
STATE-DEPENDENT FORCES IN COLD QUANTUM GASES

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CHAPTER I

Introduction

THE SUBJECT OF STUDY of this thesis is Bose–Einstein condensation, as well as associated experimental and theoretical techniques and phenomena in cold atom physics. The following chapters describe my work in a cold atom research group over the past several years, pertaining to apparatus construction, experiment, theory, and software design and development. An overarching theme is *state-dependent forces* on cold atoms. Selectively subjecting atoms to forces based on what state they are in is at the core of many phenomena in cold atom physics. As I go into in the following chapters, different types of state selectivity allow for cooling and imaging techniques that would otherwise not be possible, momentum state-selectivity is central to wave-mixing phenomena; and semiclassical models run into a problem when state-selective forces cannot be disregarded in determining the classical force that atoms modelled semiclassically ought to be subjected to.

Bose–Einstein condensates (BECs) in dilute atomic gases are superfluids that can be created in the lab at extremely low temperatures. This strange state of matter was predicted in 1925 by Bose and Einstein [1, 2], first produced experimentally in 1995 [3] in a cloud of rubidium atoms, and has since been made out of many other atoms, usually alkali metals [4–7]. In a BEC, a macroscopic sample of bosonic atoms all occupy the same quantum state, and many of the features of the single particle wavefunctions are exhibited by the cloud as a whole. Bose–Einstein condensation and cold atoms and ions more generally have rich applications in precision measurement [8, 9], quantum computation [10, 11] and quantum simulation [12, 13].

1.1 Chapter overview

Various experimental techniques are used to produce and study Bose–Einstein condensates, many of which exploit or necessitate an understanding of the quantum behaviour of the atomic systems in question. I summarise some of these techniques and detail the atomic physics principles underlying them in chapter 2.

The fields of Bose–Einstein condensation and cold atoms more generally enjoy a tight coupling between theory and experiment, not least because of the enduring usefulness and accuracy of mean-field theory. In mean-field theory, the quantum matter field operator of the atoms comprising a Bose–Einstein condensate is replaced with its expectation value at each point in space, allowing the entire multi-particle system to be modelled with little more computational complexity than that required to model a single-particle wavefunction.¹ The resulting differential equation—the Gross–Pitaevskii equation—is nonlinear and using it to propagate a condensate wavefunction in time generally requires numerical techniques rather than analytic ones. My favourite numerical methods for do-

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¹Mean field theory is accurate in the low-temperature limit, in which it has some remaining limitations—it does not for example predict the observed s-wave scattering halos when two BEC wavepackets collide [14], but it is sufficient for modelling a wide range of experiments nonetheless.

ing so (which apply more generally to numerically evolving quantum systems of all kinds) are described in chapter 3. In chapter 3 I also develop a variation on fourth-order Runge–Kutta integration which improves on one of its deficiencies for simulating quantum systems. I also present arguments that a fairly sophisticated method of discretising partial differential equations—the finite element discrete variable representation—may offer less computational efficiency than simpler methods for computing solutions of comparable accuracy to the Gross–Pitaevskii and Schrödinger wave equations.

As an experimental field, BEC research involves the construction of apparatuses capable of implementing the techniques described in chapter 2 in order to produce, control, and measure BECs. Chapter ?? describes some of the process of constructing such an apparatus, which involves a vacuum system, magnetic coils and optical systems. I present an optical layout for producing magneto-optically trapped ^{87}Rb atoms (a step on the way to condensation) that I designed and assembled as an exchange student in the group of József Fortágh at the University of Tübingen’s Physikalisches Institut.

Production, control, and measurement of cold atom systems require more than the necessary optics and magnetic sources to be installed—they must be controllable in a time-accurate way in order to execute the necessary cooling processes, manipulate the system as desired, and observe the results. Production of a condensate takes on the order of tens of seconds, requiring precisely timed pulses of laser light at specific frequencies, sweeps of magnetic field strengths, and frequency sweeps of radio and microwave radiation. This cannot all be done by human experimenters alone, and so requires computer automation of some kind. In chapter ?? I reproduce our publication on a suite of software programs, the *labscript suite*. This software leverages modern software development techniques such as object orientation, abstraction and isolation as well as older principles—such as aspects of the Unix philosophy—to produce a powerful, maintainable, extensible system for designing, running and analysing shot-based experiments on commodity hardware.

As superfluids, BECs have zero viscosity and as such can support persistent flows. In classical fluid dynamics the absence of viscosity means that a fluid cannot support vorticity,² and must be irrotational. However, fluid circulation can still occur around points of zero fluid density, known as vortices. In BECs this circulation is also quantised, in units of \hbar/m .

These quantised vortices are topological defects—the phase of the macroscopic wavefunction winds by a multiple of 2π around them, and is undefined at the center of the vortex core itself. Quantised vortices were observed in superfluid helium³ in the early 1960s [15], and in BEC in a dilute atomic gas in 1999 [16]. The formation, dynamics and decay of these vortices are believed to be important for the study of superfluid turbulence [17].

In chapter ?? I present simulations exploring the feasibility of imaging these vortices in-situ using *tracer particles*. Atoms of one kind (^{87}Rb) may become trapped in the cores of quantised vortices in a condensate of another kind (^{41}K), and if imaged in a time-resolved way, reveal the motion of these vortices. A primary concern in any implementation of such a scheme is keeping the tracer atoms cold enough that they remain trapped in the vortex cores even as they scatter light for imaging. To that end, in chapter ?? I present modelling of a novel—if impractical—laser cooling scheme for Sisyphus cooling of ^{87}Rb atoms in a 34 G magnetic field—a field strength at which ^{87}Rb and ^{41}K repel each other strongly (leading to tighter trapping in the vortex cores).

In a BEC, the wavelike behaviour of matter is apparent, unlike at higher temperatures at which atoms are well described as classical particles. Not only are atoms in a BEC wavelike, they are described by a *nonlinear* wave equation. As with nonlinear optical systems, they therefore exhibit wave-mixing behaviour whereby a number of momentum states may interact to produce additional momentum states. In chapter ?? I describe our lab’s four wave mixing experiment, which reproduced an existing result for BECs, and then our attempt at six wave mixing—a higher order effect (in the sense of perturbation

theory). We did not observe the expected six wave mixing, rather we saw *four* wave mixing despite the required resonance condition apparently being violated. This result agreed with mean-field theory simulations however, implying that the result was unlikely to be due to some experimental error.

Atoms have spin, and when other spin-projection states cannot be disregarded, mean-field theory for spinful BECs takes the form of multiple, interacting fields. This allows for richer nonlinear dynamics that a single-component BEC, with wave mixing producing new momentum states and new spin-projection states in tandem. In chapter ?? I present simulation results showing this “spin wave” mixing, although the main result is that—for ^{87}Rb at least—the effect is very small and unlikely to be experimentally observable in existing ^{87}Rb BEC experiments.

As mentioned above, at high temperatures (higher than that at which atoms Bose-condense) atoms are well described as classical particles. This is true in the sense that the wavelike nature of the atoms can be disregarded—they move through space like classical billiard balls obeying Newtonian mechanics. The internal state of the atoms, however—for example the state of an outer shell electron—may not be well modelled by classical mechanics. Even at room temperature, an electron is poorly described as a classical charged particle orbiting a nucleus. When there is *coupling* then, between this internal state of an atom and its motional state, the quantum-ness of the internal state can in some sense ‘leak’ into its motional state even if the motion is otherwise modelled well classically. The classic example of this is the Stern–Gerlach experiment [18], in which a beam of atoms splits into two beams as it passes through a magnetic field gradient. A similar situation arises for atoms in a magnetic trap—a common feature of cold atom experiments and often used in the final stage of cooling to BEC. To correctly model the losses of atoms from these traps, one needs to model the internal state of the atoms quantum-mechanically, but it is computationally expensive to also model their spatial motion using full quantum wavefunctions. We would like a way to model the atoms’ motion classically, but in such a way that it can reproduce Stern–Gerlach separation—with modelled atoms taking one or the other trajectory probabilistically, with the probabilities consistent with those of a fully quantum treatment. In chapter ?? I present such a model, one that is based on a *hidden variable* carried around with each atom being modelled, which selects one of the atom’s internal eigenstates. The apparent definiteness of the hidden variable allows the spatial motion part of the modelling to treat the atoms’ spin projection degree of freedom as if it were in a definite state, allowing the modelling to take a single, definite trajectory. The hidden variable itself is evolved using a stochastic hidden variable theory that ensures its probability of corresponding to any particular spin-projection state is consistent with the underlying quantum evolution of the atom’s internal degrees of freedom.

1.2 What's new in this thesis

[WHATS NEW IN THIS THESIS SECTION (compare to alex, lisa, martijn)]

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CHAPTER 2

Atomic physics: Experimental techniques and theory

BOSE-EINSTEIN CONDENSATES provide such a tantalising opportunity for studying quantum phenomena not only because of their interesting properties, but also because of the level of control they afford, with parameters able to be tuned and manipulated in order to investigate various regimes. Many of the same techniques which allow experimentalists such control over a BEC are also employed in the production of BEC.

The main experimental techniques used to create BEC are Doppler cooling, magneto-optical, magnetic, and dipole trapping, polarisation gradient (Sisyphus) cooling, and evaporative cooling. These were discovered, perhaps by no coincidence, in roughly the same order as they are called for in a BEC experiment, and many were first discovered along the way to the first realisations of Bose–Einstein condensation.

Section 2.1 is a whirlwind summary of these and a few other experimental techniques. Then in section 2.2, I'll present some of the theory describing superfluid flow and vortices and in BEC, which is central to the simulations of vortex tracking presented in chapter ???. The final section of this chapter, section 2.3, will construct the detailed theory for describing the internal state of a ^{87}Rb atom in magnetic and optical fields, considering only the first $L = 0$ to $L = 1$ excitation is accessible to rubidium's sole outer-shell electron. The resulting Hamiltonian describes 32 sublevels, and forms the basis of any detailed calculations of laser cooling and trapping.

2.1 Cooling, trapping, and manipulating atoms

2.1.1 Doppler cooling

Doppler cooling, demonstrated in 1978 [19] is a consequence of the simple observation that atoms see the wavelength of incident light Doppler shifted depending on their velocity. For example, the electric field of a linearly polarised optical plane wave is:

$$\mathbf{E} \propto \cos(\mathbf{k} \cdot \mathbf{r} - \omega t), \quad (2.1)$$

which for an atom moving with constant velocity \mathbf{v} can be written:

$$\mathbf{E} \propto \cos(\mathbf{k} \cdot \mathbf{v}t - \omega t) \quad (2.2)$$

$$= \cos(-\omega_{\text{eff}}t), \quad (2.3)$$

where $\omega_{\text{eff}} = \omega - \mathbf{k} \cdot \mathbf{v}$ is the effective angular frequency of the laser as seen by the atom.

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This can be used to selectively transfer momentum to only fast-moving atoms, by tuning an incident laser slightly redder than would be required for a resonant absorption. If six lasers in counterpropagating pairs orthogonal to each other surround a cloud of atoms, the atoms can be cooled close to the *Doppler limit* [20, p 58]

$$k_B T_D = \frac{\hbar\Gamma}{2} \quad (2.4)$$

¹The D₂ line, 5S_{1/2} → 5P_{3/2}, approximately 780 nm.

where Γ is the linewidth of the atomic transition. For the cooling transition used for Doppler cooling ⁸⁷Rb¹, this gives 146 μK, which is approximately a factor of a thousand too high for Bose-condensation. These atoms are also not trapped.

2.1.2 Magneto-optical and magnetic trapping

Magneto-optical trapping, first demonstrated in 1987 [21] comes from the realisation that a magnetic field can be used to *spatially* vary the detuning from resonance that the atoms in the above mentioned arrangement of lasers see. This is possible due to the Zeeman effect [22], in which the wavelengths of atomic transitions are shifted in a magnetic field.

If a field profile can be found which causes the transition to come closer to resonance as the atoms move away from a central point, then it forms a trap—atoms that stray too far from the centre will absorb more strongly and be deflected back².

The field configuration used in an anti-Helmholtz one: two coils opposite each other carrying opposing currents. The resulting magnetic field profile has a zero in the middle and increases in magnitude in all directions.

With the Doppler beams off, this magnetic field still provides a trapping potential, due to the magnetic dipole interaction:

$$V(\mathbf{r}) = -\boldsymbol{\mu} \cdot \mathbf{B}, \quad (2.5)$$

where $\boldsymbol{\mu}$ is the atomic magnetic moment, and \mathbf{B} the magnetic field. This only traps some atomic spin states, and has losses due to spin-flips [23] near the field zero.

The optimal magnetic field gradient for forming a magneto-optical trap (MOT) is lower than that required to merely hold atoms against gravity, and thus there is little if any magnetic trapping occurring in a MOT—the trapping almost entirely results from the position-dependent rate of photon absorption caused by the Zeeman shift due to the magnetic field. To transfer atoms from a MOT to a solely magnetic trap, the magnetic field gradient must be increased substantially.

2.1.3 Optical trapping

Optical trapping on the other hand relies on the *dipole force*, in which off-resonant light shifts the energy of the eigenstates of the combined atom-light system, the so called *dressed states*. This energy shift, called the *light shift*, depends on the intensity of the light, and so results in a potential that spatially varies as the intensity of the light. In the limit of large detuning (compared to Rabi frequency), the resulting energy shift for an atom in a groundstate (the shift for an excited state is the same but with the opposite sign) is [20, p 8]:

$$\Delta E = \frac{\hbar\Omega^2}{4\delta} \quad (2.6)$$

where δ is the detuning from resonance and the Rabi frequency is:

$$\Omega = \frac{E_0}{\hbar} \langle n' | e\hat{r} | n \rangle, \quad (2.7)$$

where E_0 is the amplitude of the light's electric field and $\langle n'|\hat{e}\hat{r}|n\rangle$ is the transition dipole moment between the two states in a two-level system. This transition dipole moment can be either for two specific sublevels coupled by a laser, the calculation of which is detailed in section 2.3.5, or an overall effective transition dipole moment between two manifolds containing many sublevels, if the detuning much larger compared to the spacing between the sublevels.

With the potential proportional to E_0^2 , and thus the light's intensity, the force the atom experiences is proportional to the light's intensity gradient. For this reason, the dipole force is also called the *gradient force*. The name *dipole force* comes from the fact that the force can be equivalently understood to arise from the polarisability of atoms in a light field, giving rise to a force identical to that which traps polarisable materials in optical tweezers [24].

2.1.4 Polarisation gradient cooling

Polarisation gradient cooling, also called Sisyphus cooling, was proposed in 1989 [25, 26] to explain experimentally measured cold atom cloud temperatures [27] which, at NIST in 1988, were found to be well below the expected limit obtainable by the well understood method of Doppler cooling³, one of the few examples of experiments turning out better than expected. A one dimensional theory has been developed [25] which has found remarkable agreement with three dimensional experiments [28]

One common configuration for Sisyphus cooling comprises two counterpropagating laser beams in each spatial dimension, both linearly polarised but with their polarisation angles perpendicular to one another. The optical field resulting from the two beams' superposition has regions of linear polarisation and of both helicities of circular polarisation, and varies between them on a length scale shorter than an optical wavelength.

The effect on multi-level atoms as they move from regions of one circular polarisation to another is that they are pumped alternately from one extreme of their spin-projection states to the other, alternately climbing and descending potential hills due to the dipole forces from the regions of different polarisations⁴. And so, like the Greek legend of Sisyphus⁵, who was doomed to push a rock uphill for eternity, the atoms are climbing hills repeatedly. Due to the state dependence of the strength of the dipole forces, the atoms climb steeper hills than they descend, and are thus slowed and cooled.

This type of cooling does not work in a magnetic field; the splitting of transition frequencies makes it impossible for an atom to traverse its spin manifold on one laser frequency. For this reason the Sisyphus cooling stage is performed with magnetic fields off, though a sufficiently short period is required such that the atoms can be recaptured when the trapping field is restored.

³As well as to explain other discrepancies between experiments and the theory of Doppler cooling, such as the optimal detuning of light being much greater than predicted.

⁴If you consider only one polarisation of light, its intensity varies sinusoidally in space, creating a series of potential hills and wells via the dipole force

⁵Polarisation gradient cooling is but one of a family of so called 'Sisyphus cooling' methods, all of which involve atoms repeatedly climbing potential hills.

2.1.5 Evaporative cooling

The final stage of cooling is forced RF evaporative cooling [3, 29], which decreases the temperature of the cloud by systematically removing the hottest atoms. This is performed in a magnetic trap, which as mentioned earlier, only traps certain spin states. Evaporation proceeds by using an *RF knife* to induce spin flips in the atoms. The RF frequency is chosen such that it is only resonant with atoms some distance away from the center of the trap (via the Zeeman shift). The furthest out atoms are the most energetic, possessing the energy to climb the magnetic potential the furthest. By flipping their spins, these atoms are ejected due to the magnetic field becoming anti-trapping for them.

