

Oh

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Tous les fichiers svg :

Dedication

For those who hate looking at a template with 500 lines of code and an extra 300 lines commented out.

Declaration

Acknowledgements

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1.1.3.1 The physics

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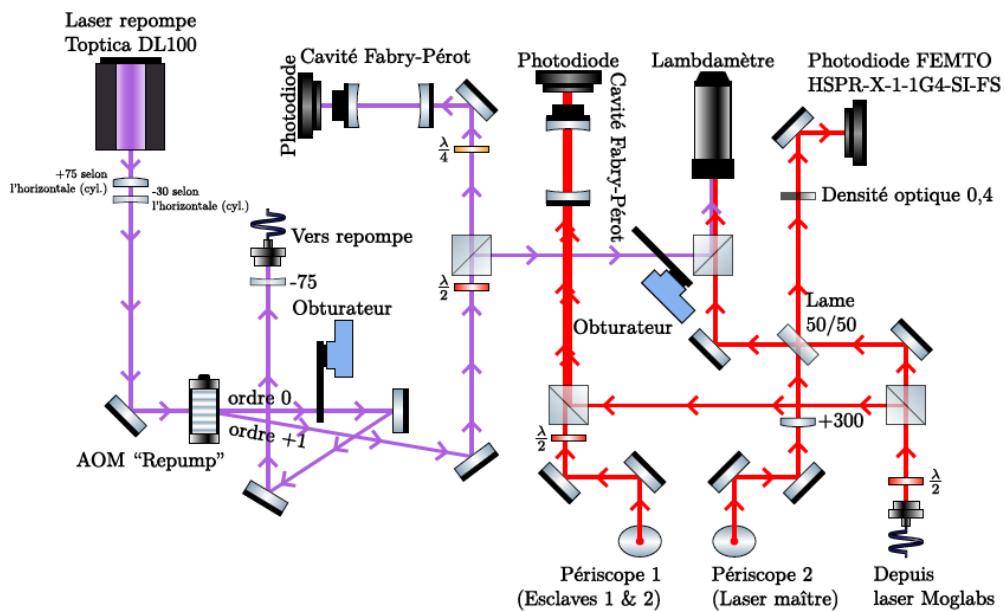


Figure 1.1: Caption

1.1.3.3 Comment on the hyperfine states (+boson 88)

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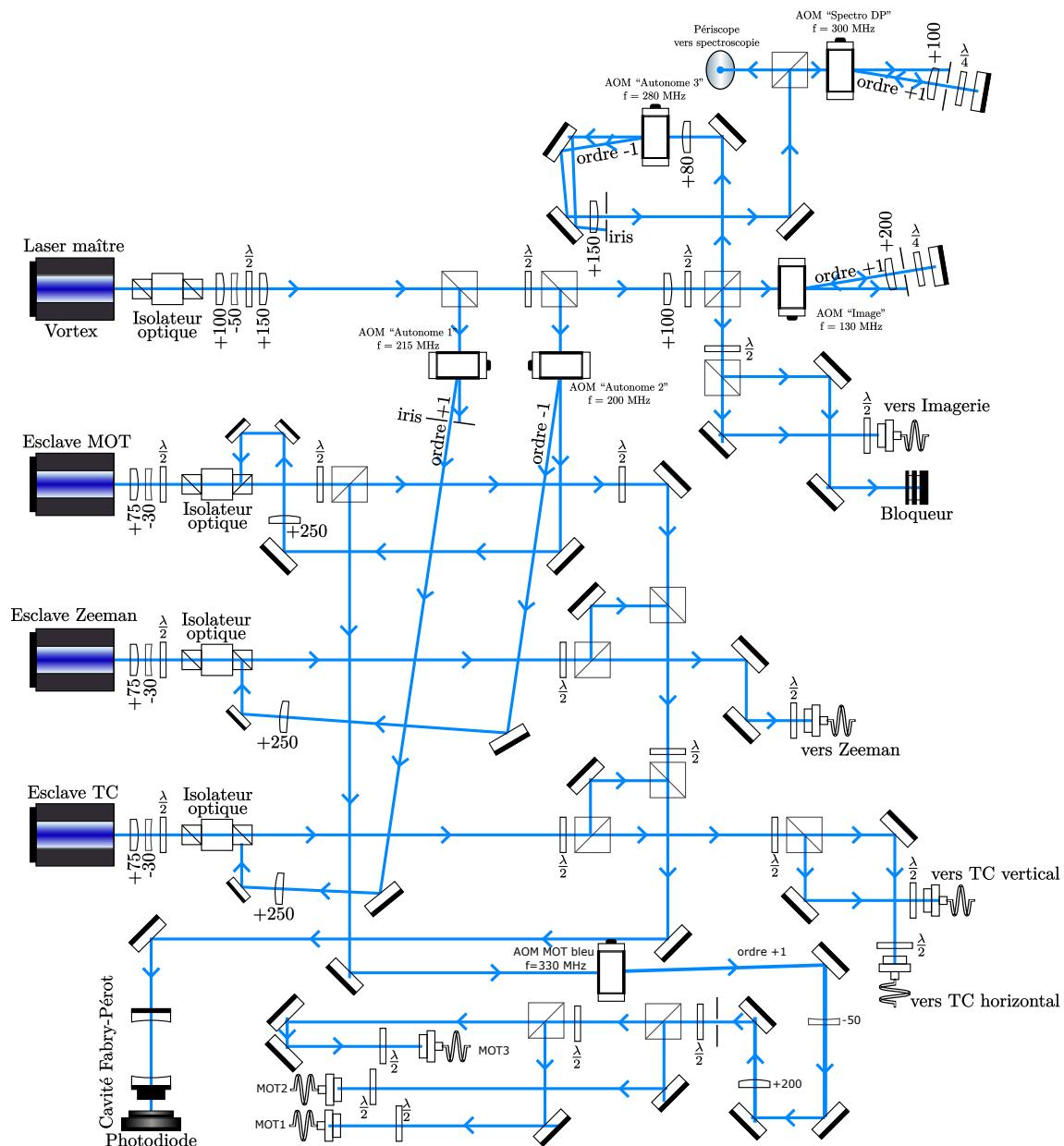


