle + |\varepsilon_2 \kappa_2\rangle$
- The (direct) double photoionization of atoms has attracted much interest during the last decades as the emission of a second electron is caused by the electron-electron interaction.
- The recent advancements of free electron lasers (FEL) has lead to intense short pulses with rather high photon energies, and which may help investigate the (direct) double photoionization process for multiply charged ions.

#### Cross sections:

- Triple-differential cross section for the double ionization of polarized atoms: If one expresses the cross section in terms of the solid angles  $\Omega_{1,2}$  of the two electrons as well as the energy  $E_1$  of one of the emitted electrons, this (five-fold differential) cross section can be written as trace over the final-state density matrix  $\rho_f$

$$\frac{d^3 \sigma}{d\Omega_1 d\Omega_2 dE_1} = c \text{Tr} \rho_f = c \text{Tr} (R \rho_i R^+) = c \text{Tr} (D \rho^\gamma \rho_i R^+),$$

where  $\rho^\gamma$  is the photon density matrix and  $c$  just a kinematically determined normalization constant. Usually, this expression is evaluated in a basis with well-defined symmetry  $\mathbb{J}$ .

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- Within the density matrix and the statistical tensor formalism, the triply differential cross section for single-photon double ionization from above can be written as (Berakdar and Kabachnik, 2005)

$$\frac{d^3 \sigma}{d\Omega_1 d\Omega_2 dE_1} = c \sum_{\alpha\alpha' \mathbb{J}\mathbb{J}' kq} \rho_{kq}^f(\alpha\mathbb{J}, \alpha'\mathbb{J}') \epsilon_{kq}^*(\alpha\mathbb{J}, \alpha'\mathbb{J}')$$

where  $\rho_{kq}^f(\alpha\mathbb{J}, \alpha'\mathbb{J}')$  is the final-state statistical tensor of the “ion + two electrons” with symmetries  $\mathbb{J}$  and  $\mathbb{J}'$  and where  $\epsilon_{kq}(\alpha\mathbb{J}, \alpha'\mathbb{J}')$  denotes the efficiency tensor of the detector system to record these final states. Here, the summation over  $\alpha, \alpha'$  refers to all other quantum numbers which remain unobserved.

- **Final-state density matrix:** The final-state density matrix of the “photoion + two free electrons” can be expressed in terms of the statistical tensors of the total initial state “atom + photon” as well as the reduced electron-photon (photoionization) amplitudes (Berakdar and Kabachnik, 2005)

$$\begin{aligned} \rho_{kq}^f(\alpha\mathbb{J}, \alpha'\mathbb{J}') &= \sum_{k_o q_o k_\gamma q_\gamma} [k_o, k_\gamma]^{1/2} \langle k_o q_o, k_\gamma q_\gamma | kq \rangle \left\{ \begin{matrix} J_o & 1 & J \\ J_o & 1 & J' \\ k_o & k_\gamma & k \end{matrix} \right\} \rho_{k_o q_o}(\alpha_o \mathbb{J}_o) \rho^\gamma(P_1, P_2, P_3) \\ &\quad \times \langle \alpha\mathbb{J} || \mathbb{O}^{(\text{photoionization})} || \alpha_o \mathbb{J}_o \rangle \langle \alpha'\mathbb{J}' || \mathbb{O}^{(\text{photoionization})} || \alpha_o \mathbb{J}_o \rangle^* . \end{aligned}$$

- Although such expressions for the (triple-differential) cross sections appear to be rather complex, they can be handle quite easily in JAC.

### 7.3.i. Double-Auger decay (DoubleAuger)

#### Process, notations & application:

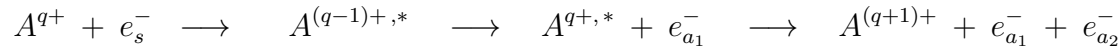
- **Simultaneous double Auger electron emission (autoionization)** of an atom or ion:  $A^{q+*} \longrightarrow A^{(q+2)+(*)} + (e_{a_1}^- + e_{a_2}^-)$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon_1 \kappa_1\rangle + |\varepsilon_2 \kappa_2\rangle$
- An autoionization of an atom may occur also via the simultaneous emission of two electrons which share the overall excess energy.

- The (direct) double Auger (DA) emission is an important second-order processes in which two electrons are ejected simultaneously from an inner-shell excited atom. A first evidence of the direct DA was found in *K*-shell photoionization experiments of neon with photon energies between 867 and 913 eV by Krause and coworkers (1965) as well as for argon.
- For the decay of a Ne 1s-hole, already Carlson and Krause (1965) found an about 8 % portion of Ne<sup>3+</sup> ions due to the simultaneous double Auger decay. This portion occurs in contrast to the  $\lesssim 1$  % prediction from a simple **shake-off** model.

### 7.3.j. Resonant excitation with sequential double autoionization (RESDA):

#### Process, notations & application:

- Resonant excitation with sequential double autoionization (RESDA) of an ion: ...



- Formal quantum notation:  $|\alpha_i \mathbb{J}_i\rangle + |\varepsilon_i \kappa_i\rangle \longrightarrow |\alpha_m \mathbb{J}_m\rangle \longrightarrow |\alpha_n \mathbb{J}_n\rangle + |\varepsilon_1 \kappa_1\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon_1 \kappa_1\rangle + |\varepsilon_2 \kappa_2\rangle$
- Lithium-like ions are perhaps the simplest ions in which a RESDA can occur. This process can be observed by measuring the ratio of the numbers of trapped ions of two neighbouring charge states as a function of electron energy within an electron beam ion trap.

### 7.3.k. Resonant excitation with direct double autoionization (REDDA):

#### Process, notations & application:

- Resonant excitation with direct double autoionization (REDDA) of an ion: ...



- Formal quantum notation:  $|\alpha_i \mathbb{J}_i\rangle + |\varepsilon_i \kappa_i\rangle \longrightarrow |\alpha_m \mathbb{J}_m\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon_1 \kappa_1\rangle + |\varepsilon_2 \kappa_2\rangle$
- In the REDDA process, an electron is resonantly captured into a (doubly) excited level that subsequently decays by the simultaneous emission of two electrons, a so-called double autoionization, and where the two electrons share the overall excess energy.

### 7.3.l. Radiative double electron capture

#### Process, notations & application:

- Radiative double electron capture (RDEC) of an ion:  $A^{q+*} + e_1^- + e_2^- \longrightarrow A^{(q-2)+} + \hbar\omega$
- Formal quantum notation:  $|\alpha_i \mathbb{J}_i\rangle + |\varepsilon_1 \kappa_1\rangle + |\varepsilon_2 \kappa_2\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + \hbar\omega(\mathbb{M})$
- The radiative double electron capture (RDEC) is the time-reversed process of the single-electron direct double photoionization; it includes the transfer of two target electrons into a bound state of the projectile under the simultaneous emission of a single photon. Indeed, RDEC can be considered as the simplest and (almost) background-free tool for studying electron-electron correlations in high- $Z$  ions.
- Experimental evidence: Simon *et al.* (2010) have presented experimental evidence of radiative double electron capture in the collision of 38 MeV  $O^{8+}$  with a carbon foil. When compared to theoretical cross sections, the experiments suggest a 5 times larger cross sections. This discrepancy may arise from the fact that no capture into excited states were included into the computations and that the electrons were considered as quasi-free in all computations.
- The (non-resonant) radiative electron capture is known to be dominant for fast collisions of heavy ions with light target atoms and shows a rather weak dependence on the inter-electron interaction between the projectile and target ions. The RDEC is a single-step process where the energy of two correlated captured electrons are converted into the energy and momentum of one emitted photon (Chernovskaya *et al.*, 2011).
- The RDEC can be described as the inverse process to double photoionization. Indeed, RDEC can be seen as a prominent tool for investigating the role of interelectronic interactions in the process of ion-atom collisions.

### 7.3.m. Generalized oscillator strengths (GOS)

#### Process, notations & application:

- Inelastic scattering of high-energetic photons or electrons The generalized oscillator strengths has been utilized to describe the inelastic scattering of high-energetic photons or fast electrons at atoms. This inelastic scattering has been discussed quite early by Bethe (1931) but has received less attention afterwards.
- Inelastic scattering of a fast electron by an atom or ion:  $A^{q+} + e_s^- \longrightarrow A^{r+} + e_s^{-'}$  with  $r \geq q$ .



- The notion of GOS has been introduced if only the effective scattering cross section (strength) is of interest and if all the other released electrons are neglected from the theoretical treatment.

### GOS expressions:

- **Fast electron impact:** The GOS for an excitation of the atom by fast electrons has been expressed as (Zhu *et al.*, 2006)

$$f(E, K) = \frac{E}{2} \frac{p_i}{p_f} K^2 \frac{d\sigma}{d\Omega} = \frac{2E}{K^2} |\langle \alpha_f \mathbb{J}_f M_f | T_1(\mathbf{K}) | \alpha_i \mathbb{J}_i M_i \rangle|^2,$$

where  $E$  is the excitation energy,  $K$  the momentum transfer,  $p_i, p_f$  the incident and scattered electron momenta, while  $\frac{d\sigma}{d\Omega}$  refers to the differential scattering cross section of the incident electron.

### 7.3.n. Electron-impact ionization (ImpactIonization)

#### Process, notations & application:

- **Electron-impact ionization** of an atom or ion:  $e_s^- + A \longrightarrow A^* + e_s'^- + e^-$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + |\varepsilon_i \kappa_i\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon_f \kappa_f\rangle + |\varepsilon_c \kappa_c\rangle$  where  $|\varepsilon_c \kappa_c\rangle$  is the initially bound *atomic* electron.
- The electron-impact ionization process is important in all (high-temperature) plasma, both in astro and plasma physics, and elsewhere.
- We refer to the three *free* electrons, that are involved in this process, as **incident, (final-) scattered and additional (-ly released) electron**, and with the corresponding indices above.

#### Cross sections:

- **Electron-impact ionization cross section:** These cross sections are obtained rather similarly as the electron-impact excitation cross sections but by allowing an **ionizing transition between fine-structure levels**  $\alpha_i \mathbb{J}_i(N) \rightarrow \alpha_f \mathbb{J}_f(N-1)$  (Fontes *et al.* 2015)

$$\sigma^{(\text{ionization})}(\alpha_i \mathbb{J}_i, \alpha_f \mathbb{J}_f) = \frac{8}{k_i^2 g_i} \sum_{\mathbb{J}_c, \mathbb{J}_t} (2J_s + 1) \sum_{\kappa_i \kappa_f \kappa_e} \int_0^{\varepsilon_i - I_p} \varepsilon' \left| \langle ((\alpha_f \mathbb{J}_f, \varepsilon_c \kappa_c) \mathbb{J}_c), (\varepsilon_f \kappa_f) \mathbb{J}_t \parallel \mathbb{V}^{(e-e)} \parallel (\alpha_i \mathbb{J}_i, \varepsilon_i \kappa_i) \mathbb{J}_t \rangle \right|^2$$

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- This cross section contains a summation over all possible values  $\mathbb{J}_t$  of the initial ion and the (incoming) incident electron in the partial wave  $|\varepsilon_i \kappa_i\rangle$ .

### 7.3.o. Electron-impact multiple ionization

#### Process, notations & application:

- **Electron-impact multiple ionization** of an atom or ion:  $e_s^- + A \longrightarrow A^* + e_s'^- + e_a^- + e_b^- + \dots$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + |\varepsilon_i \kappa_i\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon_f \kappa_f\rangle + |\varepsilon_a \kappa_a\rangle + |\varepsilon_b \kappa_b\rangle + \dots$   
where  $|\varepsilon_a \kappa_a\rangle$ ,  $|\varepsilon_b \kappa_b\rangle$ , ... refer to the initially bound *atomic* electrons.
- To describe the **multiple ionization in electron-impact processes**, one needs to consider the (electron-impact) excitation, ionization as well as several resonant excitation processes. This is usually done in the **isolated resonance approximation**, in which all the processes are treated independently. The multiple ionization is then obtained from a proper summation over individual decay pathways and branching ratios.
- **Charge state distribution (CSD) in plasma:** Electron-impact double and multiple ionization plays a central role upon the charge state distribution of a plasma if the electron temperature changes rapidly either in time or space.
- For a plasma in a **collisional-ionization equilibrium**, the CSD is determined by the balance between electron-ion recombination and electron-impact – single and multiple – ionization. Therefore, accurate electron-impact excitation and ionization cross sections are important for plasma diagnostics over a wide range of plasma (electron) temperatures, electron densities and elemental abundancies.

### 7.3.p. Three-body recombination

#### Process, notations & application:

- **Three-body recombination** of an ion:  $A^{q+} + e_i^- + e_c^- \longrightarrow A^{(q-1)+} + e_f^-$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + |\varepsilon_i \kappa_i\rangle + |\varepsilon_c \kappa_c\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon_f \kappa_f\rangle$
- The three-body recombination of an ion is the inverse process of the electron-impact ionization.

- The three-body recombination combination of an electron results in the capture of an electron by an positive ion due to energy and momentum transfer of the (captured) free electron to another free electron in the neighborhood of the atom or ion.
- **Three-body recombination:** For low-temperature afterglow and recombination-laser plasmas, the three-body collisional recombination to higher levels and deexcitation between higher levels are often dominant.

### 7.3.q. Elastic scattering of electrons and ions

#### Process, notations & application:

- **Elastic scattering of electrons and positrons** by atoms and ions of an ion:  $A^{q+} + e_i^- \longrightarrow A^{q+} + e_f^-$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + |\varepsilon_i \kappa_i\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon_f \kappa_f\rangle^+$
- **Potential scattering of relativistic electrons and positrons:** The scattering of relativistic electrons or positrons by a – real or complex – central-field  $V(r)$  is **completely described by the direct scattering amplitude  $f(\vartheta)$  and the spin-flip scattering amplitude  $g(\vartheta)$** , which are both complex functions of the polar scattering angle  $\vartheta$ .
- These scattering amplitudes are derived from the asymptotic behaviour for  $r \rightarrow \infty$  of the distorted plane-wave solutions of the Dirac equation *plus* an outgoing spherical wave.

#### Static scattering potentials:

- **Static-field approximation:** This approximation only includes the (electro-) static potential *plus* approximate local exchange potential for the interaction between the incoming or outgoing electron and the (target) ion. For electrons with kinetic energies  $E_{\text{kin}} \lesssim 10$  keV, moreover, one often includes optionally a semi-empirical correlation-polarization potential in order to account for the dynamically induced polarization of the target.
- **Local exchange potential:** For electron scattering processes, three simple analytical approximations for a local exchange potential refer to the Thomas-Fermi, Furness-McCarthy as well as the Riley-Truhlar potential (Salvat *et al.*, 2005).
- **Absorption potential:** For the calculation of elastic scattering cross sections, it is sometimes helpful to optionally include an imaginary (absorptive) part into the static potential in order to account for the coupling with inelastic channels.

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- **Exchange potential:** For in- and/or outgoing electrons, a possible exchange with a bound electron need to be taken into account as it occurs naturally if all electrons are treated together by an (antisymmetrize) Slater determinant. In general, however, these exchange terms in the Dirac-Hartree-Fock equations are difficult to deal with and are often approximated in some suitable form.

### Scattering amplitudes and cross sections:

- **Partial-wave expansion of the scattering amplitudes:** For a projectile with wave number  $k = \hbar/p$  and kinetic energy  $E$ , the direct and spin-flip scattering amplitudes are given by

$$f(\vartheta) = \frac{1}{2ik} \sum_{\ell=0}^{\infty} \{ (\ell+1) [\exp(2i\delta_{\kappa=-\ell-1}) - 1] + \ell [\exp(2i\delta_{\kappa=\ell}) - 1] \} P_{\ell}(\cos \vartheta)$$

$$g(\vartheta) = \frac{1}{2ik} \sum_{\ell=0}^{\infty} [\exp(2i\delta_{\kappa=\ell}) - \exp(2i\delta_{\kappa=-\ell-1})] P_{\ell}^1(\cos \vartheta), \quad (c\hbar k)^2 = E(E + 2mc^2),$$

and where  $P_{\ell}^1(x)$  is an associated Legendre function. The scattering amplitudes above apply for both, electrons and positrons if the sign of their charges is taken into account, and which turns an attractive potential into a repulsive potential and *vice versa*.

- Analogue to non-relativistic theory, **attractive potentials are generally associated with positive phase shifts and repulsive potentials with negative phase shifts.**
- **Elastic differential and total CS:** (Salvat *et al.*, 2005)

$$\frac{d\sigma^{(\text{elastic})}}{d\Omega} = |f(\vartheta)|^2 + |g(\vartheta)|^2, \quad \sigma^{(\text{elastic})} = \int d\Omega \frac{d\sigma^{(\text{elastic})}}{d\Omega} = 2\pi \int_0^{\pi} d\vartheta \sin \vartheta \frac{d\sigma^{(\text{elastic})}}{d\Omega}.$$

### Spin-polarization of elastically scattered electrons:

- **Sherman function:** This function describes the degree of spin polarization of electrons or positrons as function of the scattering angle  $\vartheta$  for an initially unpolarized beam

$$S(\vartheta) \equiv i \frac{f(\vartheta)g^*(\vartheta) - f^*(\vartheta)g(\vartheta)}{|f(\vartheta)|^2 + |g(\vartheta)|^2}.$$

### Applications of elastic scattering:

- **Application of differential elastic cross sections**  $\frac{d\sigma^{(\text{elastic})}}{d\Omega}$ : Accurate differential cross sections (DCS) for the elastic scattering of electrons and positrons by atoms and molecules are required in surface science, electron microscopy, electron-probe microanalysis, the design of radiation detectors, radiation protection, radiation-therapy planning and at several places elsewhere (Salvat *et al.*, 2005). Moreover, DCS for the elastic scattering by positive ions are needed also in order to describe the electron transport in plasmas.

### 7.3.r. Negative-continuum dielectronic recombination

#### Process, notations & application:

- **Negative-continuum dielectronic recombination** of an ion:  $A^{q+} + e_s^- \longrightarrow A^{(q-2)+*} + e^+$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + |\varepsilon_i \kappa_i\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon_f \kappa_f\rangle^+$
- The capture of an electron by a bare heavy nucleus may occur non-resonantly via the creation of a free-positron-bound-electron pair. This process is referred to as **negative-continuum dielectronic recombination** since it leads to the capture of (two) electrons into a bound state while a positron is released.

### 7.3.s. Nonradiative electron capture (NRC)

#### Process, notations & application:

- **Non-radiative electron capture** of an ion:  $\dots A^{q+} + e^- + Z_t \longrightarrow A^{(q-1)+, (*)} + Z'_t$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + |\varepsilon_f \kappa_f\rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle$
- The nonradiative electron capture of a quasi-free electron by a projectile ion mainly occurs if the electron and ion velocities match to each other.

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- The nonradiative capture is often described within the eikonal approximation, which enables one to estimate the capture from any filled shell of the target to any shell of the projectile.
- In high-Z targets, the nonradiative electron capture into excited states of the projectile is dominant, as predicted by the eikonal calculations and confirmed by measurements.

### Nonrelativistic capture cross sections:

- In the nonrelativistic theory, the (nonradiative electron capture) cross section of a stationary bare nucleus to capture an electron with energy  $E = \frac{m}{2} v^2$  is given by (Spitzer, 1956)

$$\sigma^{(\text{NRC, non-rel})}(E) = A \sum_{n=1}^{\infty} \frac{\varepsilon_g}{\hbar \omega_n E n^3} g_n,$$

where  $A = 32\pi/\sqrt{27} \hbar e^2/(m^2 c^3) \approx 2.11 \times 10^{-22} \text{ cm}^2$ ,  $\varepsilon_g$  is the ground-state binding energy of the hydrogenic projectile and  $\hbar \omega_n = E - \varepsilon_n$  the emitted photon energy for capture into level  $n$ . The factor  $g_n \approx 1$  is a correction factor which is typically not specified in detail for most computations.

- For  $g_n = 1$ , this cross section can be written

$$\sigma^{(\text{NRC, non-rel})}(E) = A \left( \frac{\varepsilon_g}{E} \right) \sum_{n=1}^{\infty} \frac{1}{n (n^2 + \frac{\varepsilon_g}{E})} = \frac{A \varepsilon_g}{2 E} \left\{ \psi \left( 1 + i \sqrt{\frac{\varepsilon_g}{E}} \right) + \psi \left( 1 - i \sqrt{\frac{\varepsilon_g}{E}} \right) + 2 \gamma_E \right\}$$

where  $\psi(z) = \frac{1}{\Gamma(z)} \frac{d}{dz} \Gamma(z)$  is the digamma function and  $\gamma_E \approx 0.57772$  is Euler's constant.

- Using the asymptotic form of  $\psi(z)$ , the cross sections can then be approximated as

$$\sigma^{(\text{NRC, non-rel})}(E) = A \left( \frac{\varepsilon_g}{E} \right) \left[ \gamma_E + \ln \sqrt{\frac{\varepsilon_g}{E}} \right].$$

### Velocity distributions:

- Two velocity distributions have been frequently applied to study nonradiative electron capture processes at proton and heavy-ion storage rings.
- **Maxwell distribution:**  $f(E) = \frac{1}{(2\pi kT/m)^{3/2}} \exp\left(-\frac{E}{kT}\right).$
- **'Flattened' Maxwell distribution:**  $f(E) = \frac{1}{(2\pi kT/m)} \exp\left(-\frac{E}{kT}\right) \delta(v_{\text{ex}})$   
which is obtained by suppressing one of the components of velocity.
- **Nonradiative electron-capture rates:** By using the asymptotic form of the digamma function in the nonradiative electron capture cross section above but different velocities distributions for the electrons to capture, **various approximations for the capture rates** can be derived. From the capture cross section  $\sigma^{(\text{NRC, non-rel})}(E)$ , the averaged coefficient  $\alpha_r$  can be evaluated analytically for different velocity distributions

$$\alpha_r^{(\text{NRC, non-rel})}(\text{Maxwell}) = \frac{\sqrt{2} A \varepsilon_g}{\sqrt{m, kT}} \frac{2}{\sqrt{\pi}} \left\{ \gamma_1 + \frac{\gamma_E}{2} + \Gamma\left(\frac{4}{3}\right) \gamma_2 \left(\frac{kT}{\varepsilon_1}\right)^{1/3} + \ln \sqrt{\frac{\varepsilon_1}{kT}} \right\}$$

$$\alpha_r^{(\text{NRC, non-rel})}(\text{flattened}) = \frac{\sqrt{2} A \varepsilon_g}{\sqrt{m kT}} \sqrt{\pi} \left\{ \gamma_1 + \frac{\gamma_E}{2} + \Gamma\left(\frac{5}{6}\right) \frac{\gamma_2}{\sqrt{\pi}} \left(\frac{kT}{\varepsilon_g}\right)^{1/3} + \ln \left(2 \sqrt{\frac{\varepsilon_1}{kT}}\right) \right\}.$$

### Semi-relativistic capture cross sections:

- Cross section calculations for nonradiative electron capture in relativistic heavy-ion collisions are often based on an eikonal approximation. Anholt (1985) adapted some prior form of the asymmetric eikonal theory in order to calculate total nonradiative electron capture cross sections from any filled target shell to any projectile shell by just scaling the nuclear charges  $Z_t$  and  $Z_p$  by the principal quantum numbers  $n_t$  and  $n_p$ . In general, this approximation was found to become better for high-energetic ions and can be applied confidently also for ultra-relativistic collisions.

### 7.3.t. Vacuum-electron capture

#### Process, notations & application:

- **Vacuum electron capture** of an ion:  $\dots A^{q+} + Z_t + e^- + e^+ \longrightarrow A^{(q-1)+, (*)} + Z'_t + e^+$
- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + |\varepsilon_- \kappa_- \rangle + |\varepsilon_+ \kappa_+ \rangle \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon_+ \kappa_+ \rangle$
- The vacuum electron capture process refers to the capture of an electron by fast projectiles due to the creation of an electron-positron pair and under the emission of the positron.
- In the vacuum electron capture, the target is considered to act as an perturbing potential that help create an electron-positron pair in the field of the projectile nucleus. While the electron is captured in one of the bound states (predominantly into  $1s$ ), the positron is emitted.
- The creation of an electron-positron pair from the vacuum is similar to the Coulomb ionization process of the projectile ion, except that the relative momentum of the ion is imparted into the vacuum instead of exciting an electron into a higher state of the projectile.

#### Semi-relativistic capture cross sections:

- Due to the high value of momentum transfer, which is needed to excite an electron-positron pair from the vacuum, **vacuum-electron capture takes place predominantly at small impact parameters**, where the screening of the target nucleus by the target electrons is unimportant.
- **Vacuum-capture cross sections** The vacuum-capture process involves a momentum transfer from the target nucleus to the vacuum followed by an electron capture into the projectile K-shell. Therefore, the vacuum capture cross section is  $\sim Z_t^2$ , similar to most momentum transfer processes like excitation and ionization, and also  $\sim Z_p^5$  as typical for projectile electron capture processes. For large  $\gamma$ , the cross sections increases as  $\ln \gamma$  by they vanish for  $v \ll c$  unlike for projectile ionization.
- Anholt and Becker (1987, table I) provide (reduced) vacuum capture cross sections for any combination of target-projectile and for  $\gamma \geq 10$ . For these large values of  $\gamma$ , the cross section is approximately

$$\sigma^{(\text{VAC})} = Z_p^5 Z_t^2 a \ln \left( \frac{\gamma}{\gamma_0} \right) \quad (\gamma > \gamma_0)$$

where  $a(Z_p)$  is a slowly varying function between 0.4 and 2.7 and  $\gamma_0$  is an offset. As the capture into higher shells than the K-shell varies as  $n^{-3}$ , the cross section for  $1s$  has been multiplied by 1.2 by Anholt and Becker to provide an estimate for the total capture. For  $\gamma > 10$ , these cross sections are expected to be accurate within approximately 20 %.



- The capture of an electron into the 1s shell of the projectile under the simultaneous emission of a positron has been calculated by Becker *et al.* (1987) and by Bertulani and Bauer (unpublished).

### 7.3.u. Positron-bound-electron pair annihilation with single-photon emission (PairAnnihilation1Photon)

#### Process, notations & application:

- Positron-bound-electron pair annihilation with single-photon emission from an atom or ion:  $A^{q+} + e^+ \longrightarrow A^{(q+1)+*} + \hbar\omega$
- Formal quantum notation:  $|\alpha_i \mathbb{J}_i\rangle + |\varepsilon \kappa\rangle^+ \longrightarrow |\alpha_f \mathbb{J}_f\rangle + \hbar\omega(\mathbb{M})$
- A positron can annihilate with a bound electron from an atom or ion under the emission of one, two or more photons. Here, we consider the emission of a single photons which is possible only due to the presence of the atom.

### 7.3.v. Positron-bound-electron pair annihilation with two-photon emission (PairAnnihilation2Photon)

#### Process, notations & application:

- Positron-bound-electron pair annihilation with two-photon emission from an atom or ion:  $A^{q+} + e^+ \longrightarrow A^{(q+1)+*} + \hbar\omega_1 + \hbar\omega_2$
- Formal quantum notation:  $|\alpha_i \mathbb{J}_i\rangle + |\varepsilon \kappa\rangle^+ \longrightarrow |\alpha_f \mathbb{J}_f\rangle + \hbar\omega_1(\mathbb{M}_1) + \hbar\omega_2(\mathbb{M}_2)$
- A positron can annihilate with a bound electron from an atom or ion under the emission of one, two or more photons. Here, we consider the emission of a two photons.

### 7.3.w. Positron-bound-electron pair production by a photon (PairProduction)

#### Process, notations & application:

- Positron-bound-electron pair production by a (single) photon of an atom or ion:  $A^{q+} + \hbar\omega \longrightarrow A^{(q-1)+*} + e^+$

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- **Formal quantum notation:**  $|\alpha_i \mathbb{J}_i\rangle + \hbar\omega(\mathbb{M}) \longrightarrow |\alpha_f \mathbb{J}_f\rangle + |\varepsilon \kappa\rangle^+$
- At high photon energies, the electron-photon interaction can lead to the creation of an electron-positron from which the electron is eventually bound by a multiply or highly-charged ion.

## 7.4. Other topics closely related to atomic processes

### 7.4.a. Atomic database from the literature

#### CHIANTI: A database for emission lines:

- CHIANTI is a database of assessed atomic parameters and transition rates that are needed for the simulation of line and continuum emission spectra from optically thin and collisionally-dominated plasma.
- CHIANTI was first released in 1996 and, since then, several new releases have been made available in order to expand the database and improve the quality of the data. Emphasis in developing this database has been given especially to the line identification and improvement of the reference wavelengths.
- For several ions, the wavelengths and identifications are different from and considered to be more accurate than those of the NIST database. Therefore, CHIANTI is now often applied as a reference atomic database for ions and has been included into several other atomic codes and packages.
- In its last version (Del Zanna *et al.*, 2015), the CHIANTI database includes a large amount of new data which improve simulations for the soft x-ray, extreme UV (EUV) as well as UV spectral regions. These regions are currently covered by several space missions.

#### DREAM: A database for lanthanide spectra:

- **DREAM database** The Mons group (Biemont and coworkers) have calculated various neutral and near-neutral lanthanide spectra; these data are kept within the DREAM database. In these computations, they employed the least-squares fitting procedure in the Cowan code (Cowan 1981) together with a core-polarization potential in order to obtain spectroscopic quality radiative rates.

#### NIST-XCOM:

- The XCOM database provides photon scattering data and attenuation coefficients between 1 keV and 100 GeV for all the elements of the periodic table. It also lists total cross sections, attenuation coefficients and partial interaction coefficients for selected processes, such as Compton and Rayleigh scattering, photoelectric absorption as well as pair production (Amako *et al.*, 2005).

### NIST-ESTAR:

- The ESTAR database provides (electron) stopping powers and ranges for incident electrons with energies between 10 keV and 1 GeV, and as derived from the so-called ICRU Report (Amako *et al.*, 2005).

### NIST-PSTAR:

- The PSTAR and ASTAR databases provide (proton) stopping powers and ranges for incident protons with energies between 1 keV – 10 GeV, and as derived from the so-called ICRU Report 49.

## 7.4.b. Codes which require atomic data input

### Codes from astro and plasma physics:

- The main reason for differences in the output of different codes are typically the different sources of atomic data.
- **Radiative-Collisional code based on FAC, RCF:** This code has been used to simulate steady-state plasmas under non-local thermodynamic equilibrium condition, and especially for photoionization-dominated plasmas. RCF takes almost all of the radiative and collisional atomic processes into account by rate equations in order to interpret the plasmas systematically. The Flexible Atomic Code (FAC) supplies all the atomic data that are needed for RCF in order to ensure completeness and consistency of the atomic data.
- RCF is a steady-state, collisional-radiative, optically thin model. Its rate-equation

$$\frac{dN_{i,j}}{dt} = \text{populating processes} - \text{depopulating processes} = 0$$

where  $N_{i,j}$  is the density of the j-th level in the i-th charge state. RCF includes ionization and recombination between neighboring charge states as well as the excitation and de-excitation within the same charge state. Moreover, the inverse processes are typically taken into account by the detailed-balance principle.

- GALAXY employs an average-of-configuration approximation for the electronic states as well as screened hydrogenic wave functions for both, the collisional and radiative processes. This code also employs Hartree-Dirac-Slater or Kramers cross-sections for photoionization.