The cloud is given some time to rethermalise and the knife⁶ is moved inward where it removes slightly colder atoms. This is repeated until the desired compromise of lower temperature/lower atom number is reached. Usually some method is employed to prevent atoms near the center of the trap from undergoing spin flips [23] as they move across the

⁶So called because it cuts the tail off the velocity distribution of the atom cloud.

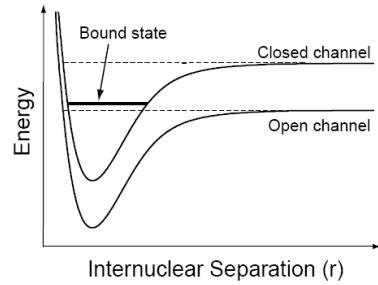


Figure 2.1: When atoms approach each other with spins aligned, they are in the *open channel*. In this channel they are unbound, but do not have enough energy to be free in the other channel - the *closed channel*. In the close range however, the atoms may have energy corresponding to a bound (molecular) state of the closed channel, a resonance which causes a divergence in the scattering length. The energy difference between the two channels can be tuned with a magnetic field and so these resonances can be induced in a wide range of situations.

field zero. The method used in our lab is to use an optical dipole trap in combination with the magnetic trap [30], such that the coldest atoms get trapped in the dipole trap which is offset from the magnetic field zero.

2.1.6 Feshbach resonances

A Feshbach resonance [31] is an enhancement of the interparticle interaction strength when when a certain magnetic field strength is applied⁷. This phenomenon was first discovered in ultracold atoms in 1998 [32], and is now a staple of cold atom experiments.

The interparticle interaction mentioned above:

$$g = \frac{2\pi\hbar^2 a}{m_r} \quad (2.8)$$

where m_r is the reduced mass of a pair of the interacting particles, is dependent on a parameter a called the *s-wave scattering length*, which characterises low energy collisions between atoms. It is sensitive not only to what species of atoms are colliding, but also to their spin states. For each combination of spins, there is a different inter-atomic potential (called a *channel*) which determines the collision dynamics (Figure 2.1).

The resulting scattering length is sensitive to any bound states of this inter-atomic potential which are near the collision energy. If the channels of different spin states are coupled via the hyperfine interaction⁸, then the scattering length is also sensitive to bound states in the channels other than the one the atoms are in when they are far from each other. Due to the Zeeman effect, the energies between the different channels can be shifted with a magnetic field, and so a bound state can be shifted close to the collision energy, which causes the scattering length to diverge.

The end result is that at certain magnetic field strengths we find that atoms are much more strongly attracted to or repelled from each other.

A particular Feshbach resonance of interest is shown in Figure 2.2, and can be used to enhance the interspecies repulsion between ^{87}Rb and ^{41}K . The use of this resonance is assumed for the simulations in chapter ?? to trap tracer particles more strongly in vortex cores of a BEC.

⁷Feshbach resonances can also be induced optically and with RF but magnetic resonances are the most commonly used.

⁸Requiring that the atoms in question have a nuclear magnetic moment.

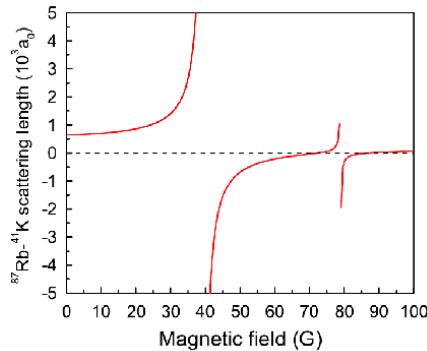


Figure 2.2: Predicted interspecies scattering length [33] as a function of magnetic field strength, for ^{41}K and ^{87}Rb both in their lowest energy hyperfine groundstate. The 35 gauss resonance is one of the main reasons for this pair of atoms being used in this project. It has a particularly low field strength and large width compared to most Feshbach resonances.

2.2 Mean field theory: The Gross–Pitaevskii equation and vortices

Bose-condensates are described well by *mean field* theory, whereby the many-body wavefunction is approximated by a product of identical single-particle wavefunctions. Indeed, that the majority of the atoms are in the same quantum state is one of the defining features of BEC. The effect of interparticle interactions is included as a nonlinear term in the Schrödinger equation for the single particle wavefunctions, known as the Gross–Pitaevskii equation:

$$\frac{\partial \Psi}{\partial t} = \left[-\frac{\hbar^2}{2m} \nabla^2 + V(\mathbf{x}) + g|\Psi|^2 \right] \Psi, \quad (2.9)$$

where g characterises the strength of the interparticle interactions⁹, and $\Psi = \sqrt{N}\Psi_{\text{single}}$ is the single-particle wavefunction scaled by the square root of the number of particles¹⁰.

In the hydrodynamic formulation of quantum mechanics [34], the flow velocity of a spatial wavefunction can be defined by considering the probability current to be a product of density and velocity. This allows us to define the superfluid velocity of a BEC as:

$$\mathbf{v} = \frac{\hbar}{m} \nabla \phi \quad (2.10)$$

where ϕ is the complex phase of the condensate wavefunction Ψ . Integrating this velocity over any closed path γ gives us the circulation:

$$C = \frac{\hbar}{m} \oint_{\gamma} \nabla \phi \cdot d\mathbf{s} \quad (2.11)$$

$$= \frac{\hbar}{m} 2\pi n. \quad n = 0, 1, 2 \dots \quad (2.12)$$

The fact that the circulation is quantised means that vorticity cannot exist in the condensate except in one-dimensional lines, about which the wavefunction's phase winds by a multiple of 2π . These topological defects are the quantised vortices that are central to the vortex tracking simulations in chapter ??.

At a vortex core, the atom density of a BEC must go to zero. This can be intuitively understood to arise from centrifugal forces, but is also required in order for the wavefunction to be continuous and single-valued across the core. This drop in density in the

⁹The nonlinear constant g is usually positive—having the effect of stabilising BECs by self-repulsion.

¹⁰Thus giving it the property that $|\Psi|^2$ is the particle density.

vicinity of a vortex core is exploited by the method simulated in chapter ?? to trap atoms within the cores.

2.3 The ^{87}Rb D line

We atomic physicists do our theory work at an intermediate level of abstraction, at which many quantities and systems of interest can be computed and simulated with accurate models using standard quantum mechanics, but with models that are not fully a-priori. Instead, the Hamiltonians we feed to the machinery of quantum mechanics encapsulate some of the details we are not interested in or that are too hard to compute, with the link between the underlying layers of reductionism and the higher layer usually provided by experimentally measured values rather than calculations from fundamental physics. In this way we can readily compute results about the atoms we are interested in by treating them as simpler systems than they actually are, with some of the underlying details encapsulated by terms in an effective Hamiltonian for the dynamics that we are interested in.

In this section I'll summarise what the ^{87}Rb D line looks like from the perspective of a cold atom physicist, building up a Hamiltonian containing all 32 sublevels of the ground and first excited state of ^{87}Rb including fine structure, hyperfine structure, interaction with a magnetic field, and optical transitions between states. This Hamiltonian is the starting point for any calculations regarding cooling, trapping, and coherent control of ^{87}Rb , and for other alkali earth metals is much the same.

Much of the details of this section are drawn from references [20, 35–38], but the reader should be aware that there are considerable conventional and notational differences between different literature sources. What is presented here is summarised and framed in a way that I think is useful to an experimentalist looking to use the theory to make concrete calculations about real systems.

2.3.1 Fine structure

The rubidium 87 D line refers to the first excitation available to the sole outer electron of ^{87}Rb . Both the ground and excited state of this transition have the same principal quantum number, but different orbital angular momentum quantum numbers. Upon closer inspection, it is not just one transition between a ground state and an excited state—there are two excited states, and the two resulting transitions are called, in order of their transition frequencies, the D₁ and D₂ lines. Thus the ground and first excited state of ^{87}Rb are actually a ground state plus two non-degenerate excited states, once we take into account fine structure. The groundstate is an S state (electronic orbital angular momentum quantum number $L = 0$), called the $S_{1/2}$ state, and the two excited states are P states ($L = 1$), one with the electron spin anti-aligned with its orbital angular momentum (resulting in total angular momentum quantum number $J = 1/2$) and one with the electron spin aligned with its orbital angular momentum ($J = 3/2$), called the $P_{1/2}$ and $P_{3/2}$ states respectively. In all of these states, ^{87}Rb 's single outer-shell electron occupies an orbital with principle quantum number $n = 5$, which for brevity we leave out of the notation. The transition between $S_{1/2}$ and $P_{1/2}$ is called the D₁ line, with experimentally measured (angular) transition frequency ω_{D_1} , and the transition between $S_{1/2}$ and $P_{3/2}$ is the D₂ line with angular transition frequency ω_{D_2} . These transition frequencies correspond to optical wavelengths of $\lambda_{\text{D}_1} \approx 795 \text{ nm}$ and $\lambda_{\text{D}_2} \approx 780 \text{ nm}$ [35].

This fine structure is treated entirely empirically for our purposes, and so our base Hamiltonian for the rubidium D line, taking into account only fine structure, is simply a statement of the experimentally measured energy differences between the states:

$$\hat{H}_{\text{fs}} = \hbar\omega_{\text{D}_2} \hat{1}_{P_{3/2}} \oplus \hbar\omega_{\text{D}_1} \hat{1}_{P_{1/2}} \oplus \hat{0}_{S_{1/2}}, \quad (2.13)$$

where $\hat{1}_{P_{3/2}}$, $\hat{1}_{P_{1/2}}$, and $\hat{0}_{S_{1/2}}$ are identity and zero operators each acting on the subspace of states within the $P_{\frac{3}{2}}$, $P_{\frac{1}{2}}$, and $S_{\frac{1}{2}}$ manifolds respectively, and \oplus is the direct sum.¹¹ The matrix representation H_{fs} of \hat{H}_{fs} in the basis in which it is diagonal (which we will call the $\{|L_J\rangle\}$ basis, since L and J are good quantum numbers¹² for specifying one of the three states, which we write with the spectroscopic notation letter— S or P —corresponding to the value of L in place of its numerical value) is

$$H_{\text{fs}} = \begin{bmatrix} \ddots & & & \\ & \begin{bmatrix} \ddots & \hbar\omega_{D_2} & \ddots \\ & \ddots & \ddots \\ & & \ddots \end{bmatrix} & & \\ & & \begin{bmatrix} \ddots & \hbar\omega_{D_1} & \ddots \\ & \ddots & \ddots \\ & & \ddots \end{bmatrix} & & \\ & & & \begin{bmatrix} \ddots & 0 & \ddots \\ & \ddots & \ddots \\ & & \ddots \end{bmatrix} \end{bmatrix}, \quad (2.14)$$

which is a block-diagonal matrix with each block also being a diagonal matrix. We have not yet specified the size of each submatrix—the size of each differs and depends on how many hyperfine and Zeeman sublevels are in that state.

This base Hamiltonian is worth singling out since the energy differences between its three states are orders of magnitude larger than any of the energy differences between hyperfine and Zeeman sublevels within them. When doing any sort of calculations or simulations then, this time-independent Hamiltonian can often be removed from the equations using an interaction picture (see Section 3.2.2), as done in Section ?? in simulating laser cooling.

2.3.2 Hyperfine structure

Within each of the $S_{1/2}$, $P_{1/2}$ and $P_{3/2}$ states, the single outer-shell electron's total angular momentum $\hat{\mathbf{J}}$ has an interaction with ^{87}Rb 's nuclear angular momentum $\hat{\mathbf{I}}$. This results in multiple discrete energy levels depending on the relative orientation of the two separate angular momenta. The interaction Hamiltonian for this hyperfine structure is¹³ [35, 39]:

$$\hat{H}_{\text{hfs}} = \frac{A_{\text{hfs}}}{\hbar^2} \hat{\mathbf{I}} \cdot \hat{\mathbf{J}} + \frac{B_{\text{hfs}}}{\hbar^2} \frac{3(\hat{\mathbf{I}} \cdot \hat{\mathbf{J}})^2 + \frac{3}{2}\hat{\mathbf{I}} \cdot \hat{\mathbf{J}} - I(I+1)J(J+1)}{2I(2I-1)J(2J-1)}, \quad (2.15)$$

where J is the total angular momentum quantum number of the electron, equal to either $\frac{1}{2}$ or $\frac{3}{2}$ depending on which state in the D line we are considering, $I = \frac{3}{2}$ is the total angular momentum quantum number of the nucleus, and A_{hfs} and B_{hfs} are empirically determined coupling constants. Here we see the boundary between the quantities we can calculate with the machinery of quantum mechanics and those that we determine empirically—this expression applies so long as J and I are good quantum numbers,¹⁴ and the two terms are the dipolar and quadrupolar interactions [39] between two angular momenta, with the coupling constants determined empirically and encapsulating details that are difficult to compute a-priori, such as relativistic effects and the exact shape of the electron orbitals given the presence of inner shell electrons. For the spherically-symmetric $S_{1/2}$ groundstate, there is no quadrupolar interaction and so B_{hfs} is only nonzero for the two P excited states.¹⁵ The values of A_{hfs} and B_{hfs} for each of the three L_J states of the D line can be found in [35].

For a given state of the D line, we can construct the matrix representation $H_{\text{hfs}}^{I \times J}$ of \hat{H}_{hfs} in the $I \times J$ basis—which we define as the basis in which both the z vector components \hat{I}_z and \hat{J}_z of $\hat{\mathbf{I}}$ and $\hat{\mathbf{J}}$ are diagonal—by constructing matrix representations $\hat{\mathbf{I}}^{I \times J}$ and $\hat{\mathbf{J}}^{I \times J}$ of the operators $\hat{\mathbf{I}}$ and $\hat{\mathbf{J}}$ in that basis and then applying the expression (2.15). The overset $I \times J$ on each of the matrices indicates that the matrix is a representation of its respective operator in the $I \times J$ basis, so named because its basis vectors can be obtained

¹¹Not to be confused with the Kronecker sum, with which it shares notation. The direct sum concatenates matrices as blocks, producing a larger, block-diagonal matrix with dimension equal to the *sum* of the dimensions of the matrices being direct-summed, whereas the Kronecker sum is the regular sum of matrices after each has been multiplied using the Kronecker-product with identity matrices with sizes of the other matrices in the sum, producing matrices with dimension equal to the *product* of those being summed.

¹²A *good quantum number* is a number that can be used to label (not necessarily uniquely) an energy eigenstate, and on which the energy eigenvalue of that state depends. Saying J and L are good quantum numbers is saying that the eigenstates of the overall Hamiltonian are also eigenstates of \hat{L}^2 and \hat{j}^2 , since eigenstates of these two operators can be specified by stating their quantum numbers L and J .

¹³Note that this expression differs from those in the cited references by a factor of $1/\hbar^2$ —this is because I define the $\hat{\mathbf{I}}$ and $\hat{\mathbf{J}}$ angular momentum operators in SI units, rather than in units of \hbar^2 .

¹⁴ J is a good quantum number so long as the hyperfine splitting is small compared to the spacing between the three states of the D line, which it is, and I is a good quantum number so long as the hyperfine splitting is small compared to the energy difference between the groundstate and the first *nuclear* excited state, which it most certainly is.

¹⁵The quadrupolar term should be explicitly excluded from numerical computations of the hyperfine splitting on the $S_{1/2}$ state, as it contains a division by zero in this case, which may lead to erroneous results even if the term is subsequently multiplied by $B_{\text{hfs}} = 0$.

¹⁶The different types of products available for matrices, vectors, operators, spaces, and their sets of basis vectors are rife with subfield-specific conventions and over-loaded notation, leading to much ambiguity. Although ‘Cartesian product’ connotes well what I mean here, it is still ambiguous. What I mean is that the $\mathcal{I} \times \mathcal{J}$ basis is the set of basis vectors $\{\mathbf{u} \otimes \mathbf{v} \mid \mathbf{u} \in \mathcal{I}, \mathbf{v} \in \mathcal{J}\}$. The notation $\mathbf{u} \otimes \mathbf{v}$, which is also ambiguous, denotes the Kronecker product of two column vectors, producing a column vector with number of elements equal to the product of the number of elements in each of the two vectors—not the ‘outer’ or ‘tensor’ product, which would produce a matrix.

¹⁷This explicit procedure for constructing the matrix representations of these operators is useful for entering into a programming language to produce programs capable of performing atomic physics calculations for arbitrary total angular momentum quantum number J , without having to explicitly enter the angular momentum operators for each value of J , which can be tedious and prone to human error.

¹⁸I will hereafter use the terms “ \mathcal{J} basis” and “ $\{|m_J\rangle\}$ basis”, and similarly for other bases, interchangeably, in the understanding that this identification of standard unit vectors with the basis kets is implied.