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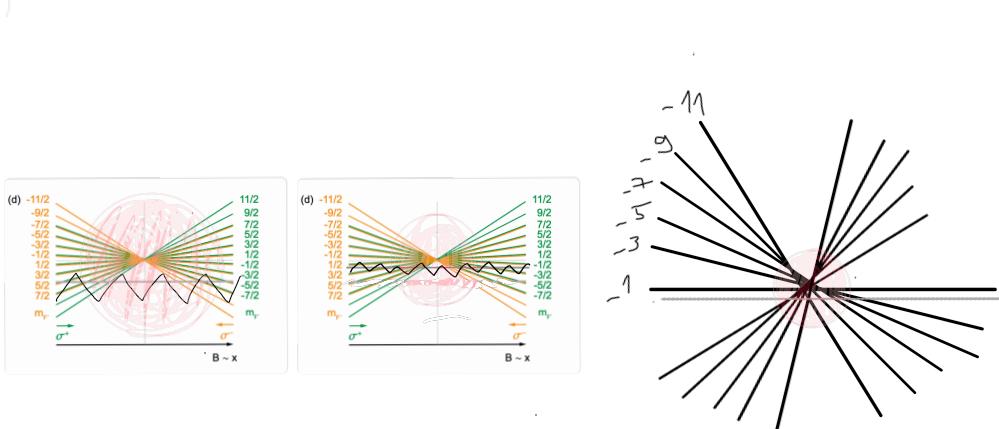


Figure 1.4: Caption

1.1.5.1 First step

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Need a stir because :



Figure 1.5: Caption

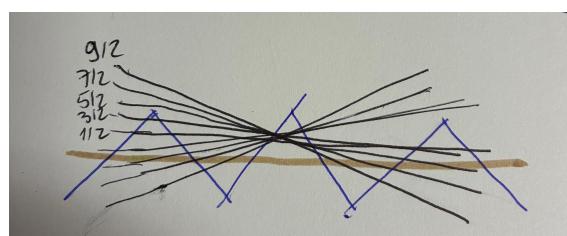


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1.1.7 Narrow MOT

cf p.43 S.Stellmer thesis

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Include images of the cloud for different I and detuning ?

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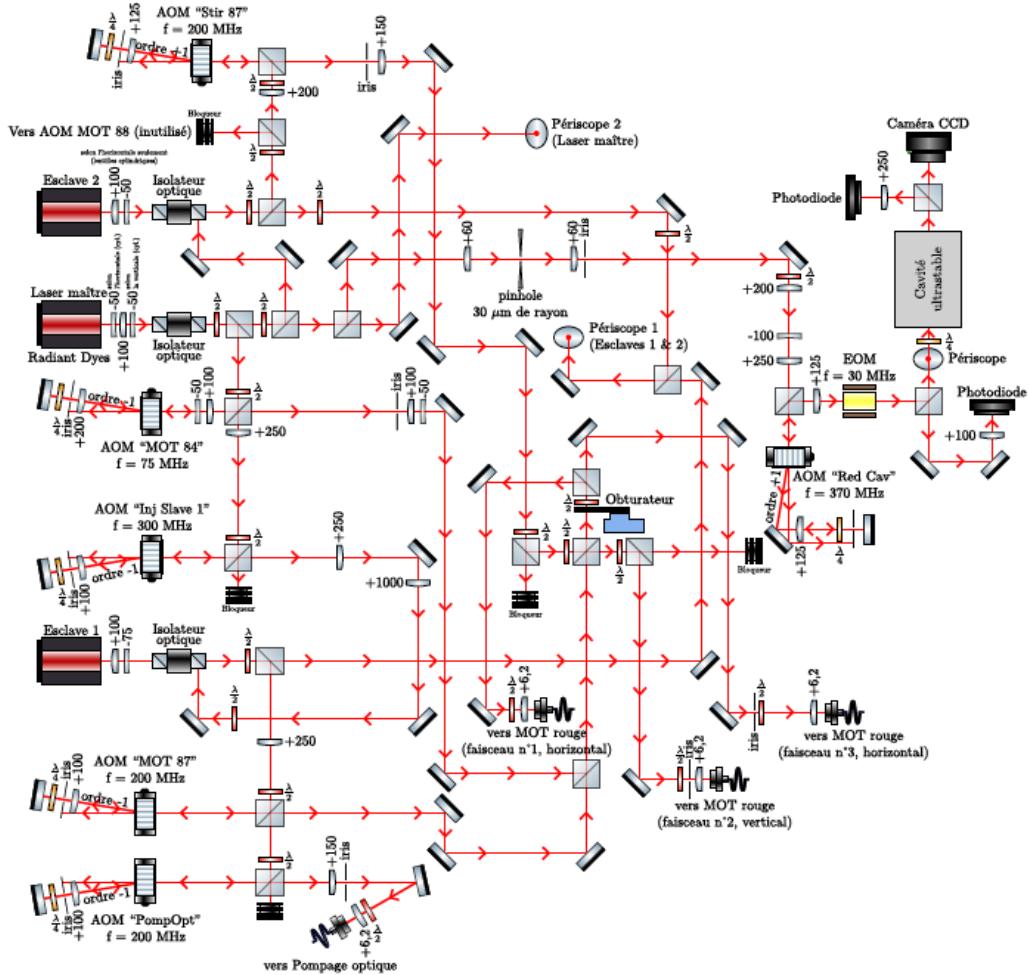