- FLYCHK uses a hydrogenic approximation in order to calculate energy levels and level populations. However, the results of GALAXY and FLYCHK largely deviate from measured ones.
- In PHICRE, the energy levels and spontaneous decay rates are taken from the NIST database, and other rate coefficients are calculated by widely used formulas.
- ATOMDB calculates the x-ray/UV spectrum of a hot, collisionally-dominated optically-thin plasma, based on knowledge of the atomic transition rates and energies of the involved ions. This code also help analyse the interplay between the different rates. The line and continuum results of ATOMDB (<http://www.atomdb.org>) apply to an optically-thin thermal plasma with astronomical abundances; cf. Anders and Grevesse (1989).
- XSTAR is a command-driven, interactive, computer program for calculating the physical conditions and emission spectra of photoionized gases; cf. <http://ascl.net/9910.008>. It may be applied in a wide variety of astrophysical contexts. In particular, XSTAR computes the reradiated light when the gas of a star absorbs energy. The user typically need to supply the shape and strength of the incident continuum, the elemental abundances in the gas as well as the density or pressure, and the thickness of the gas layer.

### 7.4.c. Radiative opacity

#### Radiative opacity computations:

- **Radiative opacity of a plasma:** The radiative opacity of a plasma is of great significance for various research fields, such as inertial confinement fusion, stellar physics as well as the development of x-ray lasers. The (so-called) **Opacity Project** calculated the atomic data for most astrophysically abundant elements and also generated a database of opacities.
- **Opacity computations generally imply the accurate knowledge of different atomic data, such as level energies, oscillator strength, photoionization cross sections for all ionic species in the plasma.** They also require simulations over a wide range of temperatures and densities, an adequate equation of state for determining the ionization fractions and level populations and information about line broadening mechanisms.
- Knowing the opacity of a plasma, the **fraction of transmitted radiation** with regard to some incident radiation of given intensity can be expressed as:

$$F(h\nu) = \exp(-\rho \eta' L), \quad \eta' = \eta'(h\nu)$$

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where  $L$  is the path length of the light travelling through the plasma,  $\rho$  the mass density and  $\eta'(h\nu)$  the radiative opacity. The fraction of transmitted radiation can be readily determined experimentally and can be obtained theoretically by taking the integral over a Gaussian function, if the full width corresponds to the spectrometer resolution.

- **Radiative opacity of iron plasma:** The radiative opacity of iron plasmas at high temperatures has been found important in astrophysics and, therefore, various experiments have been carried out to accurately measure their radiative opacity during recent decades.
- Although the fraction of iron is typically quite small in most stellar environments, especially when compared with the high abundance of hydrogen and helium, the opacity of iron plays an important role for understanding the radiative transfer in astrophysical plasmas, such as the Sun.
- Opacity computation can provide valuable information about a plasma, including the fractional distribution of the different ion stages and whether the plasma is in a local thermodynamic equilibrium (or not) as well as about the temperature gradient in plasma.
- Radiative opacity are often analysed and displayed either as **isothermal sequences for different mass densities** and as **isodensity sequences for different temperatures**.
- Barnes and Kasen (2013) applied the line data from the VALD database in order to study the opacity of the lanthanides and to demonstrate that these heavy elements may exhibit a much larger opacity. At present, however, the limitations of VALD does not allow to compute reliable  $r$ -process opacities over the wavelength and temperature range as required by realistic transport calculation. In particular, VALD contains almost no lines with wavelengths greater than  $1\ \mu\text{m}$  and, hence, cannot be applied to compute sensible infrared opacities or emissivities of the neutron-star merger.
- A detailed modelling of the radiative properties of the high- $Z$  ions will be a long-term endeavor because of the number of electrons and the shell structure of the lanthanides as well as most elements beyond.
- The atomic properties of the light elements ( $Z < 30$ ) are known reasonably well from experiment and many previous computations. In particular, Kurucz (1993) and Kurucz and Bell (1995) provide an extensive list of lines (with about 42 million lines and 500,000 line shifts), which is dominated by the lines from the iron group elements. These lists were generated by using the CATS code (Cowan 1981) and by tuning the theoretical energies semi-empirically in order to reproduce the extensive observed experimental level energies.

### Definition of detailed and mean radiative opacities:

- **Radiative opacity  $\eta(\hbar\omega)$ :** For a plasma in a local thermodynamic equilibrium at temperature  $T$  and mass density  $\rho$ , the radiative opacity  $\eta(\hbar\omega)$  is given by (Cheng and Jiaolong, 2008)

$$\rho \eta(\hbar\omega) = [\mu^{(\text{bound-bound})}(\hbar\omega) + \mu^{(\text{bound-free})}(\hbar\omega) + \mu^{(\text{free-free})}(\hbar\omega)] \left(1 - \exp\left(\frac{\hbar\omega}{k_B T}\right)\right) + \mu^{(\text{scattering})}(\hbar\omega)$$

where  $\mu^{(\text{bound-bound})}$ ,  $\mu^{(\text{bound-free})}$ ,  $\mu^{(\text{free-free})}$  and  $\mu^{(\text{scattering})}$  are the absorption coefficients due to bound-bound, bound-free, free-free and well as scattering processes, respectively.

- **Rosseland mean opacity:** Radiative transfer is mainly controlled by the Rosseland mean opacity and is often expressed in terms of the reduced photon energy  $u$

$$\frac{1}{\kappa_R} = \left[\frac{dB}{dT}\right]^{-1} \left[\int_0^\infty d\nu \frac{1}{\kappa_\nu} \frac{\partial B_\nu}{\partial T}\right] = \int_0^\infty du \frac{1}{\kappa_u} \times g(u), \quad g(u) = \frac{15}{4\pi^4} u^4 \frac{\exp(-u)}{[1 - \exp(-u)]^2}, \quad u = \frac{h\nu}{kT}.$$

- **Rosseland mean opacity:** The Rosseland opacity is defined by

$$\frac{1}{K^{(\text{Rosseland})}} = \int_0^\infty du \frac{W^{(\text{Rosseland})}}{\eta(u)}, \quad W^{(\text{Rosseland})} = \frac{15}{4\pi^4} \frac{u^4 e^{-u}}{(1 - e^{-u})^2}, \quad u = \frac{\hbar\omega}{k_B T}.$$

- **Planck mean opacity:**

$$\frac{1}{K^{(\text{Planck})}} = \int_0^\infty du W^{(\text{Planck})} [\eta(u) - \eta^{(\text{scattering})}(u)]; \quad W^{(\text{Planck})} = \frac{15}{\pi^4} \frac{u^3 e^{-u}}{(1 - e^{-u})}, \quad u = \frac{\hbar\omega}{k_B T}.$$

- **Rosseland and Planck mean opacities:** The spectrally-resolved Rosseland and Planck mean opacities for iron plasmas at different temperatures from a few eV to a few hundred eV have been found useful important in studying the evolution of stars, such as the Sun. These studies range from the stellar envelopes to their interiors.
- The Rosseland and Planck mean opacities are required for many practical applications such as radiative transfer.
- **Gray opacity:** The gray opacity value just results from the Thomson-scattering opacity,  $\kappa^{(\text{Thomson})} = 0.4(\bar{Z}/A) \text{ cm}^2/\text{g}$  and if one assumes a (fully) ionized material (medium) with  $\bar{Z} = Z$ . Here,  $A$  is atomic weight and one often applies  $\bar{Z}/A \approx 1/2$  for most elements.

### Spectral intensity and energy density:

- **Spectral intensity  $I_\lambda$ :** The (specific) light intensity describes the amount of electromagnetic radiation with wavelengths between  $\lambda \dots \lambda + d\lambda$  that passes in the time  $dt$  through the area  $dA$  and into the solid angle  $d\Omega = \sin \vartheta d\vartheta d\varphi$  (Pettini 2018)

$$W_\lambda d\lambda = I_\lambda(\vartheta, \varphi) d\lambda dt dA \cos \vartheta \sin \vartheta d\vartheta d\varphi$$

In astrophysics, the intensity is often measured in units:  $\text{erg s}^{-1} \text{cm}^{-2} \text{\AA}^{-1} \text{sr}^{-1}$ .

- **Spectral and total energy density  $u_\lambda$ :** The specific energy density refers to the energy per unit volume that is contained in the radiation field with wavelengths between  $\lambda \dots \lambda + d\lambda$ , and from which the total energy density is easily obtained

$$u_\lambda d\lambda = \frac{d\lambda}{c} \int d\Omega I_\lambda(\vartheta, \varphi) = \frac{4\pi}{c} J_\lambda d\lambda \quad \Rightarrow \quad u = \int_0^\infty d\lambda u_\lambda$$

### 7.4.d. Opacities for astrophysical matter clouds

#### Astrophysical opacities models:

- **Opacity:** The opacity of a (plasma) medium characterizes the transport of radiation through matter. Opacity computations generally require atomic data for a large number of processes involving the absorption and scattering of radiation (Seaton 1987).
- A new astrophysical opacity model (called OPAL) has been developed at the Lawrence Livermore National Laboratory (LLNL) with which it was shown that  $\Delta n = 0$  transitions results in many overlapping lines. This overlap significantly increase the opacity of the iron component of the mixtures and which was suggested also to apply for their observation of stellar pulsations.
- **Need of opacities:** Opacities are frequently applied in various research fields, such as the standard solar model (SSM), helio and astero-seismology, non-LTE 3D hydrodynamic photospheric modeling, nuclear reaction rates, solar neutrino observations, computational atomic physics as well as in plasma experiments (Mendoza 2018).
- Typically opacity computations have considered a total of up to  $10^8$  radiative lines. Since these computations cannot so easily be extended towards high- $Z$  elements, the **average-atom** description has been applied as an alternative approach.



- **Local Thermodynamic Equilibrium (LTE):** Despite the complexity of (and complications with) astrophysical plasma, one makes often use of the LTE approximation as long as the mean free path of the photons and particles is small compared to the scale over which the temperature changes significantly. This applies for instance to many stellar interiors because of the high density and temperature and, hence, quite small mean distances between collisions.

### Opacity calculations:

- Radiative opacities are required for the modeling of astrophysically relevant plasmas under local thermodynamic equilibrium (LTE) conditions. For the computation of these opacities, the atomic structure calculations should be carried out for all fine-structure levels and by including electronic correlations.
- Early opacity computations have often used data obtained from simple atomic models. However, already Simon (1982) suggested that such simple estimates are too small by factor of 2-3 for all elements heavier than helium.
- Three typical example applications for opacity calculations are: (i) The iron opacities at conditions of the solar convection zone, (ii) nickel opacities for modeling various stellar envelopes and (iii) the samarium opacities for the modeling of light curves as produced by neutron star mergers.
- In some recent work, the frequency-dependent LTE opacity of iron has been studied for the conditions of the solar convection zone. These studies revealed rather large (30-400%) differences between experiment and theory for the monochromatic opacity of an neon-like iron plasma (or nearby charge states). Until the present, the source of this not yet fully understood and remains an active area of research.
- **Opacity calculations for lanthanides or actinides:** In the previous opacity computations of the lanthanides or actinides, rather large uncertainties arose from the – experimentally and theoretically – unknown atomic states and line strengths of these complex elements. Indeed, these high- $Z$  atoms still represent a major challenge for accurate atomic calculations and often requires the use of statistical models, and which need to be calibrated to experimental data.
- **Opacities for ejecta from NS mergers:** At the near-IR/optical frequencies, the dominant source of opacity arises from many bound-bound transitions between the low-lying levels of the atoms and ions. For this *continuum* of lines, the opacity is determined by the strength and wavelength density of the lines, and which depend sensitively on the composition of the ejecta.
- **Opacities for ejecta from NS mergers:** The opacity of the ejecta from NS mergers can significantly rise if they just contain a modest fraction of elements with a partially-filled  $f$ -shell, such as the elements from the lanthanide and actinide groups.
- **Iron opacity:** When compared to previous estimations, OPAL iron monochromatic opacities at density  $\rho = 6.82 \times 10^{-5} \text{ g cm}^{-3}$  and temperature  $T = 20 \text{ eV}$  indicated large enhancements at photon energies around 60 eV, which may arise from  $\Delta n = 0$  unresolved

transition arrays in Fe, and which is supported by laboratory photoabsorption measurements.

### Opacity measurements for homogeneous plasma with uniform density:

- **Uniform plasma density:** For a uniform density, the transmission of the sample

$$T(\omega) = \frac{I(\omega)}{I_o} = \exp(-\rho L \kappa(\omega))$$

is measured where  $I(\omega)$  is the attenuated intensity of the probe (backlighter) radiation and  $I_o$  the reference intensity. This transmission is related also to opacity by the Beer-Lambert law if the reabsorption can be neglected in the homogeneous plasma. Here,  $\rho$  is the density of the material and  $L$  its thickness.

- **Iron opacity:** Bailey *et al.* (2015) measured iron opacities at electron temperatures between 1.9 – 2.3 million kelvin and electron densities between  $(0.7 - 4.0) \times 10^{22} \text{cm}^{-3}$  at the Z facility. In these measurements, the observed optical absorptions was found about  $\sim 30 - 400 \%$  higher than those predicted by currently applied opacity models.

### Neutron star mergers:

- **Neutron stars:** These stars have been postulated already short after the discovery of the neutron as the ultimate fate of massive stars, that end in a supernova. The existence of neutron stars was later shown in the 1960s following the observation of the first pulsars.
- Today, the equation of state and the distribution of neutron star masses are believed to be known reasonably well.
- The mass of the ejecta from neutron star mergers are predicted to be rather high and to reach from  $10^{-3} \dots 10^{-2} M_{\odot}$  for the overall r-process matter.
- Modern simulations of neutron star mergers take into account both, the (elemental) composition of the ejecta as well as the neutrino wind along the poles.
- **Material ejected from neutron-star mergers:** The material that is ejected during or immediately after the merger of two neutron stars may produce considerable amounts of heavy elements through the *r*-process. The subsequent radioactive decay of these (neutron-rich) nuclei can power transient electromagnetic emission similar to, but significantly dimmer than, an ordinary supernova (Kasen *et al.*, 2013). One goal of future observations of the electro-magnetic transients in gravitational-wave astronomy is to better understand this *r*-process and, hence, the astrophysical sources of gravitational waves.

- **Material ejected from neutron-star mergers:** A key in studying the matter outside of the remnant of the neutron-star merger is the release of neutron rich-matter and a cascade of nuclear reactions. These reaction gives then rise to kilonova which can be observed on the time scale of hours, days and even months due to the radioactive decay of freshly synthesized, neutron-rich elements in the ejected matter.
- **Light curves from neutron-star mergers:** Since very little is known about the optical properties of the heavy  $r$ -process elements, most previous light-curve models have adopted the opacities from the iron group elements. Barnes and Kasen (2013) considered first the opacities of lanthanides and concluded, that the opacity of the ejecta material from neutron-star mergers might be larger by several orders of magnitude. In further detail, Barnes and Kasen included these opacities into time-dependent, multi-wavelength radiative transport calculations in order to predict the broadband light curves for a range of masses and velocities of the ejecta.
- **Light curves of NS mergers:** For heavy atoms and ions, such as the lanthanides and all heavier elements, the opacities of the  $r$ -process material are expected to be order of magnitude higher than those of the iron group elements. This increase of the opacity is not only caused by the complex shell structure of the heavy atoms and ions but also by the number and diversity of available atomic species. Such an enhanced opacity then results in (predicted) light curves that are longer, dimmer, and redder than previously thought. These light curves are not sharply peaked  $t \sim 1$  day may lead to bolometric light curves  $\sim$  week(s), cf. Barnes and Kasen (2013). Indeed, simulation suggest pseudo-blackbody spectra with some broad absorption features and some peak in the infrared ( $1 \mu\text{m}$ ).
- **Light curves from NS mergers:** Tanvir *et al.* (2017) report the observation of optical and near-infrared spectra of the transient state of the ejecta from the NS merger (AT2017gfo). In particular, they have analyzed the broadband spectral energy distribution of AT2017gfo in a galaxy at  $\sim 40$  Mpc over the first 12 days, which shows a remarkable color change from blue to red. In more detail, the observed light curves exhibit both, a rapidly evolving blue and more slowly evolving red components.
- **Light curves from NS mergers:** The rise and decay of the (optical) light curves can be used together with the velocity of the ejecta in order to estimate its mass in terms of the opacity  $\kappa$

$$M \sim 5 \times 10^{-3} M_{\odot} \left( \frac{0.1 \text{ g cm}^{-2}}{\kappa} \frac{\nu}{0.1 c} \right).$$

This estimate suggests, however, that only a fraction  $10^{-3}$  of the totally released energy contributes to the kinetic energy of the ejecta.

- **Evolution of light curves:** The evolution of the light from the transient state of the NS merger has been found consistent with predictions of kilonovae which are powered by the radioactive decay of massive neutron-rich nuclides due to the  $r$ -process nucleosynthesis in the neutron-star ejecta.
- **Light curves from NS mergers:** The analysis by Tanvir *et al.* (2017) of the multi-band light curves indicates the presence of at least two emission components: one with high opacity, and a second with low opacity. The high-opacity component arises due to the *tidal part* of

the dynamical ejecta with a low electron fraction  $Y_e \leq 0.25$ .

### Gravitational-wave astronomy:

- The discovery of gravitational waves (GW) from the inspiral and coalescence of binary black holes (BH) and/or neutron stars (NS) by the Laser Interferometer Gravitational Wave Observatory (LIGO) has opened a new window on the cosmos.
- The discovery of electromagnetic signals in coincidence with a chirp at the gravitational-wave detectors leads to a better location and a much richer picture about the merger process. This may allow to extract directly information on the binary formation channels, the age of the stellar population as well as the dynamics of binary mergers.
- In general, the merger of BH binaries is not expected to produce luminous electro-magnetic emission because there is not enough baryonic matter in these systems.
- Despite of a rather large number of expected BH-BH mergers during the next few years, a better synthesis and analysis of gravitational wave and the associated electro-magnetic signals will probably come from NS-NS or BH-NS mergers, i.e. if neutron stars are involved.
- **Sources of gravitational waves (GW):** While most detected GW sources are nowadays believed to arise from black-hole binaries, neutron star mergers are another major source of gravitational waves. For these NS mergers, an electro-magnetic signal is expected to arise in two forms: gamma-ray bursts (GRB) or from kilonovae (Metzger *et al.*, 2010). The GRB are expected to be short-lived with a jet-type signal, while the kilonovae are predicted to generate relatively isotropic signals.
- **Sources of gravitational waves:** In practice, it is not easy to precisely locate the sky position of GW sources since they need to be determined by triangulating the arrival times of the GW signals from different detectors. At present, the uncertainties in this triangulation are quite large, although these uncertainties will be reduced to  $\approx 10 - 100 \text{ deg}^2$ , once Virgo in Italy and, eventually, KAGRA in Japan and LIGO-India will join the network (Metzger 2017). Initially, these uncertainties were  $\approx 850 \text{ deg}^2$  for GW15091 and with a later improvement to  $\approx 250 \text{ deg}^2$ , though
- Even an localisation of the GW signals within  $10 - 100 \text{ deg}^2$  greatly exceeds however the fields of view of most radio, optical and x-ray telescopes, and in particular the field of view of those telescopes that are sensitive enough to detect the electro-magnetic counterparts of NS-NS and BH-NS mergers.
- **Short  $\gamma$ -ray bursts (GRB):** Although these short GRB are bright and high-energetic events, only a subset of these bursts will be observed since the most intense emission from a given merger will typically not intersect the line of sight from the Earth.
- **$\gamma$ -ray bursts (GRB):** Short GRB are commonly believed to be powered by the accretion of a massive remnant disk, following some BH and/or NS mergers. These GRB are typically expected to occur already within seconds of the GW chirp, providing an unambiguous

temporal signature to the GW signal. Although such short GRB are likely the cleanest electro-magnetic counterparts to the GW signals, their measured rate will remain low within the Advanced LIGO detection volume. This rate will probably be less than one observable event per year if the GW signal arises from NS-NS mergers since the  $\gamma$ -ray emission is beamed into a narrow solid angle not seen from Earth (Metzger 2017).

- NS-NS and BH-NS mergers are also predicted to be accompanied by a more isotropic counterpart, commonly known as kilonovae. These kilonovae provide both a robust electro-magnetic counterpart to the GW chirp, and which is expected to accompany a good fraction of BH-NS mergers and essentially all NS-NS mergers (Metzger 2017).
- The single NS-NS merger event (GW170817, GRB170817A, AT2017gfo) made a big qualitative step for gravitational-wave astronomy since it showed that these mergers can be observed and are directly linked to gravitational waves, gamma-ray bursts, kilonovae as well as the generation of heavy elements (nucleosynthesis).

### Kilonovae:

- **Kilonova:** Kilonovae are day to week-long thermal, supernova-like transients which are powered by the radioactive decay of heavy, neutron-rich elements synthesized in the expanding ejecta of NS-NS or NS-BH mergers. The radioactive decay of such unstable neutron-rich nuclei powers a rapid evolution of a supernova-like transition of stellar matter that is **known as kilonova or macronova** in the literature. These kilonovae can be understood also as roughly isotropic electromagnetic counterpart to the GW signal and, thus, provide a unique and direct probe of such r-process sites (Metzger 2017).
- Kilonovae refer to optical transients that are powered by the radioactive decay of material ejected in the neutron-star merger. Simulations suggest an ejected mass of  $10^{-4} \dots 10^{-1} M_{\odot}$  due to a tidal stripping during the merger process itself or by a subsequent disk wind in course of the evolution of a post-merger remnant.
- **Kilonovae:** Kilonovae are powered by the radioactive decay of excited nuclear matter. These kilonovae will emit more isotropically (compared to the accretion disk of the merger event) and will peak later than short-GRB afterglows. Kilonovae are the best candidates for the electromagnetic counterparts of the GW detection.
- **Photon emission from kilonovae:** A detailed understanding of the photon emission from kilonovae requires (detailed) information and simulations on the composition, velocity structure and the opacity of the ejecta as well as on the transport of radiation through the ejecta.
- **Key observables of kilonovae:** Three main ingredients need to be understood to quantify the key observables of a kilonova, namely its peak time scale, the luminosity and the effective temperature. These ingredients are: (i) The mass and velocity of the ejecta from NS-NS

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or BH-NS mergers; (ii) the opacity  $\kappa$  of the expanding neutron-rich matter; and (iii) the variety of sources which contribute to the heat of the ejecta  $Q(t)$ . These ingredients are important especially on the time scale of  $t^{(\text{peak})}$ , when the ejecta is first becoming transparent.

### r-process:

- Burbidge *et al.* (1957) and Cameron (1957) realized already more than 60 years ago that approximately half of the elements heavier than iron are synthesized via the **capture of neutrons by light seed nuclei**, such as iron, in a dense neutron-rich environment, in which the time scale for neutron capture is shorter than the lifetimes for  $\beta$ -decay.
- **Opacities of the  $r$ -process** Barnes and Kasen (2013) showed that more realistic opacities of the  $r$ -process material may have a dramatic effect upon the predicted **light curves of a kilonovae**. In these simulations, they applied improved estimates of the radiative opacities of heavy elements as derived from *ab-initio* atomic-structure computations. The enhanced opacities of the lanthanides strongly affect the radioactive powered light curves from the ejecta of the neutron-star merger.
- **Production of heavy elements:** In all the suggested scenarios, the  $r$ -process nuclei are produced in two steps. (i) An initial explosion burns at high temperatures until the electrons and ions freeze-out during expansion of the ejecta but still with a high neutron-to-seed ratio, and which is followed by (ii) the rapid capture of neutrons by the seed nuclei, leading eventually to heavy nuclei.
- **Production of heavy elements:** Despite of intense research, little is known so far about the neutron-rich and in the  $r$ -process produced heavy elements. Indeed, the nuclei beyond the iron group elements (with mass numbers  $A \sim 90..100$ ) need to be formed via successive neutron capture owing to the high Coulomb barriers for other nuclear reactions with charged particles.

### r-process environments:

- **r-process:** The rapid neutron capture or (so-called)  $r$ -process is considered to be responsible for about half of the elements heavier than iron, although this process was previously attributed to core-collapse supernovae. However, recent studies of the  $r$ -process have disfavored supernovae because their conditions are not suitable to produce the heavy elements near to the *platinum peak* around  $A \propto 195$ . Instead, neutron-star mergers have received a lot of recent interest as a major  $r$ -process production site (Tanvir *et al.*, 2017).
- **r-process environment:** A current question in astrophysics is whether the mergers of neutron stars gives rise to a robust  $r$ -process environment that lead to heavy neutron-rich isotopes with mass  $A \geq 130$ , and in proportions similar to solar abundancies (Thielemann *et al.*, 2017).
- **r-process environment:** Various astrophysical sites have been suggested for the  $r$ -process of the nucleosynthesis. Apart from regular supernovas, the  $r$ -process has been associated especially with neutrino-induced processes in the outer shells of massive stars, the ejecta from

compact binary mergers as well as from so-called magneto-hydrodynamic (MHD) jet supernovae which exhibit high magnetic fields and neutron-rich jet ejecta along the poles.

- In the neutron-rich ejecta from neutron star (NS-NS) and BH-NS binaries, in particular, the so-called **r-process nucleosynthesis (rapid neutron capture)** is expected to proceed rapidly, leading to many heavy elements like gold, platinum and many others.

#### 7.4.e. Absorption spectra of distinct astrophysical objects

##### Identification of absorption spectra

- Sako *et al.* (2001) present first results from high-resolution x-ray observation of IRAS 13349+2438 with the XMM-NEWTON observatory. The observed spectrum was obtained with the Reflection Grating Spectrometer (RGS) and shows a wealth of discrete spectral features, including the (first) astrophysical detection of an unresolved transition array (UTA) due to the  $2p - 3d$  inner-shell excitations of  $M$ -shell iron ions. These spectral features could have been easily misidentified as an O VII edge, when observed with only moderate resolution spectrometers.
- Other prominent features in the spectra of Sako *et al.* (2001) were K-shell absorption lines of hydrogen- and helium-like carbon, nitrogen, oxygen and neon, and L-shell lines of Fe XVII-XX.

#### 7.4.f. Ionization equilibria in astrophysical sources

##### Astrophysical observation of EUV spectra:

- **High-quality EUV spectra:** The Solar EUV Rocket Telescope and Spectrograph (SERTS, 1992) and the Extreme Ultraviolet Explorer (EUVE, 1990) spectrometers provided high-quality EUV spectra and helped improve the diagnostics of high-temperature plasmas in astrophysical objects.
- The ability to resolve a large number of individual emission lines and to measure the accurate intensities over a wide wavelength range have allowed to analyze the temperatures, densities and abundances of elements in a variety of astrophysical sources (Brickhouse *et al.*, 1995).
- The fine structure of (line) multiplets nowadays allow to solve the full rate equations for several atoms in the plasma, including collisions among all levels and cascades. This enables one to calculate also the emissivities of weak lines that were previously omitted from the investigation.



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- **X-ray emission from astrophysical objects:** From the analysis of line intensities in exotic dense-plasma environments, several key properties of astrophysical objects may be derived that are not attainable through other observational windows.
- **Temperature of stars:** This can be defined in various different ways (Pettini 2018):
  1. **Effective temperature:** as defined in terms of the luminosity of the star and its radius;
  2. **Excitation temperature:** as given by the relative population of different excited levels of an atom or ion in the plasma, and by using Boltzmann's equation;
  3. **Ionization temperature:** as defined by the relative population of different ionisation stages of an atom, and by using Saha's equation;
  4. **Kinetic temperature:** as defined by the Maxwell-Boltzmann distribution:

$$n_v dv = n \left( \frac{m}{2\pi kT} \right)^{3/2} e^{-mv^2/2kT} 4\pi v^2 dv$$

with the number density  $n_v$  (number of particles per unit volume) of particles with speeds between  $v \dots v + dv$  and the total number density  $n$  of all particles with mass  $m$ ;

5. **Color temperature:** as the temperature of the blackbody whose spectral energy distribution resembles most closely that of the given star.

### 7.4.g. Photoionized, steady-state plasma

#### Ion distributions:

- All plasmas, that are produced in the laboratory, are generally quite far from local thermodynamic equilibrium (LTE). Therefore, the ion distribution of such laboratory plasma cannot be described by the usual Saha-Boltzmann equations, neither for selected ions nor for dealing with different charge states of the ions.
- **Z-pinch experiments:** In the laboratory, one often has a **steady-state plasma, far from equilibrium, that is formed by x-ray radiation from a Z-pinch**. In such plasma, the ambient radiation field is important in addition to the electron-collisional processes in order to determine the the excitation and ionization dynamics. **Such photoionized plasmas are believed to occur also in many astrophysical situations**; in fact, Z-pinch experiment allow for the first time to compare numerical models of photoionized plasmas with experiment (Rose *et al.*, 2004).



- The standard method for calculating the ion distribution in a non-LTE system is to solve the rate equations for all the important ionic states and by including different excitation and ionization channels due to the incident radiation. From the solutions of these rate equations, a time history of the electron temperature and density as well as the time evolution of the different charge states in the plasma can be evaluated. The **steady-state solution** is then obtained by following the time evolution until the populations do no longer change with time.
- For ion-distribution calculations in non-LTE plasma, a major simplification is achieved if only the shell populations, averaged over the distribution of the ionic states, is taken into account within the **average-atom model**. This approach reduces the number of differential equations significantly, when compared to a detailed treatment of all charge states though the shell occupancies are non-integers then.

#### 7.4.h. Plasma light sources for nanolithography

##### Development of EUV light sources:

- **Ions with several open shells:** The (correlated) electronic structure of medium and heavy ions with several open shells is both, tedious and difficult to calculate, and their complex level structure often hampers a straightforward experimental assignment.
- **EUV light sources at 13.5-nm wavelength:** In order to generate extreme ultraviolet (EUV) light at 13.5-nm wavelength by laser-produced plasmas, the emission of multiply-charged tin ions have been found promising for nanolithographic applications. For these ions, the EUV light is generated by thousands of transitions that form so-called unresolved transition arrays (UTA), and with rather little dependence on the particular charge state of the tin ions (Winderberger *et al.*, 2016).
- For multiply charged tin ions with a  $[\text{Kr}] 4d^m$   $m = 6 \dots 0$  configuration, the upper configurations for an EUV emission are  $[\text{Kr}](4p^6 4d^{m-1} 4f + 4p^6 4d^{m-1} 5p + 4p^5 4d^{m+1})$  (Winderberger *et al.*, 2016).
- Apart from application for EUV light sources, the optical transitions in medium and heavy ions with multiple open shells, such as  $\text{Sn}^{11+}$ – $\text{Sn}^{14+}$ , also represent a stringent test for *ab-initio* atomic structure calculations of complex (many-electron) spectra with non-negligible Breit contributions.

### 7.4.i. Laser-induced fluorescence spectroscopy (LIFS) in flames

#### Laser-induced fluorescence & applications:

- Laser-induced fluorescence describes the spontaneous emission from atoms or molecules following the prior excitation by laser radiation. In multi-level atoms or molecules, the laser is usually tuned to one of the absorbing transitions. From the so excited level, the atoms then decay either to some lower lying state or are re-populated by collisions to adjacent states. In LIFS, therefore, a fluorescence spectrum is observed that contains individual transition of the excited level as well as from levels nearby (Daily 1997).
- The laser-induced fluorescence signal after laser excitation can be exploited in various ways: (i) The observed spectrum provides a measure for the population of the initially excited states; (ii) it allows to extract the total number density of different species in the probe, if a clear relationship can be established between the number density of all involved quantum states; (iii) the observed spectrum (may) help deduce the ground-state distribution function and, hence, the temperature of the probe. Moreover, (iv) the pressure, temperature and velocity may be deduced from the observed line shapes under suitable circumstances.