¹⁹We’re using the standard convention of ordering the eigenkets $\{|m_J\rangle\}$ in descending order of m_J . This is at odds with the computer programming convention of looping over most indices in ascending order, and so care should be taken when constructing these matrices in a computer program.

via a Cartesian product¹⁶ of the two sets of basis vectors from the \mathcal{I} and \mathcal{J} bases for the two individual nuclear and electronic angular momentum degrees of freedom, in which \hat{I}_z and \hat{J}_z are respectively diagonal.

To construct matrix representations of angular momentum operators in the product space basis $\mathcal{I} \times \mathcal{J}$, we first need their matrix representations $\overset{\mathcal{I}}{\mathbf{I}}$ and $\overset{\mathcal{J}}{\mathbf{J}}$ in the bases of their respective subspaces, which we will write as \mathbf{I} and \mathbf{J} for brevity. We can then expand the two operators into the product space by applying a Kronecker product with an appropriate identity matrix to each:

$$\overset{\mathcal{I} \times \mathcal{J}}{\mathbf{I}} = \mathbf{I} \otimes \overset{\mathcal{J}}{\mathbb{I}} \quad (2.16)$$

$$\overset{\mathcal{I} \times \mathcal{J}}{\mathbf{J}} = \overset{\mathcal{I}}{\mathbb{I}} \otimes \mathbf{J} \quad (2.17)$$

where $\overset{\mathcal{I}}{\mathbb{I}}$ is the matrix representation of the identity operator in the \mathcal{I} basis of the nuclear spin degree of freedom, equal to a $(2I + 1) \times (2I + 1)$ identity matrix, with $\overset{\mathcal{J}}{\mathbb{I}}$ defined similarly for the electronic spin degree of freedom.

Each of the two matrices \mathbf{I} and \mathbf{J} is actually a vector of matrices, one for the angular momentum projection in each of the directions x , y and z . The procedure for constructing such matrices for arbitrary total angular momentum quantum numbers is as follows.¹⁷ I’ll show the procedure for constructing J_x , J_y and J_z only for an arbitrary J , the procedure is identical for computing the vector components of \mathbf{I} .

For a given total angular momentum quantum number J , the vector components of \mathbf{J} in the \mathcal{J} basis (the basis in which J_z is diagonal) can be constructed using the raising and lowering operators \hat{J}_+ and \hat{J}_- . Since the action of the raising and lowering operators on an eigenstate of J_z with angular momentum projection quantum number m_J is to produce an adjacent ($m_J \pm 1$) eigenstate multiplied by a constant [40, p 192], this fact can be used to compute the nonzero matrix elements of \hat{J}_+ and \hat{J}_- in the $\{|m_J\rangle\}$ basis (the basis kets of which we identify with the standard basis vectors¹⁸ for a $2J + 1$ dimensional space in order to define the set of concrete basis vectors \mathcal{J}):

$$\langle m_J + 1 | \hat{J}_+ | m_J \rangle = \hbar \sqrt{J(J+1) - m_J(m_J+1)}, \quad -J \leq m_J < J, \quad (2.18)$$

$$\langle m_J - 1 | \hat{J}_- | m_J \rangle = \hbar \sqrt{J(J+1) - m_J(m_J-1)}, \quad -J < m_J \leq J, \quad (2.19)$$

and therefore compute explicit matrices for J_+ and J_- in the $\{|m_J\rangle\}$ basis:¹⁹

$$J_+ = \begin{bmatrix} 0 \langle J | \hat{J}_+ | J-1 \rangle & 0 & \dots \\ \dots & 0 \langle J-1 | \hat{J}_+ | J-2 \rangle & 0 & \dots \\ & \ddots & \ddots & \ddots \\ & \dots & 0 \langle -J+2 | \hat{J}_+ | -J+1 \rangle & 0 \\ & & \dots & 0 \langle -J+1 | \hat{J}_+ | -J \rangle \\ & & & \ddots \\ & & & 0 \end{bmatrix}, \quad (2.20)$$

$$J_- = \begin{bmatrix} 0 & \dots \\ \langle J-1 | \hat{J}_- | J \rangle & 0 \\ 0 & \langle J-2 | \hat{J}_- | J-1 \rangle & 0 & \dots \\ \dots & 0 & \langle J-3 | \hat{J}_- | J-2 \rangle & 0 & \dots \\ & \ddots & \ddots & \ddots & \ddots \\ & \dots & 0 \langle -J | \hat{J}_- | -J+1 \rangle & 0 \end{bmatrix}, \quad (2.21)$$

both of which have nonzero values along only one non-main diagonal adjacent to the main diagonal, and which form a Hermitian conjugate pair (or indeed, a transpose pair, since all elements are real). The matrix representations of \hat{J}_x and \hat{J}_y can then be computed by rearranging the defining expressions for \hat{J}_+ and J_- :

$$\hat{J}_+ = \hat{J}_x + i\hat{J}_y, \quad (2.22)$$

$$\hat{J}_- = \hat{J}_x - i\hat{J}_y, \quad (2.23)$$

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for \hat{J}_x and \hat{J}_y , and then applying the result to our matrix representations of \hat{J}_+ and \hat{J}_- to obtain matrix representations J_x and J_y of \hat{J}_x and \hat{J}_y in the $\{|m_J\rangle\}$ basis:

$$J_x = \frac{J_+ + J_-}{2}, \quad (2.24)$$

$$J_y = \frac{J_+ - J_-}{2i}. \quad (2.25)$$

Finally, since $\{|m_J\rangle\}$ is the eigenbasis of J_z with eigenvalues $\{\hbar m_J\}$, the matrix representation of J_z in the \mathcal{J} basis is simply the diagonal matrix of eigenvalues:

$$J_z = \begin{bmatrix} \hbar J & & & \\ & \hbar(J-1) & & \\ & & \ddots & \\ & & & \hbar(-J+1) \\ & & & & -\hbar J \end{bmatrix}. \quad (2.26)$$

We can also construct the matrix representation J^2 of the total (squared) angular momentum operator \hat{J}^2 as

$$J^2 = J_x^2 + J_y^2 + J_z^2, \quad (2.27)$$

or equivalently

$$J^2 = J(J+1)\hbar^2 \mathbb{I}, \quad (2.28)$$

since every m_J state is an eigenstate of the J^2 operator with eigenvalue $J(J+1)\hbar^2$.

The above prescription can be used to produce matrix representations of angular momentum operators J_x, J_y, J_z and J^2 for any integer or half-integer total angular momentum quantum number J . The three components can be considered a vector of matrices, \mathbf{J} , for the vector angular momentum operator $\hat{\mathbf{J}}$. Below is a Python function that computes these matrices as well as the corresponding eigenvectors:

```

1 import numpy as np
2 hbar = 1.054571628e-34
3
4 def angular_momentum_operators(J):
5     """Construct matrix representations of the angular momentum operators Jx,
6     Jy, Jz and J2 in the eigenbasis of Jz for given total angular momentum
7     quantum number J. Return them, as well as the number of angular momentum
8     projection states, a list of angular momentum projection quantum numbers
9     mJ, and a list of their corresponding eigenvectors, in the same order as
10    the matrix elements (in descending order of mJ)."""
11    n_mJ = int(round(J**2 + 1))
12    mJlist = np.linspace(J, -J, n_mJ)
13    Jp = np.diag([hbar * np.sqrt(J*(J+1) - mJ*(mJ + 1)) for mJ in mJlist if mJ < J], 1)
14    Jm = np.diag([hbar * np.sqrt(J*(J+1) - mJ*(mJ - 1)) for mJ in mJlist if mJ > -J], -1)
15    Jx = (Jp + Jm) / 2
16    Jy = (Jp - Jm) / 2j
17    Jz = np.diag([hbar*mJ for mJ in mJlist])
18    J2 = Jx**2 + Jy**2 + Jz**2
19    basisvecs_mJ = [vec for vec in np.identity(n_mJ)]
20    return Jx, Jy, Jz, J2, n_mJ, mJlist, basisvecs_mJ

```

Using the above prescription to construct a matrix representation of the $\hat{\mathbf{J}}$ operator with $J = \frac{1}{2}$ for the $S_{1/2}$ and $P_{1/2}$ states, or $J = \frac{3}{2}$ for the $P_{3/2}$ state, and to construct a matrix representation of the \hat{I} operator with $I = \frac{3}{2}$, we can explicitly construct a matrix representation of the hyperfine interaction Hamiltonian for any state of the ^{87}Rb D line. Remaining is to obtain the matrix representations of the two operators in the $\mathcal{I} \times \mathcal{J}$ product space basis using (2.16) and (2.17), and then we can apply (2.15) to our matrices

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to obtain the matrix representation $H_{\text{hfs}}^{I \times J}$ of \hat{H}_{hfs} for a given J corresponding to one of the three states on the D line:

$$H_{\text{hfs}}^{I \times J} = \frac{A_{\text{hfs}}}{\hbar^2} \mathbf{I} \cdot \mathbf{J} + \frac{B_{\text{hfs}}}{\hbar^2} \frac{\frac{3}{2}(\mathbf{I} \cdot \mathbf{J})^2 + \frac{3}{2}\mathbf{I} \cdot \mathbf{J} - I(I+1)J(J+1)}{2I(2I-1)J(2J-1)}, \quad (2.29)$$

where the products of vector components within the dot products are computed with ordinary matrix multiplication. Alternatively, one can use the matrices in their individual subspaces rather than their equivalents in the product space, so long as one interprets the dot products as “Kronecker dot products”:

$$H_{\text{hfs}}^{I \times J} = \frac{A_{\text{hfs}}}{\hbar^2} \mathbf{I} \circledast \mathbf{J} + \frac{B_{\text{hfs}}}{\hbar^2} \frac{3(\mathbf{I} \circledast \mathbf{J})^2 + 3\mathbf{I} \circledast \mathbf{J} - I(I+1)J(J+1)}{2I(2I-1)J(2J-1)}, \quad (2.30)$$

where \circledast is the Kronecker dot product:

$$\mathbf{I} \circledast \mathbf{J} \equiv I_x \otimes J_x + I_y \otimes J_y + I_z \otimes J_z. \quad (2.31)$$

In the above way one can construct an explicit matrix representation of \hat{H}_{hfs} in the $I \times J$ basis for the hyperfine interaction for a given L_J state of the D line.

Column vectors in the $I \times J$ basis

Because the matrix representations \mathbf{I} and \mathbf{J} of the electron and nuclear angular momentum operators were constructed in the $\{|m_I\rangle\}$ and $\{|m_J\rangle\}$ bases of their respective subspaces, the matrices we have constructed are in the basis $I \times J$ with basis vectors:

$$I \times J = \left\{ |m_I m_J\rangle \mid |m_I\rangle \in \{|m_I\rangle\}, |m_J\rangle \in \{|m_J\rangle\} \right\}, \quad (2.32)$$

where $|m_I m_J\rangle = |m_I\rangle \otimes |m_J\rangle$. The vector representation ψ of a state vector $|\psi\rangle$ in this basis is:

$$\psi = \begin{bmatrix} \langle m_I=I m_J=J | \psi \rangle \\ \langle m_I=I m_J=J-1 | \psi \rangle \\ \vdots \\ \langle m_I=I m_J=-J | \psi \rangle \\ \langle m_I=I-1 m_J=J | \psi \rangle \\ \langle m_I=I-1 m_J=J-1 | \psi \rangle \\ \vdots \\ \langle m_I=-I m_J=-J | \psi \rangle \end{bmatrix}. \quad (2.33)$$

The hyperfine Hamiltonian is not diagonal in the $\{|m_I m_J\rangle\}$ basis. The basis in which it is diagonal—the $\{|F m_F\rangle\}$ basis—will be discussed in section 2.3.4.

2.3.3 Zeeman sublevels

The states differing only in their m_F quantum numbers—called Zeeman sublevels—are degenerate in energy with respect to the hyperfine Hamiltonian, but an external magnetic field lifts this degeneracy. The Zeeman effect [22,35] results in an energy shift proportional to the external magnetic field \mathbf{B} and to a system’s magnetic moment μ :

$$V = -\mu \cdot \mathbf{B}. \quad (2.34)$$

Our atom is a composite particle, made of a nucleus with its own intrinsic magnetic moment, an electron with its own intrinsic one as well, and a contribution from the orbital motion of the electron about the nucleus. Each magnetic moment is proportional to the angular momentum of the subsystem in question, with the proportionality constants,

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called Landé g-factors written as dimensionless multiples of $-\mu_B/\hbar$, where μ_B is the Bohr magneton.²⁰ Since J is a good quantum number so long as energy shifts are smaller than the (large) energy spacing between the three L_J states of the D line, on the level we work we don't consider the electron spin and orbital angular momenta separately, rather we encapsulate them with a single, empirically determined Landé g-factor g_J [35] for the magnetic moment of the electron in each of the three L_J states. Similarly we consider the nucleus as a single spin with an experimentally determined g_I [35], resulting in a Zeeman Hamiltonian:

$$\hat{H}_Z = -\hat{\mu} \cdot \mathbf{B} \quad (2.35)$$

$$= -(\hat{\mu}_I + \hat{\mu}_J) \cdot \mathbf{B} \quad (2.36)$$

$$= \left(\frac{g_I \mu_B}{\hbar} \hat{I} + \frac{g_J \mu_B}{\hbar} \hat{J} \right) \cdot \mathbf{B}. \quad (2.37)$$

Separate g_S and g_L values are known and can be used in two terms instead of the one containing \hat{J} above if J is not a good quantum number, but in the regime we work that is not usually the case (and if it were, the fine and hyperfine structure Hamiltonians above would also be inadequate since they assume that J is a good quantum number). If J is a good quantum number then it is more accurate to use the above expression with empirically measured g_J values, since they encapsulate QED effects and corrections due to the multi-electron structure of ^{87}Rb that are not captured by the simple Zeeman Hamiltonian with separate \hat{S} and \hat{L} terms [35].

If the energy shifts from the Zeeman effect are small compared to the hyperfine splitting, then F (the total spin quantum number, defined in the next section) is a good quantum number and a given hyperfine level can be treated as a single magnetic moment subject to the Zeeman Hamiltonian:

$$\hat{H}_{Z\text{lin}} = \frac{g_F \mu_B}{\hbar} \hat{F} \cdot \mathbf{B}, \quad (2.38)$$

where [35]

$$g_F = g_J \frac{F(F+1) - I(I+1) + J(J+1)}{2F(F+1)} + g_I \frac{F(F+1) + I(I+1) - J(J+1)}{2F(F+1)}. \quad (2.39)$$

The direction in which each Zeeman sublevel shifts in energy for small magnetic fields is depicted in Figure 2.3. Experimentally, Zeeman shifts that depart from this linear regime are not infrequently encountered, and so it is an approximation that cannot always be made.

An explicit matrix representation H_Z^J of \hat{H}_Z in the $\{|m_I m_J\rangle\}$ basis for each of the three L_J states of the D line can be constructed by applying (2.37) to the matrix representations of \hat{I} and \hat{J} in that basis:

$$H_Z^J = -\boldsymbol{\mu} \cdot \mathbf{B}, \quad (2.40)$$

where

$$\boldsymbol{\mu} = -\frac{g_I \mu_B}{\hbar} \hat{I}^J - \frac{g_J \mu_B}{\hbar} \hat{J}^I \quad (2.41)$$

$$= -\frac{g_I \mu_B}{\hbar} \hat{I} \otimes \hat{I}^J - \frac{g_J \mu_B}{\hbar} \hat{J} \otimes \hat{J}^I. \quad (2.42)$$

2.3.4 Putting it all together: the $\{|F m_F\rangle\}$ basis

So far we have described how to construct a matrix representation of the fine structure Hamiltonian for the three L_J states of the ^{87}Rb D line, as well as matrix representations of

²⁰We are using the sign convention that defines the Landé g-factor g_J for the electron as positive.

each state's hyperfine and Zeeman Hamiltonians, these latter two in the same $\{|m_I m_J\rangle\}$ basis. In the subspaces of each of the three L_J states then, we can sum together the hyperfine and Zeeman Hamiltonians to form (a matrix representation of) a Hamiltonian that takes interactions into account:

$$\overset{I \times J}{H}_{P_{3/2}} = \overset{I \times J}{H}_{\text{hfs} P_{3/2}} + \overset{I \times J}{H}_{Z P_{3/2}}, \quad (2.43)$$

$$\overset{I \times J}{H}_{P_{1/2}} = \overset{I \times J}{H}_{\text{hfs} P_{1/2}} + \overset{I \times J}{H}_{Z P_{1/2}}, \quad (2.44)$$

$$\overset{I \times J}{H}_{S_{1/2}} = \overset{I \times J}{H}_{\text{hfs} S_{1/2}} + \overset{I \times J}{H}_{Z S_{1/2}}, \quad (2.45)$$

where the subscripts $P_{3/2}$, $P_{1/2}$, and $S_{1/2}$ on the terms on the right hand side indicate that the expressions for the matrix elements of the hyperfine and Zeeman Hamiltonians are to be evaluated using the specific values of J , A_{hfs} , B_{hfs} , and g_J relevant to that state. The total Hamiltonian for the D line including fine structure, hyperfine structure and the Zeeman interaction is then

$$\hat{H} = \hat{H}_{\text{fs}} + (\hat{H}_{P_{3/2}} \oplus \hat{H}_{P_{1/2}} \oplus \hat{H}_{S_{1/2}}) \quad (2.46)$$

$$\Rightarrow \hat{H} = (\hbar\omega_{D_2} \hat{1}_{P_{3/2}} + \hat{H}_{P_{3/2}}) \oplus (\hbar\omega_{D_1} \hat{1}_{P_{1/2}} + \hat{H}_{P_{1/2}}) \oplus \hat{H}_{S_{1/2}}, \quad (2.47)$$

the matrix representation of which in the $\{|L_J m_I m_J\rangle\}$ basis is the block diagonal matrix

$$\overset{I \times J}{H}_{\text{D}^{87}\text{Rb}} = \begin{bmatrix} [\hbar\omega_{D_2} + \overset{I \times J}{H}_{P_{3/2}}] & & \\ & [\hbar\omega_{D_1} + \overset{I \times J}{H}_{P_{1/2}}] & \\ & & [\overset{I \times J}{H}_{S_{1/2}}] \end{bmatrix}, \quad (2.48)$$

where addition of scalars with matrices implies the addition of a scalar multiple of the appropriately sized identity matrix.