FIGURE 14 – Préparation des différents lasers rouges

Figure 1.7: Caption

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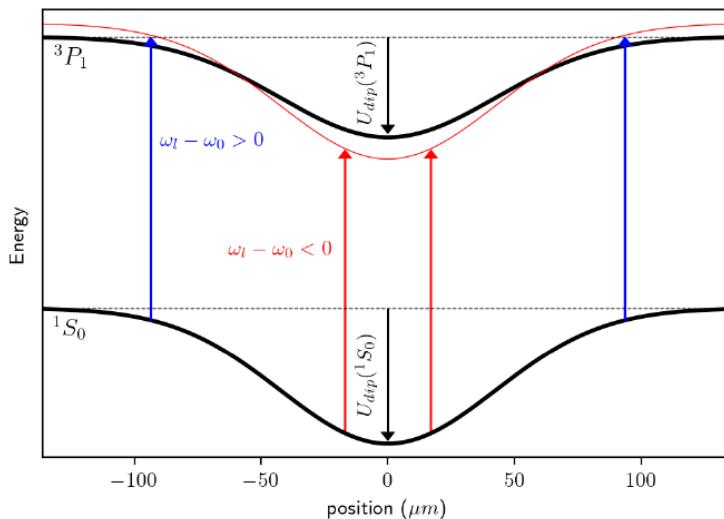


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Comment on the LS it does to each state

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Figure 1.10: Caption

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Ramsey interferometers on qudit

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blablablagtg

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2.2.1 Driving long coherence time Rabi oscillations

2.2.1.1 Rabi oscillations

Comment on what the FP could add as a longer coherence time of the qubit

2.2.1.2 Interferometer of $\text{su}(2)$ symmetry

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2.2.3 Measuring two non commuting observables

2.2.3.1 Principle

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e.g densité gaz, alimentation bobines - ζ comment faire mieux que les chiffres actuels)

Chapter 3

Engineering highly entangled system of photoassociated ^{87}Sr atoms

Engineering Dicke states

3.1 Introduction on photoassociation

3.1.1 Preamble on scattering theory

It is convenient in the case of the ^{87}Sr to approximate the electrons following the movement of the nuclei because of the mass ratio. In a Born-Oppenheimer picture we neglect the coupled terms between the kinetic energy of the electrons and the kinetic energy of the nuclei. In this manner we can study separately the radial and the angular part of the electronic movement.

As in the LANDAU-ZENER we will find the solutions of the 1D radial Schrödinger equation of the electrons

$$-\frac{\hbar^2}{2\mu} \left(\frac{1}{r^2} \frac{d^2\Psi_l}{dr^2} r^2 - \frac{l(l+1)}{r^2} + \hat{V}(r) \right) \Psi = E\Psi \quad (3.1)$$

- Ψ the eigenfunction - by approximating the scattering of the atoms as a sum of an incoming plane wave with a spherical scattered one - considered as plane wave at long distance 3.1

$$\Psi = e^{i\mathbf{k}\cdot\mathbf{z}} + f(\theta) \frac{e^{ikr}}{r} \quad (3.2)$$

$f(\theta)$ a function that will tell the absorption of the scattered wave if the collision is elastic or inelastic. θ is the angle between the vector wore by the incoming wave with the vector worn by the scattered wave.

The solution at asymptotes contains the scattering part of the phenomenon studied on this subsection. To solve this hamiltonian we will take the results of [1] that changes the solution basis to find a general solution of the hamiltonian 3.1.

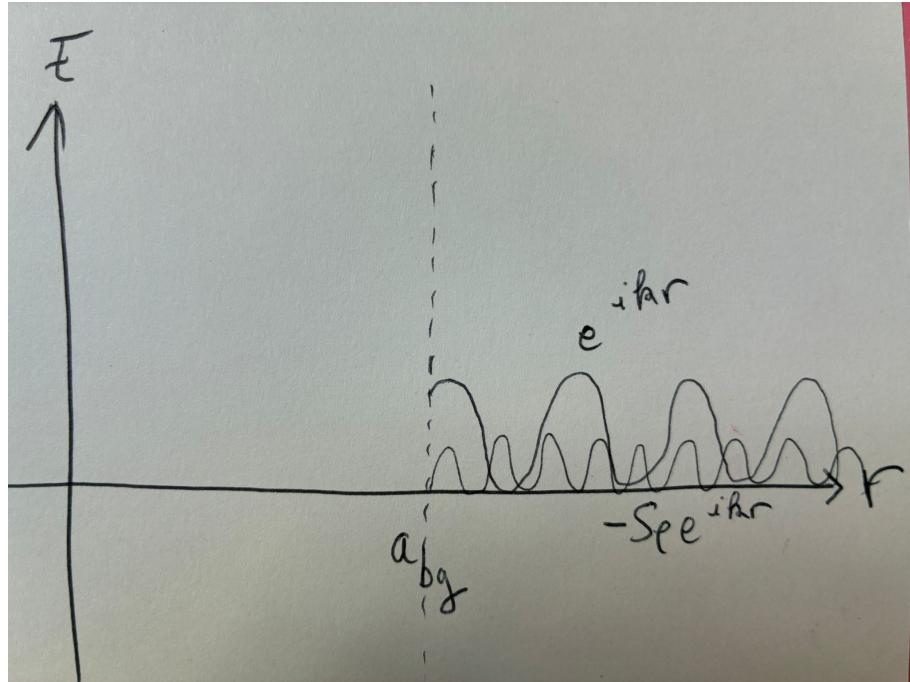


Figure 3.1: Caption

$$\Psi = \frac{a_l}{r} \sin(kr - \frac{l\pi}{2} + \delta_l) \quad (3.3)$$

with a_l a normalization factor and δ_l the phase acquired by the atoms at the distance at which they scatter.