#### Combustion processes:

- **Combustion:** Combustion is the primary propulsion technology and is very frequently applied for the generation of electricity, the heating of buildings and at many places elsewhere. Combustion is applied also for the synthesis of new materials and for chemical processing (Daily 1997). In nature, combustion is a frequent source when natural areas are changed by fire.
- **Combustion systems:** Flames and combustion systems vary widely in their size, fuel type and purpose. They range from small mm-to-cm systems, such as furnace and stove pilot flames, gas stove burners or small laboratory burners to medium-size motors and heating systems, and up to pulverized coal fire boilers, marine diesels and forest fires, each with length scales on the order of several meters (Daily 1997).
- **Performance evaluation of combustion processes:** Combustion processes are often analyzed with regard to their energy conversion efficiency, material synthesis yield, material purity and/or pollutant emissions. This analysis is often made by measuring the exhaust and without that the combustion zone is affected directly.

## 7.4.j. Synthetic spectra for laser-induced breakdown spectroscopy (LIBS)

### Laser-induced breakdown spectroscopy:

- The laser-induced breakdown spectroscopy (LIBS) has been found a versatile analytical tool with various applications.
- In this technique, a high-power laser beam is typically focused onto the surface of a sample that need to be analyzed. If the laser power exceeds a certain threshold, an optical breakdown occurs due to the formation of a laser-induced plasma. In practice, the breakdown threshold depends on the wavelength and duration of the laser pulses as well as on the analyzed media.
- For a nanosecond laser ablation, for example, the threshold varies from  $1.9 \cdot 10^8 \text{ W cm}^{-2}$  for copper up to  $10^{10} \dots 10^{11} \text{ W cm}^{-2}$  for aqueous solutions (Yaroshchik *et al.* 2006).
- During the breakdown, a plasma of highly-ionized atoms is formed with an electron density of  $10^{16} \dots 10^{20} \text{ cm}^{-3}$  for typical Nd:YAG laser-induced plasmas.

### Computation of LIBS spectra:

- Ciucci *et al.* (1999) first estimated with limited accuracy the composition of a laser-induced plasma by using computational methods instead of the (standard) calibration standards. Their computation were based on the assumption that the plasma and sample compositions are the same and that the plasma is optically thin and in local thermal equilibrium (LTE) during the observation.
- Yaroshchik *et al.* (2006) describe some software for synthesizing laser-induced breakdown spectroscopy emission spectra and to predict sample composition by using a proposed calibration-free algorithm. This code applies a database of atomic emission lines in order to create a theoretical emission spectrum for selected elements using defined plasma parameters.
- The code by Yaroshchik *et al.* (2006) is also based on the assumption of a LTE and a well-defined plasma temperature and electron density; it takes a number of characteristics atomic emission lines as input to create a synthesized theoretical emission spectrum of selected elements which are compared to the spectra of four separate compact spectrometers covering the spectral region 185–950 nm.
- Until the present, all heavy elements are typically represented in LIBS plasmas by their neutral atoms and their singly- and doubly-charged ions that emit radiation at characteristic wavelengths. From the intensity of individual wavelength, the composition of the samples are then estimated both, qualitatively and quantitatively.

### 7.4.k. X-ray absorption of solid-state materials

#### Near-edge spectroscopy of solid-state samples:

- **X-ray near-edge absorption (XANES):** XANES as well as other related spectroscopies, such as the resonant x-ray scattering spectroscopy (RXS) and x-ray magnetic circular dichroism (XMCD), are powerful spectroscopic techniques for probing the electronic structure around the absorbing atom in a fluid or solid-state probe sequences (Bunau and Joly, 2009).
- In the near-edge region, one is very sensitive to the surroundings of the absorbing atom. This makes absorption spectroscopies a very useful tool to investigate the geometric and electronic structure of the samples.
- In x-ray absorption spectroscopy, there has been a great need for codes that help reproduce the near-edge structure within a reasonable calculation time.
- These codes can be classified into two major categories:
  - i) multi-electronic but monoatomic codes which are based on some multiplet theory and which give rise to a proper parametrized description of the localized electronic state. In these codes, however, while the delocalized states can often not be described correctly.
  - ii) multi-atomic but mono-electronic codes which provide correct description of the extended states but fail to account for the highly correlated state that interact with the core hole. This applies, for example, for the calculation of the L2 and L3 edges of the 3d elements.

### 7.4.l. X-ray quantum optics

#### Phenomena in x-ray quantum optics:

- X-ray quantum optics has evolved rapidly during the last year owing to the development of novel x-ray light sources and x-ray optics.
- A few important phenomena from x-ray quantum optics include: i) x-ray parametric down conversion; ii) nuclear  $\gamma$ -ray superradiance, and others. Indeed, various concepts from conventional (near-visible) quantum optics have been suggested to study also at higher photon energies. These concepts include photon correlations and entanglement, squeezing, and others.
- There are important differences between the interaction of x-rays and optical photons because of the inner-shell atomic energy levels, nuclear resonances, or quite large energy shifts in Compton scattering as the photon energies approach the electron rest mass.

### 7.4.m. Coherent x-ray scattering at high intensities

#### X-ray radiation sources:

- In 2009, the Linac Coherent Light Source (LCLS) at the SLAC National Accelerator Laboratory was the first x-ray FEL in operation with photon energies up to 8.3 keV, up to  $2 \times 10^{12}$  photons per pulse, and with a pulse length of just a few femtoseconds.
- Nowadays, further x-ray FEL include the SPring-8 Compact SASE Source (SCSS, Japan) and the European x-ray FEL (XFEL) at DESY. This XFEL can deliver photon energy up to 14 keV and has a 50-500 times higher average brightness than the (early) LCLS.

#### Single-shot imaging:

- **Single-shot imaging:** One central goal and applications of the European XFEL is single-shot imaging of individual macromolecules by means of coherent x-ray scattering. Such an imaging should help determine the atomically resolved structure of noncrystallized nanoparticles and biomolecules. Single-shot imaging might become possible for a sufficiently high-fluence and well-focused FEL pulse in order to generate a sufficiently large number of scattered photons from each (single-molecule) sample.
- **Radiation damage:** One of the key problems in single-shot imaging is the radiation damage in course of the irradiation of the sample. Typically, a much higher flux is required than the **conventional radiation damage limit of about 200 photons/Å<sup>2</sup>**, and this may seriously degrade the scattering patterns and, hence, the possibility of determining the atomic positions within the target molecule (Son *et al.*, 2011).

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Single-shot imaging therefore requires a sufficiently short pulse duration of less than 10 fs in order to keep the atomic position effectively frozen during the irradiation by the pulse.

- **Single-shot imaging:** Even if the atomic positions are nearly frozen, the single-shot imaging will be affected however by changes in the electron density of the target. This makes it necessary to understand also the detailed ionization and relaxation dynamics in individual parts of the molecules under ultrashort and ultraintense x-ray pulses (Son *et al.*, 2011). This dynamics can be simulated by means of time-dependent rate equations in order to explore the formation and relaxation of the inner-shell vacancies.

### Ionization dynamics in short light pulses:

- **Rate equations for configuration-averaged population dynamics:** The transitions among all possible electronic configurations  $c$  of a given atom are described by a set of coupled rate equations of the form,

$$\frac{dP_c(t)}{dt} = \sum_{c' \neq c} (\Gamma_{c' \rightarrow c} P_{c'}(t) - \Gamma_{c \rightarrow c'} P_c(t))$$

where  $P_c$  is the (relative) population of the  $c$ -th configuration,  $\Gamma_{c' \rightarrow c}$  is the rate for generating the configuration  $c$  from  $c'$  and  $\Gamma_{c \rightarrow c'}$  the rate for all decay transitions from configuration  $c$  to configuration  $c'$ . In a simplified dynamical model,  $\Gamma$  can be either a (configuration-averaged) time-independent Auger or fluorescence rate, or a time-dependent photoionization rate given by  $\sigma_p J(t)$  with the photon flux  $J(t)$  of the x-ray pulse at some given time  $t$ .

- For unpolarized and linearly-polarized x-rays, the differential cross section for coherent scattering is given by (Son *et al.*, 2011)

$$\frac{d\sigma^{(\text{unpolarized})}}{d\Omega} = \alpha^4 |f^{(0)}(Q)|^2 \frac{1 + \cos^2 \vartheta}{2}, \quad \frac{d\sigma^{(\text{lin-polarized})}}{d\Omega} = \alpha^4 |f^{(0)}(Q)|^2 (1 - \cos^2 \varphi \sin^2 \vartheta),$$

where  $\varphi$  is the azimuthal angle of the scattered photon momentum with respect to the x-ray propagation and polarization axes. This differential cross section gives distributions of the scattered x-ray emission if one restricts the model to fixed electronic configurations. From measurement of the x-ray distribution, one can retrieve information about the electronic density. If an atomic cloud is exposed to an ultraintense x-ray pulse, the atomic electron density is dynamically modified. This may require to introduce a suitably averaged, time-dependent differential scattering cross section,

$$\frac{d\sigma^{(\text{averaged})}}{d\Omega}(t) = \sum_c^{\text{all configs}} P_c(t) \frac{d\sigma^{(c)}}{d\Omega}$$

where  $P_c(t)$  is the population of the  $c$ -th configuration as obtained from the rate equations above. The differential scattering cross section for the  $c$ -th configuration is evaluated from the form factor for the charge density  $\rho(r)$  of the  $c$ -configuration.

## 7.4.n. Decay of medical radioactive isotopes

### Radioactive isotopes used in clinical therapy:

- **Linear energy transfer (LET):** Most radioactive isotopes, that are used in clinical therapy, emit  $\beta$  particles as ionizing radiation. The biological effect of this radiation is often characterized by the so-called linear energy transfer (LET) and measured in units of  $[\text{keV}/\mu\text{m}]$  for the deposited energy along the particle track.
- The LET for most therapeutic  $\alpha$  emitters ranges from 25 to 230  $[\text{keV}/\mu\text{m}]$ .
- A new class of radioactive nuclides refer to the  $\alpha$  emitters  $^{149}\text{Tb}$ ,  $^{213}\text{Bi}$ ,  $^{211}\text{Po}$ ,  $^{211}\text{A}$ ,  $^{223}\text{Ra}$ ,  $^{225}\text{Ac}$ ,  $^{227}\text{Ac}$ ,  $^{226}\text{Th}$  and  $^{230}\text{U}$  and have been considered for therapy.
- The electrons and positrons emitted in the nuclear  $\beta$  decay and the internal conversion processes often have kinetic energies from tens of keV to several MeV, while their LET  $\sim 0.2 [\text{keV}/\mu\text{m}]$  is much lower than for  $\alpha$  emitters.

### 7.4.o. Configuration-averaged energies and cross sections

#### Configuration averages:

- For various applications, it is desirable to use and readily access configuration-averaged energies, rates and cross sections.
- **Configuration-averaged energies** Cowan (1981) provides a non-relativistic expression for the configuration-averaged energy as it often occurs in the computation of radiative and collisional rates (Peyrusse, 1999)

$$E_c = \sum_a N_a \langle a \rangle + \sum_{ab} N_a (N_b - \delta_{ab}) \langle a, b \rangle$$

$$\langle a \rangle = \varepsilon_a + \left\langle a \left| -V^{(\text{SCF})}(r) - \frac{2Z}{r} \right| a \right\rangle, \quad \langle a, b \rangle = \frac{1}{2} \frac{g_a}{g_a - \delta_{ab}} \left[ R^{(0)}(ab, ab) - \frac{1}{2} \sum_k \begin{pmatrix} \ell_a & k & \ell_b \\ 0 & 0 & 0 \end{pmatrix}^2 R^{(k)}(ab, ba) \right].$$

Here,  $g_a = 2(2\ell_a + 1)$  is the degeneracy of subshell  $a$ ,  $V^{(\text{SCF})}(r)$  is the self-consistent atomic potential and  $R$  is the non-relativistic Slater integral.

### 7.4.p. Mass attenuation coefficients

#### Attenuation coefficients:

- The photon mass-attenuation coefficient  $\mu/\rho$  is calculated from the density  $\rho$  of the target material and the thickness  $d$  of the material in the direction of the incident photons as

$$\frac{\mu}{\rho} = -\frac{1}{\rho d} \ln \left( \frac{N}{N_o} \right),$$

and where  $N_o$ ,  $N$  are the numbers of incident and transversing photons, respectively.



## 7.4.q. Strong-field ionization phenomena and theory

### Strong-field phenomena:

- **Tunnel ionization:** Tunnel ionization at every or every second half-cycle can be enhanced or suppressed by using a third or second harmonic field with parallel polarization.
- **Frustrated tunnel ionization:** Frustrated tunnel ionization is one of the dominant channels in strong-field ionization that results in the excitation of atoms. Theory predicts that the excitation efficiency increases with the pulse duration and decreases when going to a few-cycle regime.
- **Subcycle tracing of ionization enabled by infrared (STIER) pulses:** STIER pulses refers to a streak camera that temporally resolves strong-field ionization caused by a linearly-polarized few-cycle pulse. STIER provides insight into the sub-cycle dynamics of strong-field ionization. STIER samples the photoelectrons produced by a few-cycle laser pulse in the near-visible spectral range, for example at 735 nm (Kübel *et al.*, 2017).
- **STIER:** This method enables one to separate the electron wave packets generated at different half-cycle maxima of the visible pulse. The best separation occurs, when the visible pulse is centered around a field maximum of the IR pulse, i.e. when the signal in the STIER trace is centered around  $p_z = 0$ . In this case, the observed momentum distribution directly reflects the ionization dynamics during a single half-cycle.
- **Stabilization of atoms against ionization in intense fields:** Atomic excitation in strong optical laser fields take place even at intensities that clearly exceed saturation. This is closely related to the existence of the so-called Kramers-Henneberger (KH) atom, a bound atomic system in an intense laser field.

### Theoretical methods:

- **Time-dependent CI singles:** This one-electron formulation for describing strong-field ionization is guided by the first step in all the strong-field processes, in which an electron is pulled away from the parent ion by the laser field. The use of the single-active electron (SAE) approximation has been shown to be sufficient to predict the single-ionization production, the cutoff in the HHG spectrum as well as the ATI photoelectron spectra in different atoms.
- **Beyond the SAE:** An approximate one-electron approach, that goes beyond the SAE treatment, is time-dependent configuration interaction with single excitations (TDCIS). This is an *ab-initio* electronic-structure technique, in which the time-dependent wave function is restricted to spin-singlet conserving single-particle excitations from the ground-state determinant.

### Above-threshold ionization (ATI):

- Above-Threshold Ionization refers to the absorption by an atomic electron of more photons than are required for ionization. ATI was discovered in 1979 and is a rather universal phenomenon. This process occurs in light pulses that are short and strong enough to eliminate the influence of collisions.

#### 7.4.r. Guo-Aberg-Crasemann (GAC) theory. Strong-field ionization

### Differential GAC ionization cross sections:

- Lowest-order perturbation theory typically fails to explain the the shape of the envelope of the individual ATI photoelectron peaks. Perturbation theory also fails to predict the dependence of the strength of the ATI peaks on the intensity of the laser light, if the laser intensity exceeds  $1 \text{ TW cm}^{-1}$ .
- **Differential transition-rate formula:** Within the GAC theory, the lowest-order transition matrix element can be utilized together with the (integer) constraint  $z = m - n$ , in order to obtain a **differential transition-rate formula for the ionization of electron that have absorbed  $q = lm$  photons** (Guo *et al.*, 1989)

$$\frac{d\omega}{d\Omega} = \frac{(2m^3\omega^5)^{1/2}}{4\pi^2} |\Phi_i(\mathbf{p} - q\mathbf{k})|^2 q^2 \left(q - \frac{E_b}{\omega}\right)^{1/2} |\mathcal{J}_{q+z}|^2 |\mathcal{J}_z|^2,$$

where  $\Phi_i(\mathbf{p} - q\mathbf{k})$  is the Fourier transform of the initial-state wave function with binding energy  $E_b$ . The angular and polarization dependences of this differential transition-rate formula are determined by  $\mathbf{p} - q\mathbf{k}$  in  $\Phi_i$  as well as by the arguments of the two **elliptically polarized Bessel functions  $\mathcal{J}_t$** .

- In the nonrelativistic limit and for a large photon-number, the final state can be reduced to a single but modified Volkov state, if the ponderomotive energy is an integer multiple of the photon energy. Therefore, the scattering amplitude strictly exists only in the integer case, where it becomes a product of the original KFR amplitude and the overlap between the final plane-wave state and the Volkov state, and vanishes otherwise.

### 7.4.s. Two-color sideband formation in atomic photoionization

#### Processes considered so far:

- Photoionization by a long twisted pulse and a long NIR pulse:

$$\hbar\omega^{(\text{long, twist})} + A + \hbar\omega^{(\text{String NIR})} \longrightarrow A^+ + e_p^- (E_p, \text{sidebands})$$

- In this strong-field ionization process, the photoelectron moves after its ionization in a Volkov state with the NIR field, and this gives rise to well-defined sidebands in the photoionization probability  $W(E_p; m, \Lambda_X, \Lambda_L)$ .

### 7.4.t. Phase control of atomic photoionization in multi-color fields

#### Phase control:

- **Coherent-control schemes:** While various quantum coherent-control schemes have been developed and utilized already within the optical and IR regions, their use at high frequencies has been hampered by the lack of sufficient coherent XUV sources. An important step ahead was made by seeded FEL beams, for instance at FERMI, which provide high longitudinal coherence. This coherence of the light beams enables one to observe interference of two- and one-photon pathways in the photoelectron angular distributions following the XUV irradiation.
- **Application of phase control:** The phase control of photoprocesses has attracted much recent interest and has found promising applications, such as the control of chemical reactions, the study of biological changes, ultrafast and nonlinear optics as well as four-dimensional (4D) ultrafast electron microscopy.
- **Differential ionization probability:** In second-order perturbation theory, the angle-differential ionization probability for the emission of a photoelectron and for a transition  $i \rightarrow f$  of the photoion can be expressed as

$$\frac{dW}{d\Omega} = \frac{1}{2J_i + 1} \sum_{M_i, M_f \mu} \left| \eta U_{M_i, M_f \mu}^{(1, 2\omega)} + U_{M_i, M_f \mu}^{(2, \omega)} \right|^2$$

where  $U_{M_i, M_f \mu}^{(1, 2\omega)}$  is the first-order amplitude for ionization by the second harmonic and  $U_{M_i, M_f \mu}^{(2, \omega)}$  the second-order amplitude for two-photon ionization by the fundamental frequency. Here, the magnetic quantum numbers refer to the initial atomic state  $|J_i M_i\rangle$ , the residual ion  $|J_i M_i\rangle$  as well as the photoelectron  $|\mathbf{p} \mu\rangle$ .

## 7. Atomic processes

### ➤ Electric dipole approximation:

$$U_{M_i, M_f \mu}^{(1, 2\omega)} = -i \langle \alpha_f \mathbb{J}_f M_f, \mathbf{p} \mu^{(-)} | e \mathbb{D} | \alpha_i \mathbb{J}_i M_i \rangle T^{(1)}$$

$$U_{M_i, M_f \mu}^{(2, \omega)} = \sum_{\nu} \int \langle \alpha_f \mathbb{J}_f M_f, \mathbf{p} \mu^{(-)} | e \mathbb{D} | \alpha_{\nu} \mathbb{J}_{\nu} M_{\nu} \rangle \langle \alpha_{\nu} \mathbb{J}_{\nu} M_{\nu} | e \mathbb{D} | \alpha_i \mathbb{J}_i M_i \rangle T_{E_{\nu}}^{(2)}.$$

Here  $\alpha_{\nu}$  specifies the quantum numbers of the intermediate states  $\nu$ , while  $\sum_{\nu} \int$  denotes the infinite summation over all discrete and continuum states. The minus sign in the outgoing electron wave indicates the proper asymptotic form of the continuum wave function,  $E_{\nu}$  is the energy of the intermediate virtual atomic state. The time-dependent  $T^{(1)}$  and  $T_{E_{\nu}}^{(2)}$  were described elsewhere in the literature (Gryzlova *et al.*, 2018).

## 8. Atomic cascades

### 8.1. Cascade approaches and cascade simulations

#### Cascade approaches:

- Atomic cascades (may) often require an enormous amount of computational effort in order to generate, simulate and interpret all data as required for a given experiment. It therefore appears necessary and highly desirable to distinguish a number of predefined cascade approaches which define a hierarchy of possible approximations for a given scenario.
- Here, we briefly summarize these (predefined) approaches which can be further refined by choosing proper parameters for the given computation.
- **Decay cascade:** Such a cascade usually starts from a given set of initial levels (and together with their relative population) that may decay by various atomic processes, such as autoionization, photon emission and others, until a given number of electrons is released and/or the ions cannot further decay.
- i) **Average single-configuration approach (AverageSCA):** In this approach, only **one common set of orbitals is applied for all ionization stages of the atom**. Moreover, all electron configurations in the cascade are treated in terms of single-CSF levels, i.e. without any configuration mixing.
- In practice, the AverageSCA approach should be realistic for (almost) all atoms and ions from the periodic table, independent of how well this approach describes the underlying physics of a given autoionization cascade.
- ii) **Single-configuration approach (SCA):** This approach applies a common set of orbitals for electron configuration (cascade block) in the given cascade. stage of the atom.
- iii) **Multiple-configuration approach (MCA):** In this approach, we intent to incorporate further electron-electron correlations by configuration mixing contributions with closeby-lying configurations; the program attempts to *group* different configurations together by using physical insight into the cascade process.

## 8. Atomic cascades

- For each of these (cascade) approaches above, the mean computational effort should increase by (at least) one order of magnitude, perhaps also by several orders.
- Although the overall computational effort is quite difficult to estimate and formalize for general atomic systems, **the various cascade approaches should be clearly discernible from each other and the 'costs' of any simpler approach (than the currently chosen one) should be typically negligible.** In general, however, not all of these approaches will be feasible for a given cascade computation.

### Implementation of cascade computations:

- Formally, a cascade computation is separated into several steps which are treated independently: (i) A list of electron configurations is generated as they may occur during the decay of levels from one or several initial (hole-state) configurations; (ii) From the configurations, a list of *cascade blocks* (groups of configurations) is formed and a common SCF field and Hamiltonian matrix calculated for each of these blocks; (iii) For each pair of cascade blocks, all fine-structure transitions are calculated for all selected decay processes; (iv) A (common) list of radiative lines, Auger lines, etc. is compiled and forms the basis for all subsequent cascade simulations.

### Cascade simulations:

- **Cascade simulation:** A cascade simulation makes use of the calculated (many-electron) amplitudes, rates and cross sections in order to derive different distributions and spectra which can be compared with experiment. **Cascade computations and cascade simulations are always handled independently.**

## 8.2. In JAC implemented cascade approaches

### 8.2.a. Average single-configuration approach (AverageSCA), using a common set of orbitals (Cascade)

#### Cascade, notations & application:

- **AverageSCA approach:** This approach applies a **common set of orbitals for all steps of the cascade.** Moreover, each single configuration forms an independent multiplet (cascade block) and all configuration interactions are omitted in this simple approach. Generally, the orbitals of the initial configurations/levels are utilized also in order to compute the cross sections and rates for all subsequent steps of the

cascade. This single-CSF representation of all atomic levels is assumed to be rather equivalent to the use of *configuration-averaged rates and cross sections*, although this equivalence is far from being obvious for complex configurations.

- **General restrictions:** (a) Only electric-dipole (E1) transitions are included for all radiative decays; (b) continuum orbitals are generated at an configuration-averaged transition energy and only for partial waves with  $|\kappa| \leq 4$ .

### Computational steps of the AveragedSCA and SCA approaches:

- The following steps are carried out separately in these two approaches:
  - 1) **Cascade blocks:** Determine all single configurations that are energetically allowed in photon emission and autoionization processes from the initially chosen configurations/levels, and until a given maximum number of electrons are released. These electron configurations form a list of (so-called) cascade blocks from which the decay cascade is then built by calculating all requested decay lines for each pair of blocks.
  - 2) **Radial orbitals:** Bound radial orbitals are calculated only for the first (and uppermost) block in the cascade, i.e. the initial hole configuration with the energetically highest-lying energy levels. These bound orbitals are then applied in order to set the `basis::Basis` of all other blocks.
  - 3) **Level representation:** An individual level multiplet is calculated for each single non-relativistic electron configuration (cascade block) but without any CI, i.e. without configuration mixing. Therefore, the levels in all cascade blocks are given by a single CSF, and no explicit diagonalization of the Hamiltonian matrix is needed in this approach.
  - 4) **Decay processes & cascade steps:** The requested radiative and non-radiative decay processes are considered between each configuration pair. This is achieved by performing standard atomic computations with the corresponding submodules, such as `AutoIonization` or `PhotoEmission`. The following restrictions are applied throughout the computations: (a) Only electric-dipole (E1) transitions are included for all radiative decays; (b) continuum orbitals are generated at an configuration-averaged transition energy and only for partial waves with  $|\kappa| \leq 4$ . – This step gives rise to a list of `lines` but without that all details are printed to screen.
  - 5) **Decay lines:** For each pair of cascade blocks, the calculated decay lines are appended to common lists `linesR`, `linesA`, ... as finally returned by an instance of `Cascade.Data`.

### 8.2.b. Single-configuration approach (SCA) with individual sets of orbitals (Cascade)

#### Cascade, notations & application:

- **SCA approach:** This approach applies an **individual set of orbitals for all steps of the cascade**, while each individual electron configuration in the decay tree still forms an independent cascade block, i.e. level multiplet. For each of these blocks, an independent SCF is generated and the configuration mixing included by diagonalizing the corresponding Hamiltonian matrix. **The single-configuration approach (SCA) is expected to already provide a quite reasonable description of all strong decay pathes.**
- **General restrictions:** (a) Multipoles of the radiative transitions can be defined by the given **Settings**; (b) continuum orbitals are generated in a slightly simplified Dirac-Fock-Slater potential of the corresponding (final-level) configuration but for energies of the individual fine-structure transitions. (c) No orbital relaxation is incorporated into the amplitudes apart from the separate orbital representation (SCF) of the initial and final levels.
- See subsection 8.2.a for the individual steps of the cascade computation.

## 8.3. In JAC partly-implemented cascade approaches

### 8.3.a. User-grouped multi-configuration approach (UserMCA; Cascade)

#### Cascade, notations & application:

- **UserSCA approach:** This approach applies an **individual set of orbitals for all steps of the cascade**, which can be compiled by the user. For this, the program first provides a list of (single electron) configurations that can be either combined interactively or by some explicit calls (code) in the function . . . . For each of these user-specified cascade blocks, an independent SCF is generated and the configuration mixing included by diagonalizing the corresponding Hamiltonian matrix. **The UserMCA enables one to include further physical insight about strong configuration mixing into the computation of the cascade.**
- **General restrictions:** (a) Multipoles of the radiative transitions can be defined by the given **Settings** or directly in the code; (b) continuum orbitals are still generated in a simplified Dirac-Fock-Slater potential of the corresponding (final-level) configuration but for energies of the individual fine-structure transitions. (c) No orbital relaxation is incorporated into the amplitudes apart from the separate orbital representation (SCF) of the initial and final levels.



- See subsection 8.2.a for the major steps of the cascade computation.

### 8.3.b. Multi-configuration approach with shake transitions (ShakedMCA; Cascade)

#### Cascade, notations & application:

- **ShakedMCA approach:** This approach applies an *individual set of orbitals for all steps of the cascade*. However, these steps do include not only the nominal electron configurations as they obviously occur in the cascade but also configurations with the *replacement of additional electrons*, and which refer to shake-up or shake-down transition. The generation of these additional configurations are done via a *maximal number of electron replacements* that are considered in addition to the standard configurations.
- These *shake configurations* can be treated as individual groups of configurations or together with other configurations of the decay cascade.
- Similar as in the UserMCA, the program first provides a list of (single electron) configurations that can be either combined interactively or by some explicit calls (code) in the function . . . . For each of these user-specified cascade blocks, again, an independent SCF is then generated and the configuration mixing included by diagonalizing the corresponding Hamiltonian matrix.
- Although the UserMCA and ShakedMCA are quite similar in spirit, the incorporation of shake-up or shake-down configurations typically enlarges the effort enormously and need to be handled with great care. *The ShakedMCA enables one to include further physical insight about strong shake processes and configuration mixing into the computation of the cascade.*
- **General restrictions:** (a) Multipoles of the radiative transitions can be defined by the given **Settings** or directly in the code; (b) continuum orbitals are still generated in a simplified Dirac-Fock-Slater potential of the corresponding (final-level) configuration but for energies of the individual fine-structure transitions. (c) No orbital relaxation is incorporated into the amplitudes apart from the separate orbital representation (SCF) of the initial and final levels.

## 8.4. Further cascade approaches not yet considered in JAC

### Light pulses: shape and intensities

➤

### Rate equations

➤

## 8.5. In JAC implemented cascade simulations

### 8.5.a. Ion distribution (...; Cascade)

#### Distribution, notations & application:

➤

## 8.6. In JAC partly-implemented cascade simulations

### 8.6.a. Electron intensity spectrum (...; Cascade)

#### Spectrum, notations & application:

➤

### 8.6.b. Photon intensity spectrum (...; Cascade)

#### Spectrum, notations & application:





## 9. Field- and collision-induced atomic responses

### Atoms in intense laser pulses

- Atoms, which are irradiated by intense laser pulses, can emit radiation at multiples of the laser frequency, the (high) so-called harmonics.
- Such high-order harmonics, i.e. highly-energetic photons, can be produced by focusing either a high-frequency (excimer) laser or some low-frequency laser (Nd:glass, Ti:sapphire) into an atomic gas.
- Especially for low-frequency lasers, the harmonic spectrum is known to exhibit a characteristic behaviour: A rapid decrease of the yield for the first few harmonics, followed by a plateau and a rather abrupt cut-off.

### 9.1. Floquet theory

#### Motivation:

- Many time-dependent problems in physics are periodic to a good approximation. For such physical systems, the (time) evolution can be solved either directly by numerical integration of the (time-dependent) Schrödinger, and making use of the time-dependent Hamiltonian.
- **Floquet formalism:** For a periodic evolution, time-dependent problems can be transformed also into a time-independent problem by applying the Floquet formalism, a very desirable transformation as time-independent problems can often be solved more efficiently.
- While, in applying the Floquet formalism, often the Lindblad and Bloch-Redfield master equations are adapted, we note that these equations are derived for a time-independent Hamiltonian. In many cases, however, these standard master equations have still been found useful for solving time-dependent problems.

### Floquet ansatz:

➤ Vector potential of a linearly-polarized electric field:

$$\mathbf{A}(t) = \epsilon A_o \sin(\omega t); \quad A_o = -\frac{c E_o}{\omega} \quad \rightsquigarrow \quad \mathbf{E}(t) = \epsilon E_o \cos(\omega t).$$

➤ Floquet ansatz for the wave function:

$$\psi(t) = \exp(-i E t) \sum_{n=-\infty}^{\infty} \exp(-i n \omega t) \psi_n, \quad E = \Re(E) - i\Gamma/2$$

where the functions  $\psi_n$  are often called the harmonic components, while  $E$  is the complex quasi-energy of the Siegert state  $\psi(t)$  and  $\Gamma$  the total ionization rate of this state.

## 9.2. In JAC implemented responses

None.

## 9.3. In JAC partly-implemented responses

None.