While $\overset{I \times J}{H}_{\text{D}^{87}\text{Rb}}$ is a block diagonal matrix, each of the three submatrices is not diagonal, since m_I and m_J are not good quantum numbers for the hyperfine interaction

(though they are good quantum numbers for the Zeeman interaction and hence $\overset{I \times J}{H}_{\text{D}^{87}\text{Rb}}$ becomes approximately diagonal at high magnetic field). Once one has constructed

$\overset{I \times J}{H}_{\text{D}^{87}\text{Rb}}$ or its submatrices, there are two other bases one might consider transforming the submatrices into depending on the circumstances. One is the $\Sigma\mathcal{F} = \{|F m_F\rangle\}$ basis²¹, in which the matrix representation of the hyperfine interaction is diagonal. The matrix representation of the Zeeman Hamiltonian is also approximately diagonal in the $\{|F m_F\rangle\}$ basis, so long as one is in the linear Zeeman regime. Furthermore, the transition dipole moments for optical transitions are most easily calculated in the $\{|L_J F m_F\rangle\}$ basis, as we'll see in the next subsection. For these reasons the $\{|F m_F\rangle\}$ basis is the most commonly used and referred to.

The $\{|F m_F\rangle\}$ basis is defined as the simultaneous eigenbasis of the \hat{F}^2 and \hat{F}_z operators, which are the total (squared) and z component of the total angular momentum operator $\hat{F} = \hat{\mathbf{I}} + \hat{\mathbf{J}}$, the matrix forms F^2 and F_z of which in the $|m_I m_J\rangle$ basis can be

²¹The name $\Sigma\mathcal{F}$ refers to the fact that this is the basis that lends itself to the interpretation of the total space being the direct sum of subspaces, each of which has a well defined F quantum number. So if $I \times J$ is a basis for a product space with $(2I+1) \times (2J+1)$ dimensions, then $\Sigma\mathcal{F}$ is a basis for a sum space with $\sum_i (2F_i + 1)$ dimensions, where F_i ranges from $|I - J|$ to $I + J$ in integer steps. The product and sum bases have the same dimensionality and span the same space, so the difference is only in the identification and composition of their component subspaces.

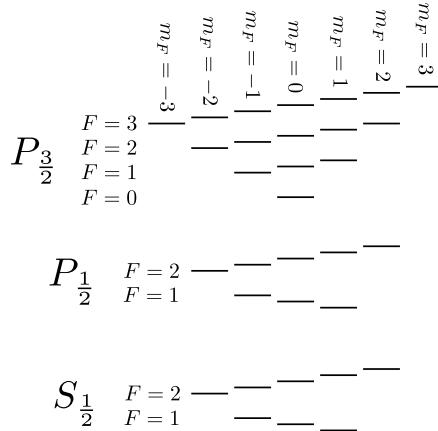


Figure 2.3: The 32 states of the rubidium ^{87}D line, ordered vertically by energy (not to scale) in a small magnetic field. At zero magnetic field, Zeeman sublevels sharing a common F quantum number and within the same hyperfine level are degenerate. At small magnetic fields this degeneracy is lifted, with state energies shifting in the directions depicted. However, F is no longer a good quantum number at nonzero magnetic field, as most of the non-degenerate sublevels are actually equal to linear combinations of two states of different F quantum numbers. Nevertheless at low magnetic fields the states are labelled using F anyway, since one F state dominates the linear combination, and at higher magnetic fields the states are labelled using an index α , equal to the value of F of the state that would dominate the superposition if the field were smoothly reduced to zero. m_F remains a good quantum number at nonzero magnetic field however, and so at all fields a state can be specified by the numbers L, J, α (equal to F at small field) and m_F .

constructed from the matrix forms of the individual angular momentum operators:

$$\begin{aligned} \mathbf{F} &= \mathbf{I} + \mathbf{J} \\ &= \mathbf{I} \otimes \mathbb{I} + \mathbb{I} \otimes \mathbf{J} \end{aligned} \tag{2.49}$$

$$\begin{aligned} \mathbf{F}_z &= \mathbf{I}_z + \mathbf{J}_z \\ &= \mathbf{I}_z \otimes \mathbb{I} + \mathbb{I} \otimes \mathbf{J}_z \end{aligned} \tag{2.50}$$

$$\mathbf{F}^2 = \mathbf{F} \cdot \mathbf{F} \tag{2.51}$$

The $\{|Fm_F\rangle\}$ basis allows the eigenstates of the hyperfine interaction to be labelled with F and m_F quantum numbers. For a state of the ^{87}Rb D line with electron total angular momentum quantum number J , there are $1 + J + I - |(I - J)|$ hyperfine levels, with F quantum numbers running from $|I - J|$ to $I + J$. Within each hyperfine level there are $2F + 1$ degenerate states with different m_F quantum numbers ranging from $-F$ to F . This results in a total of 32 possible states for the rubidium D line, a schematic of which is shown in Figure 2.3.

To transform each submatrix between the $\{|m_I m_J\rangle\}$ and $\{|F m_F\rangle\}$ bases, we use a unitary matrix whose elements are Clebsch–Gordan coefficients, each defined as the inner product of a $|F m_F\rangle$ state with a $|m_I m_J\rangle$ state (written in full as $|I m_I J m_J\rangle$), and

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calculable as [35]

$$\langle I m_I J m_J | F m_F \rangle = (-1)^{I-J+m_F} \sqrt{2F+1} \begin{pmatrix} I & J & F \\ m_I & m_J & -m_F \end{pmatrix} \quad (2.52)$$

where the object in parentheses is a Wigner 3-j symbol. The Clebsch–Gordan coefficients are real: $\langle F m_F | I m_I J m_J \rangle = \langle I m_I J m_J | F m_F \rangle$, and are zero unless $m_I + m_J = m_F$. Given that the possible range of the F quantum number is from $|I - J|$ to $I + J$, and the possible range of m_F quantum numbers is from $-F$ to F for each F (both in integer steps), an explicit construction of the unitary matrix U_{CG} of Clebsch–Gordan coefficients, in the convention where the column vectors in the $\{|F m_F\rangle\}$ basis have F running from its highest value to lowest from top to bottom, and m_F also running from highest to lowest within each value of F , (omitting I and J for brevity) is:

$$U_{CG} = \begin{bmatrix} \langle m_I=I m_J=J | F=I+J m_F=F \rangle & \dots & \langle m_I=I m_J=J | F=I+J m_F=-F \rangle & \dots & \dots & \langle m_I=I m_J=J | F=|I-J| m_F=-F \rangle \\ \vdots & & & & & \vdots \\ \langle m_I=-I m_J=J | F=I+J m_F=F \rangle & & & & \ddots & \\ \vdots & & & & & \vdots \\ \langle m_I=-I m_J=-J | F=I+J m_F=F \rangle & & & \dots & & \langle m_I=-I m_J=-J | F=|I-J| m_F=-F \rangle \\ \vdots & & & & & \end{bmatrix}. \quad (2.53)$$

This the matrix that takes vectors from the $\{|F m_F\rangle\}$ basis to the $\{|m_I m_J\rangle\}$ basis, so its Hermitian conjugate U_{CG}^\dagger is needed for the inverse transformation (which is equal to the transpose since the matrix elements are real). Within the convention we're using to order the matrix elements, the vector representation ψ of a state vector $|\psi\rangle$ in the $\{|F m_F\rangle\}$ basis of one of the L_J states is:

$$\psi = \begin{bmatrix} \langle F=I+J m_F=F | \psi \rangle \\ \langle F=I+J m_F=F-1 | \psi \rangle \\ \vdots \\ \langle F=I+J m_F=-F | \psi \rangle \\ \langle F=I+J-1 m_F=F | \psi \rangle \\ \langle F=I+J-1 m_F=F-1 | \psi \rangle \\ \vdots \\ \langle m_F=|I-J| m_F=-F | \psi \rangle \end{bmatrix}. \quad (2.54)$$

Each submatrix of the total Hamiltonian for the D line can be transformed into its $\{|F m_F\rangle\}$ basis by using a unitary matrix of Clebsch–Gordan coefficients with the appropriate value of J , yielding a total Hamiltonian for the rubidium 87 D line in the $\{|J F m_F\rangle\}$ basis:

$$H_D^{87\text{Rb}} = \begin{bmatrix} \left[\begin{smallmatrix} \hbar\omega_{D_2} + U_{CG\ 3/2}^\dagger H_{P_{3/2}} U_{CG\ 3/2} \end{smallmatrix} \right] & \left[\begin{smallmatrix} \hbar\omega_{D_1} + U_{CG\ 1/2}^\dagger H_{P_{1/2}} U_{CG\ 1/2} \end{smallmatrix} \right] & \left[\begin{smallmatrix} U_{CG\ 1/2}^\dagger H_{S_{1/2}} U_{CG\ 1/2} \end{smallmatrix} \right] \\ \left[\begin{smallmatrix} \hbar\omega_{D_2} + H_{P_{3/2}} \end{smallmatrix} \right] & \left[\begin{smallmatrix} \hbar\omega_{D_1} + H_{P_{1/2}} \end{smallmatrix} \right] & \left[\begin{smallmatrix} H_{S_{1/2}} \end{smallmatrix} \right] \end{bmatrix} \quad (2.55)$$

$$= \begin{bmatrix} \left[\begin{smallmatrix} \hbar\omega_{D_2} + H_{P_{3/2}} \end{smallmatrix} \right] & \left[\begin{smallmatrix} \hbar\omega_{D_1} + H_{P_{1/2}} \end{smallmatrix} \right] & \left[\begin{smallmatrix} H_{S_{1/2}} \end{smallmatrix} \right] \\ \left[\begin{smallmatrix} \hbar\omega_{D_2} + \Sigma_F^P_{3/2} \end{smallmatrix} \right] & \left[\begin{smallmatrix} \hbar\omega_{D_1} + \Sigma_F^P_{1/2} \end{smallmatrix} \right] & \left[\begin{smallmatrix} \Sigma_F^P_{1/2} \end{smallmatrix} \right] \end{bmatrix} \quad (2.56)$$

where the additional subscripts on the unitary matrices indicates the value of J used in its construction. Armed with the $\{|F, m_F\rangle\}$ basis, we have little remaining reason to use the $\{|m_I m_J\rangle\}$ basis any more.

Here is a recap of what we have taken into account with our model of the Rubidium D line:

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- The two lowest quantum numbers $L = 0$ and $L = 1$ of the electron's orbital angular momentum. This yields the S groundstate and P excited state.
- Fine structure: the two possible orientations of the electron's spin with respect to its orbital angular momentum. This splits the P excited state into two states, $P_{1/2}$ and $P_{3/2}$, and leaves the S groundstate as the single state $S_{1/2}$. The energies of these three states are determined entirely empirically—without any modelling of the fine structure.
- Hyperfine structure: The possible orientations of the electron's total angular momentum with respect to the nuclear angular momentum. This splits each state so far into $1 + J + I - |(I - J)|$ hyperfine states. The hyperfine interaction is treated semi-empirically, using an analytic form of the hyperfine interaction but with empirically determined coupling constants within each of the three L_J states of the D line.
- Zeeman effect: the possible orientation of the angular momenta $\hat{\mathbf{I}}$ and $\hat{\mathbf{j}}$ —or at low field their sum $\hat{\mathbf{F}}$ —onto an external magnetic field. The Zeeman Hamiltonian is modelled analytically, but with empirically determined Landé g factors for each of the three L_J states of the D line.

The above considerations allow us to construct a Hamiltonian, as a concrete matrix representation, for all 32 possible states of the rubidium D line in a fixed magnetic field, in a basis in which it is diagonal for the case of zero magnetic field. In a nonzero magnetic field, one can analytically diagonalise the $S_{1/2}$ and $P_{1/2}$ blocks of this matrix using the Breit–Rabi formula [36, p. 347; 41], but the same cannot be done for the entire D line, and diagonalisation must be performed numerically for the $P_{3/2}$ block. At nonzero magnetic field, F ceases to be a good quantum number, but states are nonetheless labelled using a variable α , defined as the value of F a state *would* have if the magnetic field were reduced to zero. m_F remains a good quantum number however, and thus at nonzero magnetic field the $\{|\alpha m_F\rangle\}$ basis, is the one in which the Hamiltonian is diagonal, with the transformation between it and the $\{|F m_F\rangle\}$ basis requiring numerical diagonalisation.

If the magnetic field is static, then this Hamiltonian $\hat{H}_{\text{D}^{87}\text{Rb}}$ is a time independent Hamiltonian, to which additional time-dependent Hamiltonians can be added to take into account optical or RF transitions, detailed below. In a dynamic magnetic field, only \hat{H}_{hfs} is time-independent, in which case the Zeeman Hamiltonian will also be part of the time-dependent part of the Hamiltonian. This has implications for the use of an interaction picture (described in section 3.2.2), in which the time independent part of a Hamiltonian can be analytically removed from the equations of motion, which can make further computations more tractable both analytically and numerically.

2.3.5 Optical dipole transitions

Optical transitions due to laser fields appear as off-diagonal matrix elements in the $\{|F m_F\rangle\}$ basis²². A given laser of a certain polarisation may give rise to nonzero transition matrix elements between several different pairs of states, depending on selection rules.

The potential the outer electron of rubidium is subject to in a classical electric field is the electric dipole potential [20, 36]:

$$\hat{H}_{\text{d}} = -\hat{\mathbf{d}} \cdot \mathbf{E}, \quad (2.57)$$

where \mathbf{E} is the classical electric field and $\hat{\mathbf{d}}$ is the electric dipole operator:

$$\hat{\mathbf{d}} = -e\hat{\mathbf{r}}, \quad (2.58)$$

²²The matrix elements can be obtained in the $\{|\alpha m_F\rangle\}$ basis at nonzero magnetic field via the numerically computed transformation mentioned in the previous subsection

which gives the electric dipole of the atom in terms of the electron's charge $-e$ and its position operator (with respect to the nucleus) $\hat{\mathbf{r}}$. A single matrix element of \hat{H}_d for coupling an initial state $|L_J F m_F\rangle$ to a final state $|L'_J F' m'_F\rangle$ is therefore:

$$\langle L'_J F' m'_F | \hat{H}_d | L_J F m_F \rangle = \langle L'_J F' m'_F | e \hat{\mathbf{r}} \cdot \mathbf{E} | L_J F m_F \rangle, \quad (2.59)$$

which we will abbreviate as

$$\langle n' | \hat{H}_d | n \rangle = \langle n' | e \hat{\mathbf{r}} \cdot \mathbf{E} | n \rangle, \quad (2.60)$$

²³The choice of the symbol n is unrelated to the principal quantum number, which does not vary between the states of the D line.

writing the initial state $|L_J F m_F\rangle = |n\rangle$ and final state $|L'_J F' m'_F\rangle = |n'\rangle$, encapsulating the set of quantum numbers as a single index²³ for the state.