Continuity of the wavefunction in every region should be respected. It gives

$$\delta_l \propto k^{2l+1} \quad (3.4)$$

For s-waves ($l = 0$), this express as

$$a = -\lim_{k \rightarrow 0} \delta_0/k \quad (3.5)$$

To understand how the atoms scatter we are going to express the scattered part of 3.2 as $\Psi - e^{ikz}$ in a spherical basis where the solutions are a sum on all the partial waves.

$$\Psi \simeq \frac{i}{2kr} \sum_{l=0}^{\infty} (2l+1) ((-1)^l e^{-ikr} - S_l e^{ikr}) \quad (3.6)$$

with $S_l = e^{2i\delta_l}$ the scattering amplitude.

For a **conservative hamiltonian** $|S_l| = 1$ meaning $\delta_l \in \mathbb{R}$. In this case the wavefunction is well a sum of incoming and outgoing waves of same amplitude.

The elastic cross section representing the surface where the scope of the potential is effective writes as

$$\sigma_{el} = \frac{\pi}{k^2} \sum_{l=0}^{\infty} (2l+1) |1 - S_l|^2 \quad (3.7)$$

But we see in equation 3.6 that $|S_l|$ needs to be < 1 to allow dissipation of the wavefunction. Defining δ_l as a complex number satisfies this condition.

In this case the inelastic collisions cross section reads as

$$\sigma_{inel} = \frac{\pi}{k^2} \sum_{l=0}^{\infty} (2l+1) (1 - |S_l|^2) \quad (3.8)$$

In ultracold gases, only s-wave atoms can collide because they do not have enough kinetic energy to pass the potential barrier of higher angular momentum states so $l = 0$ and the scattering amplitude is

$$S_0 = e^{2i\delta_0} \simeq 1 + 2i\delta_0 \quad (3.9)$$

By using the result of 3.5

$$1 + 2i\delta_0 \simeq 1 - 2ika \quad (3.10)$$

Because δ_0 is complex, a direct result of these equations is that we can write a general scattering length as

$$a = a_{bg} + ia'' \quad (3.11)$$

At low temperatures we can re-express equations 3.7 and 3.8 as

$$\sigma_{el} = 4\pi|a|^2 \quad (3.12)$$

$$\sigma_{inel} = 4\pi|a''|^2 \quad (3.13)$$

A usual case of the elastic collisions is atoms in the ground state. For inelasticity photoassociation is a good example of dissipative process where two atoms in the ground state are coupled to a molecular state by a laser.

3.1.2 What is photoassociation

Photoassociating two atoms consists in bounding two colliding atoms with light that occurs mostly with two-body and three body losses : it means adding a dissipative part on the interaction.

In a first step we can express the equation evolution of a pair of atoms without dissipation in the basis $\{| \psi_g \rangle, | \psi_e \rangle\}$ of a non coupled closed channel of free atoms $\psi_g = |g, g\rangle$ and their first molecular state virtual channel $\psi_e = |g, e\rangle$

$$i\hbar \frac{d}{dt} \begin{pmatrix} \psi_g \\ \psi_e \end{pmatrix} = \begin{pmatrix} -\frac{\hbar^2}{2\mu} \nabla + V_g & 0 \\ 0 & -\frac{\hbar^2}{2\mu} \nabla + V_e \end{pmatrix} \begin{pmatrix} \psi_g \\ \psi_e \end{pmatrix}. \quad (3.14)$$

In the virtual channel we see in 3.2 the presence of the internal states of the molecule. It is clear that the atoms get more bounded with the increase of the potential amplitude.

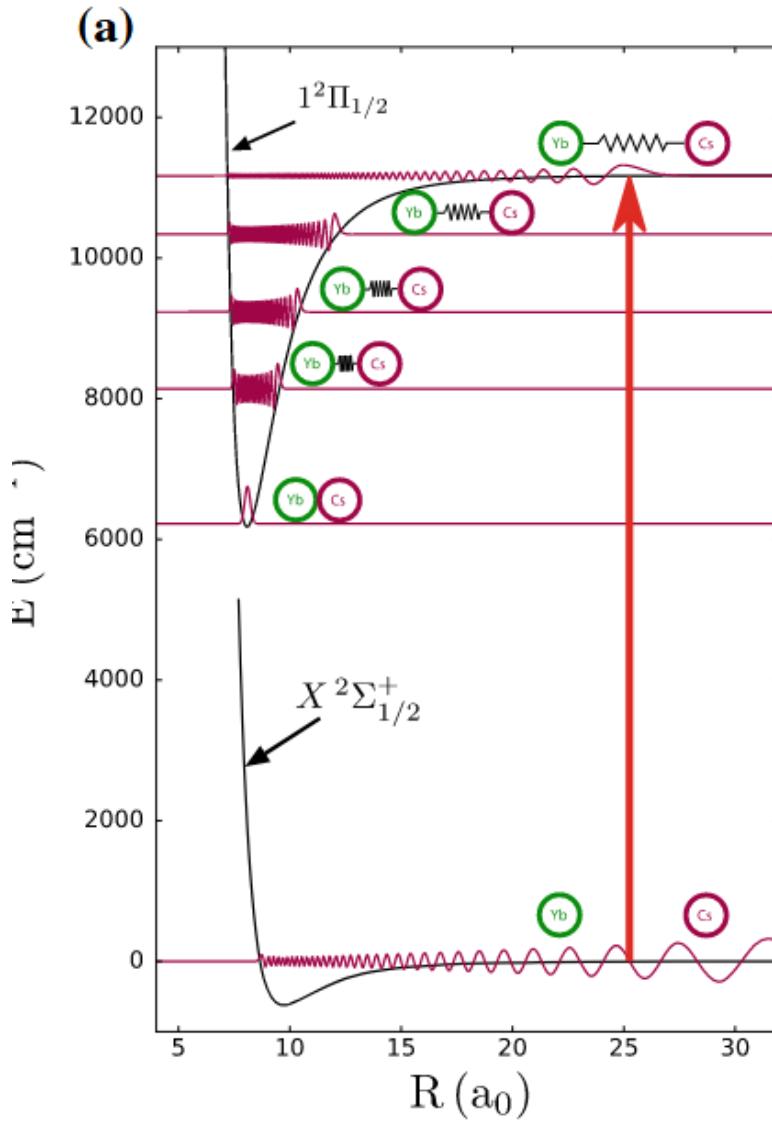


Figure 3.2: Caption

Then at resonance $\Delta = \omega_{at} - \omega_L = 0$ with a bounded molecular state a dissipative term appears on the evolution of the system because of the molecular spontaneous emission linewidth γ_{mol} .