## 9.4. Further responses not yet considered in JAC

### 9.4.a. Collisional-radiative (CR) models

#### Basic CR models:

- Collisional radiative models are sometimes called **zero-dimensional plasma models**. They are often applied in order to calculate atomic state distribution functions for one or more atomic species as function of the particle densities and temperatures. Most CR models assume that the radiative absorption and emission as well as electron-collision processes dominate the plasma and that all other (ion-atom, three-atom, ...) processes are negligible.
- A CR model can be used to calculate the effective emissivity or the effective ionization and recombination rates as function of electron temperature and density. These rates and results can then be applied in **plasma transport model** for determining local ionization, recombination and radiation effects.
- CR models help determine atomic level population densities of atoms in different plasma environments. Calculations are typically based on the CR recombination and ionization rate coefficients, radiative energy losses as well as on various individual rate coefficients.
- Different CR models can be distinguished due to the – analytic or numerical – treatment of the rate coefficients as well as the number and type of processes, for which such rate coefficients are taken into account.
- The CR models have been found useful for the interpretation of experimental observations, for instance in order to derive the electron density and temperature. In addition, CR models can be applied to **determine effective ionization and recombination rates as well as the effective emissivity of plasma as function of its electron temperature and density**.
- **Collisional-radiative (CR) models:** These 0-dimensional plasma models are typically applied in order to compute **atomic charge-state distributions for – (ionic) mixtures of – one or several elements**, and in particular as function of particle densities and temperatures.

#### Approximations:

- CR models are typically built on the level densities  $n(p)$  for a series of levels  $p_1, p_2, \dots, p_N$ . In general, the relation between the temporal and spatial distribution of the level density  $n(p)$  is given by the diffusion equation

$$\left( \frac{\partial n(p)}{\partial t} \right)_{\text{CR}} = \frac{\partial n(p)}{\partial t} + \nabla \cdot (n(p) \omega(p))$$

## 9. Field- and collision-induced atomic responses

Often, several atomic levels are simply combined into a single (effective) level  $p_e$ , and where an equal distribution is assumed for all of its physical levels:  $n(p) = n(p_e)/g(p_e)$  with  $g(p_e)$  being the statistical weight of the effective level  $p_e$ .

➤ If changes in the **level density are dominated by radiative and collisional (de-) excitation processes**, most CR models simply assume

$$\left( \frac{\partial n(p)}{\partial t} \right)_{\text{CR}} = 0 = P(p) - n(p) D(p)$$

and just solve a matrix equation. Here,  $P(p)$  is the collisional and radiative production rate, and  $D(p)$  the (so-called) *destructive* factor. The two functions  $P(p)$ ,  $D(p)$  can be written in terms of different rate coefficients for electron-impact excitation, ionization as well as for radiative and three-body recombination.



# 10. Time-evolution of many-electron atomic state functions and density matrices

## 10.1. Time-dependent approximations of many-electron states

### Time-dependent methods:

- **Time-dependent theory:** The response of multi-electron atoms to strong laser fields with intensity  $\sim 10^{14}$  W/cm<sup>2</sup> can be typically described only by explicitly time-dependent approaches. For short and intense light pulses, for example, the ionization rate is known to change rapidly during the pulse, and this makes the use of just a monochromatic field incorrect in the theoretical description. For laser intensities  $\gtrsim 10^{16}$  W/cm<sup>2</sup>, moreover, the electric field becomes even stronger than the inter-atomic fields, leading to a breakdown of perturbation theory. Then, the time-dependent Schrödinger equation need to be integrated explicitly in order to properly deal with the shape and intensity of the light pulses, although the *limits* of perturbation theory are not yet well understood in this time and intensity domain.
- **Light sources for studying electron dynamics in atoms:** New light sources and technologies facilitate nowadays various experiments with intense attosecond extreme-ultraviolet (XUV) pulses in order to resolve in time the ejection of electrons from the atomic target.
- **New coherent light sources:** Recently developed light sources include attosecond lasers, high-order-harmonic generation (HHG) sources or free-electron lasers. These novel light will pave the way for studying the interaction of matter with intense femtosecond or even sub-femtosecond radiation pulses in the visible and XUV regime.
- The development of sub-femtosecond radiation pulses enables one to explore the electron dynamics in atoms *or* the (combined) electron and nuclei dynamics in molecules with a high degree of temporal and spatial resolution.
- **Need for new time-dependent theoretical models:** To support modern experiments with intense atto-second pulses, new time-dependent models and methods need to be developpe beyond the current standard schemes, such as the single-active-electron approximation.

## 10. Time-evolution of many-electron atomic state functions and density matrices

- **Time-dependent Hartree-Fock (TDHF) method:** In this method, the time evolution of the electronic orbitals is accounted for during the interaction with an external field or collision partner by following the (time-dependent) mean field of the total electron density in the given potential. This method therefore properly treats the nuclear interaction and the *mean field* of the electrons but averages over all electronic correlations. In practice, such a mean-field approximation significantly reduces the dimension of the  $N$ -particle quantum system from  $3N$  spatial coordinates to just three coordinates for the *mean* density. **The TDHF method is an initial value problem that typically starts from the wave function of the static ground state of the system.**
- **Time-dependent RAS method with space partition:** In this method, the wave functions are expanded in terms of Slater determinants with explicitly time-dependent CI coefficients. The Slater determinants are usually based on two type of time-independent orbitals: (i) Orbitals that are just defined in an inner region close to the atomic nucleus and (ii) orbitals in an outer region with zero-contribution to the inner region (Miyagi and Madsen, 2017).
- **Space partition methods:** A partitioning of space has been widely utilized for describing atom-electron collisions since the very first advent of the R-matrix theory. In all these models, a configuration-interaction (CI) expansion of the many-electron wave function is typically made in some inner region in order to deal with electronic correlations. In the outer region, in contrast, the electrons are supposed to move independently of each other in just a long-range effective potential, and without that further exchange and correlation contributions are taken into account. This piecewise definition of the wave function then requires to make sure that it remains continuously differentiable across the surface between the two regions.
- **Time-dependent MCHF (TD-MCHF) method:** This method is presently seen as the most accurate and widely used technique in order to describe the many-electron dynamics in explicitly time-dependent potentials, such as the laser-induced dynamics of  $N$ -electron atoms. When compared to the space-partition methods, the TD-MCHF method has the great advantage that, for a sufficiently large (radial) domain of the orbitals, all single- and multi-electron interactions (continua) can be taken into account, at least in principle.
- **Time-evolution with separable (atomic) potentials (TESP):** Galstyan *et al.* (2018) describe a new and **computationally inexpensive method for studying the interaction of complex atoms or molecules with intense ultrashort laser pulses within the single active electron approximation.** With this method, they have analyzed the photoionization of water molecules from the highest occupied molecular orbital and calculated, in particular, the ionisation yield for different orientations of the molecule with regard to the polarization of the incident pulse.
- **TESP method:** This method is an approximation since the exact solution would require to incorporate an infinite number of separable potentials in the representation of the kernel. This method is therefore especially useful for pulses with high frequency where only the lowest (few) levels play a crucial role for the ionization of the system.

- **Central idea of the TESP method:** The main idea of this method is to work in **momentum space and to replace the kernel of the Coulomb potential by a sum of  $N$  symmetric separable potentials which each support just a single bound state of the system**. The use of separable potentials enables one to reduce the  $N$ -dimensional time-dependent Schrödinger equation (TDSE) to a system of  $N$  one-dimensional Volterra integral equations which only depend on time. This method is known also by Separable Potentials for Atoms and Molecules (SPAM).
- **Weaknesses of the TESP method:** The separable potential is not necessarily unique and does usually not represent the correct asymptotic behaviour of the potential. Moreover, the method is not gauge invariant and the question about proper gauges will need to be explored in further detail.

### Time-evolution of ASF:

- A reliable description of the time-dependent (TD) many-electron dynamics has attracted recent interest in order to theoretically support a real-time analysis and control of the ultrafast electronic and nuclear dynamics of atoms and molecules in intense laser pulses.
- A computational and conceptual simple approach to the many-electron dynamics is the TD configuration-interaction singles (TD-CIS) method in which only singly excited configurations, relative to the Hartree-Fock ground state, are taken into account in the configuration interaction expansion.
- Miyagi and Madsen (2013) make use of the TD-CIS method by performing an optimization of the orbitals as well as the expansion coefficients at each time step in the evolution. This separate optimization makes it possible to construct a reasonable accurate wave function with just a relatively small number of electronic configurations. For the optimization, Miyagi and Madsen (2013) utilized the Dirac-Frenkel-McLachlan TD variational principle.

## 10.2. Time-dependent statistical tensors

### Motivation:

- In physics, Liouville's theorem is known from classical statistical and Hamiltonian mechanics, and it asserts that the phase-space distribution function is constant along the trajectory of the system.
- Similarly, the quantum-mechanical analog of Liouville equation describes the time evolution of a mixed state. The quantum form can be obtained by the standard replacements of classical variables with quantum operators and Poisson brackets with commutators.

## 10. Time-evolution of many-electron atomic state functions and density matrices

➤ The Liouville equation for the atomic density matrix  $\rho$  reads

$$\frac{\partial \rho}{\partial t} = -\frac{i}{\hbar} [\mathbb{H}, \rho] - \Gamma \rho, \quad \mathbb{H} = \mathbb{H}^{(\text{atom})} + \mathbb{H}^{(e-\gamma)}$$

where the first term in the Hamiltonian describes the (inner-) atomic and the atom-photon interactions, and the second term the relaxation matrix  $\Gamma$ . This relaxation matrix accounts for the atomic spontaneous decay due to either radiative or nonradiative (Auger) processes of the atom.

### 10.3. Time-integration of statistical tensors

#### Motivation:

➤ Various methods can be applied to solve the coupled first-order ODEs for the time-evolution of the statistical tensors. Apart from those methods which solve the equations with given accuracy (but typically require the evaluation of the rhs at any given time  $t$ ), **we are mainly interested in shooting or predictor-corrector methods** which make only use of previous solutions of  $\rho$  and  $d\rho/dt = f$  on an equidistant time grid.

#### Numerical methods:

➤ If we use a step-size  $h$  and define  $\rho(t_j) = \rho_j$ ,  $\frac{d\rho}{dt}(t_j) = f_j$ , the following integration formulas can be applied.

➤ **Euler's methods:**

$$\rho_{j+1} = \rho_j + h f_j$$

➤ **Adams-Bashford methods:**

$$\rho_{j+1} = \rho_j + h \left[ \frac{3}{2} f_j - \frac{1}{2} f_{j-1} \right]$$

$$\rho_{j+1} = \rho_j + h \left[ \frac{23}{12} f_j - \frac{16}{12} f_{j-1} + \frac{5}{12} f_{j-2} \right]$$

$$\rho_{j+1} = \rho_j + h \left[ \frac{55}{24} f_j - \frac{59}{24} f_{j-1} + \frac{37}{24} f_{j-2} - \frac{9}{24} f_{j-3} \right]$$

## 10.4. Time evolution of statistical tensors. Formalism

### 10.4.a. Liouville equation for the atomic density matrix

#### Vector potential and state multipoles:

- In the electron-photon interaction, the vector potential of the electromagnetic field is typically chosen to fulfill  $\text{div} \mathbf{A}(\mathbf{r}_i, t) = 0$ .
- When the electron-photon interaction Hamiltonian  $\mathbb{H}^{(e-\gamma)}$  is applied upon the density operator  $\rho$ , we shall consider only processes with one photon.
- Moreover, we assume a vector potential for a monochromatic field with frequency  $\omega$  and wave vector  $\mathbf{k} = (\vartheta_k, \varphi_k)$  ( $|\mathbf{k}| = \frac{\omega}{c}$ ) with regard to the coordinates in the laboratory system.
- **Envelope of the vector potential:** For a pulse with (real) envelope  $f(t)$ , the vector potential then includes two terms

$$\mathbf{A}(\mathbf{r}, t) = \mathbf{A}(\mathbf{r}) f(t) e^{-i\omega t} + \mathbf{A}^*(\mathbf{r}) f^*(t) e^{i\omega t}.$$

## Liouville equation in the representation of well-defined angular momenta:

➤ **Liouville equation:** In an atomic basis with well-defined total angular momenta and parity, the Liouville equation can be re-written as

$$\begin{aligned} \frac{\partial}{\partial t} \langle \alpha \mathbb{J}_\alpha M_\alpha | \rho | \beta \mathbb{J}_\beta M_\beta \rangle &= -\frac{i}{\hbar} (E_\alpha - E_\beta) \langle \alpha \mathbb{J}_\alpha M_\alpha | \rho | \beta \mathbb{J}_\beta M_\beta \rangle \\ &- \frac{i}{\hbar} \sum_{\nu \mathbb{J}_\nu M_\nu} \left( \langle \alpha \mathbb{J}_\alpha M_\alpha | \mathbb{H}^{(e-\gamma)} | \nu \mathbb{J}_\nu M_\nu \rangle \langle \nu \mathbb{J}_\nu M_\nu | \rho | \beta \mathbb{J}_\beta M_\beta \rangle - \langle \alpha \mathbb{J}_\alpha M_\alpha | \rho | \nu \mathbb{J}_\nu M_\nu \rangle \langle \nu \mathbb{J}_\nu M_\nu | \mathbb{H}^{(e-\gamma)} | \beta \mathbb{J}_\beta M_\beta \rangle \right) \\ &- \sum_{\gamma \mathbb{J}_\gamma M_\gamma} \langle \alpha \mathbb{J}_\alpha M_\alpha | \Gamma | \gamma \mathbb{J}_\gamma M_\gamma \rangle \langle \gamma \mathbb{J}_\gamma M_\gamma | \rho | \beta \mathbb{J}_\beta M_\beta \rangle \end{aligned}$$

### 10.4.b. Time-dependent statistical tensors of atomic lines

#### Time-dependent statistical tensors:

- **Statistical tensors:** With the substitutions of the statistical tensors from above and the notations  $[ab...c] \equiv (2a+1)(2b+1)...(2c+1)$ , the Liouville equation can be written as Grum-Grzhimailo, 2012)

$$\begin{aligned}
 i\hbar \dot{\rho}_{kq}(\alpha \mathbb{J}_\alpha, \beta \mathbb{J}_\beta) &= -2\pi \sqrt{\frac{\hbar}{c k}} \sum_{\nu \mathbb{J}_\nu K Q} \\
 &\left[ f(t) e^{-i(\omega - \omega_{\alpha\nu})t} \sum_{LM\lambda} g_\lambda (-1)^{J_\alpha + J_\beta + K + L} [LK]^{1/2} \langle KQ, LM | kq \rangle \left\{ \begin{matrix} J_\alpha & J_\beta & k \\ K & L & J_\nu \end{matrix} \right\} D_{M\lambda}^L(\varphi_k, \vartheta_k, 0) \left( \langle \alpha \mathbb{J}_\alpha \| \mathbb{T}_L^{(m)} \| \nu \mathbb{J}_\nu \rangle + i\lambda \langle \alpha \mathbb{J}_\alpha \| \mathbb{T}_L^{(e)} \| \nu \mathbb{J}_\nu \rangle \right) \tilde{\rho}_{KQ}(\nu \mathbb{J}_\nu, \beta \mathbb{J}_\beta) \right. \\
 &- f^*(t) e^{i(\omega - \omega_{\nu\alpha})t} \sum_{LM\lambda} g_\lambda^* (-1)^{J_\beta + k + J_\nu + L} [Lk]^{1/2} \langle kq, LM | KQ \rangle \left\{ \begin{matrix} J_\alpha & J_\beta & k \\ K & L & J_\nu \end{matrix} \right\} D_{M\lambda}^{L*}(\varphi_k, \vartheta_k, 0) \left( \langle \nu \mathbb{J}_\nu \| \mathbb{T}_L^{(m)} \| \alpha \mathbb{J}_\alpha \rangle - i\lambda \langle \nu \mathbb{J}_\nu \| \mathbb{T}_L^{(e)} \| \alpha \mathbb{J}_\alpha \rangle \right) \tilde{\rho}_{KQ}(\nu \mathbb{J}_\nu, \beta \mathbb{J}_\beta) \\
 &+ f(t) e^{-i(\omega - \omega_{\nu\beta})t} \sum_{LM\lambda} g_\lambda (-1)^{J_\alpha + J_\beta + K} [LK]^{1/2} \langle KQ, LM | kq \rangle \left\{ \begin{matrix} J_\alpha & J_\beta & k \\ L & K & J_\nu \end{matrix} \right\} D_{M\lambda}^L(\varphi_k, \vartheta_k, 0) \left( \langle \nu \mathbb{J}_\nu \| \mathbb{T}_L^{(m)} \| \beta \mathbb{J}_\beta \rangle + i\lambda \langle \nu \mathbb{J}_\nu \| \mathbb{T}_L^{(e)} \| \beta \mathbb{J}_\beta \rangle \right) \tilde{\rho}_{KQ}(\alpha \mathbb{J}_\alpha, \nu \mathbb{J}_\nu) \\
 &- f^*(t) e^{i(\omega - \omega_{\beta\nu})t} \sum_{LM\lambda} g_\lambda^* (-1)^{J_\alpha + K + J_\nu} [Lk]^{1/2} \langle kq, LM | KQ \rangle \left\{ \begin{matrix} J_\alpha & J_\beta & k \\ L & K & J_\nu \end{matrix} \right\} D_{M\lambda}^{L*}(\varphi_k, \vartheta_k, 0) \left( \langle \beta \mathbb{J}_\beta \| \mathbb{T}_L^{(m)} \| \nu \mathbb{J}_\nu \rangle - i\lambda \langle \beta \mathbb{J}_\beta \| \mathbb{T}_L^{(e)} \| \nu \mathbb{J}_\nu \rangle \right) \tilde{\rho}_{KQ}(\alpha \mathbb{J}_\alpha, \nu \mathbb{J}_\nu) \Big] \\
 &- i\hbar \sum_{\gamma \mathbb{J}_\gamma} G_{kq}^{KQ}(\alpha \mathbb{J}_\alpha, \beta \mathbb{J}_\beta, \gamma \mathbb{J}_\gamma) e^{i\omega_{\alpha\gamma}t} \tilde{\rho}_{KQ}(\gamma \mathbb{J}_\gamma, \beta \mathbb{J}_\beta)
 \end{aligned}$$

## 10.5. Observables to be derived from time-dependent statistical tensors

### Motivation:

- If we solve the set of time-dependent (Liouville) equations for the state multipoles, we can obtain the atomic statistical tensors for
- (a) discrete atomic states  $\rho_{kq}(\alpha \mathbb{J}_\alpha; \beta \mathbb{J}_\beta)$       (b) for atomic states with one electron in continuum  $\rho_{kq}(\alpha_f \mathbb{J}_f, \varepsilon \kappa : \alpha \mathbb{J}_\alpha, \alpha'_f \mathbb{J}'_f, \varepsilon' \kappa' : \alpha \mathbb{J}_\alpha),$
  - (c) *nondiagonal* elements  $\rho_{kq}(\alpha_f \mathbb{J}_f, \varepsilon \kappa : \alpha \mathbb{J}_\alpha, \beta \mathbb{J}_\beta)$     and     $\rho_{kq}(\alpha_f \mathbb{J}_f, \varepsilon \kappa : \alpha \mathbb{J}_\alpha, \alpha'_f \mathbb{J}'_f, \varepsilon' \kappa' : \beta \mathbb{J}_\beta).$



# 11. Semiempirical estimates

## 11.1. In JAC implemented estimates for atomic properties and data

None.

## 11.2. In JAC partly-implemented estimates for atomic properties and data

### 11.2.a. Ionization energies

- There are numerous tabulations for the first and second ionization potentials of (almost) all elements from the periodic table.
- A simple parametrization for the ionization of inner-shell electrons has been worked out by Thayer Watkins for elements from the  $n\ell$ -shell, although it seems not to be very reliable

$$E_b = 10.42924 \left[ (Z/n)^2 + 1.371650145 \right]^2 - 16.76869944 = -11.0053 + 10.31048 (Z/n)^2 + 28.83251 (Z/n) - 2.88439 \ell$$

- In JAC, an estimate for the binding energy of an  $n\ell$ -electron in the neutral atom can be obtained from `JAC.estimate("ionization potential: inner-shell", shell::Shell, Z::Int64)`.

### 11.2.b. Atomic radii and susceptibilities

#### Role of atomic radii:

- **Atomic radii:** These radii are sometimes applied in order to estimate a number of size-dependent properties of isolated atoms from the periodic table of elements, such as the diamagnetic part of the atomic susceptibility, the atomic polarizability or the chemical hardness. For a good number of atoms, indeed, the estimated global hardness and atomic polarizability closely agree with available experimental data (Ghosh, 2002). In practice, however, these **atomic radii can hardly be verified experimentally since there is no well-defined quantum mechanical notion for the size of atoms or ions.**
- **Atomic radii:** The concept of atomic and ionic radii has been found helpful for understanding, explaining or even predicting various physico-chemical properties of atoms, ions and molecules. This is quite different to the use of such radii in crystal chemistry where they are used in order to estimate and to explain the binding lengths in crystal lattices.
- **Atomic radii:** Ghosh (2002) applied computed theoretical radii in order to estimate a number of size-dependent physical properties of isolated atoms: (1) the diamagnetic part of the magnetic susceptibility  $\chi^{(\text{diamagnetic})}$  for all elements up to xenon, (2) the atomic polarizability  $\alpha$  and (3) the global hardness  $\eta$  for all elements of the periodic table.

### 11.2.c. Weak-field ionization of effective one-electron atoms

#### Property & notations:

- **Field ionization:** by a quasi-static external field:  $A^{(*)} + F \mathbf{e}_z \longrightarrow A^+ * e^-$ .
- **Formal quantum notation:**  $|\alpha \mathbb{J}\rangle \longrightarrow \dots$
- Not yet implemented.
- In JAC, all estimates on the field-ionization rates of atoms are always based on a non-relativistic, single-electron approach. No attempt has (yet) been made to generalize this approximation towards real many-electron atoms and ions with their fine-structure.

### Motivation:

- Ionization of atoms by an external electric field is a fundamental atomic process that plays an important role, for instance, in the formation of charge-state distributions in (dilute) plasma under strong fields.
- **Vice versa:** The observations of temporal variations in the line intensities of ions, that are produced by field ionization, can help determine spectroscopically the macroscopic-field distribution in plasma.
- Since the field-ionization probability grows (very) rapidly with the electric field strength  $F$ , the field strength can be derived from accurate measurements of the ionization probabilities.

### Further information:

- **Ionization probabilities are often calculated within the WKB approximation, i.e. in the weak-field limit.** In general, field ionization probability depends critically on the projection of the total orbital angular momentum upon the direction of the field.
- In most field-ionization probability computations, a quasi-static and homogeneous electric field is assumed as it occurs, for example, for macroscopic-scale fields in plasma. This assumption is in contrast to many local fields that may occur at the atomic scale.
- For the case of a single electron with binding energy  $E$  in the valence shell ( $n\ell$ ) and outside of closed shells otherwise, the quasi-classical escape rate is given by (Fisher *et al.*, 1998)

$$\Gamma(E, \ell, m) = B_{n\ell}^2 \frac{(2\ell + 1)(\ell + m)!}{2^{m+1} k^m m! (\ell - m)!} \left( \frac{2k^2}{F} \right)^{2Z/k - m - 1} \exp \left( -\frac{2k^3}{3F} \right)$$

where  $Z$  is the charge of the parent ion as seen by the electron are large  $r$ ,  $k = \sqrt{2|E|}$  is the modulus of the (electron's) wave vector and  $B_{n\ell}$  is the amplitude that characterizes the outgoing electron wave.

### 11.2.d. Electron-impact ionization. Cross sections

#### Semi-empirical electron-impact excitation and ionization cross sections:

- **Temperature dependence of cross sections and rate coefficients:** Vriens and Smeets (1980) analyzed how the electron-impact excitation, ionization and de-excitation cross sections and rate coefficients depend on the electron energy and temperature as well a few other parameters. They cover a rather wide range of electron-impact energies from *sudden* up to *adiabatic* collisions. In particular, Vriens and Smeets (1980) constructed a set of simple analytical formulas for the electron-impact cross sections and rate coefficients as well as for the total depopulation and three-body recombination, although it appears rather difficult to follow this paper in detail, and especially to which entities the given formulas refer and when they are valid.
- **Electron-impact cross sections for CR simulations:** CR models have been applied in the literature in order to simulate different nonequilibrium plasmas; these models typically require the cross sections or rates for all major electronic transfer processes. While different approximate cross-section and rate formulas are available from the literature, these formulas often have a limited range of validity, and there occur serious discrepancies among the different approximations. Vriens and Smeets (1980) here provide practical formulas that are accurate over a reasonable wide range of primary energies and which are applicable for different transitions.
- Two characteristic energies of electron-impact processes refer to the binding energy  $E^{(\text{binding})}$  of the (least-bound) valence electron and to the kinetic energy of the incident electrons  $E^{(\text{impact})}$ .
- **Two regimes for electron-impact processes:** (i) In the sudden regime of rather high impact energies,  $E \gtrsim 4 E^{(\text{binding})}$ , there is generally good agreement between the binary-encounter and the Bethe-Born approximation. (ii) For low collision energies  $E^{(\text{binding})} \lesssim E \lesssim 4 E^{(\text{binding})}$  the binary-encounter cross sections are larger (up to  $\sim 40\%$ ) than the classical three-body as obtained by Monte-Carlo trajectory calculations.

#### Kolbenstvedt model for the electron-impact ionization:

- Haque and coworkers (2010) explored a so-called Kolbenstvedt model for the electron-impact ionization of the  $K$ -,  $L$ - and  $M$ -shells of neutral atoms and ions.

## 11.3. Further estimates on atomic properties, not yet considered in JAC

### 11.3.a. Electron and positron stopping powers (StoppingPower)

#### Property & notations:

- **Stopping power:** due to the penetration of electrons and positrons through matter:  $A^{(*)} \dots$

#### Motivation:

- **Stopping powers of matter for electrons are important** for many applications involving energy deposition. In radiation physics, chemistry, biology and medicine, it is important to have simple but accurate estimates about the stopping power of energetic electrons in various media for (Gümüs, 2005).
- Until the present, no simple and practical model exists for the stopping power of electrons with energies below 10 keV.

#### Approximate stopping power formulas:

- The stopping power of materials arise from two types of processes, namely collisional and radiative processes. The collisions of the incident particles with the target material is most important and mainly arises from collisions between the incident particles and the atomic electrons.

$$\frac{dE}{dx} = \left( \frac{dE}{dx} \right)^{(\text{collision})} + \left( \frac{dE}{dx} \right)^{(\text{radiative})}$$

- For incoming electrons, a modified collisional stopping power formula for incoming electrons can be written as given by Sugiyama (1985) or Rohrlich and Carlson (1954)
- **Full Bethe-Bloch formula:**

$$-\frac{dE}{dx} = \left( \frac{e^2}{4\pi\epsilon_o} \right) \frac{4\pi z^2 N_A Z \rho}{mc^2 \beta^2 A} \left[ \ln \left( \frac{2mc^2 \beta^2}{I} \right) - \ln(1 - \beta^2) - \beta^2 \right]$$

### 11.3.b. Stark broadening of spectral lines in plasma

#### Motivation:

- **Stark broadening of spectral lines:** This broadening has been found very important in (so-called) DA and DB white dwarf atmospheres.
- Apart from white dwarfs, Stark broadening of spectra is also an important pressure-broadening mechanism for various types of other stars.
- Various semi-empirical calculations of the Stark widths and shifts have been performed in the impact approximation by using Griems (1968) formula.
- For selected lines, it was shown that Stark broadening can change the observed widths by 10-45 %.

#### Semi-empirical approximations:

- **Lorentzian profile of isolated lines:** For isolated lines  $|\alpha_i \mathbb{J}_i\rangle \rightarrow |\alpha_f \mathbb{J}_f\rangle$  with level energies  $E_i$  and  $E_f$ , the line profile is often assumed to be Lorentzian, shifted by the energy  $d$  and with the (total) line widths  $\Gamma$  due to the (plasma) Stark broadening.

$$F(\omega) = \frac{(\Gamma/2\pi)}{(\omega - \omega_{if} - d)^2 + (\Gamma/2)^2}, \quad \omega_{if} = \frac{E_i - E_f}{\hbar}.$$

where

- **Energy shift and (total) widths of the Stark-broadened line:** Both, the energy shift  $d$  as well as the (total) widths  $\Gamma$  can be expressed in terms of the velocity distribution  $f(v)$  of the electrons in the plasma, its density  $N^{(\text{elec})}$ , the impact parameter  $b$  of the incident electrons as well as the elastic and inelastic (electron-impact excitation) cross sections for **excitations of the initial and final levels to neighbored levels**.

$$\Gamma = N^{(\text{elec})} \int dv v f(v) (\sigma^{(\text{elastic})} + \sigma^{(\text{inelastic})}); \quad d = N^{(\text{elec})} \int dv v f(v) \int_{R_1}^{R_{\text{Debye}}} db 2\pi b \sin(\phi_p)$$

and where  $\phi_p$  is a **plasma-specific phase shift**. Further details on the Stark broadening of spectral lines are given by Dimitrijevic and Sahel-Brechot (1996).

- In 1968, Griem suggested simple semiempirical formulas for the Stark line shift and widths in impact approximation which is based on some original formula by Baranger (1958) as well as an effective Gaunt factor by Seaton (1962) and Regemortel (1962). With these formulas,

the Stark linewidths and Stark line shifts can be obtained by:

$$\frac{\Gamma}{[a.u.]} = 8 \left(\frac{\pi}{3}\right)^{3/2} \frac{\hbar}{m a_o} N_e \left(\frac{E_H}{k T}\right)^{1/2} \left[ \sum_{i'} |\langle i' | r | i \rangle|^2 g_{se} \left(\frac{E}{\Delta E_{i'i}}\right) + \sum_{f'} |\langle f' | r | f \rangle|^2 g_{se} \left(\frac{E}{\Delta E_{f'f}}\right) \right]$$

$$\frac{d}{[a.u.]} = -8 \left(\frac{\pi}{3}\right)^{3/2} \frac{\hbar}{m a_o} N_e \left(\frac{E_H}{k T}\right)^{1/2} \left[ \sum_{i'} \left(\frac{\Delta E_{i'i}}{|\Delta E_{i'i}|}\right) |\langle i' | r | i \rangle|^2 g_{sh} \left(\frac{E}{\Delta E_{i'i}}\right) + \sum_{f'} \left(\frac{\Delta E_{f'f}}{|\Delta E_{f'f}|}\right) |\langle f' | r | f \rangle|^2 g_{sh} \left(\frac{E}{\Delta E_{f'f}}\right) \right]$$

In these formulae,  $E_H$  is the hydrogen ionization energy,  $N_e$  is the free-electron density of the perturber levels,  $T$  is the electron temperature,  $E = 3/2 kT$  the mean energy of the perturbing electron and  $g_{se}$ ,  $g_{sh}$  are the effective Gaunt factors which are calculated and tabulated by Griem. These Gaunt factors are slowly varying functions of  $x_{j'j}$  where  $x_{j'j} = E/\Delta E_{j'j}$  is the energy difference between a perturbing level  $j'$  and the perturbed initial/upper (final/lower) level  $j$ .

- The atomic matrix elements need often to be obtained from *ab-initio* Hartree-Dirac-Fock computations, and the summation over the levels is obtained by making use of the table by Moore (1958).