Plane waves

The electric field of a plane wave of amplitude E_0 , polarisation unit vector $\hat{\mathbf{e}}$ and angular frequency ω can be written:

$$\mathbf{E}(\mathbf{r}, t) = \frac{E_0}{2} (\hat{\mathbf{e}} e^{i(\mathbf{k} \cdot \mathbf{r} - \omega t)} + \hat{\mathbf{e}}^* e^{-i(\mathbf{k} \cdot \mathbf{r} - \omega t)}), \quad (2.61)$$

where $*$ is complex conjugation, which in the case of a real valued polarisation unit vector $\hat{\mathbf{e}}$ does nothing, making the above equal to a cosine wave $E_0 \hat{\mathbf{e}} \cos(\mathbf{k} \cdot \mathbf{r} - \omega t)$ as expected for a linearly polarised plane wave. However, (2.61) also allows for arbitrary circular or elliptical polarisation, if $\hat{\mathbf{e}}$ is allowed to have complex components. This comes with the caveat, however, that the amplitude E_0 is a misnomer in these cases and no longer corresponds to any actual amplitude. For example, a circularly polarised plane wave propagating in the z direction constructed using complex polarisation vector $\hat{\mathbf{e}} = \mp(\hat{x} \pm i\hat{y})/\sqrt{2}$ and “amplitude” E_0 actually has an electric field vector with constant magnitude $E_0/\sqrt{2}$. The “amplitude” E_0 in (2.61) is actually equal to $\sqrt{2}E_{\text{rms}}$ and thus only equal to the peak electric field strength in the case of linear polarisation. Using the intensity of the beam²⁴ $I = \epsilon_0 c E_{\text{rms}}^2$ instead, which is less ambiguous and more relatable to experimentalists [37], resolves this potential confusion:²⁵

$$\mathbf{E}(\mathbf{r}, t) = \frac{1}{2} \sqrt{\frac{2I}{\epsilon_0 c}} (\hat{\mathbf{e}} e^{i(\mathbf{k} \cdot \mathbf{r} - \omega t)} + \hat{\mathbf{e}}^* e^{-i(\mathbf{k} \cdot \mathbf{r} - \omega t)}). \quad (2.62)$$

The spherical basis

Consider one of the three polarisation unit vectors $\hat{\mathbf{e}}_q$ for which the photon has a well defined angular momentum projection in the z direction:

$$\hat{\mathbf{e}}_{\pm} = \mp(\hat{x} \pm i\hat{y})/\sqrt{2}, \quad (2.63)$$

$$\hat{\mathbf{e}}_0 = \hat{z}, \quad (2.64)$$

where $q = \pm 1$ is denoted with the subscript \pm . These are the basis vectors of the spherical tensor basis [36], which is the basis in which it is easiest to compute matrix elements, as well as the one in which it is easiest to think about which transitions a given laser will drive. Light with polarisation vector $\hat{\mathbf{e}}_{\pm}$ is called σ^{\pm} polarised, and excites groundstates to excited states that differ in m_F quantum number by ± 1 . A single beam with this polarisation would have to be propagating in the z direction, with left or right hand circular polarisation. Light with polarisation vector $\hat{\mathbf{e}}_0$ is called π polarised, and drives transitions between states of equal m_F quantum numbers. A single beam with this polarisation would have to be propagating with a \mathbf{k} vector in the x, y plane and be linearly polarised in the z direction. However, lasers with propagation directions and polarisation

vectors different from these three configurations can be decomposed into the spherical basis. For example, light with linear polarisation vector in the x, y plane drives both σ^+ and σ^- transitions.

The spherical basis is defined somewhat differently than how we are used to for ordinary real-valued vectors. The components A_q of a vector \mathbf{A} in the spherical basis are defined as the coefficients of *complex conjugates* of the basis vectors:

$$\mathbf{A} = \sum_q A_q \hat{\boldsymbol{\epsilon}}_q^* = \sum_q (-1)^q A_q \hat{\boldsymbol{\epsilon}}_{-q}, \quad (2.65)$$

with the second equality resulting from the property of the basis vectors that

$$\hat{\boldsymbol{\epsilon}}_q^* = (-1)^q \hat{\boldsymbol{\epsilon}}_{-q}, \quad (2.66)$$

and the dot product in terms of the spherical basis components of two vectors is:

$$\mathbf{A} \cdot \mathbf{B} = \sum_q (-1)^q A_q B_{-q}. \quad (2.67)$$

The definition (2.65) has the counter-intuitive consequence that the unit vectors themselves, written in terms of their components $\hat{\boldsymbol{\epsilon}}_q = ((\varepsilon_q)_-, (\varepsilon_q)_0, (\varepsilon_q)_+)$ in the spherical basis are:

$$\hat{\boldsymbol{\epsilon}}_- = (0, 0, -1), \quad (2.68)$$

$$\hat{\boldsymbol{\epsilon}}_0 = (0, 1, 0), \quad (2.69)$$

$$\hat{\boldsymbol{\epsilon}}_+ = (-1, 0, 0), \quad (2.70)$$

that is, the unit vector $\hat{\boldsymbol{\epsilon}}_q$ has a component of $(-1)^q$ in the $-q$ position, and zeros elsewhere.

Considering a plane wave with polarisation vector equal to one of the spherical basis vectors $\hat{\boldsymbol{\epsilon}}_q$, we have the electric field:

$$\mathbf{E}_q(\mathbf{r}, t) = \frac{1}{2} \sqrt{\frac{2I}{\varepsilon_0 c}} (\hat{\boldsymbol{\epsilon}}_q e^{i(\mathbf{k} \cdot \mathbf{r} - \omega t)} + \hat{\boldsymbol{\epsilon}}_q^* e^{-i(\mathbf{k} \cdot \mathbf{r} - \omega t)}), \quad (2.71)$$

which, removing the complex conjugation using (2.66), can be written:

$$\mathbf{E}_q(\mathbf{r}, t) = \frac{1}{2} \sqrt{\frac{2I}{\varepsilon_0 c}} (\hat{\boldsymbol{\epsilon}}_q e^{i(\mathbf{k} \cdot \mathbf{r} - \omega t)} + (-1)^q \hat{\boldsymbol{\epsilon}}_{-q} e^{-i(\mathbf{k} \cdot \mathbf{r} - \omega t)}). \quad (2.72)$$

The dipole approximation

Substituting (2.72) into (2.59) results in the matrix element of the dipole Hamiltonian for a q polarised plane wave:

$$\langle n' | \hat{H}_d(q, I) | n \rangle = \frac{1}{2} \sqrt{\frac{2I}{\varepsilon_0 c}} \left(\langle n' | e \hat{\mathbf{r}} \cdot \hat{\boldsymbol{\epsilon}}_q e^{i(\mathbf{k} \cdot \mathbf{r} - \omega t)} | n \rangle + (-1)^q \langle n' | e \hat{\mathbf{r}} \cdot \hat{\boldsymbol{\epsilon}}_{-q} e^{-i(\mathbf{k} \cdot \mathbf{r} - \omega t)} | n \rangle \right). \quad (2.73)$$

The *dipole approximation* is the approximation that the spatial extent of the electron's orbital is much smaller than the wavelength of the light, and thus that the factors of $e^{\pm i\mathbf{k} \cdot \mathbf{r}}$ are approximately constant over the integral and can be taken outside it, with \mathbf{r} taken to be the expectation value of the atom's position²⁶. For optical wavelengths this is a good

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²⁶The resulting classical position \mathbf{r} of the atom should not be confused with the electron's position operator $\hat{\mathbf{r}}$ with respect to the nucleus. Strictly speaking, (2.73) should have \mathbf{r} in the exponents replaced with a quantum operator with a different name to signify that it is the absolute position of the electron rather than the position relative to the nucleus, but I suspect this would cause more confusion than it

approximation, yielding our matrix element in the dipole approximation:

$$\langle n' | \hat{H}_d(q, I) | n \rangle = \frac{1}{2} \sqrt{\frac{2I}{\varepsilon_0 c}} \left(\langle n' | e \hat{r} \cdot \hat{\epsilon}_q | n \rangle e^{i(\mathbf{k} \cdot \mathbf{r} - \omega t)} + (-1)^q \langle n' | e \hat{r} \cdot \hat{\epsilon}_{-q} | n \rangle e^{-i(\mathbf{k} \cdot \mathbf{r} - \omega t)} \right). \quad (2.74)$$

Evaluating the dot products using (2.67), we get $\hat{r} \cdot \hat{\epsilon}_q = \sum_p (-1)^p \hat{r}_p (\epsilon_q)_{-p} = \hat{r}_q$ and $\hat{r} \cdot \hat{\epsilon}_{-q} = \sum_p (-1)^p \hat{r}_p (\epsilon_{-q})_{-p} = \hat{r}_{-q}$, where \hat{r}_q are the components of the position operator \hat{r} in the spherical basis, the position space representations of which are proportional to the $\ell = 1$ spherical harmonics and the radial coordinate r :

$$\langle r, \theta, \phi | \hat{r}_\pm | r, \theta, \phi \rangle = \mp \frac{r}{\sqrt{2}} \sin \theta e^{\pm i\phi}, \quad (2.75)$$

$$\langle r, \theta, \phi | \hat{r}_0 | r, \theta, \phi \rangle = r \cos(\theta). \quad (2.76)$$

We now have our final expression for the matrix element $\langle n' | \hat{H}_d(q, I) | n \rangle$ of the dipole Hamiltonian for laser intensity I , angular frequency and wavenumber ω and \mathbf{k} respectively, and complex polarisation vector $\hat{\epsilon}_q$, in the dipole approximation in terms of matrix elements of $e \hat{r}_q$, the calculation of which is detailed below:

$$\langle n' | \hat{H}_d(q, I) | n \rangle = \frac{1}{2} \sqrt{\frac{2I}{\varepsilon_0 c}} \left(\langle n' | e \hat{r}_q | n \rangle e^{i(\mathbf{k} \cdot \mathbf{r} - \omega t)} + (-1)^q \langle n' | e \hat{r}_{-q} | n \rangle e^{-i(\mathbf{k} \cdot \mathbf{r} - \omega t)} \right). \quad (2.77)$$

A further approximation often used is to approximate the atom as stationary, replacing the $\mathbf{k} \cdot \mathbf{r}$ from the exponents with a constant (usually zero if the exact phase of the optical field is not of interest). Alternately, if the atom is moving at an approximately constant velocity, the $\mathbf{k} \cdot \mathbf{r}$ term can be absorbed into ω resulting in an effective Doppler-shifted angular frequency as in section 2.1.1. We will leave the $\mathbf{k} \cdot \mathbf{r}$ terms where they are, but note that they are often absent in most literature sources for the preceding reasons.

The rotating wave approximation

If the plane wave's angular frequency ω is close to the resonant angular frequency ω_{D_1} or ω_{D_2} of the D₁ or D₂ line, then one of the two exponential terms is oscillating at a frequency much further from resonance (approximately 2ω) than the other, and thus does not contribute significantly to coupling between states. Discarding it is known as the *rotating wave approximation* (RWA), so called because in an interaction picture (section 3.2.2) where the evolution of the base Hamiltonian $\hat{H}_{D^{87}\text{Rb}}$ is factored out—the so called *rotating frame*—this fact is more readily apparent and a formal approximation can be made showing that the off-resonant term averages to zero over time scales much shorter than any other dynamics. Which term is discarded depends on whether the transition being considered is from a groundstate to excited state, or vice-versa. Our matrix element $\langle n' | \hat{H}_d(q, I) | n \rangle$ in the RWA becomes:

$$\langle n' | \hat{H}_d(q, I) | n \rangle \stackrel{\text{RWA}}{\approx} \begin{cases} \frac{1}{2} \sqrt{\frac{2I}{\varepsilon_0 c}} \langle n' | e \hat{r}_q | n \rangle e^{i(\mathbf{k} \cdot \mathbf{r} - \omega t)}, & E_{n'} > E_n \\ \frac{(-1)^q}{2} \sqrt{\frac{2I}{\varepsilon_0 c}} \langle n' | e \hat{r}_{-q} | n \rangle e^{-i(\mathbf{k} \cdot \mathbf{r} - \omega t)}, & E_{n'} < E_n \end{cases}. \quad (2.78)$$

The spherical basis components \hat{r}_q of the position operator \hat{r} are not Hermitian operators, instead obeying the following relation [36]:

$$\hat{r}_q^\dagger = (-1)^q \hat{r}_{-q}, \quad (2.79)$$

$$\Rightarrow \langle n | e \hat{r}_q | n' \rangle^* = (-1)^q \langle n' | e \hat{r}_{-q} | n \rangle, \quad (2.80)$$

allowing us to simplify (2.78) to:

$$\langle n' | \hat{H}_d(q, I) | n \rangle \stackrel{\text{RWA}}{\approx} \begin{cases} \frac{\hbar}{2} \Omega_{n'n}(q, I) e^{i(\mathbf{k} \cdot \mathbf{r} - \omega t)}, & E_{n'} > E_n \\ \frac{\hbar}{2} \Omega_{nn'}^*(q, I) e^{-i(\mathbf{k} \cdot \mathbf{r} - \omega t)}, & E_{n'} < E_n \end{cases}, \quad (2.81)$$

where $\Omega_{n'n}(q, I)$ is the *Rabi frequency* of the $n \rightarrow n'$ transition for polarisation q and intensity (I),

$$\Omega_{n'n}(q, I) \equiv \hbar \sqrt{\frac{2I}{\epsilon_0 c}} \langle n' | e\hat{r}_q | n \rangle, \quad (2.82)$$

with the symmetry that $\Omega_{n'n}(q, I) = \Omega_{nn'}^*(q, I)$. Because of this symmetry, the electric dipole Hamiltonian for a plane wave in the rotating wave approximation is Hermitian, and all its matrix elements can be computed considering only transitions from groundstate to excited state (with the Rabi frequencies for the reverse transitions being simply the complex conjugates of those of the forward transitions).

The rotating wave approximation is accurate when $\omega - \omega_0 \gg |\omega + \omega_0|$, which is the case for transitions being driven close to resonance. When far off-resonant light is being used however—for example to impart optical dipole forces as in optical traps—it may not be a good approximation, and the full matrix elements (2.77) of the dipole Hamiltonian may need to be considered.

Dipole matrix elements

We now move on to how to compute the dipole matrix elements

$$\langle n' | e\hat{r}_q | n \rangle \equiv \langle L'_J F' m'_F | e\hat{r}_q | L_J F m_F \rangle. \quad (2.83)$$

Using these matrix elements, one can compute Rabi frequencies for matrix elements of the dipole Hamiltonian (2.81) in the rotating wave approximation, or to calculate the full matrix elements (2.77) of the dipole Hamiltonian.

In the spherical basis we've been working in, most of the integrals required to compute these matrix elements can be treated analytically, and there are a series of reductions that can be made to ultimately reduce the matrix elements down to a product of an analytically computable coefficient depending only on angular momentum quantum numbers, and an experimentally measurable ‘reduced’ dipole matrix element.

Following [35], the dependence on the photon angular momentum projection quantum number q and atomic angular momentum projection quantum numbers m_F and m'_F can be encapsulated in a Clebsch–Gordan coefficient using the Wigner–Eckart theorem, allowing us to write:

$$\langle L'_J F' m'_F | e\hat{r}_q | L_J F m_F \rangle = \langle F' m'_F | F 1 m_F q \rangle \langle L'_J F' | e\hat{r} | L_J F \rangle \quad (2.84)$$

$$= (-1)^{F-1+m'_F} \sqrt{2F+1} \begin{pmatrix} F & 1 & F' \\ m_F & q & -m'_F \end{pmatrix} \langle L'_J F' | e\hat{r} | L_J F \rangle, \quad (2.85)$$

where the object in parentheses is a Wigner 3-j symbol, and $\langle L'_J F' | e\hat{r} | L_J F \rangle$ is the ‘reduced’ matrix element for the transition from $|L_J F\rangle$ to $|L'_J F'\rangle$ free of dependence on angular momentum projection quantum numbers. Due to the fact that the Clebsch–Gordan coefficient is zero unless $m_F + q = m_F$, these matrix elements are zero unless the final and initial states are being coupled by a laser of the correct polarisation to conserve angular momentum projection in the z direction, one of the selection rules of optical transitions.²⁷

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²⁷Observe equation (2.77) however, and note that for the case of σ^\pm light, both $q = \pm 1$ dipole matrix elements are present, and thus σ^\pm light couples a groundstate to both $m'_F = m_F \pm 1$ excited states, with one being far off resonance and discarded by the rotating wave approximation for the case of near-resonant light.

The final reduction is to express this reduced matrix element as yet another analytic coefficient multiplied by a reduced matrix element that doesn't depend on the F quantum numbers:

$$\langle L'_J, F' | e\hat{r} | L_J, F \rangle = (-1)^{F+J'+1+I} \sqrt{(2F+1)(2J'+1)} \begin{Bmatrix} J' & J & 1 \\ F & F' & I \end{Bmatrix} \langle L'_J | e\hat{r} | L_J \rangle, \quad (2.86)$$

where I is the nuclear spin quantum number, the object in curly braces is a Wigner 6-j symbol, and $\langle L'_J | e\hat{r} | L_J \rangle$ is a reduced matrix element for the transition from $|L_J\rangle$ to $|L'_J\rangle$ free of dependence on F quantum numbers.