The evolution of the atom number evolves

$$\frac{dN}{dt} = -\beta Nn - \Gamma N \quad (3.15)$$

where the 2-body losses express as $\beta \propto \gamma_{mol}$ (we are going to comment more in the next section). In comparison with [2]), $\beta = \frac{K_2 V_e}{2}$, where V_e is the effective volume of the cloud $(2\pi)^{3/2} \sigma_x \sigma_y \sigma_z$ for a density Gaussian distribution and Γ the one-body losses rate. This latter is limited by the vacuum quality and the lifetime of the atoms in the trap.

Adding the inelastic collisions of the PA process gives the following effective evolution equation

$$i\hbar \frac{d}{dt} \begin{pmatrix} \psi_g \\ \psi_e \end{pmatrix} = \begin{pmatrix} -\frac{\hbar^2}{2\mu} \nabla + V_g & \Omega/2 \\ \Omega/2 & -\frac{\hbar^2}{2\mu} \nabla + V_e + \hbar\Delta - i\frac{\gamma_{mol}}{2} \end{pmatrix} \begin{pmatrix} \psi_g \\ \psi_e \end{pmatrix}. \quad (3.16)$$

The eigenvector of the ground state $|\psi_g\rangle$ is then modified by this coupling. It also means that the phase δ_l (see 3.3) of the free state wavefunction is modified so the scattering length. The real part of a encodes two-body losses rate whereas its imaginary part changes the collisions properties. In this case we can write an effective scattering length enclosing the background scattering length a_{bg} standing for elastic collisions without any coupling, δa a modification of it because of the coupling and still the inelastic part of it

$$a = a_{bg} + \delta a + ia'' \quad (3.17)$$

In the case of spectroscopy we usually start with the least bounded state to have the atomic state reference. From that we sweep the atom frequency to address the molecular states red detuned from the atomic line. The quantity that will testify there is two-body losses is the atom number remaining in the trap after a photoassociation pulse. The molecule leaves the trap by acquiring kinetic energy from the binding energy.

PA is mostly used in Feschbach resonance field because determining the exact position of the vibrational states enables by changing the magnetic field to tune the scattering length of the atoms and thus the interactions in the system.

3.1.3 Parameters for PA study

In this subsection we are going to admit the expression of the two-body losses rate we introduced above but we will take a look on its dependencies.

This rate writes as

$$K_2 = \frac{1}{k} \frac{\frac{1}{4} \gamma_{stim} \gamma_{mol}}{\Delta^2 + (\frac{1}{2}(\gamma_{stim} + \gamma_{mol})^2)} \quad (3.18)$$

γ_{stim} is linear with the intensity of the PA laser, proportionnal to the 3D density n , and the probability to transit from ψ_g to ψ_e that we call **Franck-Condon factor**

$$F_{FC} = \left| \int_0^{\infty} \psi_g(r) \psi_e(r) dr \right|^2 \quad (3.19)$$

It depends directly with the overlap of the wavefunctions of the initial and final states. **The Condon radius** is the distance between the atoms for which this factor is maximum which also means that in a classical approach the atoms spend the most time at this position.

In the molecular community it is conventional to characterize the strength of the photoassociation rate with the **optical length**

$$K_2(I) = \frac{4\pi\hbar}{\mu} l_{opt}(I) \quad (3.20)$$

μ the reduced mass of the two atoms. From this formula we can say that the major experimental changes enabling the increase of the PA rate is the density of the cloud and the intensity if the PA beam as we are going to test in the experimental results.

Exciting an atom from a ground state to a single excited state is straightforward and quite easy to do precisely. It is not as easy for molecules because there exist a big number molecular potentials and internal states for each potential. These internal and external accesible states of the molecule are presented in the following via the introduction of the molecular quantum numbers.

The good quantum numbers

We write the total angular momentum of the molecule $T = R + F = R + I + J$, F being the total spin of the two atoms $f_1 + f_2$, $J = j_1 + j_2$ is the angular momentum of the molecule that describes the spin-orbit coupling, and R the rotational angular momentum of the molecule.

We define the projection of the different angular momentum on the inter-nuclear axis as

$$\Omega = \Lambda + \Sigma \quad (3.21)$$

with Λ the projection of the orbital momentum of the molecule on the internuclear axis and Σ the projection of the momentum spin.

M_T is the projection of the total angular momentum T onto a defined quantization axis.