### 11.3.c. Atomic electron-momentum densities

#### Definition & notations:

- For a given atomic level  $|\alpha\mathbb{J}\rangle$ , the (radial electron) momentum density  $I(p; \alpha\mathbb{J})$  and the one-electron momentum density  $\Pi_{n\kappa}(p; \alpha\mathbb{J})$  and is given by (Koga and Thakkar, 1996)

$$I(p; \alpha\mathbb{J}) = 4\pi p^2 \Pi(p; \alpha\mathbb{J}), \quad \Pi(p; \alpha\mathbb{J}) = \frac{1}{4\pi p^2} \sum_{i=1}^N \langle \psi_\alpha | \delta(p - p_i | \psi_\alpha \rangle.$$

- The radial one-electron momentum density  $I$  is used to calculate the corresponding **moments of the momentum density**

$$\langle p^k \rangle = \int_0^\infty dp p^k I(p; \alpha\mathbb{J}), \quad -2 \leq k \leq 4$$

- These moments are often applied in:

## 11. Semiempirical estimates

- density functional theory: since  $\langle p \rangle / \pi$  is close to the Dirac-Slater exchange energy;
- x-ray cristallography: since  $\langle p^3 \rangle$  is roughly proportional to the initial value of the so-called Patterson function
- Compton profiles: since  $J(q) = \frac{1}{2} \int_{|q|}^{\infty} dp p^{-1} I(p)$  is twice the peak height of the isotropic Compton profile.



# 12. Beams of light and particles

## 12.1. Helmholtz wave equation

### Propagation of light in free space:

- **Helmholtz wave equation:** If we consider the propagation of light in free space, the Maxwell equations simplify and the spatial structure of their solution satisfy the well-known **Helmholtz wave equation**, an elliptic partial differential equation

$$\nabla^2 u(\mathbf{r}) + k^2 u(\mathbf{r}) = 0.$$

This equation follows from the (full) time-dependent wave equations due to the **separation of variables**.

- **Harmonic solution of the Helmholtz equation:**

$$u(\mathbf{r}) = C_1 e^{i\mathbf{k}\cdot\mathbf{r}} + C_2 e^{-i\mathbf{k}\cdot\mathbf{r}}, \quad k = |\mathbf{k}| = \frac{\omega}{c}.$$

- The solution in time is a linear combination of sine and cosine functions with angular frequency  $\omega$ , while the form of the spatial solution of the Helmholtz equation depends rather sensitively on the given boundary conditions.
- **Vector potential of optical beams:** There exist various expressions to represent the vector potential of optical beams of different kind and spatial structure, and which are solutions of the Helmholtz equation. These vector potentials represent different spatial structures and phases.
- Below, several of these solution are provided for different beams in terms of the (scalar) amplitudes of the corresponding light fields.
- **Synonym notions of vortex beams:** helically phased light beams; optical vortex beams with a field dependence  $e^{im\varphi}$ , ... .

### Paraxial approximation to the Helmholtz equation:

- **Paraxial approximation:** The Helmholtz equation still represents a complicated partial differential equation; in the science of optics, one often has well-defined paraboloidal waves or Gaussian beams in which the field change along the propagation direction is small compared to the change in perpendicular direction;  $\partial u / \partial z \ll \partial u / \partial x, \partial u / \partial y$ .
- For these conditions, the Helmholtz equation simplifies to

$$\nabla_{\perp}^2 u(\mathbf{r}) + 2ik \frac{\partial u}{\partial z}(\mathbf{r}) = 0.$$

- **Conditions for the paraxial approximation:** The restrictions upon the variation of the amplitude function  $u$  is often also written in the form

$$\left| \frac{\partial u}{\partial z} \right| \ll |k u| \quad \text{and} \quad \left| \frac{\partial^2 u}{\partial z^2} \right| \ll |k^2 u|.$$

- The paraxial approximation is obtained if the ansatz  $\psi(\mathbf{r}) = u(\mathbf{r}) e^{ikz}$  is used in the Helmholtz equation and if the term  $\sim \frac{\partial^2 u}{\partial z^2}$  is neglected, when compared with  $\sim \frac{\partial \psi}{\partial z}$ .

## 12.2. Light beams

### Light beams:

- **Characterization and properties of beams:** In studying light-matter interactions, light beams play a major practical role and, hence, their characterization and properties are very important in order to understand the details of the interaction.
- **Beams from applications:** Light beams are usually formed when the em field from a given source is projected by filters and blends into a beam. In some lighting devices, lamps or parabolic reflector are used to produce an artificially light beam with a more or less large divergence (for example, car headlights, spotlights, ect.).
- **Experimental beam parameters:** The geometry and behavior of a Gaussian beam are governed by a set of beam parameters which are defined slightly different for different kinds of beams but which are often used to characterize the beams. These beam parameters include:

- **Beam width:** There are over five definitions of beam width.
- **Beam quality:** Quantified by the beam quality parameter, M2.
- **Beam divergence:** This is a measure of how the beam spreads with distance with regard to the focus plane.
- **Beam profile:** A beam profile is the 2D intensity plot of a beam at a given location along the beam path. A Gaussian or flat-top profile is often desired for many applications. The beam profile indicates nuisance due to high-order spatial modes in a laser cavity as well as hot spots in the beam.
- **Beam astigmatism:** A beam is called **astigmatic** when the vertical and horizontal focus of the beam is placed at different locations along the beam path.
- **Beam wander or jitter:** The amount by which the centroid or peak value of the beam profile moves with time.

### Twisted (vortex) light beams:

- Charles Darwin (1932; 1887-1962) was one of the first who worked on **light with a spatial structure of the phase**.
- For circularly-polarized plane-wave light, the photons carry a projection of the angular momentum  $\Delta L_z = \pm\hbar$  upon the propagation direction.
- **Polarization state of a light beam:** In 1892, Poincare showed that the state of polarization of a light beam can be described as a point on the surface of a unit sphere now known as the Poincare sphere.

#### 12.2.a. Gaussian beams

### Characerization and use:

- **Gaussian beam:** In optics, a Gaussian beam is a beam of electromagnetic radiation whose **transverse electric field amplitude and, hence, its intensity distributions are approximately described by some Gaussian distribution**.
- Both, planes waves and Gaussian beams are solutions of the Maxwell equation *and* the homogenous wave equation they both generally possess a **polarization** but can be readily described in terms of complex amplitudes.

## 12. Beams of light and particles

- **Complex-valued electric field amplitude  $u(\mathbf{r})$ :** In a Gaussian beam, the electric field amplitude

$$u(\rho, z) = u_0 \frac{w_0}{w(z)} \exp\left(\frac{-\rho^2}{w^2(z)}\right) \exp\left(-ikz - ik\frac{\rho^2}{2R(z)} + i\zeta(z)\right)$$

obeys the paraxial Helmholtz equation. This amplitude is sufficient in order to describe the properties of the beam since the electric field and magnetic field propagate together and obey the same wave equation.

- **Characteristic parameters:** In the definition of the electric-field amplitude above, we have

- $\rho$  ... radial distance from the center axis of the beam;
  - $z$  ... the axial distance from the beam's narrowest point (the **beam waist**),
  - $k = \frac{2\pi}{\lambda}$  ... the wave number (in radians per meter),  $E_0 = |E(0, 0)|$
  - $w_0 = w(0)$  ... the waist size,
  - $w(z)$  ... the radius at which the field amplitude drop to  $1/e$  and, hence, the intensity to  $1/e^2$  at axis with regard to  $w_0 = w(0)$ ,
  - $R(z)$  ... the radius of curvature of the beam's wavefronts,
  - $\zeta(z)$  ... the Gouy phase shift, an extra contribution to the phase that is seen in Gaussian beams.
- In fact, many laser emit beams with roughly a Gaussian profile; these beams are often called the TEM<sub>00</sub> fundamental modes of the laser's optical resonator. **Gaussian beams form a widespread model in laser physics.**

## 12.2.b. Vortex beams. Characterization and properties

### General remarks:

- The recent years have seen a tremendous effort in generating and manipulating vortex beams together with a good number of new or promising applications, including the capture, manipulation and transport of nanoscopic particles (optical tweezer), high-resolution microscopy or for data transmission.

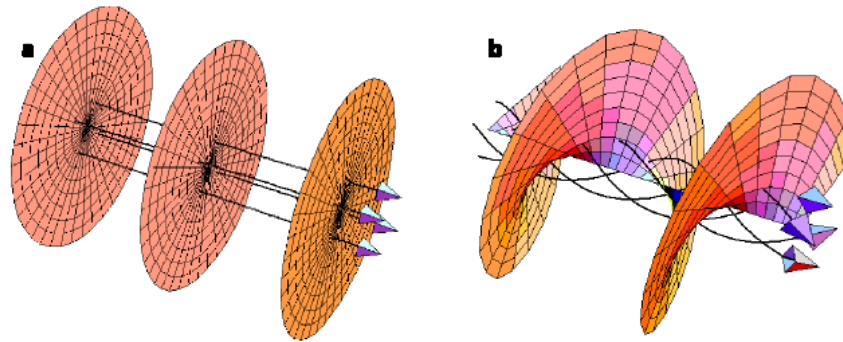


Figure 12.1.: Laser beams usually have planar wavefronts with wavevectors parallel to the beam axis. Beams with helical wavefronts have wavevectors which spiral around the beam axis and give rise to an orbital angular momentum; from Padgett and Allen (2000).

- Twisted photons carry both, spin angular momentum (SAM) as well as orbital angular momentum (OAM) along their propagation direction.
- **Topological charge** The  $z$ -projection of the OAM,  $m$ , is often called the topological charge or the **winding number of the beam**. This charge describes the number of the  $2\pi$  windings of the helical phase around the vortex in one wavelength.
- **Synonym notions of vortex beams:** helically phased light beams; optical vortices with a field dependence  $e^{im\varphi}$ .
- Optical vortex beams are well-known and have been routinely used for the last 20 years; a first demonstration dates back to Durnin *et al.* (1987).
- For optical vortices, charge values of  $m \leq 5000$  have been realized (Shen *et al.*, 2013).
- Reviews on optical vortex beams are given by Allen *et al.* (1999); Franke-Arnold *et al.* (2008), Zhan (2009) and Yao and Padgett (2011).

### Characterization of vortex beams:

- A vortex state of light carries a well-defined projection of the orbital angular momentum (OAM) with respect to the beam axis.
- **Wave functions:**  $\sim e^{im\varphi}$  are eigenfunctions of  $\ell_z = \frac{\partial}{\partial \varphi}$ .
- **Photon field:** For a quantized radiation field, each particle also carries a projection  $m\hbar$  of the OAM, that is part of the total angular momentum. In general, the spin and orbital angular momentum cannot be separated.
- In the paraxial approximation, the spin and OAM longitudinal components of the vortex state can be separated from each other. Beyond the paraxial limit, however, the spin-orbit interaction does not allow a simple separation of the spin and orbital angular momentum.

### Properties of vortex beams:

- For all points in the beam, the ratio between the azimuthal and  $z$ -components of the momentum is found to be  $\ell/k r$ .
- The linear momentum of each photon is given by  $\hbar k$ ; since the azimuthal component of the wavevector is  $m/r$  and is independent of the wavelength, the orbital angular momentum per photon is

$$p = \hbar k \quad \longrightarrow \quad p_\varphi = \hbar k \phi \quad \longrightarrow \quad \ell = r \times p_\varphi = r \times \frac{\hbar k}{r} = m\hbar$$

- However, OAM  $\neq$  SAM, despite the rather similar illustrations in the literature; in particular, there exist also vortex states for spin-less waves, such as acoustic vortex waves.

### 12.2.c. Vortex beams. Generation

#### Generation of vortex beams with different photon energies:

- For intense, short-wavelength FEL radiation, optical methods can often not be applied due to strong limitations and difficulties in the fabrications of optical surfaces (Terhalle *et al.* 2011, Peele *et al.* 2002).
- **Twisted EUV beams:** Ribic and coworkers (2014) proposed an efficient scheme in order to generate intense coherent and twisted OAM beams in the EUV region. These beams can be realized by means of a seeded FEL, i.e. by using a (seeding) laser pulse with a transverse, staircaselike phase. For this phase pattern, in particular, the diffraction and mode selection drive the radiation profile towards a dominant OAM mode at saturation.
- **Twisted x-ray beams:** Hemsing and coworkers proposed two clever approaches for generating FEL vortex beams within the x-ray region:
  - (i) By using the interaction of an electron beam with a seeded laser within a helical undulator (Hemsing *et al.*, 2011);
  - (ii) by using a so-called echo-enabled harmonic generation (EEHG) scheme, in which two seed lasers and two magnetic chicanes are utilized in order to produce harmonic microbunches of an electron beam with a corkscrew distribution.

#### Generation of vortex beams by pitch-fork holograms:

- The superposition of a regular phase profile with some proper grating already results in a fork hologram that help generate vortex beams.
- Fork-like or pitch-fork holograms are known to diffract the light and to generate beams with OAM. Such holograms can be generated also dynamically by using a spatial light modulator and readily controlled by a computer.
- A pitchfork hologram is obtained adding a twisting and an oblique phase. The width of six  $2\pi$  phase ramps is highlighted.

#### Generation of vortex beams by spiral phase plates:

- A (plane-wave) light beam can acquire a non-zero OAM by crossing a spiral phase plate with an inhomogeneous thickness.
- **Spiral phase plates (SPP):** Such a SPP is a transparent dielectric plate with a thickness that varies as a smooth ramp; this variation in the thickness adds a phase shift to an incident field that increases linear with the azimuthal angle  $\varphi$ .
- **Light field:** If a SPP is placed into the waist of a Laguerre-Gaussian beam with well-defined quantum numbers,  $u_{\ell p}^{\text{LG}}$ , the complex light amplitude just behind the plate can be described by means of a phase-plate operator  $S(m_s, \alpha)$

$$\langle r, \varphi | S(m_s, \alpha = 0) | l, p \rangle = u_{lp}^{\text{LG}}(r, \varphi) e^{im_s \varphi},$$

## 12. Beams of light and particles

and where  $m_s$  is the phase shift per unit angle of the given SPP.

- SPP help explore high-dimensional entanglement with just two detector, while 6 detectors are required by other popular methods in order to prove the entanglement of the OAM degree of freedom for two entangled photons (Oemrawsingh *et al.*, 2004).
- The figure below shows the transformation of a  $\text{TEM}_{00}$  mode into a helical beam.

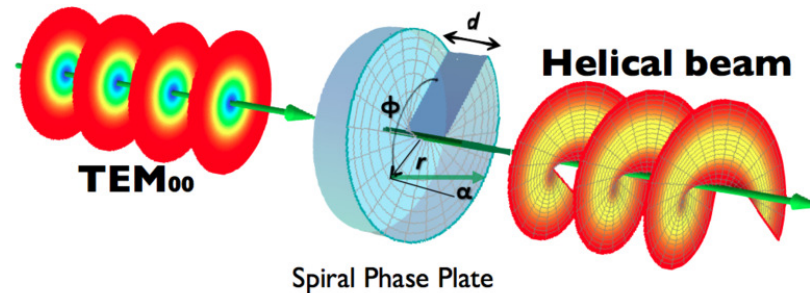


Figure 12.2.: Experimental scheme for generating light orbital angular momentum with spiral phase plates. Taken from Beijersbergen *et al.* (1994).

### Generation of vortex beams by q-plates:

- **q-plate:** A (so-called) q-plate is a device that is realized by means of liquid crystals, polymers or sub-wavelength gratings; this device exploits a change of sign if the polarization of the incoming light is modified.

### Generation of vortex beams by spatial light modulators:

- xxxxx A spatial light modulator (SLM) is an object that imposes some form of spatially varying modulation on a beam of light. A simple example is an overhead projector transparency. Usually when the phrase SLM is used, it means that the transparency can be controlled by a computer. In the 1980s, large SLMs were placed on overhead projectors to project computer monitor contents to the screen. Since then more modern projectors have been developed where the SLM is built inside the projector. These are commonly used in meetings of all kinds for presentations.



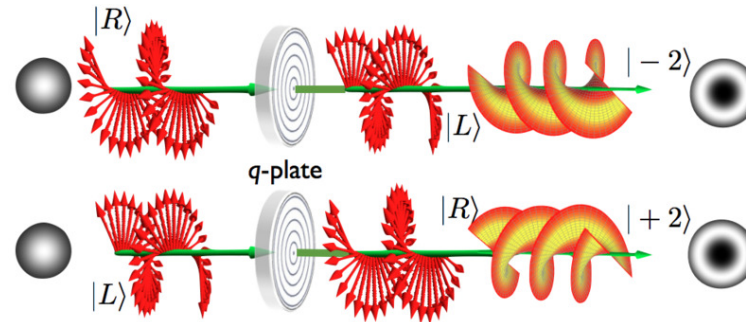


Figure 12.3.: The q-plate effect for left and right-hand circular polarizations.

### Generation of vortex beams by cylindrical mode converters:

- A Hermite-Gaussian beam can be converted into a Laguerre-Gaussian beam with well-defined OAM by using an **astigmatic** system with two well-aligned cylindrical lenses. In this scheme, these lenses are placed at some specific distance in order to introduce a well-defined relative phase between the horizontal and the vertical Hermite-Gaussian beams.

### Generation of vortex beams by helical undulators:

- X-ray vortex beams with photon energies  $\hbar\omega \approx 100$  eV have been generated in a helical undulator by Bahrtdt *et al.* (2013) and Hemsing *et al.* (2013).
- Another interesting approach has been suggested by Ribic and coworkers (2014). These authors proposed to modify the phase mask and to manipulate directly the seeding laser in order to change the transverse properties of the FEL light.

## 12.2.d. Hermite-Gaussian beams

### Characterization and use:

- **Hermite-Gaussian modes**  $HG_{mn}$ : have a rectangular symmetry and are described by means of two mode indices  $m$  and  $n$ ; these mode indices give the number of nodes in the  $x$  and  $y$  directions respectively.

## 12. Beams of light and particles

- **Hermite-Gaussian modes:** These modes are a convenient description for the output of lasers whose cavity design is not radially symmetric but asymmetric in horizontal and vertical directions.

### 12.2.e. Laguerre-Gauss beams

#### Characterization and use:

- The Laguerre-Gaussian (LG) beams carry an intrinsic OAM like the intrinsic SAM for beams with a circular polarization (Allen *et al.* 1992, 2003). This appeared originally rather surprising as the OAM of LG beams arise from the helical phase distribution about the beam axis, and which coincides also with a phase singularity (Nye and Berry, 1974).
- **Laguerre-Gaussian modes  $u_{\ell p}^{\text{LG}}$ :** For a Laguerre-Gaussian beam with well-defined magnetic and radial quantum numbers  $(m, p)$ , the complex amplitude of the light field in the waist of the beam is defined as in the polar representation

$$u_{\ell p}^{\text{LG}}(r, \varphi) = \langle r, \varphi | m, p \rangle = R_{mp}(r) \Phi_m(\varphi) \quad \text{with} \quad \Phi_m(\varphi) = \frac{e^{im\varphi}}{\sqrt{2\pi}}$$

- **Laguerre-Gaussian modes  $u_{\ell p}^{\text{LG}}$ :** If the problem is cylindrically symmetric, these modes are the natural solution of the paraxial wave equation. They are often written in cylindrical coordinates by using Laguerre polynomials.

### 12.2.f. Bessel beams

#### Characterization and use:

- **Bessel beams:** A Bessel beam refers generally to an electromagnetic field whose amplitude is given by a Bessel function of the first kind. 5
- Although these beams are an idealization, true **Bessel beam are monochromatic and non-diffractive and, hence, do not not diffract and spread out** when they propagate. This behaviour is quite in contrast to other typical optical (or sound) waves if they are focused at some spot.
- Bessel beams are known to be **self-healing if they are obstructed at some point, i.e. they reform to a Bessel beam again further down the beam axis.**

- Pure Bessel modes generally carry a well-defined OAM; in particular, they typically include three components of the topological charge since only the (projection of the) total angular momentum is conserved for these beams. For this reason, the  $z$ -projections of the orbital and spin angular momentum are not defined independently.
- **Bessel modes  $|\kappa, m\rangle$ :** Bessel beams can be written also as a superposition over plane waves with well-defined (fixed) transverse momentum, longitudinal momentum as well as circular polarization

$$\psi(\mathbf{r}) = e^{-i\omega t + i k_z z} |\kappa, m\rangle, \quad |\kappa, m\rangle \propto e^{im\varphi} J_m(\kappa \rho)$$

- **Vector potential:** The vector potential of the Bessel beams is often written as Fourierintegral over Bessel states of different projections of the orbital angular momentum  $a_{\kappa m}(\mathbf{k}_\perp)$ ,

$$\mathbf{A} = e^{-i\omega t} \frac{1}{(2\pi)^2} \int d\mathbf{k}_\perp e^{-i\mathbf{k}\cdot\mathbf{r}} a_{\kappa m}(\mathbf{k}_\perp) \mathbf{e}_{\mathbf{k},\lambda}, \quad a_{\kappa m}(\mathbf{k}_\perp) = \sqrt{\frac{2\pi}{\kappa}} (-i)^m e^{-im\varphi_{\mathbf{k}}} \delta(k_\perp - \kappa)$$

$$\mathbf{e}_{\mathbf{k},\lambda} = \frac{1}{\sqrt{2}} \begin{pmatrix} \cos \vartheta_k \cos \varphi_k - i\lambda \sin \varphi_k \\ \cos \vartheta_k \sin \varphi_k + i\lambda \cos \varphi_k \\ -\sin \vartheta_k \end{pmatrix} \quad \tan \vartheta_k = \frac{\kappa}{k_z}$$

and where  $\mathbf{e}_{\mathbf{k},\lambda}$  are known as polarization (unit) vector and  $\vartheta_k$  as the opening angle of the Bessel beam.

- **Plane-wave expansion:** Of course, every monochromatic plane-waves can be expanded also in terms of Bessel waves as they are both just two different bases in order to describe wave packets

$$e^{i(\mathbf{k}\cdot\mathbf{r} - \omega t)} = \sum_{m=-\infty}^{\infty} e^{-i\omega t + i k_z z} i^m e^{-im\varphi_{\mathbf{k}}} |\kappa, m\rangle$$

- **Electron-photon interaction operator for Bessel beams:** For (a Bessel beam of) twisted photons, the transition operator reads as:

$$V^{(\text{tw})} = \alpha \mathbf{A}_{\varkappa m_\gamma k_z \lambda}(\mathbf{r}) \mathbf{p}.$$

$$\mathbf{A}_{\varkappa m_\gamma k_z \lambda}(\mathbf{r}) = \int \mathbf{e}_{\mathbf{k}\lambda} e^{i\mathbf{k}\cdot\mathbf{r}} a_{\varkappa m_\gamma}(\mathbf{k}_\perp) e^{-i\mathbf{k}_\perp \cdot \mathbf{b}} \frac{d^2 k_\perp}{(2\pi)^2}, \quad a_{\varkappa m_\gamma}(\mathbf{k}_\perp) = (-i)^{m_\gamma} e^{im_\gamma \varphi_{\mathbf{k}}} \sqrt{\frac{2\pi}{k_\perp}} \delta(k_\perp - \varkappa).$$

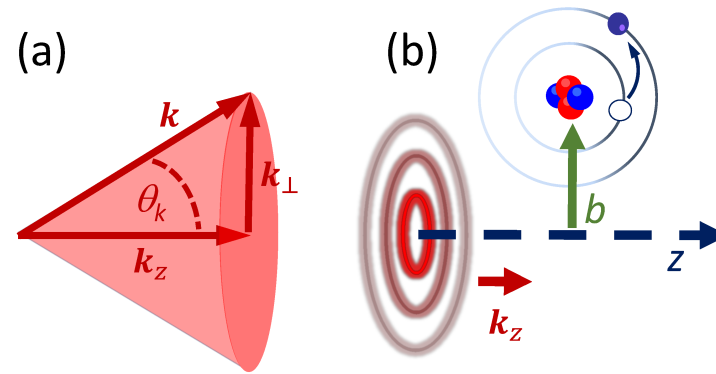


Figure 12.4.: (a) In momentum representation, the twisted light can be seen as a coherent superposition of plane waves. These plane waves with wavevectors  $\mathbf{k}$  are all lying on a cone with (polar) opening angle  $\vartheta_k = \arctan(|\mathbf{k}_\perp|/k_z)$  and with polarization vectors  $\mathbf{e}_{\mathbf{k},\lambda}$  which are perpendicular to  $\mathbf{k}$ . (b) For a given atom, the absorption amplitude depends on the position of the atom in a plane perpendicular to the propagation direction. The position of the target atom with regard to the beam axis is specified by the impact parameter (vector)  $\mathbf{b}$  (c-05-photobeam-interaction.eps).

In this expression, the factor  $e^{-i\mathbf{k}_\perp \cdot \mathbf{b}}$  specifies the position of a target atom within the incident wave-front, and which refers to the complex spatial structure of the Bessel beam.

## 12.2.g. Airy beams

### Characterization and use:

- **Airy beams:** Similar to Bessel beams, the cross section of an ideal Airy beam has an area of principal intensity and with a series of adjacent, less luminous areas up to infinity. Airy beams can be derived from the Airy integral that was first developed by George Biddell Airy in the 180s in order to explain the optical caustics in a rainbow and elsewhere.
- An Airy beam preserves its shape of intensity during propagation and forms a parabolic curve, quite analogue to the trajectory of a free projectile as seen in a plane perpendicular to the propagation.
- Airy beam are freely accelerating, i.e. they bend to form a parabolic arc as they propagate (Berry and Balazs, 1979; Siviloglou *et al.*, 2007), although the linear momentum remains conserved for these beams since the intensity centroid of an Airy beam forms a straight line.

- Airy beam do also not diffract, i.e. they do not spread out after focussing.

### 12.2.h. Necklace ring beams

#### Characterization and use:

- **Necklace ring beams:** These beams refer to optical vector solitons that consist of incoherently coupled self-trapped ‘necklace’ beams; these beams can carry zero, integer or even some fractional angular momentum.
- Necklace ring beams exhibit a self-stabilizing mutual attraction between different beam components and a quasistable propagation for much larger distances than the corresponding scalar vortex solitons

### 12.2.i. Light beams with non-integer OAM

#### Characterization and use:

- **Beams with non-integer OAM:** These beams refer to structurally propagation invariant light beams carrying non-integer orbital angular momentum (OAM) using Hermite-Laguerre-Gaussian (HLG) modes.

### 12.2.j. Vector beams

#### Characterization and use:

- **Vector beams:** The vector beams need to be described in the product space of the optical SAM and OAM subspaces and are characterized by higher-order Stokes parameters.
- Vector beams exhibit a spatially variant polarization.
- Cylindrical vector beams are axially symmetric solution to the full electromagnetic (vector) wave equation.
- **Generation of vector beams:** These beams can be generated via different active and passive methods.

### 12.2.k. Traktor beams

#### Characterization and use:

- **Tractor beam:** These beams refer to optical fields that help attract one object to another from a distance.
- The concept of tractor beams originates from the novel *Spacehounds of IPC* by E. E. Smith (1931) and by re-coining an earlier ‘attractor beam’, in contrast to other, ‘repulser beams’.
- Tractor beam are defined to use (a negative) nonconservative radiation pressure, that differs from a gradient force, and which is directed towards the source. This generally requires a continuous redirection of momentum flux (Sukhov and Dogariu, 2010, 2011; Chen *et al.* 2011).
- Traktor beams are based on the maximization of forward scattering of light due to interference of different radiation multipoles. It was shown especially that an simultaneous excitation of multipoles is necessary in the particle in order to realize a negative (pulling) optical force.

### 12.2.l. Polarization radiation

#### Characterization and use:

- In electrodynamics, one generally distinguishes **two classes of emitted radiation: bremsstrahlung and polarization radiation**. While bremsstrahlung is produced by accelerated charges, polarization radiation can be emitted by a uniformly moving charge, although only in the presence of a medium.
- In dependence on the given medium or target geometry, one distinguishes various forms of polarization radiation: Cherenkov radiation, transition radiation, diffraction radiation as well as Smith-Purcell radiation.
- **Polarization radiation:** Electromagnetic radiation can be emitted not only by accelerated charges but also due to the motion of magnetic moments and higher multipoles. Until the present, however, this (so-called) polarization radiation has never been verified experimentally for any frequency of light.
- The main experimental difficulty for approving this radiation refers to the (very) small current of a magnetic moment that is suppressed by  $\sim 10^{-5}$  w.r.t. a corresponding charge current.
- A pure quasi-classical treatment of the polarization radiation is inconsistent since quantum corrections are of the same order for this type of radiation phenomena.

## 12.2.m. Manipulation of optical beams

### Optical elements:

- **Dove prism:** This is a (type of) reflective prism that can be used to invert a given image. These prisms are often shaped from a truncated right-angle prism.
- If a beam enters the prism at one of the sloped faces and parallel to the longitudinal axis, it undergoes total internal reflection at the inside of the longest (bottom) face, and then leaves the prism at the opposite sloped face. Therefore, any image that passes through the prism is flipped (mirrored) and also also inverted but not laterally transposed.
- **Hologram:** A hologram usually refers to a three-dimensional image of an object that is created by holography, i.e. an optical technique for recording and reconstructing the amplitude and phase distributions of a coherent wave. Holography is often used to produce three-dimensional images or holograms.
- **Half-wave plate:** A wave plate, sometimes known as a retarder, designates an optical device that alters the polarization state of a light wave. While a half-wave plate shifts the polarization direction of linearly polarized light, a quarter-wave plate converts linearly-polarized light into circularly-polarized light and vice versa. Quarter-wave plate can be used also in order to produce elliptically-polarized light.

### Mach-Zehnder interferometers:

- These interferometers make use of two separate beam splitters in order to split and recombine an input beam but from which the beam can emerge and detected at two outputs. The optical path lengths in the two arms of the interferometer may be nearly identical or may be different in order to introduce an extra delay path.
- If a beam is sent through the interferometer, the optical powers at the two outputs depends on in practice sensitively on the the precise difference in optical arm lengths as well as on the wavelength of the light.
- The Mach-Zehnder interferometer was first developed by Ludwig Mach and Ludwig Zehnder.

### 12.2.n. Optical forces of vortex beams

#### Optical forces:

- :
- A positive radiation force is relatively intuitive and occurs in the backscattering or absorption of the forward-directed momentum of a beam; it was first reported already by Nichols and Hull (1903).

### 12.2.o. Application of optical (vortex) beams

#### Remarks:

- Different applications of optical vortex beams and the angular momentum of light have been envisaged and are currently explored in research laboratories, although no real commercial application are yet known.
- **Fields of applications:** Cold atoms, trapped nanoparticles, micro manipulation and detection of spinning objects, remote measurement of the rotation of mesoscopic particles, entanglement and quantum information, optical data transmission, interaction with matter, microscopy and astrophysics.
- **X-ray magnetic circular dichromism** By using different OAM states in the measurements, the dipole and quadrupole contributions to the circular dichroism signal can be separated from each other.

#### Optical tweezer:

- **Optical tweezer:** These tweezers usually refer to single-beam gradient force trap that are realized by some highly focused laser beam. These focused beams provide an attractive or repulsive force on nano- or mesoscopic dielectric objects.
- Optical tweezers have been utilized to manipulate the orientation of particles or particle aggregates.
- In optical tweezers, the forces are typically on the order of piconewtons; they depends on the **refractive index** of the object.



### Optical phasors:

- **Phasor:** A complex constant in **complex wave notation** that encapsulates the amplitude; sometimes known also as sinor or complexor.

### Applications in quantum information theory:

- The orbital-angular momentum (OAM) of light has recently emerged as a promising candidate for quantum and classical information systems.
- In OAM beams, the spatial degrees of freedom provide a high-dimensional alphabet to quantum information processing which enables one to use qunits instead of qubits.
- **Quantum information encoding,** Higher-dimensional quantum information encoding has been discussed in the literature for possible future applications in quantum cryptography or quantum computations. The discrete, unbounded state-space of OAM promises in particular vastly enhanced data rates as well as an increased tolerance to eavesdropping in quantum communication.