This final reduced matrix element can be linked to experimentally known numerical values via the linewidths Γ_{D_1} and Γ_{D_2} of D₁ and D₂ lines [36, eq. 7.3.7.4]:

$$\Gamma = \frac{\omega_0^3}{3\pi\epsilon_0\hbar c^3} \frac{2J_g + 1}{2J_e + 1} |\langle L'_g | e\hat{r} | L_{J_e} \rangle|^2, \quad (2.87)$$

where Γ and ω_0 are the linewidth and resonant (angular) frequency respectively of the transition, L'_g is the groundstate ($S_{1/2}$) and L_{J_e} is the excited state (either $P_{1/2}$ or $P_{3/2}$) of the transition. Thus the reduced matrix elements for transitions from excited to groundstates can be expressed in terms of known constants. For the reduced matrix elements for transitions from ground to excited states instead, one cannot simply take the complex conjugates of the excited-to-ground reduced matrix elements—the reduced matrix elements are defined instead such that they have the following property [36, eq. 7.3.5.1]:

$$\langle L_J | e\hat{r} | L'_{J'} \rangle = (-1)^{J-J'} \sqrt{\frac{2J'+1}{2J+1}} \langle L'_{J'} | e\hat{r} | L_J \rangle^*. \quad (2.88)$$

Putting these together gives the following reduced matrix elements for the D₁ and D₂ lines in both directions:

$$\langle S_{1/2} | e\hat{r} | P_{1/2} \rangle = \sqrt{\frac{3\pi\epsilon_0\hbar c^3\Gamma_{D_1}}{\omega_{D_1}^3}} \quad (2.89)$$

$$\langle S_{1/2} | e\hat{r} | P_{3/2} \rangle = \sqrt{\frac{6\pi\epsilon_0\hbar c^3\Gamma_{D_2}}{\omega_{D_2}^3}} \quad (2.90)$$

$$\langle P_{1/2} | e\hat{r} | S_{1/2} \rangle = \langle S_{1/2} | e\hat{r} | P_{1/2} \rangle^* \quad (2.91)$$

$$\langle P_{3/2} | e\hat{r} | S_{1/2} \rangle = -\frac{1}{\sqrt{2}} \langle S_{1/2} | e\hat{r} | P_{3/2} \rangle^* \quad (2.92)$$

The first two of these, having only their absolute values related to the linewidths and frequencies via (2.87), have an undetermined phase factor which we have taken to be equal to +1. Since only relative phases are important, if we limit ourselves to only one of the D₁ or D₂ lines at a time, all dipole matrix elements, being proportional to these reduced dipole matrix elements, will have the correct relative phases. However, we are imposing a fixed phase on *both* the reduced dipole matrix elements for the D line. We can verify that a relative phase factor of +1 between the two reduced dipole matrix elements is the correct choice by decomposing [35] the reduced matrix elements into yet another Wigner-6j symbol and a reduced dipole matrix element $\langle L' = \frac{1}{2} | e\hat{r} | L = 1 \rangle$, common to both fine structure lines:

$$\langle L'_J | e\hat{r} | L_J \rangle = (-1)^{J+L+S+1} \sqrt{(2J+1)(2L'+1)} \begin{Bmatrix} L' & L & 1 \\ J & J' & S \end{Bmatrix} \langle L' | e\hat{r} | L \rangle, \quad (2.93)$$

where $S = 1/2$ is the electron spin. This expression can be used to verify the ratio:

$$\frac{\langle S_{1/2} | e\hat{r} | P_{3/2} \rangle}{\langle S_{1/2} | e\hat{r} | P_{1/2} \rangle} = \sqrt{2}, \quad (2.94)$$

confirming that the choice of relative phase factor +1 between the two L_J reduced dipole matrix elements is the correct one.

We have now arrived at the ground floor of reductionism for most experimental atomic physicists when it comes to optical transitions, with the values of the transition dipole matrix elements resting on the foundation of empirically measured reduced dipole matrix elements, and calculable by traversing this section in reverse.

2.3.6 Magnetic dipole transitions

After all the complexity of optical dipole transitions, magnetic dipole transitions—for either RF transitions between Zeeman sublevels, or microwave transitions between hyperfine states—are relatively simple. The matrix elements in the $\{|F m_F\rangle\}$ basis for a perturbing magnetic field $\mathbf{B}(t)$ are simply those of the Zeeman Hamiltonian (2.40):

$$\langle F' m'_F | \hat{H}_Z | F m_F \rangle = -\langle F' m'_F | \hat{\mu} \cdot \mathbf{B} | F m_F \rangle, \quad (2.95)$$

which for a magnetic plane wave with linear polarisation $i \in \{x, y, z\}$ and amplitude B_0 , and in the dipole approximation, is:

$$\langle F' m'_F | \hat{H}_Z(i, B_0) | F m_F \rangle = -\langle F' m'_F | \hat{\mu}_i | F m_F \rangle B_0 \cos(\mathbf{k} \cdot \mathbf{r} - \omega t). \quad (2.96)$$

The magnetic dipole matrix elements $\langle F' m'_F | \hat{\mu}_i | F m_F \rangle$ can be computed using the known matrix form (2.41) of the magnetic moment operator μ_i in the $\{|m_I m_J\rangle\}$ basis, and transforming it into the $\{|F m_F\rangle\}$ basis using a unitary of Clebsch–Gordan coefficients as in (2.52).

As with optical dipole transitions, the cosine can be written as a sum of exponentials, revealing co-rotating and counter-rotating terms:

$$\langle F' m'_F | \hat{H}_Z(i, B_0) | F m_F \rangle = -\frac{B_0}{2} \langle F' m'_F | \hat{\mu}_i | F m_F \rangle (e^{i(\mathbf{k} \cdot \mathbf{r} - \omega t)} + e^{-i(\mathbf{k} \cdot \mathbf{r} - \omega t)}), \quad (2.97)$$

allowing for a rotating wave approximation and definition of a Rabi frequency as with optical dipole transitions:

$$\langle F' m'_F | \hat{H}_Z(i, B_0) | F m_F \rangle \stackrel{\text{RWA}}{\approx} \begin{cases} \frac{\hbar}{2} \Omega_{F' m'_F F m_F}(i, B_0) e^{i(\mathbf{k} \cdot \mathbf{r} - \omega t)}, & E_{F' m'_F} > E_{F m_F} \\ \frac{\hbar}{2} \Omega_{F m_F F' m'_F}^*(i) e^{-i(\mathbf{k} \cdot \mathbf{r} - \omega t)}, & E_{F' m'_F} < E_{F m_F} \end{cases}, \quad (2.98)$$

where $\Omega_{F' m'_F F m_F}(i, B_0) = -\hbar B_0 \langle F' m'_F | \hat{H}_Z(i, B_0) | F m_F \rangle$.

The required linear combinations of the above matrix elements can be used to compute matrix elements for RF waves of different phases, or circular polarisations. It is also often practical, when numerically modelling magnetic dipole transitions, particularly between Zeeman sublevels, to simply include the magnetic field vector as a function as time and evaluate the entire matrix H_Z directly in whichever basis one is working, rather than decomposing it into fixed frequency waves. This is particularly useful for frequency sweeps and other complex magnetic field sequences, which may not be good candidates for the rotating wave approximation, or may not be well represented well as single-frequency waves at all. This is most practical for transitions between Zeeman

sublevels, since the energy splittings can be on the order of MHz and therefore well sampled by numerics with nanosecond to microsecond sized timesteps, which are feasible to numerically compute. Microwave transitions, being 6.8 GHz in ^{87}Rb would require sub-nanosecond timesteps to simulate in this direct way, and so is more taxing but still within the range of possibility, as opposed to optical transitions which are in the hundreds of THz and impractical to simulate without rotating wave approximations and the assumption of monochromatic waves as treated in the previous subsection.

2.3.7 Summary

We are now finally at the end of our assembly of the Hamiltonian for the ^{87}Rb D line in the presence of magnetic fields and monochromatic plane waves. Armed with this Hamiltonian and the ability to compute its coupling terms for various lasers, the atomic physicist can predict and simulate much of what is needed for the basics of laser cooling, trapping, and coherent control.

Quantum mechanics on a computer

THIS CHAPTER COMPRISES a summary of some of the methods used by cold atom physicists to compute numerical results pertaining to cold atom systems. Many a problem in quantum mechanics is not analytically solvable, especially when the real world of experimental physics rears its ugly head, violating theorists' assumptions of simplicity left and right. In particular, atomic physics experiments are time-dependent, with each run of an experiment generally proceeding in stages. Lasers may turn on and off, magnetic fields may vary in magnitude and direction, RF pulses may be chirped to reliably induce particular transitions [42]. Much of the numerical computation performed by researchers in cold atom physics groups such as ours are accordingly of the time-dependent variety, and are fairly literal simulations of specific experiments that may be carried out in the lab.

Here I explain some fundamentals of representing quantum mechanical problems on a computer and then present my favourite algorithms for propagating state vectors and wavefunctions in time. To spoil the surprise: my favourite timestepping algorithms are the 4th-order split-step method and (unoriginally) 4th-order Runge–Kutta, and my favourite method of evaluating spatial derivatives is moderate (6th or so) order finite differences. I think Fourier transforms for spatial derivatives are overrated, and I show that the finite element discrete variable representation (FEDVR) is, despite appearances, actually less computationally efficient than simple finite differences for producing equally accurate solutions to the spatial Schrödinger equation. I also mention some methods of finding groundstates and other stationary states, and in section 3.6 I present a modification to 4th-order Runge–Kutta that enables it to take larger timesteps for certain problems.

3.1 From the abstract to the concrete: neglect, discretisation and representation

To numerically simulate a quantum mechanical system, one must evolve a state vector in time according to the Schrödinger equation:

$$i\hbar \frac{d}{dt} |\psi(t)\rangle = \hat{H}(t) |\psi(t)\rangle. \quad (3.1)$$

To do this on a computer, one must first decide which degrees of freedom are to be simulated. We necessarily neglect many degrees of freedom as a matter of course; which ones can be neglected is warranted by the specific situation and we do it so often we barely notice. For example, simulating a single component Bose–Einstein condensate entails neglecting the internal degrees of freedom of the atoms—as well as reducing the atom–light interaction to a simple potential such as an optical dipole trap or magnetic dipole

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interaction (neglecting the quantum degrees of freedom in the electromagnetic field). We may ignore one or more spatial degrees of freedom as well, say, if we are simulating an experiment in which the condensate is confined to one or two dimensions [43–45] by way of a tight trapping potential in one or more directions. Or, when simulating laser cooling [SEE SECTION ON LASER COOLING SIMS IN ATOMIC PHYSICS CHAPTER], we may care very much about the electronic state of the atom, but treat its motional state classically. In these cases we are essentially imposing the assumption that the system will only occupy one state with respect to those degrees of freedom ignored (the condensate will remain in lowest excitation level in the direction of the tight trap; the atoms will remain in one specific Zeeman sublevel), or we are assuming those degrees of freedom can be treated classically (the electromagnetic field is well described by classical electromagnetism; the atoms' motional state is described well by Newtonian mechanics). Which degrees of freedom can be neglected and which cannot requires knowledge of the situation at hand, often informed by best-practices of the research community in question and ultimately justified by experiment.¹

¹A classic example in the cold atom community of neglected degrees of freedom leading to *disagreement* with experiment is the discovery of polarisation gradient cooling (PGC), the explanation for which requires consideration of Zeeman sublevels of the atoms. The experiment that discovered PGC [46] was designed to measure the effect of Doppler cooling, which does not involve Zeeman sublevels, and it was not until afterwards that theorists determined [25] that transitions between Zeeman sublevels cannot be neglected and indeed are crucial in explaining the lower than predicted temperatures observed.

²The D-line of rubidium 87 has an energy gap of 0.6 eV, requiring a temperature of ≈ 650 K or higher in order for the Boltzmann factor $e^{-\frac{\Delta E}{k_B T}}$ describing the thermal occupation of the excited state to exceed 1×10^{-12} .

Once the degrees of freedom are known, one must decide on a basis in which to represent them concretely. The basis often cannot be complete, since for many degrees of freedom this would require an infinite number of basis states—for example the electronic state of an atom contains a countably infinite number of states, and a spatial wavefunction in free space has an uncountable number of states (one for each position in \mathbb{R}^3). For the internal state of an atom, therefore, we restrict ourselves to only the states we expect can become non-negligibly occupied, given the initial conditions and transitions involved. For example, at low temperature we can expect atoms to be almost completely in their electronic ground states, since energy gaps between ground and excited states are large compared to the thermal energy scale $k_B T$.² We need only include the small number of excited states that might become occupied as a result of optical transitions present in the situation being simulated. This can still be a large number of states if one is studying Rydberg atoms [47, 48] or using ultrafast (and therefore broad-band) laser pulses [49–51], but is otherwise fairly small. For example, including both the D₁ and D₂ lines of Rubidium 87, with all hyperfine levels and Zeeman sublevels gives 32 states (see section [LASER COOLING SIMS IN ATOMIC PHYSICS CHAPTER]).

For spatial degrees of freedom, we usually limit ourselves firstly to a finite region of space (we don't expect the Bose–Einstein condensate to have much probability amplitude on the Moon or anywhere else outside the vacuum system), and then we need to discretise the region of space remaining. To do this one can either discretise space on a grid, or use a set of orthogonal basis functions, and sometimes these can be equivalent, as we will soon see.

Once the degrees of freedom and basis vectors have been chosen, the state vector is then represented on a computer as an array of complex numbers, giving the coefficients of each basis vector required to represent a particular state vector. Matrix elements of the Hamiltonian in the same basis must be calculated, and the Schrödinger equation can then be written:

$$i\hbar \frac{d}{dt} \langle n|\psi(t)\rangle = \sum_m \langle n|\hat{H}(t)|m\rangle \langle m|\psi(t)\rangle, \quad (3.2)$$

or in standard matrix/vector notation (without Dirac notation):

$$i\hbar \frac{d}{dt} \psi_n(t) = \sum_m H_{nm}(t) \psi_m(t) \quad (3.3)$$

$$\iff i\hbar \frac{d}{dt} \psi(t) = H(t) \psi(t), \quad (3.4)$$

where $\psi_n(t) = \langle n|\psi(t)\rangle$, $H_{nm}(t) = \langle n|\hat{H}(t)|m\rangle$ and $\psi(t)$ and $H(t)$ are the vector and matrix with components and elements $\{\psi_n(t)\}$ and $\{H_{nm}\}$ respectively. This is now

something very concrete that can be typed into a computer. Programming languages generally don't know about Dirac kets and operators, and so everything that is to be computed must be translated into matrices and vectors in specific bases. This may seem so obvious as to not be worth mentioning, but was nonetheless a stumbling block in my own experience of getting to grips with quantum mechanics. Once realising that every operator has a matrix representation in some basis, at least in principle (including differential operators), and that every ket is just a list of vector components in some basis (including ones representing spatial wavefunctions), similarly at least in principle, expressions dense in bras and kets become much more concrete as the reader has a feel for exactly how they would type it into a computer. Without a good feel for the mapping between operator algebra and the actual lists of numbers that these objects imply, doing quantum mechanics on paper can seem like an exercise in abstract mumbo-jumbo.

3.2 Solution to the Schrödinger equation by direct exponentiation

As an example of something seemingly abstract being more concrete than first appearances, it is sometimes said that the 'formal' solution to the Schrödinger equation (3.1) is:

$$|\psi(t)\rangle = e^{-\frac{i}{\hbar} \int_{t_0}^t \hat{H}(t') dt'} |\psi(t_0)\rangle. \quad (3.5)$$

Saying that this is the 'formal' solution rather than just 'the solution' is presumably intended to emphasise that the arithmetic operations involved in (3.5) might not make immediate sense for the types of mathematical objects they are operating on, and that we have to be careful in defining the operations such that they produce a result that is not only sensible, but also the solution to the Schrödinger equation. If both $\hat{H}(t)$ and $|\psi(t)\rangle$ were single-valued functions of time rather than an operator valued function of time (the values at different times of which don't necessarily commute) and a vector valued function of time, then we would have no problem. However, (3.5) as written with operators and vectors is ambiguous, and we need to elaborate on it in order to ensure it is correct. I will come back to this after considering a simpler case.