For the rotational quantum number there is a parity of the total wavefunction that needs to be fulfilled. In addition there are only s-wave collisions in cold gases then the part of the wavefunction should be odd the spin one antisymmetric

$$|\Psi\rangle = (|\phi_e\rangle_A |\phi_g\rangle_B + (-1)^R |\phi_e\rangle_B |\phi_g\rangle_A) \otimes |\chi\rangle \quad (3.22)$$

$|\phi_e\rangle$ and $|\phi_g\rangle$ are the electronic wavefunctions of the two atoms A and B in ground and excited state respectively, $|\chi\rangle$ is the spin wavefunction. It gives accessible values $R = 0, 2$ or $4\dots$

For $^1S_0 - ^3P_1$ molecular state there is a strong spin-orbit coupling that couples its related momentum to the internuclear axis (3.3), quantified by Ω and described as Hund's case (c)

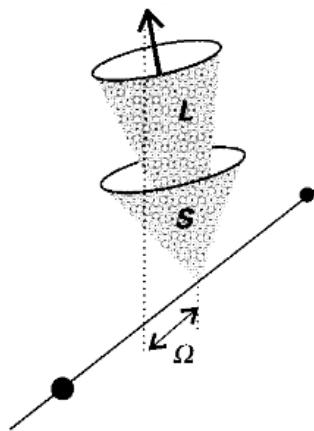


Fig. 11.4 A vector diagram for Hund's case (c), in which the spin-orbit coupling is very strong.

Figure 3.3: hund

The experiment presented in the following section is photoassociation in the $F = 11/2$ hyperfine state. We will focus in this state to simplify the discussion.

Say we have one atom in the 1S_0 state and one in the 3P_1 state: we can have a total spin of $J = 0, 1$ or 2 . It gives possible values of total atomic angular momentum F :

$$\text{For } J = 0, \quad F = \frac{9}{2} \quad (3.23)$$

$$\text{For } J = 1, \quad F = \frac{7}{2}, \frac{9}{2}, \frac{11}{2} \quad (3.24)$$

$$\text{For } J = 2, \quad F = \frac{5}{2}, \frac{7}{2}, \frac{9}{2}, \frac{11}{2}, \frac{13}{2} \quad (3.25)$$

We have one atom with $f_1 = 9/2$ from the 1S_0 state and one in the 3P_1 state with

$f_2 = 11/2$. The possible values of total angular momentum of the molecule are

$$\text{For } R = 0, \quad T = 1, 2, 3 \dots 10 \quad (3.26)$$

$$\text{For } R = 2, \quad T = 1, 2 \dots 12 \quad (3.27)$$

$$\text{For } R = 4, \quad T = 1, 2 \dots 14 \quad (3.28)$$

The initial $T^i = i_1 + i_2$ with $R = 0$ because we are in s-wave collisions. In respect with the parity $p = +1$ the possible values of T^i are $T^i = 0, 2, 4 \dots$

As for atoms there are selection rules for transitions from a free state to a molecular state fixed by the angular momentum of the electrons. $\Delta J = 0, 1$ gives $\Delta T = 0, 1$ with $T = 0 \rightarrow T = 0$ forbidden.

In the end the accessible transitions in the first molecular excited state of the $f_1 = 11/2$ are

$$\text{For } R = 0, \quad T^f = 1, 2, 3 \dots 11 \quad (3.29)$$

$$\text{For } R = 2, \quad T^f = 1, 2 \dots 13 \quad (3.30)$$

$$\text{For } R = 4, \quad T^f = 1, 2 \dots 15 \quad (3.31)$$

$$\dots \quad (3.32)$$

There are 981 combinations of accessible (R, T^f) molecular states

Another element to describe the molecule is the symmetry of the orbital wavefunction. For an even wavefunction the exchange of the two atoms A and B does not change the sign of the wavefunction, as in left figures 3.4. We say the potential is **gerade**. The probability of finding electrons in between the two atoms is high compared to the **ungerade** potential where the odd wavefunction cancels the probability of finding electrons in between.

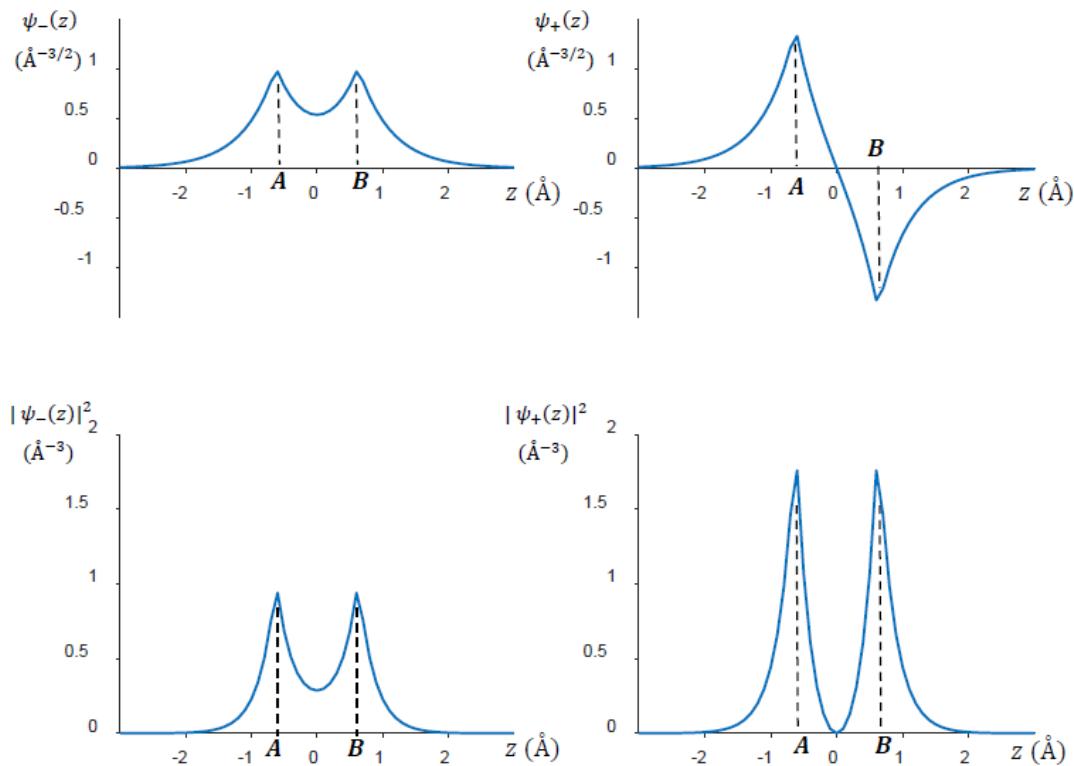


Figure 8.5 : Evolution le long de l'axe internucléaire Oz ($x=0, y=0$) des fonctions d'onde liante $|\psi_{-}\rangle$ et antiliante $|\psi_{+}\rangle$ et des densités de probabilité de présence de l'électron associées $|\psi_{-}|^2$ et $|\psi_{+}|^2$. La distance internucléaire est fixée à $R = 1,25$ Å (minimum de E_{-}).