## 12.3. Electron beams

### Electron beams:

- Electron beams are quite frequently applied in physics; such beams may refer, for instance, to some stream of electrons as obtained from a betatron. Electron beams are generally generated either by heat (thermionic emission), bombardment of surfaces with charged atoms or particles (secondary electron emission), or by means of strong electric fields (field emission).
- Electrons may be collimated by holes and slits. Because they are electrically charged, they can also be deflected, focused or accelerated by electric and magnetic fields.

### Twisted electron beams:

- Apart from twisted (vortex) solutions of the Dirac equation, a quasis-relativistic wave equation for twisted relativistic electron in arbitrary electric and magnetic fields can be derived by using a Foldy-Wouthuysen transformation.

## 12. Beams of light and particles

- The interaction of twisted (vortex) electrons with electric and magnetic fields can be described quite readily, based on Lorentz transformations.

### 12.3.a. Gaussian electron beams

#### Characterization and use:

- It has been demonstrated experimentally that a Gaussian-profile model better describes the (equilibrium) electron beam than a beam with a uniform-profile model.

### 12.3.b. Vortex electron beams

#### General remarks:

- Electron vortex beams were first realized experimentally by Uchida and Tonomura (2010), Verbeeck, Tian and Schattschneider (2010) and McMorran *et al.* (2011). Uchida and Tonomura, in particular, generated an electron vortex beams with (electron) energy  $E = 300$  keV, a projection of the OAM with  $m \lesssim 100$  and by focusing the electron beam to a spot size of  $1.2$  Å.
- Vortex electrons can generally carry a rather large intrinsic projection of the orbital angular momentum (OAM)  $m \sim 100$  with respect to their propagation direction. Therefore, the magnetic moment of vortex electrons  $\mu \approx m \mu_B$  is similarly large, when compared with the Bohr magneton  $\mu_B$  due to the spin motion of the electron.
- The magnetic moment of twisted electrons due to their OAM comes in addition to the known magnetic moment due to the spin motion.
- Since Vortex electrons are sensitive to external magnetic fields, they can be utilized to probe atomic magnetism and, thus, open novel opportunities for electron microscopy.
- While vortex photon beams have been created and routinely used for several decades already, the vortex states of electrons and charged particles is a presently emerging field.

### 12.3.c. Generation of vortex electron beams

#### Generation of vortex electron beams with different energies:

- Indeed, several new ideas were proposed in the literature in order to create electron vortex beams experimentally.

#### Generation of vortex electron beams by undulators:

- **Electrons in a helical modulator:** cf. Figure 12.3.c

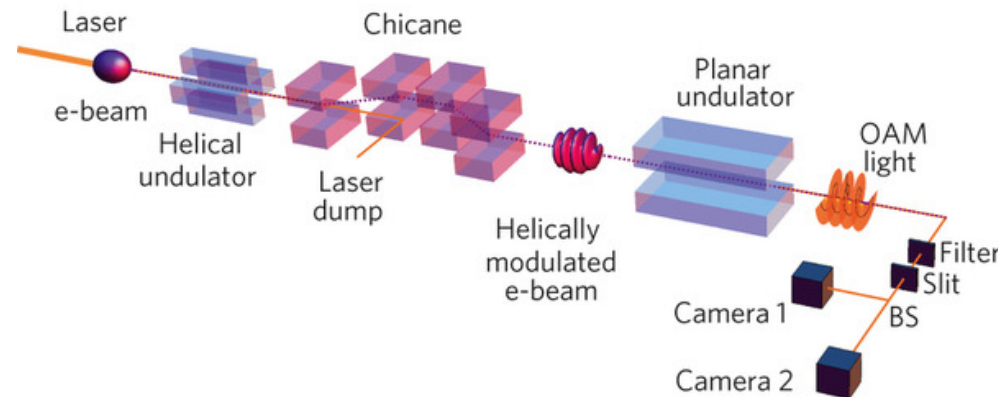


Figure 12.5.: An unmodulated relativistic electron beam interacts with a linearly polarized laser in a helical undulator, which gives the electrons an energy kick that depends on their position in the focused laser beam.

#### Generation of vortex electron beams by transmission electron microscopes:

- Voloch-Bloch *et al.* (2013) generated Airy electron beams by using a holographic technique within a transmission electron microscope. Both, self-bending and self-healing features were observed in the experiments if the beam was obstructed by a small object during its propagation.

### Generation of vortex electron beams by holograms:

- Grillo *et al.* (2014) generated an electron Bessel beam by the diffraction of electrons at a nanoscale phase hologram. This hologram imposed a conical phase structure upon the electron wave-packet. The so obtained beams propagated for 0.6 m without measurable spreading and could reconstruct also its intensity distributions, if these beams were partly obstructed by an obstacle.

#### 12.3.d. Bessel electron beams

##### Characterization and use:

- Quasi-Bessel beams can be obtained within a good approximation, while it is not possible to generate an ideal Bessel beam because they are not normalizable in their total intensity,
- Wave function of a Bessel electron: In cylindrical coordinates  $r, \varphi, z$ , the wave function of electrons in a Bessel beam is given by

$$\psi(r, \varphi, z; t) = J_m(\kappa r) e^{i m \varphi} e^{-i(\omega t - k_z z)},$$

where  $J_m(x)$  is the  $m$ -th order Bessel function of the first kind and  $\kappa$  and  $k_z$  denote the transverse and longitudinal components of the wave vector, respectively.

- Probability-density distribution of Bessel electron beams: This wave function gives rise to a probability-density distribution  $\mathcal{P} = \psi \psi^+$  that is independent of  $z$  and  $t$ :

$$\mathcal{P}(r, \varphi, z; t) = \mathcal{P}(r, \varphi, 0; 0) = [J_m(\kappa r)]^2.$$

- This shows that the probability-density distribution is stationary in the transverse plane and independent of where this plane is taken.
- The energy of the electron  $E = \hbar \omega$  defines the de Broglie wave length and the modulus of the wave vector by

$$k^2 = k_r^2 + k_z^2 = \frac{2m\omega}{\hbar} = \left( \frac{2\pi}{\lambda_{dB}} \right)^2.$$

- For Bessel beams with a nonzero  $m$  value, the electrons possess a non-uniform helical phase front with  $m$  dislocations in the phase, and where the handedness of the phase front is defined by the sign of  $m$ .
- **Current density of electron Bessel beam:** This current density circulates azimuthally in the transverse plane and introduces an (projection of the total) orbital angular momentum of  $m\hbar$  per electron.
- An Bessel electron beam with a nonzero  $m$  value possess exhibit a shape with multiple rings and with a null probability density at the origin.
- Bessel electron beams can be represented also as a coherent superposition of conical plane waves along a closed ring (circle).

### 12.3.e. Airy electron beams

#### Characterization and use:

- Airy electron beams can be generated by diffraction of electrons through a nanoscale hologram, which imprints a cubic phase modulation on wavefunction of the electrons in the transverse plane.

### 12.3.f. Application of twisted electron beams

#### Remarks:

- The scattering of twisted electrons can provide insights into the magnetic structure of the target materials.
- Beams with non-zero OAM helped explore the vacuum Faraday effect as well as Larmor and Gouy rotations.
- Bessel electron beams have been used in electron microscopy.



# 13. Symbolic evaluation of expressions from Racah's algebra

## 13.1. Racah's algebra in atomic and many-body physics

### 13.1.a. Advantages of using Racah's algebra

#### Goals and use of Racah's algebra:

- The theories of angular momentum and spherical tensor operators play a significantly role in atomic and nuclear physics, and at several places elsewhere. These two theories lead to algebraic expressions which are usually written in terms of generalized Clebsch-Gordan coefficients and/or Wigner  $3n-j$  symbols as well as the Wigner rotation matrices and spherical harmonics. Although the evaluation and simplification of such expressions is in principle a straightforward task, it can become extremely cumbersome, if more complex systems or physical scenarios are considered.
- Indeed, the study of open-shell atoms and nuclei quickly raises the question how the rotational symmetry of (closed) systems, i.e. the conservation of angular momentum, can efficiently be exploited in understanding many-particle systems. After the pioneering work by Wigner in the late thirties, Racah (1941, 1942, 1943) developed a powerful machinery, known as Racah algebra, to deal with such systems, and which is one of the fundamental concepts in the (quantitative) treatment of many-particle systems.
- In many-particle physics, an explicit re-coupling of the angular momenta is often required due to the use and sequence of different coupling schemes. In general, such recoupling transformations are written in terms of recoupling coefficients that need to be evaluated over and over again.
- **Spin-angular integration of many-particle matrix elements** In atomic and nuclear structure theory, the evaluation and spin-angular integration of many-particle matrix elements is typically based on standard quantities like the matrix elements of the unit tensor, the (reduced) coefficients of fractional parentage as well as a number of other reduced matrix elements for different products of creation and annihilation operators. These quantities arise rather frequently in both, configuration interaction approaches and in the derivation of perturbation expansions for many-particle systems by using symmetry-adapted configuration state functions.

### 13. Symbolic evaluation of expressions from Racah's algebra

- **Theory of angular momentum** In the treatment of quantum many-particle systems, the theory of angular momentum offers two crucial advantages: (i) The reduction of complicated many-electron matrix elements to a rather small number of *standard quantities* and (ii) an elegant and very powerful calculus which help simplify and evaluate sophisticated expressions.
- **Racah algebra techniques:** Owing to these advantages, the techniques from the theory of angular momentum, shortly known also as Racah algebra techniques (Racah 1941, 1942, 1943), have been utilized in a large number of applications and in quite different field of many-particle physics (Wigner 1959, Varshalovich *et al.* 1988).
- **Racah expressions:** With the symbolic evaluation of expressions from Racah's algebra, we wish to focus on the algebraic transformation of (Racah) expressions such as:

$$\begin{aligned} \text{Racahexpr} := & \sum_{j_1, j_2, l_1, \dots} (-1)^{2j_1-j_2+\dots} j_1^{3/2} [j_2] \dots \begin{pmatrix} \cdot & \cdot & j_1 \\ \cdot & \cdot & \cdot \end{pmatrix} \begin{pmatrix} j_1 & j_2 & \cdot \\ \cdot & \cdot & \cdot \end{pmatrix} \left\{ \begin{matrix} \cdot & j_3 & \cdot \\ j_1 & \cdot & \cdot \\ J & \cdot & j_2 \end{matrix} \right\} \dots \\ & \times \int d\Omega_1 Y_{\ell_1 m_1}(\Omega_1) Y_{\ell_2 m_2}(\Omega_2) \int d\beta d_{p_3 q_3}^{j_3}(\beta) d_{p_4 q_4}^{j_4}(\beta') \dots \end{aligned}$$

Here, the Clebsch-Gordan coefficients need not to shown explicitly in this expression since, apart from an additional phase, they are equivalent to the Wigner 3- $j$  symbols. Up to the present, we (must) also drop the rotation matrices and spherical harmonics which are not (yet) supported in JAC. We introduce this notation of **Racahexpr** with the intention for providing an (internal) data type which facilitates automatic manipulations and which is flexible enough to support a wide range of applications.

- **Racah expressions:** Such expressions may generally include any number of Wigner  $n$ - $j$  symbols of different kind as well as (various integrals over) the spherical harmonics and Kronecker and triangular deltas. Of course, the complexity of such Racah expressions increases rapidly as more Wigner symbols are involved in the product terms.
- The symbolic evaluation of typical expressions from Racah's algebra is naturally based on the knowledge of a large set of sum rules that may include rules with a (multiple) summations over dummy indices (Varshalovich *et al.* 1988). For complex and lengthy Racah expressions, moreover, the algebraic simplification can often be considerably accelerated if the graphical rules due to Yutsis *et al.* (1962) are taken into account.
- **Sum and orthogonality rules:** To obtain a simplification for complex Racah expressions, a large variety of sum and orthogonality rules have been implemented from the monograph by Varshalovich *et al.* (1988).
- In the past decades, various techniques in simplifying expressions from Racah's algebra have been developed, based either on graphical methods or on the explicit knowledge of special values and sum rules, as they can be found in some standard form in the literature. The direct application of these rules is however often laborious due to a large number of symmetric forms of the Wigner and related symbols.



- **Graphical loop rules:** These loop rules are typically used in order to find out about and to simplify those parts in a recoupling coefficient (or generally in any Racah expression) that belong together. The implementation of graphical rules even allows to easily simplify recoupling coefficients which include several ten angular momenta to an (completely equivalent) sum of products of Wigner 6- $j$  and/or 9- $j$  symbols, multiplied by proper weights.
- **Algebraic simplifications:** The simplification of expressions from Racah's algebra critically depends how easily **all equivalent symmetric forms of these expression are recognized internally**. Obviously, the symmetry of a Racah expression as a whole is closely related to the symmetries of all the Wigner  $3n-j$  symbols which are involved in the expression.
- **Classical symmetries of the Wigner symbols:** Apart from the classical symmetries of the Wigner symbols, there is an extended range of symmetries due to Regge (1958); these symmetries are however of minor practical importance.
- Although the mathematical background of angular momentum theory is today quite well understood, the treatment and simplification of typical expansions as they naturally arise by using Racah's algebra, is often very laborious.
- In JAC, we shall facilitate the symbolic evaluation of expressions from Racah's algebra and, in particular, for (complex) expressions for which the known algebraic and graphical methods start to become tedious and prone to making errors.

### 13.1.b. Frequently applied symbols and functions from Racah's algebra

#### Table of symbols and functions from the theory of angular momentum:

- The following symbols and functions from the **theories of angular momentum and irreducible tensor operators** occur very frequently in the computation and analysis of atomic and many-body systems

13. Symbolic evaluation of expressions from Racah's algebra

Symbol	Designation	Calls in JAC
$\begin{pmatrix} a & b & c \\ m_a & m_b & m_c \end{pmatrix}$	Wigner 3- $j$ symbol	<code>Wigner_3j()</code>
$\begin{Bmatrix} a & b & c \\ d & e & f \end{Bmatrix}$	Wigner 6- $j$ symbol	<code>Wigner_6j()</code>
$\begin{Bmatrix} a & b & c \\ d & e & f \\ g & h & i \end{Bmatrix}$	Wigner 9- $j$ symbol	<code>Wigner_9j()</code>
$\left\{ \begin{array}{cccc} a & b & c & d \\ e & f & g & h \\ i & j & k & l \end{array} \middle  s \right\}$	Wigner 12- $j$ symbol of kind $s = 1, 2$	Not implemented.
$\begin{Bmatrix} - & a_2 & a_3 & a_4 \\ b_1 & - & b_3 & b_4 \\ c_1 & c_2 & - & c_4 \\ d_1 & d_2 & d_3 & - \end{Bmatrix}$	Sharp's symbol (Sharp 1955)	Not implemented.
$\langle a \ m_a, \ b \ m_b \mid c \ m_c \rangle$	Clebsch–Gordan coefficient	<code>ClebschGordan()</code>
$W(abcd; ef)$	Racah's $W$ coefficient	Not implemented.

Symbol	Designation	Calls in JAC
$d_{mm'}^j(\beta)$	Wigner $d_{mm'}^j(\beta)$ rotation matrix	<code>Wigner_d()</code>
$D_{mm'}^j(\alpha, \beta, \gamma)$	Wigner's $D$ -function	Not implemented.
$U_{mm'}^j(\omega; \Theta, \Phi)$	Rotation matrix $U(\omega)$	Not implemented.
$Y_{\ell m}(\vartheta, \varphi)$	Spherical harmonic	Not implemented.
$\{\mathbf{Y}_{l_1}(\vartheta_1, \varphi_1) \otimes \mathbf{Y}_{l_2}(\vartheta_2, \varphi_2)\}_{LM}$	Bipolar spherical harmonic	Not implemented.
$\{\mathbf{Y}_{l_1}(\vartheta_1, \varphi_1) \otimes \{\mathbf{Y}_{l_2}(\vartheta_2, \varphi_2) \otimes \mathbf{Y}_{l_3}(\vartheta_3, \varphi_3)\}_{l_{23}}\}_{LM}$	Tripolar spherical harmonic	Not implemented.
$\mathbf{Y}_{jm}^l(\vartheta, \varphi)$	Vector spherical harmonic	Not implemented.
$\Omega_{jm}^l(\vartheta, \varphi)$	Spinor spherical harmonic	Not implemented.
$Y_{jm}^{ls}(\vartheta, \varphi)$	Tensor spherical harmonic	Not implemented.

### Brief explanation of symbols and functions:

- **Wigner  $3n-j$  symbols ( $n = 1, 2$ , and  $3$ ):** The Wigner  $3n-j$  symbols are all related to the (re-) coupling) of angular momenta between different coupling schemes. For  $n = 1, 2$ , and  $3$ , these symbols form the basic data types in the symbolic evaluation of expressions from Racah's algebra. Indeed, the Wigner  $3n-j$  symbols frequently arise in (almost) all applications of the theory of angular momentum.
- **$12-j$  symbols of first and second type. Sharp's symbol:** Wigner  $3n-j$  symbols of higher order ( $n \geq 4$ ) are rarely used in applications as their complexity increases rapidly with  $n$  and as several *kinds* of these symbols need to be distinguished (Varshalovich *et al.* 1988). In particular, there are two kinds of  $12-j$  symbols, called the *first* and *second* kind or  $12-j(1)$  and  $12-j(2)$  symbols, respectively. These symbols are often written as

$$\left\{ \begin{array}{cccc|c} a_1 & a_2 & a_3 & a_4 & s \\ b_1 & b_2 & b_3 & b_4 & \\ c_1 & c_2 & c_3 & c_4 & \end{array} \right\},$$

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where  $s = 1, 2$  selects the *kind*. Instead of the  $12-j(2)$  symbol of *second kind*, *Sharp's symbol* (Sharp 1955) is sometimes used that exhibits a slightly higher symmetry. For  $n \geq 4$ , the properties of the  $3n-j$  symbols are still a research topic in modern group theory.

- **Clebsch-Gordan coefficients:** The Clebsch-Gordan or vector coupling coefficients appear naturally as Fourier coefficients in the *re-coupling* of angular momenta; they are closely related to the Wigner  $3-j$  symbols. In JAC, we make use of the Condon-Shortley phase convention (Condon and Shortley 1935)

$$\langle j_1 m_1, j_2 m_2 | j_3 m_3 \rangle = (-1)^{j_1 - j_2 + m_3} [j_3]^{1/2} \begin{pmatrix} j_1 & j_2 & j_3 \\ m_1 & m_2 & -m_3 \end{pmatrix}.$$

- **Racah's W coefficients:** The Racah W coefficients,  $W(abed; cf) = (-1)^{a+b+d+e} \begin{Bmatrix} a & b & c \\ d & e & f \end{Bmatrix}$ , are basically equivalent to the Wigner  $6-j$  symbols and were mainly applied in the earlier literature on the recoupling of angular momenta. In the JAC program, we only use the Wigner  $6-j$  symbols in all symbolic evaluations.
- **Wigner's D-function  $D_{mm'}^j(\alpha, \beta, \gamma)$ :** The D-functions are often required for the transformation of wave functions, if either the system or the coordinates are rotated

$$|\Psi_{jm'}(\vartheta', \varphi', \sigma')\rangle = \sum_m |\Psi_{jm}(\vartheta, \varphi, \sigma)\rangle D_{mm'}^j(\alpha, \beta, \gamma).$$

Here  $\vartheta, \varphi$  and  $\vartheta', \varphi'$  are the polar angles in the initial and the rotated system, while  $\sigma(\sigma')$  denote the corresponding spin variables. More general, these functions occur in the transformation of any irreducible spherical tensor of rank  $j$ . Therefore, the D-functions occur very frequently in the treatment of scattering processes from elementary particles up to molecules and clusters, and by including capture, transfer, and emission processes. The Wigner D-functions also fulfill a large number of symmetry properties (Varshalovich *et al.* 1988) and are expressed most readily in terms of the Euler angles and the (so-called) Wigner rotation matrix  $d_{mm'}^j(\beta)$ , see below.

- **Rotation matrix  $d_{mm'}^j(\beta)$ :** This (real) matrix describes the rotation of any spherical tensor by the angle  $\beta$  around a given axis, say, the quantization axis of the system. Several explicit representations of this real function exist in terms of  $\sin(\beta/2)$  or  $\cos(\beta/2)$ , the hypergeometric function, or various other polynomials. In JAC, the rotation matrix  $d_{mm'}^j(\beta)$  is one of the basic data structures in the algebraic manipulation of expressions from Racah's algebra.
- **Spherical harmonics:** The spherical harmonics are applied in many fields of physics. In quantum mechanics, they form an important basis for classifying the one- and many-particle states since they are known to be simultaneous eigenfunctions of one component and of the square of the orbital angular momentum operator  $-\mathbf{i}\mathbf{r} \times \nabla$ . In many-particle physics, the properties of these functions (completeness, orthogonality, ...) are frequently utilized to represent and evaluate the spin-angular part of the corresponding matrix elements analytically.

- **Spherical harmonics**  $Y_{\ell m}(\theta, \phi)$ : These functions form a complete and orthonormal set on the unit sphere, and are therefore widely used in classical and quantum physics. The spherical harmonics frequently appear in the representation of wave functions for a wide range of physical systems, in the evaluation of the corresponding (quantum) matrix elements, and at many places elsewhere. In JAC, these functions are also one of the basic data structures in the algebraic manipulation of expressions from Racah's algebra.

### 13.1.c. Symmetries of the Wigner $3n-j$ symbols, rotation matrices and spherical harmonics

#### Symmetries of the Wigner $3n-j$ symbols:

- **Symmetries of the Wigner  $3-j$  symbols:** This symbol has the following symmetries with regard to a permutation of columns and with regard to a change of signs in the projections of all angular momenta

$$\begin{aligned} \begin{pmatrix} j_1 & j_2 & j_3 \\ m_1 & m_2 & m_3 \end{pmatrix} &= \begin{pmatrix} j_2 & j_3 & j_1 \\ m_2 & m_3 & m_1 \end{pmatrix} = \begin{pmatrix} j_3 & j_1 & j_2 \\ m_3 & m_1 & m_2 \end{pmatrix} = (-1)^{j_1+j_2+j_3} \begin{pmatrix} j_1 & j_3 & j_2 \\ m_1 & m_3 & m_2 \end{pmatrix} \\ &= (-1)^{j_1+j_2+j_3} \begin{pmatrix} j_2 & j_1 & j_3 \\ m_2 & m_1 & m_3 \end{pmatrix} = (-1)^{j_1+j_2+j_3} \begin{pmatrix} j_3 & j_2 & j_1 \\ m_3 & m_2 & m_1 \end{pmatrix} \\ \begin{pmatrix} j_1 & j_2 & j_3 \\ m_1 & m_2 & m_3 \end{pmatrix} &= (-1)^{j_1+j_2+j_3} \begin{pmatrix} j_1 & j_2 & j_3 \\ -m_1 & -m_2 & -m_3 \end{pmatrix} \end{aligned}$$

- **Classical symmetries:** When combined, these symmetry relations give rise to 12 formally different  $3-j$  symbols with the same absolute value. There are **additional symmetries known due to Regge (1958)** which are most easily explained in terms of Regge symbols (Varshalovich *et al.* 1988).
- A similar distinction between classical symmetries and additional symmetric forms due to Regge can be made also for the  $6-j$  symbols.
- **Number of symmetries of the Wigner  $3n-j$  symbols:** The table below lists the number of classical symmetries vs. the overall number of symmetric forms known for the Wigner symbols. In order to simplify Racah algebra expressions, it is crucial to exploit these symmetries; they are therefore all incorporated in the program as it is explained below.
- For practical purposes, however, the classical symmetries are much more important and the distinction is therefore kept. In our notation, the **classical symmetries are a subset of the Regge symmetries**.

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- Number of classical symmetries of Wigner 3- $j$  symbols vs. the symmetries known due to Regge (1958)

3- $j$ symbol	classical symmetries	Regge symmetries (Regge 1958)
3- $j$	12	72
6- $j$	24	144
9- $j$	72	—

### Definition and symmetries of the Wigner $D$ -functions and rotation matrices:

- **Wigner  $D$ -functions  $D_{pq}^j(\alpha, \beta, \gamma)$ :** The Wigner  $D$ -functions are defined as the matrix elements of the rotation operator  $\hat{R}(\alpha, \beta, \gamma)$  in  $\mathcal{R}^3$  and are often parametrized in terms of the three Euler angles  $\alpha$ ,  $\beta$ , and  $\gamma$ . In particular, these functions arise not only in the transformation of tensor components under the rotation of the coordinates but also as the eigenfunctions of the spherical top.
- **Definition of the  $D_{pq}^j(\alpha, \beta, \gamma)$  and  $d_{pq}^j(\beta)$  rotation matrices:** The Wigner  $D$ -function is defined as a product of three functions that each depend on just a single Euler angle and with the so-called Wigner rotation matrix  $d_{pq}^j(\beta)$ :

$$D_{pq}^j(\alpha, \beta, \gamma) = e^{-ip\alpha} d_{pq}^j(\beta) e^{-iq\gamma}.$$

- **Euler angles  $(\alpha, \beta, \gamma)$ :** These angle can be utilized to characterize all possible rotations of a given (coordinate) system if the following range of values are considered:

$$0 \leq \alpha \leq 2\pi, \quad 0 \leq \beta \leq \pi, \quad 0 \leq \gamma \leq 2\pi.$$

- **Symmetries of the  $D_{pq}^j$  functions and Wigner rotation matrices  $d_{pq}^j$ :** The Wigner rotation matrix  $d_{pq}^j(\beta)$  satisfies the following basic symmetries (Varshalovich *et al.*, 1988, eq. 4.4.1):

$$\begin{aligned}
d_{pq}^j(\beta) &= (-1)^{p-q} d_{-p-q}^j(\beta) &= (-1)^{p-q} d_{qp}^j(\beta) &= d_{-q-p}^j(\beta) \\
&= d_{qp}^j(-\beta) &= (-1)^{q-p} d_{qp}^j(\beta) &= (-1)^{-j-p} d_{-pq}^j(\beta - \pi) &= (-1)^{-j-2p-q} d_{p-q}^j(\beta - \pi) \\
&= (-1)^{-2j} d_{pq}^j(\beta + 2\pi) &= (-1)^{-2j} d_{pq}^j(\beta - 2\pi) &= (-1)^{j-q} d_{p-q}^j(\beta + \pi) &= (-1)^{j-2q-p} d_{-pq}^j(\beta + \pi).
\end{aligned}$$

The symmetry properties of the Wigner  $D$ -function are listed by Varshalovich *et al.* (1988, Eq. 4.2.2) and are taken into account since the  $D$ -function are internally handled in terms of the Wigner rotation matrix.

### Symmetries of the spherical harmonics:

➤ **Symmetries:** The spherical harmonics satisfy the following basic symmetries:

$$\begin{aligned} Y_{\ell m}(\vartheta, \varphi) &= (-1)^m e^{2im\varphi} Y_{\ell, -m}(\vartheta, \varphi) = (-1)^m Y_{\ell m}(-\vartheta, \varphi) = e^{2im\varphi} Y_{\ell, -m}(-\vartheta, \varphi) = (-1)^m Y_{\ell, -m}(\vartheta, -\varphi) \\ &= e^{2im\varphi} Y_{\ell m}(\vartheta, -\varphi) = Y_{\ell, -m}(-\vartheta, -\varphi) = (-1)^m e^{2im\varphi} Y_{\ell m}(-\vartheta, -\varphi) . \end{aligned}$$

### 13.1.d. Expansions of the Wigner $3n-j$ symbols, rotation matrices and spherical harmonics

#### Explicit formulas for the computation of the Wigner $3n-j$ symbols:

➤  **$\Delta$  symbol:**

$$\Delta(a, b, c) = \left[ \frac{(a+b-c)! (a-b+c)! (-a+b+c)!}{(a+b+c+1)!} \right]^{1/2} .$$

➤ **Wigner  $3-j$  symbol:** The numerical value of a Wigner  $3-j$  symbol can be calculated by the expression (Racah, 1942) and has a non-zero value only if the arguments of all factorials are non-negative integers

$$\begin{aligned} \begin{pmatrix} j_1 & j_2 & j_3 \\ m_1 & m_2 & m_3 \end{pmatrix} &= \delta_{m_1+m_2+m_3,0} (-1)^{j_1-j_2-m_3} \Delta(j_1, j_2, j_3) [(j_1-m_1)! (j_1+m_1)! (j_2-m_2)! (j_2+m_2)! (j_3-m_3)! (j_3+m_3)!]^{1/2} \\ &\times \sum_l \left[ \frac{(-1)^l}{l! (j_1+j_2-j_3-l)! (j_1-m_1-l)! (j_2+m_2-l)! (j_3-j_2+m_1+l)! (j_3-j_1-m_2+l)!} \right] . \end{aligned}$$

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➤ **Wigner 6- $j$  symbol:** Edmonds (1957) displays the following expression for the computation of a Wigner 6- $j$  symbol

$$\left\{ \begin{matrix} j_1 & j_2 & j_3 \\ l_1 & l_2 & l_3 \end{matrix} \right\} = \Delta(j_1, j_2, j_3) \Delta(j_1, l_2, l_3) \Delta(l_1, j_2, l_3) \Delta(l_1, l_2, j_3) \sum_l \left[ \frac{(-1)^l (\ell + 1)!}{(l - j_1 - j_2 - j_3)! (l - j_1 - l_2 - l_3)!} \right. \\ \left. \times \frac{1}{(l - l_1 - j_2 - l_3)! (l - l_1 - l_2 - j_3)! (j_1 + j_2 + l_1 + l_2 - l)! (j_2 + j_3 + l_2 + l_3 - l)! (j_3 + j_1 + l_3 + l_1 - l)!} \right].$$

➤ **Wigner 9- $j$  symbol:** The Wigner 9- $j$  symbol with numeric arguments is usually calculated by the known sum rule over three Wigner 6- $j$  symbols and is **zero unless the arguments in each row and column satisfy the triangular relation**

$$\left\{ \begin{matrix} j_{11} & j_{12} & j_{13} \\ j_{21} & j_{22} & j_{23} \\ j_{31} & j_{32} & j_{33} \end{matrix} \right\} = \sum_j (-1)^{2j} \left\{ \begin{matrix} j_{11} & j_{21} & j_{31} \\ j_{32} & j_{33} & j \end{matrix} \right\} \left\{ \begin{matrix} j_{12} & j_{22} & j_{32} \\ j_{21} & j & j_{23} \end{matrix} \right\} \left\{ \begin{matrix} j_{13} & j_{23} & j_{33} \\ j & j_{11} & j_{12} \end{matrix} \right\}.$$

### Explicit formulas and expansions of the Wigner rotation matrices $d_{pq}^j(\beta)$ :

➤ Below, we display six explicit expansions of the Wigner rotation matrices in alphabetical order.