If the Hamiltonian is time-independent, then (3.5) reduces to

$$|\psi(t)\rangle = e^{-\frac{i}{\hbar} \hat{H} \Delta t} |\psi(t_0)\rangle, \quad (3.6)$$

where $\Delta t = t - t_0$. Given the matrix representation H of \hat{H} and vector representation $\psi(t_0)$ of $|\psi(t_0)\rangle$ in a particular basis, this can now be directly typed into a computer as the matrix multiplication:

$$\psi(t) = U(t, t_0) \psi(t_0), \quad (3.7)$$

where

$$U(t, t_0) = e^{-\frac{i}{\hbar} H \Delta t} \quad (3.8)$$

is the (matrix representation of the) unitary evolution operator for time evolution from the initial time t_0 to time t , and is computed using a matrix exponential of $-\frac{i}{\hbar} H \Delta t$. Exponentiation of matrices is defined via the Taylor series of the exponential function:

$$e^A = \sum_{n=0}^{\infty} \frac{A^n}{n!}, \quad (3.9)$$

which reduces matrix exponentiation to the known operations of matrix multiplication and addition. However, any linear algebra programming library worth the bytes it occupies will have a matrix exponentiation function that should be used instead, as there are

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other methods of computing matrix exponentials that are more computationally efficient and numerically stable, such as the Padé approximant [52]. It should be noted that there is no known ‘best’ matrix exponential algorithm, all make compromises and perform poorly for certain types of matrices [53].

3.2.1 Matrix exponentiation by diagonalisation

Regardless of which method is used, matrix exponentiation is computationally expensive. It can be sped up however if a diagonalisation of H is known, since if

$$H = QDQ^\dagger, \quad (3.10)$$

³Note that the diagonals of D are the eigenvalues of H , and the columns of Q are its eigenvectors.

⁴The reason for this is clear from the Taylor series definition of matrix exponentiation, since matrix multiplication and addition can both be performed elementwise for diagonal matrices.

⁵Such as simulating the internal state of a large number of atoms, or evolving a spinor Bose–Einstein condensate by exponentiating the Zeeman Hamiltonian with a spatially varying magnetic field.

where D is a diagonal matrix and Q is a unitary matrix³, then

$$e^{-\frac{i}{\hbar}H\Delta t} = Qe^{-\frac{i}{\hbar}D\Delta t}Q^\dagger. \quad (3.11)$$

This is simple to evaluate because the exponentiation of a diagonal matrix can be performed by exponentiating each diagonal matrix element individually⁴

Even if a diagonalisation of H is not analytically known, numerically diagonalising H (using a linear algebra library function or otherwise) can form the basis for writing your own matrix exponentiation function, if needed. I found this necessary for efficiently exponentiating an array of matrices in Python, since the `scipy` and `numpy` scientific and numeric libraries at the present time lack matrix exponentiation functions that can act on arrays of matrices. Writing a `for` loop in an interpreted language such as Python to exponentiate the matrices individually in many cases is unacceptably slow, so for these cases⁵ I use a function such as the one below:

```

1 import numpy as np
2 from numpy.linalg import eigh
3
4 def expmh(M):
5     """Compute exp(M), where M, shape (... , N, N) is an array of N by N
6     Hermitian matrices, using the diagonalisation method. Made this function
7     because scipy's expm can't take an array of matrices as input, it can only
8     do one at a time."""
9
10    # Diagonalise the matrices:
11    evals, evecs = eigh(M)
12
13    # Now we compute exp(M) = Q exp(D) Q^\dagger where Q is the matrix of
14    # eigenvectors (as columns) and D is the diagonal matrix of eigenvalues:
15
16    Q = evecs
17    Q_dagger = Q.conj().swapaxes(-1, -2) # Only transpose the matrix dimensions
18    exp_D_diags = np.exp(evals)
19
20    # Compute the 3-term matrix product Q*exp_D_diags*Q_dagger using the
21    # einsum function in order to specify which array axes of each array to
22    # sum over:
23    return np.einsum('...ik,...k,...kj->...ij', Q, exp_D_diags, Q_dagger)

```

Matrix diagonalisation (using singular value decomposition or QR decomposition) has computational time complexity $\mathcal{O}(n^3)$, where n is the number of rows/columns in the (square) matrix. Matrix multiplication is (in practice) $\mathcal{O}(n^3)$ and exponentiating a diagonal matrix is only $\mathcal{O}(n)$, so matrix exponentiation of a Hermitian matrix via numerical diagonalisation has total cost $\mathcal{O}(n^3)$. This compares to the Padé approximant, which is also $\mathcal{O}(n^3)$ [53]. So the numerical diagonalisation method is not any worse in terms of computational resources required.

On the other hand, if an analytic diagonalisation is already known, it would seem that exponentiation is just as slow, since the computational cost of matrix multiplication

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alone is the same order in n as that of numerical diagonalisation. This is true—so there are only constant factors to be saved in computer time by using an analytic diagonalisation in order to exponentiate a matrix using (3.11). However if one's aim—as is often the case—is to ultimately compute

$$\psi(t) = e^{-\frac{i}{\hbar} H \Delta t} \psi(t_0), \quad (3.12)$$

for a specific $\psi(t_0)$, then one needs only matrix-vector multiplications and not matrix-matrix multiplications in order to evaluate

$$\psi(t) = U e^{-\frac{i}{\hbar} D \Delta t} U^\dagger \psi(t_0), \quad (3.13)$$

from right-to-left, reducing the computational cost to $\mathcal{O}(n^2)$ compared to evaluating it left-to-right.

Whilst matrix exponentiation is a way to efficiently evolve systems with time-independent Hamiltonians, if you only exponentiate a matrix once, you don't much care about the time complexity of doing so. It is mostly of interest because the real power of these exponentiation methods is as a building block for methods of approximate solutions to the Schrödinger equation in the case of time *dependent* Hamiltonians, as we will see in Section 3.2.3.

3.2.2 The interaction picture

Although time-dependent Hamiltonians are much harder to deal with than time-independent ones, often a Hamiltonian can be decomposed into the sum a time-independent and a time-dependent part:

$$\hat{H}(t) = \hat{H}_0 + \hat{H}_1(t), \quad (3.14)$$

in this case, the Schrödinger equation can be transformed into a form that is often more computationally tractable and analytically useful, by treating part of the evolution due to the time-independent part analytically. This transformation is called an *interaction picture* [54, p. 317], or sometimes, particularly in atomic physics, a *rotating frame*, due to the fact that for a spin- $\frac{1}{2}$ system in a constant magnetic field, time evolution due to \hat{H}_0 entails literal rotation of the spin vector in three dimensional space.

To analytically remove the effect of time evolution due to \hat{H}_0 , we write an ansatz for the state vector $|\psi(t)\rangle$ as an *interaction picture state vector* $|\psi_I(t)\rangle$ that has undergone time evolution due to \hat{H}_0 for time t :

$$|\psi(t)\rangle = e^{-\frac{i}{\hbar} \hat{H}_0 t} |\psi_I(t)\rangle \quad (3.15)$$

Substituting this into the Schrödinger equation yields an equation of motion for the interaction picture state vector $|\psi_I(t)\rangle$:

$$i\hbar \frac{d}{dt} e^{-\frac{i}{\hbar} \hat{H}_0 t} |\psi_I(t)\rangle = [\hat{H}_0 + \hat{H}_1(t)] e^{-\frac{i}{\hbar} \hat{H}_0 t} |\psi_I(t)\rangle \quad (3.16)$$

$$\Rightarrow i\hbar \frac{d}{dt} |\psi_I(t)\rangle = e^{\frac{i}{\hbar} \hat{H}_0 t} \hat{H}_1(t) e^{-\frac{i}{\hbar} \hat{H}_0 t} |\psi_I(t)\rangle \quad (3.17)$$

$$\Rightarrow i\hbar \frac{d}{dt} |\psi_I(t)\rangle = \hat{H}_I(t) |\psi_I(t)\rangle, \quad (3.18)$$

where $\hat{H}_I(t) = e^{\frac{i}{\hbar} \hat{H}_0 t} \hat{H}_1(t) e^{-\frac{i}{\hbar} \hat{H}_0 t}$ is an *interaction picture Hamiltonian*. Thus the interaction picture defines a time-dependent unitary transformation, with the operator for transforming from the Schrödinger to interaction picture being:

$$\hat{U}_I(t) = e^{\frac{i}{\hbar} \hat{H}_0 t}, \quad (3.19)$$

$$(3.20)$$

and state vectors and operators transforming as:

$$|\psi_I(t)\rangle = \hat{U}_I |\psi(t)\rangle \quad (3.21)$$

$$\hat{A}_I(t) = \hat{U}_I \hat{A} \hat{U}_I^\dagger, \quad (3.22)$$

which creates time dependence for operators that did not have time dependence in the Schrödinger picture.

Usually, calculations are performed in a basis in which \hat{H}_0 is diagonal, allowing the unitary transformation to be applied using only scalar exponentials rather than matrix exponentials:

$$\hat{U}(t) = \sum_n e^{\frac{i}{\hbar} E_n t} |n\rangle\langle n| \quad (3.23)$$

$$\Rightarrow \langle n|\psi_I(t)\rangle = e^{\frac{i}{\hbar} E_n t} \langle n|\psi(t)\rangle \quad (3.24)$$

where $\{|n\rangle\}$ and $\{E_n\}$ are the eigenkets and eigenvalues of \hat{H}_0 .

Use of an interaction picture can remove short timescales from the evolution of a state vector, with the original Schrödinger picture state vectors recoverable via analytic transformations, though this is often not necessary since one often only cares about the magnitude of a component $\langle n|\psi\rangle$ of $|\psi\rangle$ —which is preserved by the transformation (3.24)—and not its phase.

In atomic physics calculations an interaction picture can be used to remove the large energy scales (and hence small timescale) separating the three fine structure states of the D_2 line (Section 2.3.1), or it can be used to remove the entire hyperfine interaction Hamiltonian, and even both the hyperfine and Zeeman Hamiltonians (Sections 2.3.2 and 2.3.3) if a time-independent magnetic field is being used. This removes the necessity of small numerical timesteps needed by some numerical methods to resolve very fast phase oscillations due to these time-independent Hamiltonians, as well as aiding human interpretation of the results which would otherwise have more interesting or relevant phase changes drowned out by fast phase oscillations from \hat{H}_0 .

In Section 3.6 I present a method that exploits the use of an *instantaneous* interaction picture, that is, an interaction picture whose time independent Hamiltonian is based on a time *dependent* Hamiltonian evaluated at a specific moment in time. In this way, the use of an interaction picture can be extended to aid calculations for a certain class of time dependent Hamiltonians as well.

3.2.3 Time-ordered exponentials and time-ordered products

As hinted to earlier, the solution (3.5) is not the whole picture. It can only be taken at face value if the Hamiltonian at each moment in time commutes with itself at all other times (we will see shortly why this is). If it does—that is, if $[\hat{H}(t'_1), \hat{H}(t'_2)] = 0$ for all $t'_1, t'_2 \in [t_0, t]$, then (3.5) is the solution to the Schrödinger equation, and once represented in a specific basis can be written

$$\psi(t) = U(t, t_0)\psi(t_0), \quad (3.25)$$

with

$$U(t, t_0) = e^{-\frac{i}{\hbar} \int_{t_0}^t H(t') dt'}, \quad (3.26)$$

i.e. with an integral in the exponent rather than a simple multiplication by a time interval. Since matrix addition can be performed elementwise, so can the integral in the exponent, yielding a matrix which once exponentiated will give the evolution operator $U(t, t_0)$ for the solution to the Schrödinger equation. If the Hamiltonian at each moment in time

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does not commute with itself at all other times, however, then the unitary evolution operator for the solution to the Schrödinger equation is instead given by the following *time-ordered exponential* [55, p. 193]:

$$U(t, t_0) = T \left\{ e^{-\frac{i}{\hbar} \int_{t_0}^t H(t') dt'} \right\}. \quad (3.27)$$

In this expression, T denotes the *time-ordering operator*. The time ordering operator reorders terms within products that contain a time parameter (for us, the time parameter is the argument of the matrix-valued function H), such that the value of the time parameter is smallest in the rightmost term, largest in the leftmost term, and monotonically increasing right-to-left in between. For example:

$$T\{H(4)H(1)H(2)H(5)H(3)\} = H(5)H(4)H(3)H(2)H(1). \quad (3.28)$$

We see now why the time-ordering can be neglected when $H(t)$ commutes with itself at all times. When it does, all possible reorderings of a product of copies $H(t)$ are already equal, and so a time-ordering operator leaves the actual value of the product unchanged.

Despite appearances, this time-ordered exponential is perfectly concretely defined via the definitions of all the operations involved that we have described so far, and can—with some effort—be typed into a computer and evaluated directly. Even though this is not how I have evaluated time-ordered exponentials in my simulations of atomic systems, I'll quickly elaborate on this just to emphasise the concreteness of all these operations.

“What products is T reordering?” you might ask, as (3.27) doesn't appear to contain any products of $H(t)$. On the contrary, it does, since exponentiation is defined by its Taylor series, and so

$$U(t, t_0) = 1 + T \left\{ \sum_{n=1}^{\infty} \frac{1}{n!} \left[-\frac{i}{\hbar} \int_{t_0}^t H(t') dt' \right]^n \right\} \quad (3.29)$$

$$= 1 + \sum_{n=1}^{\infty} \frac{1}{n!} \left(-\frac{i}{\hbar} \right)^n T \left\{ \left[\int_{t_0}^t H(t') dt' \right]^n \right\}. \quad (3.30)$$

Each term in this series contains the n^{th} power (and hence a product) of an integral of $H(t)$. The time ordering operator doesn't allow us to evaluate each term by computing the matrix integral once and then raising it to a power—to do so would violate time-ordering since each integral involves evaluating $H(t)$ at all times. Instead we have to write each product of integrals as the integral of a product:

$$U(t, t_0) = 1 + \sum_{n=1}^{\infty} \frac{1}{n!} \left(-\frac{i}{\hbar} \right)^n \int_{t_0}^t dt'_1 \int_{t_0}^{t'_1} dt'_2 \cdots \int_{t_0}^{t'_{n-1}} dt'_n \cdots T \{ H(t'_1) H(t'_2) \cdots H(t'_n) \}, \quad (3.31)$$

from which we can see exactly which product of matrices the time ordering operator is acting on.

Now we are close to seeing one might evaluate $U(t, t_0)$ numerically by summing each term in the Taylor series up to some order set by the required accuracy. For the n^{th} term, one needs to evaluate an n -dimensional integral over n time coordinates, with each coordinate having the same limits of integration. This can be computed in the usual way an integral is numerically computed,⁶ with the minor change that each time the integrand is evaluated, the terms within it must be re-ordered to respect the required time-ordering. Alternatively, the integration region can be restricted to the region in which the terms

⁶By sampling the integrand on a uniform grid, using a quadrature method, or Monte-Carlo integration [56] which is widely used for high dimensional integrals such as these.

are already time-ordered, and then the total integral inferred by symmetry, which gives:

$$U(t, t_0) = 1 + \sum_{n=1}^{\infty} \left(-\frac{i}{\hbar}\right)^n \int_{t_0}^t dt'_n \cdots \int_{t_0}^{t'_3} dt'_2 \int_{t_0}^{t'_2} dt'_1 H(t'_n) \cdots H(t'_2)H(t'_1). \quad (3.32)$$

This is now a perfectly concrete expression, with each term comprising an integral over an n -simplex⁷ of a product of n matrices.

This expression for the unitary evolution operator is called the Dyson series [57]. It is not generally used for time-dependent simulations, though it is the basis for time-dependent perturbation theory, and sees use in high energy physics [57, 58] for computing transition amplitudes between incoming and outgoing waves in scattering problems (in which U is called the S-matrix). In these problems, H is an interaction Hamiltonian containing terms for all particle interactions being considered. Accordingly, the integrand for the n^{th} term, being a product of n copies of H evaluated at different times, contains one term for each possible sequence of particle interactions. The integral itself can be considered a sum of transition amplitudes over all possible times that each interaction could have occurred. Indeed, each term in the Dyson series corresponds to a number of Feynman diagrams with n nodes [58].

The Dyson series isn't really suited to time-dependent simulations, though perturbation theory is useful for approximate analytics. For one, the series must be truncated at some point, and the result won't be a U that is actually unitary.⁸. Also, we are typically interested in the intermediate states, not just the final state of a system or an average transition rate, as one might want to compute in a scattering problem.

In any case, usually when solving the Schrödinger equation by exponentiation we use the following, alternate expression for a time-ordered exponential:

$$U(t, t_0) = T \left\{ e^{-\frac{i}{\hbar} \int_{t_0}^t H(t') dt'} \right\} \quad (3.33)$$

$$= \lim_{N \rightarrow \infty} \prod_{n=N-1}^0 e^{-\frac{i}{\hbar} H(t_n) \Delta t}, \quad (3.34)$$

where here $\Delta t = (t - t_0)/N$ and $t_n = t_0 + n\Delta t$. Note that the product limits are written in the reverse of the usual order—this is important in order to produce terms with smaller n on the right and larger n on the left of the resulting product. You can convince yourself that (3.32) is equivalent to (3.34) this by replacing the integral in the exponent with a sum—as per the Riemann definition of an integral—and expanding the exponential according to its Taylor series. Expanding each exponential in (3.34) as a Taylor Series and collecting terms of equal powers of H then reveals that the two Taylor series are identical.