Figure 3.4: u-g symmetry

3.1.4 Internal energy states

3.1.4.1 WKB approximation

A semi-classical approach enables to find the eigenvectors of atoms in a molecular potential. We separate the interaction zones considering the condition $\frac{1}{2\pi}|\frac{d\lambda}{dx}| \ll 1$

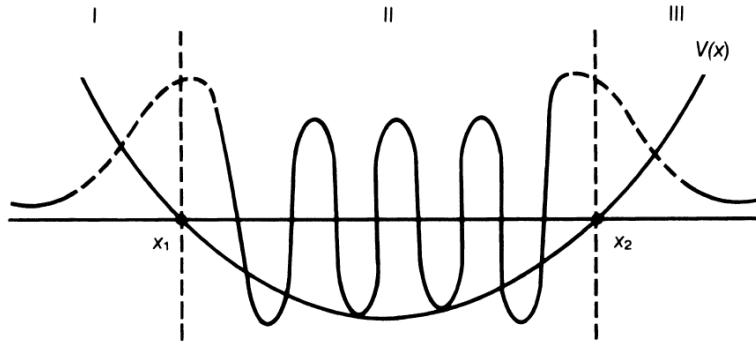


FIGURE 2.1. Schematic diagram for behavior of wave function $u_E(x)$ in potential well $V(x)$ with turning points x_1 and x_2 .

Figure 3.5: Caption

WKB approximation is a semi-classical approach where plane waves $\psi(x) = \frac{C}{\sqrt{k(x)}} e^{i \int_0^\infty k(x) dx}$ are solutions of the 1D Schrödinger equation for two particles of eigenenergie E in a potential $V(x)$

$$\frac{-\hbar^2}{2\mu} \frac{d^2\psi}{dx^2} + V(x)\psi = E\psi \quad (3.33)$$

giving the expression of the wavevector k

$$k = \frac{p}{\hbar} = \sqrt{2\mu(E - V(x))} \quad (3.34)$$

The condition of validity $\frac{d\lambda}{dx} \ll 1$ tells that the atoms must follow adiabatically the potential to adapt their movement to its variation. It means at really small x of region I this approximation is not valid.

At position $x = x_1$ and x_2 the momentum diverges which means that the plane wave approximation is not valid anymore. In the region I and III the wave is evanescent with a phase $k = i\sqrt{2\mu(|E - V(x)|)}$. It is then usual to average the potential as linear with x around these turning points.

By reflection with the left barrier of the potential at $x = x_1$ the stationary wave must come back with a 2π at x_2 . By continuity of the solutions in $[0, +\infty[$ the phase tends to

$$\int_{x_1}^{x_2} k(x) dx = (n + \frac{1}{2})\pi \quad (3.35)$$

Then this formula enables to identify the solutions E that satisfies the condition $(n + \frac{1}{2})\pi$. These values are what we call the vibrational states of a molecular potential.

3.1.4.2 Leroy-Bernstein theory

While the queue of the potential C_n/r^n dominates the potential interaction we can approximate

3.1.5 External energy states

As seen in the subsection 3.1.1 in a molecular picture the hamiltonian express as in [2] by

$$\hat{H} = \frac{\hat{p}_r^2}{2\mu} + \frac{\hbar^2}{2\mu r^2} R(R+1) + \hat{V}_{BO} + \hat{H}_{HF} \quad (3.36)$$

with

$$\hat{V}_{BO} = \frac{-C_6}{r^6} \left(1 - \frac{\sigma^6}{r^6} \right) - s \frac{C_3^\Omega}{r^3} \quad (3.37)$$

$$\hat{H}_{hf} = A(\mathbf{i}_1 \cdot \mathbf{j}_1) + B \frac{3(\mathbf{i}_1 \cdot \mathbf{j}_1)^2 + \frac{3}{2}(\mathbf{i}_1 \cdot \mathbf{j}_1) - i_1(i_1+1)j_1(j_1+1)}{2i_1(2i_1-1)2j_1(2j_1-1)} \quad (3.38)$$

$s = \pm 1$ for gerade and ungerade potentials respectively. The hyperfine term of the atom in 1S_0 is zero because its electronic angular momentum is zero. Only the hyperfine interaction of the atom with i_1, j_1 in 3P_1 counts.

3.2 About photoassociation on other species

3.2.1 Mass scaling (${}^{88}Sr$)

${}^{88}Sr$ is a bosonic isotope of strontium with no hyperfine structure. The photoassociation spectrum is then simpler than the fermionic ones : only two potentials in the first excited molecular state 0_u and 1_u . Results have been presented in [3] where their fitting model gives the following potential expression

$$\hat{V}_{BO} \sim \frac{-C_{6,\Lambda}}{r^6} - \frac{C_3}{r^3} + \frac{C_J}{r^2} \quad (3.39)$$

for the two accessible potentials $\Lambda = 0, 1$.

Mass scaling method presented in [4] enables to predict the position of the vibrational states of the different isotopes of strontium.