➤ **Asymptotic expansion of  $d_{pq}^j$  for  $j \gg 1$ :** There are three asymptotics expansion of Wigner rotation matrices as displayed by Varshalovich *et al.* (1988, Eqs. 4.18.1–4). For  $j \gg 1$ , for instance, the Wigner rotation matrix can be written as

$$d_{pq}^j(\beta) \approx \xi_{pq} \sqrt{\frac{s!(s+\mu+\nu)!}{(s+\mu)!(s+\nu)!}} \sqrt{\frac{2}{\pi s}} \frac{\cos \left[ \left( s + \frac{\mu+\nu+1}{2} \right) \beta - \frac{\pi}{4} (2\mu+1) \right]}{\sqrt{\sin \beta}} + O\left(\frac{1}{J^{\frac{3}{2}}}\right),$$

and where  $\mu$ ,  $\nu$ , and  $s$  are related to  $p$ ,  $q$ , and  $j$  as follows

$$\mu = |p - q|, \quad \nu = |p + q|, \quad s = j - \frac{1}{2}(\mu + \nu), \quad \xi_{pq} = \begin{cases} 1 & \text{if } q \geq p, \\ (-1)^{q-p} & \text{if } q < p. \end{cases}$$

➤ **Asymptotic expansion of  $d_{pq}^j$  for  $J \rightarrow \infty$  and  $\beta \rightarrow 0$ :** A second expansion for  $j\beta < \infty$  is given in terms of the Bessel functions  $J_n(x)$  by

$$d_{pq}^j(\beta) \approx J_{p-q}(j\beta)$$



➤ Asymptotic expansion of  $d_{pq}^j$  for small variations of the rotation axis:

$$d_{pq}^j(\beta) \approx \frac{\xi_{pq}}{\mu!} \sqrt{\frac{(s+\mu+\nu)!(s+\mu)!}{s!(s+\nu)!}} \left(\frac{\beta}{2}\right)^\mu \left\{ 1 - \frac{2s(s+\mu+\nu+1) + \nu(\mu+1)}{2(\mu+1)} \left(\frac{\beta}{2}\right)^2 + \dots \right\}, \quad \beta \rightarrow 0$$

$$d_{pq}^j(\beta) \approx \frac{(-1)^s \xi_{pq}}{\nu!} \sqrt{\frac{(s+\mu+\nu)!(s+\nu)!}{s!(s+\nu)!}} \left(\frac{\pi-\beta}{2}\right)^\mu \left\{ 1 - \frac{2s(s+\mu+\nu+1) + \mu(\nu+1)}{2(\nu+1)} \left(\frac{\pi-\beta}{2}\right)^2 + \dots \right\}, \quad \beta - \pi \rightarrow 0$$

➤ Clebsch-Gordan series expansion of products of two  $d_{pq}^j(\beta)$  with equal arguments: Products of two Wigner rotation matrices with equal angular arguments can be expanded in terms of their uncoupled and coupled representations, i.e. by means of the Clebsch-Gordan coefficient. When expressed in terms of the Wigner 3- $j$  symbols, the product of two Wigner rotation matrices (with equal angular arguments) are given by

$$d_{p_1 q_1}^{j_1}(\beta) d_{p_2 q_2}^{j_2}(\beta) = \sum_{JPQ} (-1)^{-2J+2j_2-p_2-q_2} (2j_2+1) \begin{pmatrix} J & j_2 & j_2 \\ P & p_2 & -p_2 \end{pmatrix} \begin{pmatrix} J & j_2 & j_2 \\ Q & q_2 & -q_2 \end{pmatrix} d_{PQ}^J(\beta)$$

in terms of two Wigner 3- $j$  symbols and one Wigner rotation matrix  $d_{PQ}^J(\beta)$ . For  $\beta = \frac{\pi}{2}$ , in particular, the magnetic quantum numbers can be replaced by  $j_1 = j_2 = j$ ,  $p_1 = p_2 = p$  and  $q_1 = q_2 = q$ , and this gives rise to the expansion

$$\left(d_{pq}^j\left(\frac{\pi}{2}\right)\right)^2 = \sum_{J=0,2,4,\dots} (-1)^{J+2j+2p} (2j+1) \frac{(J-1)!!}{J!!} \begin{pmatrix} J & j & j \\ 0 & p & p \end{pmatrix} \begin{pmatrix} J & j & j \\ 0 & q & q \end{pmatrix}.$$

Such expansion relates the Wigner rotation matrices to the Wigner  $n$ - $j$  symbols. These two expansions are displayed in Varshalovich *et al.* (1988, Eq. 4.6.1 and Eq. 4.16.9), although expressed in terms of Clebsch-Gordan coefficients.

## Expansions of the spherical harmonics:

➤ Varshalovich *et al.* (1988, section 5.2) shows several explicit formulas for computing the spherical harmonic  $Y_{\ell m}(\vartheta, \varphi)$  with given integers  $\ell$  and  $m$ , and where the summation runs over all terms with non-negative factorials.

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➤ Expansion of  $Y_{\ell m}(\vartheta, \varphi)$  in terms of  $\sin \vartheta$ : Varshalovich *et al.* (1988, Eq. 5.2.17)

$$Y_{\ell m}(\vartheta, \varphi) = e^{im\varphi} \sqrt{\frac{2\ell+1}{4\pi (\ell+m)! (\ell-m)!}} \begin{cases} \sum_{s=|m|, |m|+2, \dots}^{\ell} (-1)^{\frac{s+m}{2}} \frac{(\ell+s)!}{(s+m)!! (s-m)!!} \cdot \frac{(\ell+m)!! (\ell-m)!!}{(\ell+s)!! (\ell-s)!!} (\sin \vartheta)^s & \text{if } \ell - m \text{ is even,} \\ \cos \vartheta \sum_{s=|m|, |m|+2, \dots}^{\ell-1} (-1)^{\frac{s+m}{2}} \frac{(\ell+s)!}{(s+m)!! (s-m)!!} \cdot \frac{(\ell+m-1)!! (\ell-m-1)!!}{(\ell+s-1)!! (\ell-s-1)!!} (\sin \vartheta)^s & \text{if } \ell - m \text{ is odd.} \end{cases}$$

➤ Expansion of  $Y_{\ell m}(\vartheta, \varphi)$  in terms of  $\sin \vartheta$  and  $\cos \vartheta$ : Cf. Varshalovich *et al.* (1988, Eq. 5.2.19) with a polynomial in  $\cos \vartheta$

$$Y_{\ell m}(\vartheta, \varphi) = e^{im\varphi} \sqrt{\frac{2\ell+1}{4\pi} \frac{(\ell-m)!}{(\ell+m)!}} (\sin \vartheta)^m \sum_{\substack{s \\ \ell+m-s \text{ even}}} (-1)^{\frac{\ell+m-s}{2}} \frac{(\ell+m+s-1)!!}{(\ell-m-s)!!} \frac{(\cos \vartheta)^s}{s!}$$

## 13.2. In JAC implemented symbolic evaluations of Racah algebra expressions

### 13.2.a. Strategies in the symbolic evaluation of expressions from Racah's algebra

#### Strategies for simplification:

- For the successful simplification of expressions from Racah's algebra, one need often to exploit the orthogonality relations as well as a variety of important sum rules. The literature about Racah's algebra is therefore devoted to a large extent in order to find and to prove proper relations among the Wigner symbols and/or to compile them in some applicable form.
- **Standard presentation of orthogonality and sum rules:** The standard presentation of the different orthogonality and sum rules, as they are displayed below, gives a first impression how complex expressions can be rewritten in a simpler form. In practice, however, these standard forms are often not of great help since one first need to recognize the equivalence of some part of a Racah algebra expression, for instance with one side of a given sum rule, before this special rule can be applied.
- A numerical evaluation of a Racah expression is of course not possible as long as the input angular momentum quantum numbers are specified numerically. On the other hand, the main strength of using the techniques of Racah's algebra techniques is that such expressions can often be simplified considerably by algebraic transformations.

- Understanding the symmetries and properties of Wigner symbols, three different strategies for the simplification can be derived quite easily:
  - (i) **Use of known special values.** This replaces a Wigner  $3n-j$  symbol by a (much) simpler expression that, in particular, does not contain any implicit summation. For this strategy, each Wigner  $3n-j$  symbol can be analysed and perhaps replaced independently by some special-value rules.
  - (ii) **Use of orthogonality properties.**
  - (iii) **Use of sum rules.**
- The main obstacle in applying the strategies (ii) and (iii) is that these strategies require a careful analysis of the Racah expression as a whole. Since all orthogonality relations and sum rules include summations over formal quantum numbers, all summation variables not only have to be in the correct position in the Wigner  $3n-j$  symbols, but they must also contribute to a correct phase and weight of the overall expression. Moreover, the same variables may not occur in other Wigner symbols of the Racah expression which are not part of the selected rule.
- These difficulties in the simplification process are further enhanced by the large number of equivalent forms of the total Racah expression due to the symmetries of the Wigner  $3n-j$  symbols. Therefore, in order to simplify a Racah expression by a given sum rule, in practice one has to start with a certain part of the expression and then try to identify equivalence with some relation by means of the various symmetries of the  $3n-j$  symbols.
- Once the equivalence has been proven, this part of the Racah expression can be replaced by a corresponding simpler term. In this context, **simplification of a Racah expression always means to 'reduce' the number of summation indices and/or the number of Wigner  $3n-j$  symbols.**
- **Aim of symbolic transformations:** The aim of these symbolic transformation is to obtain an *optimum* summation formula in the sense of a minimal number of Wigner  $6-j$  symbols and/or summation variables.
- **Results of the simplification process:** The results of the simplification process will be provided as Racah expressions and may thus immediatly be used for further derivations and calculations within the theory of angular momentum.
- This rather cumbersome procedure of running through all symmetric forms of a Racah expression and of identifying algebraic equivalent parts makes the simplification of such expressions very suitable for symbolic evaluations. An efficient scheme to perform simplifications of Racah expressions by means of the steps (i-iii) is aimed to be realized by the symbolic evaluations of the JAC tools.
- A major difficulty for the simplification of general Racah expressions arise from the large number of symmetric forms of the total expression which is the direct result of the symmetries of the Wigner  $3n-j$  symbols. Apart from the so-called classical symmetries of the Wigner  $n-j$  symbols, there are additional symmetries known due to Regge for the  $3-j$  and  $6-j$  symbols; cf. section 13.1.c. These symmetries need to be exploited in order to simplify general Racah algebra expressions. In JAC, we shall keep the distinction between the classical symmetries

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and those due to Regge since the classical ones are by far more important for all practical purposes; as mentioned above, we consider the classical symmetries to be a subset of the Regge symmetries.

## 13.3. In JAC partly-implemented symbolic evaluations of Racah algebra expressions

### 13.3.a. Recursive relations for the Wigner $3n-j$ symbols and rotation matrices

#### Recursion relations of the Wigner $3-j$ symbols:

- Four recursion relations of the Wigner  $3-j$  symbols are listed by Rotenberg *et al.* (1959, Eqs. 1.45–1.48), and where we use  $J = j_1 + j_2 + j_3$ .
- **Half-integer recursion:** This recursion relation decreases two  $j$ -values by  $1/2$

$$\begin{aligned} & [(J+1)(J-2j_1)]^{1/2} \begin{pmatrix} j_1 & j_2 & j_3 \\ m_1 & m_2 & m_3 \end{pmatrix} \\ &= [(j_2+m_2)(j_3-m_3)]^{1/2} \begin{pmatrix} j_1 & j_2-1/2 & j_3-1/2 \\ m_1 & m_2-1/2 & m_3+1/2 \end{pmatrix} - [(j_2-m_2)(j_3+m_3)]^{1/2} \begin{pmatrix} j_1 & j_2-1/2 & j_3-1/2 \\ m_1 & m_2+1/2 & m_3-1/2 \end{pmatrix}. \end{aligned}$$

- **Louck's (1958) half-integer recursion:**

$$\begin{aligned} (j_2+m_2)^{1/2} (2j_3+1) \begin{pmatrix} j_1 & j_2 & j_3 \\ m_2-m_3 & -m_2 & m_3 \end{pmatrix} &= - [(J-2j_1)(J+1)(j_3+m_3)]^{1/2} \begin{pmatrix} j_1 & j_2-1/2 & j_3-1/2 \\ m_2-m_3 & -m_2+1/2 & m_3-1/2 \end{pmatrix} \\ &- [(J-2j_3)(J-2j_2+1)(J+1)(j_3-m_3+1)]^{1/2} \begin{pmatrix} j_1 & j_2-1/2 & j_3+1/2 \\ m_2-m_3 & -m_2+1/2 & m_3-1/2 \end{pmatrix}. \end{aligned}$$

➤ **Integer recursion:** A decrease of a single  $j$ -values by 1 is made with

$$\begin{aligned}
 & [(J+1)(J-2j_1)(J-2j_2)(J-2j_3+1)]^{1/2} \begin{pmatrix} j_1 & j_2 & j_3 \\ m_1 & m_2 & m_3 \end{pmatrix} \\
 &= [(j_2-m_2)(j_2+m_2+1)(j_3+m_3)(j_3+m_3-1)]^{1/2} \begin{pmatrix} j_1 & j_2 & j_3-1 \\ m_1 & m_2+1 & m_3-1 \end{pmatrix} - 2m_2 [(j_3+m_3)(j_3-m_3)]^{1/2} \begin{pmatrix} j_1 & j_2 & j_3-1 \\ m_1 & m_2 & m_3 \end{pmatrix} \\
 &\quad - [(j_2+m_2)(j_2-m_2+1)(j_3-m_3)(j_3-m_3-1)]^{1/2} \begin{pmatrix} j_1 & j_2 & j_3-1 \\ m_1 & m_2-1 & m_3+1 \end{pmatrix}.
 \end{aligned}$$

➤ **Magnetic recursion:** A recursion relation with regard to the magnetic quantum numbers reads:

$$\begin{aligned}
 & - [(j_3+m_1+m_2+1)(j_3-m_1-m_2)]^{1/2} \begin{pmatrix} j_1 & j_2 & j_3 \\ m_1 & m_2 & -m_3+1 \end{pmatrix} \\
 &= [(j_1+m_1+1)(j_1-m_1)]^{1/2} \begin{pmatrix} j_1 & j_2 & j_3 \\ m_1 & m_2+1 & -m_3 \end{pmatrix} + [(j_2+m_2+1)(j_2-m_2)]^{1/2} \begin{pmatrix} j_1 & j_2 & j_3 \\ m_1+1 & m_2 & -m_3 \end{pmatrix}.
 \end{aligned}$$

### Recursion relations of the Wigner 6- $j$ symbols:

➤ No explicit recurrence relation has yet been considered for the Wigner 6- $j$  symbols.

### Recursion relations of the Wigner 9- $j$ symbols:

➤ A single recursion relation among the Wigner 9- $j$  symbols is shown by Rotenberg *et al.* (1959, Eq. 3.24).

$$\begin{pmatrix} s & s & 1 \\ l_1 & l_2 & L \\ j_1 & j_2 & L \end{pmatrix} = \frac{l_1(l_1+1) + j_2(j_2+1) - j_1(j_1+1) - l_2(l_2+1)}{2[s(s+1)(2s+1)L(L+1)(2L+1)]^{1/2}} \begin{pmatrix} s & s & 0 \\ l_1 & l_2 & L \\ j_1 & j_2 & L \end{pmatrix}.$$

### Recursion relations of the spherical harmonics:

➤ **Recursion relations:** Cf. Varshalovich *et al.* (1988, Eqs. 5.7.1–2)

$$\begin{aligned} \cos \vartheta Y_{\ell m}(\vartheta, \varphi) &= \sqrt{\frac{(\ell - m + 1)(\ell + m + 1)}{(2\ell + 1)(2\ell + 3)}} Y_{\ell+1, m}(\vartheta, \varphi) + \sqrt{\frac{(\ell - m)(\ell + m)}{(2\ell - 1)(2\ell + 1)}} Y_{\ell-1, m}(\vartheta, \varphi) \\ - 2m \cos \vartheta Y_{\ell m}(\vartheta, \varphi) &= \sqrt{l(l+1) - m(m+1)} e^{-i\varphi} Y_{\ell, m+1}(\vartheta, \varphi) + \sqrt{l(l+1) - m(m-1)} e^{i\varphi} Y_{\ell, m-1}(\vartheta, \varphi). \end{aligned}$$

### 13.3.b. Special values of the Wigner $3n-j$ symbols and rotation matrices

#### Special values of the Wigner $3-j$ symbols:

➤ Edmonds (1957, appendix 2) lists the following 20 special values for  $3-j$  symbols

$$\begin{pmatrix} j_1 & j_2 & j_3 \\ 0 & 0 & 0 \end{pmatrix} = (-1)^{J/2} \left[ \frac{(j_1 + j_2 - j_3)!(j_1 + j_3 - j_2)!(j_2 + j_3 - j_1)!}{(j_1 + j_2 + j_3 + 1)!} \right]^{1/2} \frac{(J/2)!}{(J/2 - j_1)!(J/2 - j_2)!(J/2 - j_3)!}$$

$$\begin{pmatrix} j_1 & j_2 & j_3 \\ 0 & 0 & 0 \end{pmatrix} = 0 \quad \begin{cases} \text{if } J = j_1 + j_2 + j_3 \text{ is even} \\ \text{if } J = j_1 + j_2 + j_3 \text{ is odd} \end{cases}$$

$$\begin{pmatrix} J + 1/2 & J & 1/2 \\ M & -M - 1/2 & 1/2 \end{pmatrix} = (-1)^{J-M-1/2} \left[ \frac{J - M + 1/2}{(2J + 2)(2J + 1)} \right]^{1/2}$$

$$\begin{pmatrix} J + 1 & J & 1 \\ M & -M - 1 & 1 \end{pmatrix} = (-1)^{J-M-1} \left[ \frac{(J - M)(J - M + 1)}{(2J + 3)(2J + 2)(2J + 1)} \right]^{1/2}$$

$$\begin{pmatrix} J + 1 & J & 1 \\ M & -M & 0 \end{pmatrix} = (-1)^{J-M-1} \left[ \frac{2(J + M + 1)(J - M + 1)}{(2J + 3)(2J + 2)(2J + 1)} \right]^{1/2}$$

$$\begin{pmatrix} J & J & 1 \\ M & -M - 1 & 1 \end{pmatrix} = (-1)^{J-M} \left[ \frac{2(J - M)(J + M + 1)}{(2J + 2)(2J + 1)(2J)} \right]^{1/2}$$

$$\begin{pmatrix} J & J & 1 \\ M & -M & 0 \end{pmatrix} = (-1)^{J-M} \frac{M}{[(2J + 1)(J + 1)J]^{1/2}}$$

$$\begin{pmatrix} J + 3/2 & J & 3/2 \\ M & -M - 3/2 & 3/2 \end{pmatrix} = (-1)^{J-M+1/2} \left[ \frac{(J - M - 1/2)(J - M + 1/2)(J - M + 3/2)}{(2J + 4)(2J + 3)(2J + 2)(2J + 1)} \right]^{1/2}$$

$$\begin{pmatrix} J + 3/2 & J & 3/2 \\ M & -M - 1/2 & 1/2 \end{pmatrix} = (-1)^{J-M+1/2} \left[ \frac{3(J - M + 1/2)(J - M + 3/2)(J + M + 3/2)}{(2J + 4)(2J + 3)(2J + 2)(2J + 1)} \right]^{1/2}$$

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$$\begin{aligned}
\begin{pmatrix} J+1/2 & J & 3/2 \\ M & -M-3/2 & 3/2 \end{pmatrix} &= (-1)^{J-M-1/2} \left[ \frac{3(J-M-1/2)(J-M+1/2)(J+M+3/2)}{(2J+3)(2J+2)(2J+1)(2J)} \right]^{1/2} \\
\begin{pmatrix} J+1/2 & J & 3/2 \\ M & -M-1/2 & 1/2 \end{pmatrix} &= (-1)^{J-M-1/2} \left[ \frac{(J-M+1/2)}{(2J+3)(2J+2)(2J+1)(2J)} \right]^{1/2} (J+3M+3/2) \\
\begin{pmatrix} J+2 & J & 2 \\ M & -M-2 & 2 \end{pmatrix} &= (-1)^{J-M} \left[ \frac{(J-M-1)(J-M)(J-M+1)(J-M+2)}{(2J+5)(2J+4)(2J+3)(2J+2)(2J+1)} \right]^{1/2} \\
\begin{pmatrix} J+2 & J & 2 \\ M & -M-1 & 1 \end{pmatrix} &= (-1)^{J-M} 2 \left[ \frac{(J+M+2)(J-M+2)(J-M+1)(J-M)}{(2J+5)(2J+4)(2J+3)(2J+2)(2J+1)} \right]^{1/2} \\
\begin{pmatrix} J+2 & J & 2 \\ M & -M & 0 \end{pmatrix} &= (-1)^{J-M} \left[ \frac{6(J+M+2)(J+M+1)(J-M+2)(J-M+1)}{(2J+5)(2J+4)(2J+3)(2J+2)(2J+1)} \right]^{1/2} \\
\begin{pmatrix} J+1 & J & 2 \\ M & -M-2 & 2 \end{pmatrix} &= (-1)^{J-M+1} 2 \left[ \frac{(J-M-1)(J-M)(J-M+1)(J+M+2)}{(2J+4)(2J+3)(2J+2)(2J+1)(2J)} \right]^{1/2} \\
\begin{pmatrix} J+1 & J & 2 \\ M & -M-1 & 1 \end{pmatrix} &= (-1)^{J-M+1} 2 (J+2M+2) \left[ \frac{(J-M+1)(J-M)}{(2J+4)(2J+3)(2J+2)(2J+1)(2J)} \right]^{1/2} \\
\begin{pmatrix} J+1 & J & 2 \\ M & -M & 0 \end{pmatrix} &= (-1)^{J-M+1} 2M \left[ \frac{6(J+M+1)(J-M+1)}{(2J+4)(2J+3)(2J+2)(2J+1)(2J)} \right]^{1/2} \\
\begin{pmatrix} J & J & 2 \\ M & -M-2 & 2 \end{pmatrix} &= (-1)^{J-M} \left[ \frac{6(J-M-1)(J-M)(J+M+1)(J+M+2)}{(2J+3)(2J+2)(2J+1)(2J)(2J-1)} \right]^{1/2} \\
\begin{pmatrix} J & J & 2 \\ M & -M-1 & 1 \end{pmatrix} &= (-1)^{J-M} (1+2M) \left[ \frac{6(J+M+1)(J-M)}{(2J+3)(2J+2)(2J+1)(2J)(2J-1)} \right]^{1/2} \\
\begin{pmatrix} J & J & 2 \\ M & -M & 0 \end{pmatrix} &= (-1)^{J-M} \frac{2[3M^2 - J(J+1)]}{[(2J+3)(2J+2)(2J+1)(2J)(2J-1)]^{1/2}}
\end{aligned}$$



### Special values of the Wigner 6- $j$ symbols:

➤ Edmonds (1957, appendix 2, table 5) displays the following 19 special values for 6- $j$  symbols; in this list, moreover, we use  $s = a + b + c$ .

$$\left\{ \begin{matrix} a & b & c \\ 1 & c-1 & b-1 \end{matrix} \right\} = (-1)^s \left[ \frac{s(s+1)(s-2a-1)(s-2a)}{(2b-1)2b(2b+1)(2c-1)2c(2c+1)} \right]^{1/2}$$

$$\left\{ \begin{matrix} a & b & c \\ 1 & c-1 & b \end{matrix} \right\} = (-1)^s \left[ \frac{2(s+1)(s-2a)(s-2b)(s-2c+1)}{2b(2b+1)(2b+2)(2c-1)2c(2c+1)} \right]^{1/2}$$

$$\left\{ \begin{matrix} a & b & c \\ 1 & c-1 & b+1 \end{matrix} \right\} = (-1)^s \left[ \frac{(s-2b-1)(s-2b)(s-2c+1)(s-2c+2)}{(2b+1)(2b+2)(2b+3)(2c-1)2c(2c+1)} \right]^{1/2}$$

$$\left\{ \begin{matrix} a & b & c \\ 1 & c & b \end{matrix} \right\} = (-1)^{s+1} \frac{2[b(b+1) + c(c+1) - a(a+1)]}{[2b(2b+1)(2b+2)2c(2c+1)(2c+2)]^{1/2}}$$

$$\left\{ \begin{matrix} a & b & c \\ 3/2 & c-3/2 & b-3/2 \end{matrix} \right\} = (-1)^s \left[ \frac{(s-1)s(s+1)(s-2a-2)(s-2a-1)(s-2a)}{(2b-2)(2b-1)2b(2b+1)(2c-2)(2c-1)2c(2c+1)} \right]^{1/2}$$

$$\left\{ \begin{matrix} a & b & c \\ 3/2 & c-3/2 & b-1/2 \end{matrix} \right\} = (-1)^s \left[ \frac{3s(s+1)(s-2a-1)(s-2a)(s-2b)(s-2c+1)}{(2b-1)2b(2b+1)(2b+2)(2c-2)(2c-1)2c(2c+1)} \right]^{1/2}$$

$$\left\{ \begin{matrix} a & b & c \\ 3/2 & c-3/2 & b+1/2 \end{matrix} \right\} = (-1)^s \left[ \frac{3(s+1)(s-2a)(s-2b-1)(s-2b)}{2b(2b+1)(2b+2)(2b+3)} \right]^{1/2} \left[ \frac{(s-2c+1)(s-2c+2)}{(2c-2)(2c-1)2c(2c+1)} \right]^{1/2}$$

$$\left\{ \begin{matrix} a & b & c \\ 3/2 & c-3/2 & b+3/2 \end{matrix} \right\} = (-1)^s \left[ \frac{(s-2b-2)(s-2b-1)(s-2b)}{(2b+1)(2b+2)(2b+3)(2b+4)} \right]^{1/2} \left[ \frac{(s-2c+1)(s-2c+2)(s-2c+3)}{(2c-2)(2c-1)2c(2c+1)} \right]^{1/2}$$

$$\left\{ \begin{matrix} a & b & c \\ 3/2 & c-1/2 & b-1/2 \end{matrix} \right\} = (-1)^s \frac{[2(s-2b)(s-2c) - (s+2)(s-2a-1)] [(s+1)(s-2a)]^{1/2}}{[(2b-1)(2b)(2b+1)(2b+2)(2c-1)2c(2c+1)(2c+2)]^{1/2}}$$

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$$\begin{aligned}
\left\{ \begin{matrix} a & b & c \\ 3/2 & c-1/2 & b+1/2 \end{matrix} \right\} &= (-1)^s \frac{[(s-2b-1)(s-2c)-2(s+2)(s-2a)]}{[2b(2b+1)(2b+2)(2b+3)]^{1/2}} \frac{[(s-2b)(s-2c+1)]^{1/2}}{[2c(2c+1)(2c+2)(2c-1)]^{1/2}} \\
\left\{ \begin{matrix} a & b & c \\ 2 & c-2 & b-2 \end{matrix} \right\} &= (-1)^s \left[ \frac{(s-2)(s-1)s(s+1)(s-2a-3)}{(2b-3)(2b-2)(2b-1)2b(2b+1)} \right]^{1/2} \left[ \frac{(s-2a-2)(s-2a-1)(s-2a)}{(2c-3)(2c-2)(2c-1)2c(2c+1)} \right]^{1/2} \\
\left\{ \begin{matrix} a & b & c \\ 2 & c-2 & b-1 \end{matrix} \right\} &= (-1)^s 2 \left[ \frac{(s-1)s(s+1)(s-2a-2)(s-2a-1)}{(2b-2)(2b-1)2b(2b+1)(2b+2)} \right]^{1/2} \left[ \frac{(s-2a)(s-2b)(s-2c+1)}{(2c-3)(2c-2)(2c-1)2c(2c+1)} \right]^{1/2} \\
\left\{ \begin{matrix} a & b & c \\ 2 & c-2 & b \end{matrix} \right\} &= (-1)^s \left[ \frac{6s(s+1)(s-2a-1)(s-2b-1)(s-2a)(s-2b)}{(2b-1)2b(2b+1)(2b+2)(2b+3)} \right]^{1/2} \left[ \frac{(s-2c+1)(s-2c+2)}{(2c-3)(2c-2)(2c-1)2c(2c+1)} \right]^{1/2} \\
\left\{ \begin{matrix} a & b & c \\ 2 & c-2 & b+1 \end{matrix} \right\} &= (-1)^s 2 \left[ \frac{(s+1)(s-2a)(s-2b-2)(s-2b-1)(s-2b)}{2b(2b+1)(2b+2)(2b+3)(2b+4)} \right]^{1/2} \left[ \frac{(s-2c+1)(s-2c+2)(s-2c+3)}{(2c-3)(2c-2)(2c-1)2c(2c+1)} \right]^{1/2} \\
\left\{ \begin{matrix} a & b & c \\ 2 & c-2 & b+2 \end{matrix} \right\} &= (-1)^s \left[ \frac{(s-2b-3)(s-2b-2)(s-2b-1)(s-2b)}{(2b+1)(2b+2)(2b+3)(2b+4)(2b+5)} \right]^{1/2} \left[ \frac{(s-2c+1)(s-2c+2)(s-2c+3)(s-2c+4)}{(2c-3)(2c-2)(2c-1)2c(2c+1)} \right]^{1/2} \\
\left\{ \begin{matrix} a & b & c \\ 2 & c-1 & b-1 \end{matrix} \right\} &= (-1)^s 4 \frac{[(a+b)(a-b+1)-(c-1)(c-b+1)]}{[(2b-2)(2b-1)2b(2b+1)(2b+2)]^{1/2}} \left[ \frac{s(s+1)(s-2a-1)(s-2a)}{(2c-2)(2c-1)2c(2c+1)(2c+2)} \right]^{1/2} \\
\left\{ \begin{matrix} a & b & c \\ 2 & c-1 & b \end{matrix} \right\} &= (-1)^s 2 \frac{[(a+b+1)(a-b)-c^2+1]}{[(2b-1)2b(2b+1)(2b+2)(2b+3)]^{1/2}} \left[ \frac{6(s+1)(s-2a)(s-2b)(s-2c+1)}{(2c-2)(2c-1)2c(2c+1)(2c+2)} \right]^{1/2} \\
\left\{ \begin{matrix} a & b & c \\ 2 & c-1 & b+1 \end{matrix} \right\} &= (-1)^s 4 \frac{[(a+b+2)(a-b-1)-(c-1)(b+c+2)]}{[2b(2b+1)(2b+2)(2b+3)(2b+4)]^{1/2}} \left[ \frac{(s-2b-1)(s-2b)(s-2c+1)(s-2c+2)}{(2c-2)(2c-1)2c(2c+1)(2c+2)} \right]^{1/2} \\
\left\{ \begin{matrix} a & b & c \\ 2 & c & b \end{matrix} \right\} &= (-1)^s \frac{2}{[(2b-1)2b(2b+1)(2b+2)(2b+3)]^{1/2}} \frac{[3X(X-1)-4b(b+1)c(c+1)]}{[(2c-1)2c(2c+1)(2c+2)(2c+3)]^{1/2}}
\end{aligned}$$

and where we use  $s = a + b + c$  and  $X = b(b+1) + c(c+1) - a(a+1)$  in the last relation.

**Special values of the Wigner 9- $j$  symbols:**

➤ One special value for a Wigner 9- $j$  symbol reads as

$$\begin{Bmatrix} a & b & 0 \\ c & d & 0 \\ e & f & 0 \end{Bmatrix} = \delta_{ab} \delta_{cd} \delta_{ef} \frac{\delta(a, c, e)}{[a, c, e]^{1/2}}.$$

➤ Another special value for the Wigner 9- $j$  symbols is given by Varshalovich *et al.* (1988, Eq. 10.9.1)

$$\begin{Bmatrix} a & b & c \\ d & e & f \\ g & h & 0 \end{Bmatrix} = \delta_{cf} \delta_{gh} \frac{(-1)^{b+c+e+f}}{[c, g]^{1/2}} \begin{Bmatrix} a & b & c \\ e & d & g \end{Bmatrix}.$$

**13.3.c. Orthogonality and sum rules of the Wigner 3 $n$ - $j$  symbols****Remarks:**

- A rather comprehensive set of sum rules for the Wigner  $n$ - $j$  symbols is given by Varshalovich *et al.* (1988, chapter 12). These rules are most easily be ordered by the number of Wigner  $n$ - $j$  symbols that contribute to a given rule. For the case of an equal numbers of such symbols, we further classify these rules by the type of the Wigner symbols.
- **Orthogonality relations:** The Wigner  $n$ - $j$  symbols fulfill a variety of orthogonality relations. Two such relations are known for the Wigner 3- $j$  symbols as well as one orthogonality relation for each, the Wigner 6- $j$  and 9- $j$  symbols. Because these orthogonalities have formally the same structure like all other sum rules of the Wigner  $n$ - $j$  symbols, they are all treated together in the JAC program.