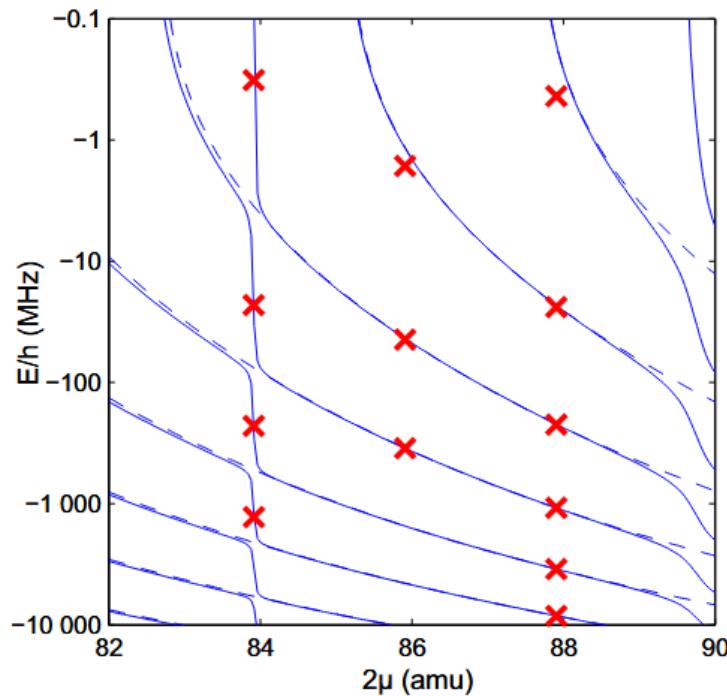


Figure 3.6: Caption

3.2.2 Ytterbium: hfs

As the ^{87}Sr the fermionic ^{173}Yb has a hyperfine structure that makes its molecular potential landscape more complex than many known species.

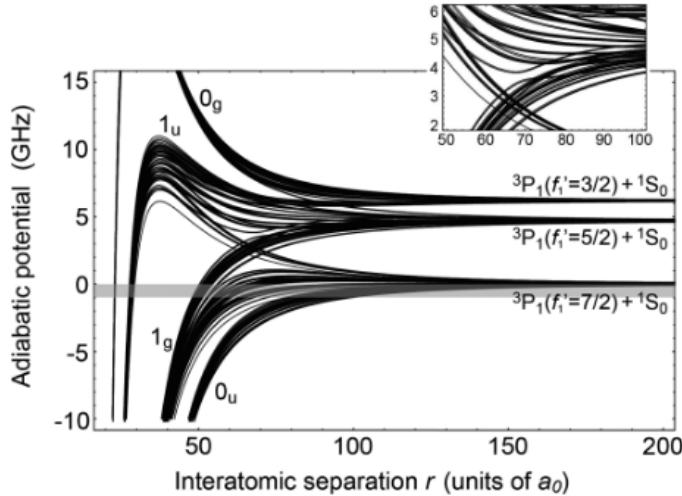


FIG. 2. Adiabatic molecular potentials for a $^{173}\text{Yb}_2$ dimer in the $^1\text{S}_0 + ^3\text{P}_1$ channel as functions of the interatomic separation r . The molecular potentials for 205 different (T, F, R) configurations are displayed, which are accessible via PA from the initial s -wave colliding atoms in the $^1\text{S}_0 + ^1\text{S}_0$ channel. At large r , the potentials converge to three asymptotic branches which correspond to excited atomic states with hyperfine numbers of $f'_1 = 3/2$, $5/2$, and $7/2$. Some of the potentials have a local minimum (inset), possibly hosting purely long-range bound states [14]. The energy offset is adjusted to the $f'_1 = 7/2$ asymptote. The shaded region indicates the spectral range of our measurements.

Figure 3.7: Caption

^{171}Yb In their fitting model they consider the same interactions as in our case (expressed in equation 3.36) but without the hyperfine term.

3.3 Experimental setup

3.4 88Sr Results

The s-wave collisions imposes that for fermions -by parity of the total wavefunction being antisymmetric- only atoms with even angular momentum can collide meaning their spin is in a singlet state as in the case of ^{87}Sr . For bosons it is the contrary : from an even orbital wavefunction, the spin wavefunction should be antisymmetric to collide. Lopt, power broadening, thermal broadening...

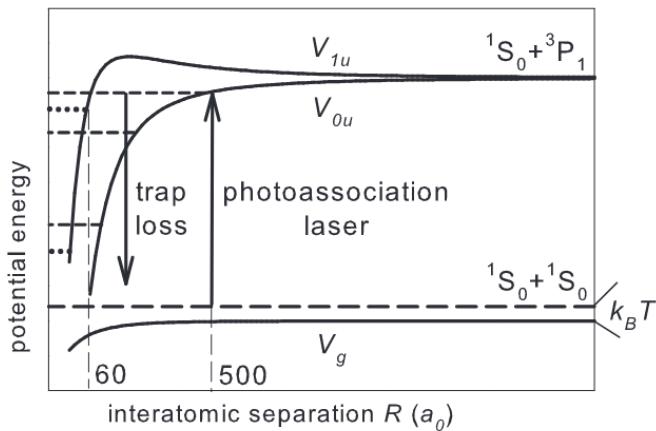


Figure 3.8: Caption

3.4.1 Technical issues of inhabilitation of photoassociation

3.4.1.1 Laser width

3.5 87Sr molecules

The s-wave collisions imposes that for fermions -by parity of the total wavefunction being antisymmetric- only atoms with even angular momentum can collide meaning their spin is in a singlet state as in the case of ^{87}Sr . Lopt questions sur nb quantique / choix de pompage optique

3.5.1 Physical sources of inhabilitation of photoassociation

3.5.1.1 On $F = 9/2$: predissociation

3.5.1.2 Coupling to more energetic state from the IR

3.5.1.3 Node of wavefunction for some vibrational states

3.5.2 Energy landscape of 87Sr-87Sr molecules

Conclusion

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Appendix A

Algorithms