### Sum rules with one Wigner 3- $j$ , 6- $j$ or 9- $j$ symbol:

➤ **Sum rules with one 3- $j$  symbol:** Varshalovich *et al.* (1988, Eq. 12.1.2) shows one sum rule with one 3- $j$  symbol

$$\sum_m (-1)^{j-m} \begin{pmatrix} j & j & j' \\ m & -m & m' \end{pmatrix} = [j]^{1/2} \delta_{m'0} \delta_{j'0}.$$

➤ **Sum rules with one 6- $j$  symbol:** Rotenberg *et al.* (1959, Eqs. 2.9–10) show sum rules with one 6- $j$  symbol. A more general form is given by Varshalovich *et al.* (1988, Eqs. 12.2.3–4)

$$\sum_X [X] \begin{Bmatrix} a & b & X \\ a & b & c \end{Bmatrix} = (-1)^{2c} \delta(a, b, c)$$

$$\sum_X (-1)^{a+b+X} [X] \begin{Bmatrix} a & b & X \\ b & a & c \end{Bmatrix} = [a, b]^{1/2} \delta_{c0}.$$

➤ **Sum rules with one 9- $j$  symbol:** Varshalovich *et al.* (1988, Eqs. 12.2.5–6) give two sum rules with one 9- $j$  symbol

$$\sum_X [X] \begin{Bmatrix} a & b & e \\ c & d & f \\ e & f & X \end{Bmatrix} = \frac{\delta_{bc}}{[b]} \delta(a, b, e) \delta(b, d, f)$$

$$\sum_X (-1)^{a+b+c+d-X} [X] \begin{Bmatrix} a & b & e \\ c & d & f \\ f & e & X \end{Bmatrix} = \frac{\delta_{ad}}{[a]} \delta(d, b, e) \delta(a, c, f).$$

### Sum rules with two 3- $j$ , 6- $j$ and/or 9- $j$ symbols:

- **Orthogonality of two 3- $j$  symbols:** Two orthogonality relations for 3- $j$  symbols are given by Rotenberg *et al.* (1959, Eqs. 1.13–14) as well as by Varshalovich *et al.* (1988, Eqs. 12.1.3–4)

$$\sum_{j_3 m_3} (2j_3 + 1) \begin{pmatrix} j_1 & j_2 & j_3 \\ m_1 & m_2 & m_3 \end{pmatrix} \begin{pmatrix} j_1 & j_2 & j_3 \\ m'_1 & m'_2 & m_3 \end{pmatrix} = \delta(m_1, m'_1) \delta(m_2, m'_2)$$

$$\sum_{m_1 m_2} \begin{pmatrix} j_1 & j_2 & j_3 \\ m_1 & m_2 & m_3 \end{pmatrix} \begin{pmatrix} j_1 & j_2 & j'_3 \\ m_1 & m_2 & m'_3 \end{pmatrix} = \frac{\delta(j_3, j'_3) \delta(m_3, m'_3)}{(2j_3 + 1)} \delta(j_1, j_2, j_3).$$

- **Orthogonality of two 6- $j$  symbols:** An orthogonality relation for the Wigner 6- $j$  symbols is given by Rotenberg *et al.* (1959, Eq. 2.6) as well as by Varshalovich *et al.* (1988, Eq. 12.2.7)

$$\sum_X [X] \begin{Bmatrix} a & b & X \\ c & d & p \end{Bmatrix} \begin{Bmatrix} c & d & X \\ a & b & q \end{Bmatrix} = \frac{\delta_{pq}}{[p]} \delta(a, d, p) \delta(b, c, p).$$

- **Sum rules with two 6- $j$  symbols:** Rotenberg *et al.* (1959, Eq. 2.7) as well as Varshalovich *et al.* (1988, Eq. 12.2.8) gives a sum rule for two 6- $j$  symbols

$$\sum_X (-1)^{p+q+X} [X] \begin{Bmatrix} a & b & X \\ c & d & p \end{Bmatrix} \begin{Bmatrix} c & d & X \\ b & a & q \end{Bmatrix} = \begin{Bmatrix} c & a & q \\ d & b & p \end{Bmatrix}.$$

- **Sum rules with two 6- $j$  symbols:** Varshalovich *et al.* (1988, Eq. 12.2.15) show a sum rule including a triple summation over a 6- $j$  symbol

$$\sum_{XYZ} [X, Y, Z] \begin{Bmatrix} X & Y & Z \\ a & b & c \end{Bmatrix} \begin{Bmatrix} X & Y & Z \\ a & b & c \end{Bmatrix} = [a, b, c].$$

- **Sum rules with one 6- $j$  and one 9- $j$  symbol:** Varshalovich *et al.* (1988, Eqs. 12.2.9–10) show two sum rules for one 6- $j$  symbol and one

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9- $j$  symbol, and with  $R = a + b + c + d + e + f + p + q$

$$\sum_X [X] \begin{Bmatrix} a & f & X \\ d & q & e \\ p & c & b \end{Bmatrix} \begin{Bmatrix} a & f & X \\ e & b & s \end{Bmatrix} = (-1)^{2s} \begin{Bmatrix} a & b & s \\ c & d & p \end{Bmatrix} \begin{Bmatrix} c & d & s \\ e & f & q \end{Bmatrix}$$

$$\sum_X (-1)^{R+X} [X] \begin{Bmatrix} a & f & X \\ d & q & e \\ p & c & b \end{Bmatrix} \begin{Bmatrix} a & f & X \\ b & e & s \end{Bmatrix} = (-1)^{2s} \begin{Bmatrix} p & q & s \\ e & a & d \end{Bmatrix} \begin{Bmatrix} p & q & s \\ f & b & c \end{Bmatrix}.$$

#### Sum rules for three Wigner 3- $j$ , 6- $j$ and/or 9- $j$ symbols:

- **Sum rules with three 3- $j$  symbols:** One such rule is given by Rotenberg *et al.* (1959, Eq. 2.20) as well as by Varshalovich *et al.* (1988, Eq. 2.1.6), and with  $S = l_1 + l_2 + l_3 + n_1 + n_2 + n_3$

$$\sum_{n_1 n_2 n_3} (-1)^S \begin{pmatrix} j_1 & l_2 & l_3 \\ m_1 & n_2 & -n_3 \end{pmatrix} \begin{pmatrix} l_1 & j_2 & l_3 \\ -n_1 & m_2 & n_3 \end{pmatrix} \begin{pmatrix} l_1 & l_2 & j_3 \\ n_1 & -n_2 & m_3 \end{pmatrix} = \begin{pmatrix} j_1 & j_2 & j_3 \\ m_1 & m_2 & m_3 \end{pmatrix} \begin{Bmatrix} j_1 & j_2 & j_3 \\ l_1 & l_2 & l_3 \end{Bmatrix}.$$

- **Sum rules with two 3- $j$  symbols and one 6- $j$  symbol:** One such rule is given by Rotenberg *et al.* (1959, Eq. 2.19) as well as by Varshalovich *et al.* (1988, Eq. 12.1.5)

$$\sum_{l_3 n_3} (-1)^{j_3 + l_3 + m_1 + n_1} [l_3] \begin{Bmatrix} j_1 & j_2 & j_3 \\ l_1 & l_2 & l_3 \end{Bmatrix} \begin{pmatrix} l_1 & j_2 & l_3 \\ n_1 & m_2 & n_3 \end{pmatrix} \begin{pmatrix} j_1 & l_2 & l_3 \\ m_1 & n_2 & -n_3 \end{pmatrix} = \sum_{m_3} \begin{pmatrix} j_1 & j_2 & j_3 \\ m_1 & m_2 & m_3 \end{pmatrix} \begin{pmatrix} l_1 & l_2 & j_3 \\ n_1 & n_2 & -m_3 \end{pmatrix}.$$

- **Sum rules with three 6- $j$  symbols:** There is a sum rule for three 6- $j$  symbols from Biedenharn and Elliot; see also Rotenberg *et al.* (1959, Eq. 2.8) or Varshalovich *et al.* (1988, Eq. 12.2.18), and with  $R = a + b + c + d + e + f + p + q + r$

$$\sum_X (-1)^{R+X} [X] \begin{Bmatrix} a & b & X \\ c & d & p \end{Bmatrix} \begin{Bmatrix} c & d & X \\ e & f & q \end{Bmatrix} \begin{Bmatrix} e & f & X \\ b & a & r \end{Bmatrix} = \begin{Bmatrix} p & q & r \\ e & a & d \end{Bmatrix} \begin{Bmatrix} p & q & r \\ f & b & c \end{Bmatrix}.$$

- **Sum rules with three 6- $j$  symbols:** Another sum rules corresponds to the representation of the 9- $j$  symbol in terms of 6- $j$  symbols; see Rotenberg *et al.* (1959, Eq. 3.1) and Varshalovich *et al.* (1988, Eq. 12.2.19)

$$\sum_X (-1)^{2X} [X] \begin{Bmatrix} a & b & X \\ c & d & p \end{Bmatrix} \begin{Bmatrix} c & d & X \\ e & f & q \end{Bmatrix} \begin{Bmatrix} e & f & X \\ a & b & r \end{Bmatrix} = \begin{Bmatrix} a & f & r \\ d & q & e \\ p & c & b \end{Bmatrix}.$$

### Sum rules with four Wigner 3- $j$ , 6- $j$ and/or 9- $j$ symbols:

- **Sum rules with three 3- $j$  symbols:** Rotenberg *et al.* (1959, Eq. 2.18) as well as Varshalovich *et al.* (1988, Eq. 12.1.8) display the representation of a 6- $j$  symbol in terms of four 3- $j$  symbols, and with  $S = l_1 + l_2 + l_3 + n_1 + n_2 + n_3$

$$\begin{aligned} \begin{Bmatrix} j_1 & j_2 & j_3 \\ l_1 & l_2 & l_3 \end{Bmatrix} &= \sum_{m_1 m_2 m_3 n_1 n_2 n_3} (-1)^S \begin{pmatrix} j_1 & j_2 & j_3 \\ m_1 & m_2 & m_3 \end{pmatrix} \begin{pmatrix} j_1 & l_2 & l_3 \\ m_1 & n_2 & -n_3 \end{pmatrix} \begin{pmatrix} l_1 & j_2 & l_3 \\ -n_1 & m_2 & n_3 \end{pmatrix} \begin{pmatrix} l_1 & l_2 & j_3 \\ n_1 & -n_2 & m_3 \end{pmatrix} \\ &= \sum_{m_1 m_2 n_1 n_2 n_3} (-1)^S [j_3] \begin{pmatrix} j_1 & j_2 & j_3 \\ m_1 & m_2 & m_3 \end{pmatrix} \begin{pmatrix} j_1 & l_2 & l_3 \\ m_1 & n_2 & -n_3 \end{pmatrix} \begin{pmatrix} l_1 & j_2 & l_3 \\ -n_1 & m_2 & n_3 \end{pmatrix} \begin{pmatrix} l_1 & l_2 & j_3 \\ n_1 & -n_2 & m_3 \end{pmatrix}. \end{aligned}$$

- Another sum rule for four 3- $j$  symbols is given by Varshalovich *et al.* (1988, Eq. 12.1.9).

$$\begin{aligned} \sum_{n_r n_p n_q n_s} (-1)^{p-n_p+q-n_q+r-n_r+s-n_s} \begin{pmatrix} p & a & q \\ n_p & n_a & -n_q \end{pmatrix} \begin{pmatrix} q & b & r \\ n_q & n_b & -n_r \end{pmatrix} \begin{pmatrix} r & s & p \\ n_r & n_s & -n_p \end{pmatrix} \begin{pmatrix} s & c & d \\ -n_s & n_c & n_d \end{pmatrix} \\ = \begin{Bmatrix} a & b & s \\ r & p & q \end{Bmatrix} \sum_{n_s} \begin{pmatrix} a & s & b \\ n_a & n_s & n_b \end{pmatrix} \begin{pmatrix} d & s & c \\ n_d & -n_s & n_c \end{pmatrix}. \end{aligned}$$

- A useful relation between 3- $j$  and 9- $j$  symbols is due to de-Shalit; cf. Rotenberg *et al.* (1959, Eq. 3.21) as well as Varshalovich *et al.* (1988, Eq. 12.1.11)

$$\sum_{MNRS} \begin{pmatrix} a & e & g \\ n_a & M & N \end{pmatrix} \begin{pmatrix} b & f & h \\ n_b & R & S \end{pmatrix} \begin{pmatrix} e & f & c \\ M & R & n_e \end{pmatrix} \begin{pmatrix} g & h & d \\ N & S & n_x \end{pmatrix} = \sum_{XY} [X] \begin{pmatrix} a & b & X \\ n_a & n_b & Y \end{pmatrix} \begin{pmatrix} X & c & d \\ Y & n_e & n_x \end{pmatrix} \begin{Bmatrix} a & b & X \\ e & f & c \\ g & h & d \end{Bmatrix}.$$

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Only this (last) sum rule for four Wigner 3- $j$  symbols need to be implemented explicitly because the other two above can be simplified in a two-step evaluation using other sum rules.

#### Sum rules with five Wigner 3- $j$ , 6- $j$ and/or 9- $j$ symbols:

➤ **Sum rules with five 3- $j$  symbols:** Varshalovich *et al.* (1988, Eqs. 12.1.12–14) give three sum rules involving products of five 3- $j$  symbols.

$$\begin{aligned}
 & \sum_{n_p n_q n_r n_s n_t n_u n_v} (-1)^{p-n_p+q-n_q+r-n_r+s-n_s+t-n_t+u-n_u+v-n_v} \begin{pmatrix} r & t & q \\ n_r & n_t & -n_q \end{pmatrix} \begin{pmatrix} q & p & u \\ n_q & n_p & -n_u \end{pmatrix} \begin{pmatrix} u & a & v \\ n_u & n_a & -n_v \end{pmatrix} \begin{pmatrix} v & s & r \\ n_v & n_s & -n_r \end{pmatrix} \begin{pmatrix} s & p & t \\ -n_s & -n_p & -n_t \end{pmatrix} \\
 &= (-1)^{2u} \frac{\delta_{uv} \delta_{a0} \delta_{na0}}{[u]^{1/2}} \left\{ \begin{matrix} q & p & u \\ s & r & t \end{matrix} \right\} \\
 & \sum_{n_p n_q n_r n_s n_t n_u} (-1)^{p-n_p+q-n_q+r-n_r+s-n_s+t-n_t+u-n_u} \begin{pmatrix} p & a & q \\ n_p & n_a & -n_q \end{pmatrix} \begin{pmatrix} q & r & t \\ n_q & n_r & -n_t \end{pmatrix} \begin{pmatrix} t & s & p \\ n_t & -n_s & -n_p \end{pmatrix} \begin{pmatrix} s & c & u \\ n_s & n_c & -n_u \end{pmatrix} \begin{pmatrix} u & b & r \\ n_u & n_b & -n_r \end{pmatrix} \\
 &= \begin{pmatrix} a & b & c \\ n_a & n_b & n_c \end{pmatrix} \left\{ \begin{matrix} a & b & c \\ u & s & r \end{matrix} \right\} \left\{ \begin{matrix} a & p & q \\ t & r & s \end{matrix} \right\} \\
 & \sum_{n_p n_q n_r n_s n_t n_u} (-1)^{p-n_p+q-n_q+r-n_r+s-n_s+t-n_t+u-n_u} \begin{pmatrix} p & a & q \\ n_p & n_a & n_q \end{pmatrix} \begin{pmatrix} p & a & q \\ n_p & n_a & n_q \end{pmatrix} \begin{pmatrix} q & t & r \\ -n_q & -n_t & -n_r \end{pmatrix} \begin{pmatrix} r & b & s \\ n_r & n_b & n_s \end{pmatrix} \\
 & \times \begin{pmatrix} s & p & u \\ -n_s & -n_p & -n_u \end{pmatrix} \begin{pmatrix} u & c & t \\ n_u & n_c & n_t \end{pmatrix} = (-1)^{r+b+s} \begin{pmatrix} a & b & c \\ -n_a & -n_b & -n_c \end{pmatrix} \left\{ \begin{matrix} a & b & c \\ p & s & u \\ q & r & t \end{matrix} \right\}
 \end{aligned}$$



### Sum rules with six Wigner 3- $j$ , 6- $j$ and/or 9- $j$ symbols:

- **Sum rules with six 3- $j$  symbols:** A sum rule corresponds to the representation of the 9- $j$  symbol in terms of six 3- $j$  symbols; see Rotenberg *et al.* (1959, Eq. 3.1).

$$\begin{pmatrix} a & b & c \\ d & e & f \\ g & h & j \end{pmatrix} = \sum_{ABCDEFGHIJ} \begin{pmatrix} a & b & c \\ A & B & C \end{pmatrix} \begin{pmatrix} d & e & f \\ D & E & F \end{pmatrix} \begin{pmatrix} g & h & j \\ G & H & J \end{pmatrix} \begin{pmatrix} a & d & g \\ A & D & G \end{pmatrix} \begin{pmatrix} b & e & h \\ B & E & H \end{pmatrix} \begin{pmatrix} c & f & j \\ C & F & J \end{pmatrix}.$$

### Graphical loop rules for the Wigner 3- $j$ symbols:

- **Loop rules for one Wigner 3- $j$  symbol:** Several loop rules ( $n$ -loops) can be implemented separately for algebraic evaluations. A loop for one Wigner 3- $j$  symbol is given by Varshalovich *et al.* (1988, Eq. 12.1.2)

$$\sum_m (-1)^m \begin{pmatrix} j & j & J \\ m & -m & M \end{pmatrix} = (-1)^j [j] \delta_{J0} \delta_{M0}.$$

- **Loop rules for two Wigner 3- $j$  symbol:** Varshalovich *et al.* (1988, Eq. 12.1.3) also displays a loop for two Wigner 3- $j$  symbols

$$\sum_{m_1, m_2} \begin{pmatrix} j_1 & j_2 & j_3 \\ m_1 & m_2 & m_3 \end{pmatrix} \begin{pmatrix} j_1 & j_2 & j'_3 \\ m_1 & m_2 & m'_3 \end{pmatrix} = \frac{1}{2j_3 + 1} \delta_{j_3, j'_3} \delta_{m_3, m'_3} \delta(j_1, j_2, j_3).$$

- **Loop rules for three Wigner 3- $j$  symbol:** Varshalovich *et al.* (1988, Eq. 12.1.6) also displays a loop for three Wigner 3- $j$  symbols.

$$\begin{aligned} \sum_{m_1, m_2, m_3} (-1)^{-m_1 - m_2 - m_3} \begin{pmatrix} j_1 & j_4 & j_2 \\ m_1 & m_4 & -m_2 \end{pmatrix} \begin{pmatrix} j_2 & j_5 & j_3 \\ m_2 & m_5 & -m_3 \end{pmatrix} \begin{pmatrix} j_3 & j_6 & j_1 \\ m_3 & m_6 & -m_1 \end{pmatrix} \\ = (-1)^{-j_1 - j_2 - j_3 + j_4 + j_5 + j_6} \begin{pmatrix} j_4 & j_5 & j_6 \\ m_4 & m_5 & m_6 \end{pmatrix} \begin{Bmatrix} j_4 & j_5 & j_6 \\ j_3 & j_1 & j_2 \end{Bmatrix}. \end{aligned}$$

## 13.4. Symbolic evaluations of Racah algebra expressions not yet considered in JAC

### 13.4.a. Orthogonality, completeness and integral representation of the Wigner rotation matrices:

#### Unitarity and completeness of the rotation matrices:

- **Unitarity condition:** This condition for the Wigner rotation matrices can be written either in terms of the Wigner  $D$ -function (Varshalovich *et al.*, 1988, Eqs. 4.1.6) or, equivalently, by means of the Wigner rotation matrices

$$\sum_{p=-j}^j d_{pq}^j(\beta) d_{pr}^j(\beta) = \delta_{qr}, \quad \sum_{q=-j}^j d_{pq}^j(\beta) d_{rq}^j(\beta) = \delta_{pr}.$$

- **Orthogonality and normalization condition for the rotation matrices:** The orthogonality and normalization condition in terms of  $d_{pq}^j(\beta)$  is shown by Varshalovich *et al.* (1988, Eqs. 4.10.6 and 4.11.7)

$$\int_0^{2\pi} d\beta \sin \beta d_{pq}^j(\beta) d_{pq}^k(\beta) = \frac{2}{2j+1} \delta_{jk}$$

$$\int_0^{2\pi} d\alpha e^{-i(p_1-p_2)\alpha} \int_0^{2\pi} d\gamma e^{-i(q_1-q_2)\gamma} \int_0^\pi d\beta \sin \beta d_{p_1 q_1}^{j_1}(\beta) d_{p_2 q_2}^{j_2}(\beta) = \frac{8\pi^2}{2j_2+1} \delta_{j_1 j_2} \delta_{p_1 p_2} \delta_{q_1 q_2}.$$

- **Completeness condition of the rotation matrices:** cf. Varshalovich *et al.* (1988, Eq. 4.10.7) with double domain proportional to  $16\pi^2$

$$\sum_{j=0, \frac{1}{2}, 1, \dots}^{\infty} \sum_{p=-j}^{+j} \sum_{q=-j}^{+j} (2j+1) e^{+i[(\alpha_1-\alpha_2)p+(\gamma_1-\gamma_2)q]} d_{pq}^j(\beta_1) d_{pq}^j(\beta_2) = 16\pi^2 \delta(\alpha_1 - \alpha_2) \delta(\cos \beta_1 - \cos \beta_2) \delta(\gamma_1 - \gamma_2).$$

### Integrals involving Wigner $D_{pq}^j$ functions or rotation matrices $d_{pq}^j$ :

- The Wigner rotation matrices fulfill various integration rules as well as integration with invariant summation rules, and where some of these integrals involve products of more than one Wigner  $D$ -function. Whenever possible, we express the Clebsch-Gordan series as a product of two Wigner 3- $j$  symbols. Most of the following integration rules require solid angle integration either over the standard domain

$$\int d\Omega \equiv \int_0^{2\pi} d\alpha \int_0^{2\pi} d\gamma \int_0^\pi d\beta \sin(\beta)$$

or over an extended domain with

$$\int d\Omega \equiv \int_0^{4\pi} d\alpha \int_0^{2\pi} d\gamma \int_0^\pi d\beta \sin(\beta) \quad \text{or} \quad \int d\Omega \equiv \int_0^{2\pi} d\alpha \int_0^{4\pi} d\gamma \int_0^\pi d\beta \sin(\beta) .$$

$$0 \leq \alpha \leq 4\pi, \quad 0 \leq \beta \leq \pi, \quad 0 \leq \gamma \leq 2\pi, \quad \text{or} \quad 0 \leq \alpha \leq 2\pi, \quad 0 \leq \beta \leq \pi, \quad 0 \leq \gamma \leq 4\pi,$$

For this redefinition of the domain of integration, the factors  $8\pi^2$  must typically be replaced by  $16\pi^2$ .

- **Integrals with one  $D_{pq}^j(\alpha, \beta, \gamma)$ :** Varshalovich *et al.* (1988, Eqs. 4.11.1 and 4.11.6) displays the following normalization integral as well as the corresponding special case for  $\alpha = \gamma = 0$  and/or  $p = q = 0$

$$\int d\Omega D_{pq}^j(\alpha, \beta, \gamma) = \int d\Omega e^{-i(p\alpha + q\gamma)} d_{pq}^j(\beta) = 8\pi^2 \delta_{j0} \delta_{p0} \delta_{q0}, \quad \text{if } j \text{ is integer}$$

$$\int_0^\pi d\beta \sin \beta d_{00}^j(\beta) = 2 \delta_{j0}.$$

### 13.4.b. Sum rules for spherical harmonics

#### Products of two and more spherical harmonics:

➤ Sum rules for two spherical harmonics:

$$\sum_{m=-\ell}^{\ell} (-1)^m Y_{\ell m}(\vartheta, \varphi) Y_{\ell, -m}(\vartheta, \varphi) = \sum_{m=-\ell}^{\ell} |Y_{\ell m}(\vartheta, \varphi)|^2 = \frac{2\ell + 1}{4\pi}$$

$$\sum_{l=0}^{\infty} \sum_{m=-\ell}^{\ell} (-1)^m m Y_{\ell, -m}(\vartheta, \varphi) Y_{\ell m}(\vartheta, \varphi) = 0$$

$$\sum_{l=0}^{\infty} \sum_{m=-\ell}^{\ell} (-1)^m m^2 Y_{\ell, -m}(\vartheta, \varphi) Y_{\ell m}(\vartheta, \varphi) = \frac{l(l+1)(2\ell+1)}{8\pi} \sin^2 \vartheta$$

$$\sum_{l=0}^{\infty} \sum_{m=-\ell}^{\ell} (-1)^m Y_{\ell, -m}(\vartheta, \varphi) Y_{\ell m}(\vartheta', \varphi') = \delta(\varphi - \varphi') \delta(\cos \vartheta - \cos \vartheta') = \delta(\varphi - \varphi') \frac{\delta(\vartheta - \vartheta')}{|\sin \vartheta|}.$$

➤ The Clebsch-Gordan expansion of two spherical harmonics: A product of two spherical harmonics with the same angles can be written as shown by Varshalovich *et al.* (1988, Eq. 5.6.9)

$$Y_{\ell_1 m_1}(\vartheta, \varphi) Y_{\ell_2 m_2}(\vartheta, \varphi) = \sum_{L, M} \left[ \frac{(2\ell_1 + 1)(2\ell_2 + 1)}{4\pi(2L + 1)} \right]^{1/2} \langle \ell_1 0, \ell_2 0 | L 0 \rangle \langle \ell_1 m_1, \ell_2 m_2 | L M \rangle Y_{LM}(\vartheta, \varphi)$$

➤ The Clebsch-Gordan expansion of  $n$  spherical harmonics: Using this relation  $n - 1$  times, one can easily derive a more general expression

$$\prod_{i=1}^n Y_{\ell_i m_i}(\vartheta, \varphi) = \sum_{L_n} B_{L_n} Y_{L_n M_n}(\vartheta, \varphi)$$

$$B_{L_n} = \frac{1}{\sqrt{(4\pi)^{n-1}(2L_n + 1)}} \sum_{\substack{L_1, L_2, \dots, L_n \\ M_1, M_2, \dots, M_n}} \prod_{i=1}^n \sqrt{2\ell_i + 1} \langle L_{i-1} 0, \ell_i 0 | L_i 0 \rangle \langle L_{i-1} M_{i-1}, \ell_i m_i | L_i M_i \rangle.$$

- **Irreducible tensor product of two spherical harmonics:** An irreducible tensor product of two spherical harmonics with the same angles can be written as (Varshalovich *et al.* 1988, Eq. 5.6.14)

$$\{Y_{\ell_1}(\vartheta, \varphi) \otimes Y_{\ell_2}(\vartheta, \varphi)\}_{LM} = \left[ \frac{(2\ell_1 + 1)(2\ell_2 + 1)}{4\pi(2L + 1)} \right]^{1/2} \langle \ell_1 0, \ell_2 0 | L 0 \rangle Y_{LM}(\vartheta, \varphi)$$

- **Irreducible tensor product of  $n$  spherical harmonics:** Using this relation  $n - 1$  times, one finds a more general expression (Varshalovich *et al.* 1988, Eq. 5.6.16):

$$\left\{ \cdot \left\{ \{Y_{\ell_1}(\vartheta, \varphi) \otimes Y_{\ell_2}(\vartheta, \varphi)\}_{L_2} \otimes Y_{\ell_3}(\vartheta, \varphi) \right\}_{L_3} \cdot \otimes Y_{\ell_n}(\vartheta, \varphi) \right\}_{L_n M_n} = \frac{1}{\sqrt{(4\pi)^{n-1}(2L_n + 1)}} \prod_{i=1}^n \left[ \sqrt{2\ell_i + 1} \langle L_{i-1} 0, \ell_i 0 | L_i 0 \rangle \right] Y_{L_n M_n}(\vartheta, \varphi).$$

- **Multipole expansions by using spherical harmonics:** The spherical harmonics are not only orthonormal to each other but they also form a complete set of functions for the expansion of any function  $f(\Theta, \Phi)$  that satisfies the given condition:

$$\int_0^{2\pi} d\Phi \int_0^\pi d\Theta \sin \Theta |f(\Theta, \Phi)|^2 < \infty : \quad f(\Theta, \Phi) = \sum_{l=0}^{\infty} \sum_{m=-l}^l \tilde{f}_{lm} Y_{lm}(\Theta, \Phi), \quad \tilde{f}_{lm} \equiv \int_0^{2\pi} d\varphi \int_0^\pi d\vartheta \sin \vartheta Y_{lm}^*(\vartheta, \varphi) f(\vartheta, \varphi).$$

### 13.4.c. Integrals involving spherical harmonics

#### Integrals involving spherical harmonics:

- **Orthonormality of the spherical harmonics:**

$$\int_0^{2\pi} d\varphi \int_0^\pi d\vartheta \sin \vartheta Y_{\ell m}^*(\vartheta, \varphi) Y_{\ell' m'}(\vartheta, \varphi) = \delta_{\ell \ell'} \delta_{m m'}.$$

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➤ **Further integrals over  $4\pi$ :** Cf. Varshalovich *et al.* (1988, Eqs. 5.9.1 and 5.9.5):

$$\int_0^{2\pi} d\varphi \int_0^\pi d\vartheta \sin \vartheta Y_{\ell m}(\vartheta, \varphi) = \sqrt{4\pi} \delta_{\ell 0} \delta_{m 0}$$

$$\int_0^{2\pi} d\varphi \int_0^\pi d\vartheta \sin \vartheta Y_{\ell_1 m_1}(\vartheta, \varphi) Y_{\ell_2 m_2}(\vartheta, \varphi) Y_{\ell_3 m_3}(\vartheta, \varphi) = \sqrt{\frac{(2\ell_1 + 1)(2\ell_2 + 1)(2\ell_3 + 1)}{4\pi}} \begin{pmatrix} \ell_1 & \ell_2 & \ell_3 \\ 0 & 0 & 0 \end{pmatrix} \begin{pmatrix} \ell_1 & \ell_2 & \ell_3 \\ m_1 & m_2 & m_3 \end{pmatrix}.$$

### Derivatives of spherical harmonics:

➤ **Derivatives of spherical harmonics:** Cf. Varshalovich *et al.* (1988, Eq. 5.8.5) and with the associated Legendre polynomials  $P_l^m(\cos \vartheta)$

$$\frac{d}{dx} \int_{g(x)}^{h(x)} dy f(y) = \frac{dh(x)}{dx} f(h(x)) - \frac{dg(x)}{dx} f(g(x)) \quad \text{chain rule}$$

$$\frac{d}{d\vartheta} Y_{\ell m}(\vartheta, \varphi) = \frac{1}{2} \sqrt{\ell(\ell+1) - m(m+1)} Y_{\ell, m+1}(\vartheta, \varphi) e^{-i\varphi} - \frac{1}{2} \sqrt{\ell(\ell+1) - m(m-1)} Y_{\ell, m-1}(\vartheta, \varphi) e^{i\varphi}.$$

$$\frac{d}{d\varphi} Y_{\ell m}(\vartheta, \varphi) = i m Y_{\ell m}(\vartheta, \varphi).$$

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