In any case, (3.34) paints an intuitive picture of solving the Schrödinger equation: one evolves the initial state vector in time by evolving it according to constant Hamiltonians repeatedly over small time intervals⁹. This has the desirable property that all intermediate state vectors are computed at the intermediate steps, meaning one can study the dynamics of the system and not just obtain the final state. This is of course useful for comparison with experiments, plenty of which involve time-dependent data acquisition and not just post-mortem analysis of some evolution.

Numerically, we can't actually take $N \rightarrow \infty$, or equivalently $\Delta t \rightarrow 0$, and so we instead choose a Δt smaller than the timescale of any time-dependence of H , and step through time using

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⁷A simplex is the generalisation of a triangle to higher dimensions, i.e. a 3-simplex is a tetrahedron.

⁸Although unitarity is not often a strict requirement - we also frequently solve (3.3) directly with fourth order Runge–Kutta, which is also not unitary.

⁹This is the usual definition of the solution to a differential equation, which is why I prefer to think of the product formula (3.34) as the *definition* of the solution to the Schrödinger equation, and the time-ordered exponential as merely a shorthand notation for it.

$$\psi(t_{n+1}) = e^{-\frac{i}{\hbar}H(t_n)\Delta t} \psi(t_n) + \mathcal{O}(\Delta t^2) \quad (3.35)$$

$$\Rightarrow \psi(t) = \left(\prod_{n=N-1}^0 e^{-\frac{i}{\hbar}H(t_n)\Delta t} \right) \psi(t_0) + \mathcal{O}(\Delta t). \quad (3.36)$$

Thus the case of a time-dependent Hamiltonian reduces to repeated application of the solution (3.7) for a time-independent Hamiltonian, and is (globally) accurate to order Δt . Note that the entire expression can be evaluated right-to-left to propagate the initial state vector in time without explicitly computing the overall unitary, which ensures the computational complexity is $\mathcal{O}(n^2)$ in the size of the system (when the exponentiation is performed via an analytic or pre-computed diagonalisation) rather than $\mathcal{O}(n^3)$.

3.2.4 The operator product/split-step method

Here I'll review decompositions similar to (3.35), but which use variable timesteps to achieve an accuracy to higher order in Δt . I'll present them at the same time as addressing another problem, which is that Hamiltonians are often not in a simple enough form to be exponentiated efficiently at all, making (3.35) difficult to evaluate. Often Hamiltonians are a sum of non-commuting operators (such as kinetic and potential terms), with time dependence such that any diagonalisation of the overall Hamiltonian at one point in time will not diagonalise it at another point in time, with the system size large enough for numerical diagonalisation to be prohibitively expensive. In these cases, we can use methods called *split-step* or *operator product* methods, which allow one to approximately exponentiate the entire Hamiltonian based on having exact diagonalisations of its component terms. In the case of time-dependent Hamiltonians, this allows us to avoid rediagonalising at every timestep whenever the time-dependence can be expressed as scalar coefficients multiplying time-independent operators:

$$\hat{H}(t) = \alpha(t)\hat{H}_1 + \beta(t)\hat{H}_2 + \dots, \quad (3.37)$$

since multiplication of a matrix by a scalar merely scales its eigenvalues, leaving its eigenbasis unchanged.

There is little downside to having an only approximate exponentiation of the Hamiltonian when the timestepping is already only approximate, so long as we ensure that neither source of error is much greater than the other. To this end I'll show split-step methods that have (global) error $\mathcal{O}(\Delta t)$, $\mathcal{O}(\Delta t^2)$ and $\mathcal{O}(\Delta t^4)$. In section 3.4, we'll see how this method applies to the case of a spatial wavefunction obeying the Gross–Pitaevskii equation.

First order split-step

Say we have (the matrix representation of) a Hamiltonian that is the sum of two non-commuting terms:¹⁰

$$H(t) = H_1(t) + H_2(t). \quad (3.38)$$

The unitary for the solution to the Schrödinger equation is as before:

$$U(t, t_0) = T \left\{ e^{-\frac{i}{\hbar} \int_{t_0}^t H(t') dt'} \right\}, \quad (3.39)$$

which, without loss of exactness, we can split into N equal time intervals of size Δt and write

$$U(t, t_0) = \prod_{n=N-1}^0 U(t_{n+1}, t_n), \quad (3.40)$$

¹⁰From here on, component terms of the Hamiltonian will be written with general time dependence, even though it is understood that these methods are mostly useful for the case where that time dependence can be written as scalar coefficients multiplying otherwise time-independent matrices.

where again $t_n = t_0 + n\Delta t$ and $\Delta t = (t - t_0)/N$, and where

$$U(t_{n+1}, t_n) = T \left\{ e^{-\frac{i}{\hbar} \int_{t_n}^{t_{n+1}} H(t') dt'} \right\}. \quad (3.41)$$

Evaluating the integral using a one-point rectangle rule gives

$$U(t_{n+1}, t_n) = e^{-\frac{i}{\hbar} H(t_n) \Delta t + \mathcal{O}(\Delta t^2)}, \quad (3.42)$$

in which we were able to drop the time ordering operator because $H(t)$ is only evaluated at a single time. Using the Taylor series definition of the exponential, we can take the error term out of the exponent and write

$$U(t_{n+1}, t_n) = e^{-\frac{i}{\hbar} H(t_n) \Delta t} + \mathcal{O}(\Delta t^2). \quad (3.43)$$

So far all we've done is justify (3.35). Now we'll acknowledge that H is a sum of two terms and write:

$$U(t_{n+1}, t_n) = e^{-\frac{i}{\hbar} (H_1(t_n) + H_2(t_n)) \Delta t} + \mathcal{O}(\Delta t^2). \quad (3.44)$$

Using the Baker–Campbell–Hausdorff formula [55, p. 158], we can expand the exponential as:

$$e^{-\frac{i}{\hbar} (H_1(t_n) + H_2(t_n)) \Delta t} = e^{-\frac{i}{\hbar} H_1(t_n) \Delta t} e^{-\frac{i}{\hbar} H_2(t_n) \Delta t} e^{\frac{1}{2\hbar^2} [H_1(t_n), H_2(t_n)] \Delta t^2 + \mathcal{O}(\Delta t^3)} \quad (3.45)$$

$$\Rightarrow U(t_{n+1}, t_n) = e^{-\frac{i}{\hbar} H_1(t_n) \Delta t} e^{-\frac{i}{\hbar} H_2(t_n) \Delta t} + \mathcal{O}(\Delta t^2). \quad (3.46)$$

So we see that the ‘commutation error’ in (3.46) caused by treating the two terms as if they commute is of the same order in Δt as the ‘integration error’ in (3.43) caused by treating the overall Hamiltonian as time-independent over one timestep, resulting in a method with local error $\mathcal{O}(\Delta t^2)$ and global error $\mathcal{O}(\Delta t)$; the first order split-step method:

$$U(t, t_0) = \prod_{n=N-1}^0 U_1(t_{n+1}, t_n) + \mathcal{O}(\Delta t), \quad (3.47)$$

where

$$U_1(t_{n+1}, t_n) = e^{-\frac{i}{\hbar} H_1(t_n) \Delta t} e^{-\frac{i}{\hbar} H_2(t_n) \Delta t}. \quad (3.48)$$

Using the diagonalisation method of matrix exponentiation discussed in section 3.2.1, this unitary U_1 can be used to step a state vector through time with only matrix-vector multiplications and scalar exponentiation, where separate diagonalisations of H_1 and H_2 are either analytically known or pre-computed, even though a diagonalisation of the total Hamiltonian H may not be feasible.

Crucially, if H_1 and H_2 act on different subspaces of the overall Hilbert space, with respective dimensionalities n_1 and n_2 , that is, $H_1(t) = \tilde{H}_1(t) \otimes \mathbb{I}_{n_2}$ and $H_2(t) = \mathbb{I}_{n_1} \otimes \tilde{H}_2(t)$, where \mathbb{I}_m is the $m \times m$ identity matrix, then the matrix-vector products involved in applying U_1 to a state vector can be much faster ($\mathcal{O}(n_1^2 n_2 + n_1 n_2^2)$ rather than $\mathcal{O}(n_1^2 n_2^2)$) than if the Hamiltonian weren't split into two terms, even if an analytic diagonalisation of the total Hamiltonian were available:

$$U_1(t_{n+1}, t_n) = \left(e^{-\frac{i}{\hbar} \tilde{H}_1(t_n) \Delta t} \otimes \mathbb{I}_{n_2} \right) \left(\mathbb{I}_{n_1} \otimes e^{-\frac{i}{\hbar} \tilde{H}_2(t_n) \Delta t} \right). \quad (3.49)$$

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By applying (3.46) recursively for the case where either H_1 or H_2 is itself the sum of two further terms, (3.48) immediately generalises to the case where H is a sum of an arbitrary number of non-commuting terms:

$$U_1(t_{n+1}, t_n) = \prod_{m=1}^M e^{-\frac{i}{\hbar} H_m(t_n) \Delta t}, \quad (3.50)$$

where $H(t) = \sum_{m=1}^M H_m(t)$.

Second and fourth order split-step

By using a two-point trapezoid rule for the integral in (3.41), evaluating $H(t)$ at the beginning and end of the timestep, one can instead obtain an integration error of $\mathcal{O}(\Delta t^3)$ per step:

$$U(t_{n+1}, t_n) = e^{-\frac{i}{\hbar}(H(t_n) + H(t_{n+1})) \frac{\Delta t}{2}} + \mathcal{O}(\Delta t^3) \quad (3.51)$$

$$= e^{-\frac{i}{\hbar}(H_1(t_n) + H_2(t_n) + H_1(t_{n+1}) + H_2(t_{n+1})) \frac{\Delta t}{2}} + \mathcal{O}(\Delta t^3). \quad (3.52)$$

Then, applying the Baker–Campbell–Hausdorff formula once more, and replacing $H_1(t_{n+1})$ (and similarly for $H_2(t_{n+1})$) wherever it appears in a commutator with the Taylor series $H_1(t_n) + \dot{H}_1(t_n)\Delta t + \mathcal{O}(\Delta t^2)$, one can show that if the individual exponentials are ordered in the following way, then the remaining commutation error is also $\mathcal{O}(\Delta t^3)$:

$$U(t_{n+1}, t_n) = e^{-\frac{i}{\hbar} H_2(t_{n+1}) \frac{\Delta t}{2}} e^{-\frac{i}{\hbar} H_1(t_{n+1}) \frac{\Delta t}{2}} e^{-\frac{i}{\hbar} H_1(t_n) \frac{\Delta t}{2}} e^{-\frac{i}{\hbar} H_2(t_n) \frac{\Delta t}{2}} + \mathcal{O}(\Delta t^3). \quad (3.53)$$

This gives the second order split-step method:

$$U(t, t_0) = \prod_{n=N-1}^0 U_2(t_{n+1}, t_n) + \mathcal{O}(\Delta t^2), \quad (3.54)$$

where

$$U_2(t_{n+1}, t_n) = e^{-\frac{i}{\hbar} H_2(t_{n+1}) \frac{\Delta t}{2}} e^{-\frac{i}{\hbar} H_1(t_{n+1}) \frac{\Delta t}{2}} e^{-\frac{i}{\hbar} H_1(t_n) \frac{\Delta t}{2}} e^{-\frac{i}{\hbar} H_2(t_n) \frac{\Delta t}{2}}, \quad (3.55)$$

or, for an arbitrary number of terms in the Hamiltonian,

$$U_2(t_{n+1}, t_n) = \left(\prod_{m=M}^1 e^{-\frac{i}{\hbar} H_m(t_{n+1}) \frac{\Delta t}{2}} \right) \left(\prod_{m=1}^M e^{-\frac{i}{\hbar} H_m(t_n) \frac{\Delta t}{2}} \right). \quad (3.56)$$

U_2 can be concisely written in terms of U_1 as:

$$U_2(t_{n+1}, t_n) = U_1^\dagger(t_n + \frac{\Delta t}{2}, t_{n+1}) U_1(t_n + \frac{\Delta t}{2}, t_n), \quad (3.57)$$

which is to say it is simply two applications of U_1 , each of duration half a total timestep, and with the multiplication order of the exponentials reversed in one compared to the other. Two shortcuts are immediately apparent when computing U_2 or its action on a vector: firstly, if there is a time-independent term in the Hamiltonian, it should be assigned to H_1 , so that the innermost two exponentials in (3.55) can be collapsed into one; secondly the final exponential of each timestep is identical to the first exponential of the next timestep, and so these can also be collapsed together (being split apart only at points in time when one wishes to sample the state vector).

The fourth order split-step method is much more difficult to derive, with a large number of commutators needing to cancel exactly to ensure the local error cancels out up to fourth order in Δt . It can be stated in terms of the second order split-step method as [59, p. 7; 60, 61]:

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$$U(t, t_0) = \prod_{n=N-1}^0 U_4(t_{n+1}, t_n) + \mathcal{O}(\Delta t^4), \quad (3.58)$$

where

$$\begin{aligned} U_4(t_{n+1}, t_n) = & U_2(t_{n+1}, t_{n+1} - p\Delta t) \\ & \times U_2(t_{n+1} - p\Delta t, t_{n+1} - 2p\Delta t) \\ & \times U_2(t_{n+1} - 2p\Delta t, t_n + 2p\Delta t) \\ & \times U_2(t_n + 2p\Delta t, t_n + p\Delta t) \\ & \times U_2(t_n + p\Delta t, t_n), \end{aligned} \quad (3.59)$$

where $p = 1/(4 - 4^{1/3})$. U_4 comprises five applications of U_2 with timesteps $p\Delta t, p\Delta t, (1 - 4p)\Delta t, p\Delta t$, and $p\Delta t$ respectively. The innermost timestep is backwards in time, since $(1 - 4p) < 0$. But this is no problem, one simply evaluates the expression for U_2 with a negative timestep exactly as written, so long as one reads Δt when written within the above expressions for $U_1(t_f, t_i)$ and $U_2(t_f, t_i)$ as referring to $t_f - t_i$, i.e. the difference between the two arguments to the expression, not its absolute value, and not to the timestep of the method in which it is embedded.

Parallelisability and other speedups for banded matrices

Expressions such as (3.55) don't at first glance appear particularly easy to parallelise, that is, to be evaluated in such a way as to leverage the computing power of multiple CPU cores, GPU cores, or multiple computers. A series of Hamiltonian terms must be exponentiated one by one and multiplied together. Whilst each exponential factor could be evaluated independently, and then all of them multiplied together before being applied to a state vector, this explicit construction of U is very costly compared to merely computing its action on a particular state vector, since the latter allows each term to act only in its specific subspace of the overall Hilbert space, and avoids having to pay the $\mathcal{O}(n^3)$ cost of matrix-matrix multiplication.

When one or more terms in the Hamiltonian has a matrix representation which is banded however, that is, all its entries further than a finite number of elements away from the main diagonal are zero, then those terms may be written as a sum of block diagonal matrices, for example:

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where zero-valued matrix elements outside the band are omitted, and those within the band replaced with dots. In this example, we first take the original 16×16 matrix A , and create four 4×4 nonoverlapping submatrices whose diagonals lie along A 's main diagonal. These submatrices don't quite cover all of A 's nonzero elements, however, so we expand each submatrix (except the final one) from its bottom-right by two elements, making it a 6×6 matrix. The submatrices are no longer non-overlapping, so we take every second one and put it in a matrix B , every other one and put it in a matrix C such that B and C are both block diagonal, divide the elements they have in common by two so we don't double count them, and then declare that $A = B + C$. In the general case, the amount of overlap between the submatrices required to encompass all of A is equal to the bandwidth b of A (in this case $b = 2$ since A has two non-main diagonals on each side of its main diagonal).

Having split a term in the Hamiltonian into two terms, we can simply apply the split operator method as normal, with the same order accuracy in Δt as before. But now the matrices that we wish to exponentiate and have act on a vector are *block diagonal*. This means that the exponentiation of each block can be applied (using the diagonalisation method) separately to the vector, independently of each other block. This enables the application of the exponentials to be performed in parallel—each block submatrix modifies different elements of the vector. So one might store different parts of the state vector on different nodes on a cluster computer, and compute $e^{-\frac{i}{\hbar} C \Delta t} \psi$ in parallel. Then some data exchange between nodes would be necessary before applying $e^{-\frac{i}{\hbar} B \Delta t}$ to the result (See Figure 3.1 for a schematic of how this works).

Whilst this specific banded matrix A is small for the sake of example, in general of course one might have a matrix of any size, allowing for B and C to contain more than two submatrices each. The submatrices have a minimum size of twice the bandwidth b of A , in order to cover all elements of A whilst only sharing elements with their nearest neighbour submatrices (any smaller and they would share elements with their next nearest neighbour submatrices as well, complicating things somewhat). But the only maximum is the size of A itself. So what is the optimal submatrix size? Although the split-step method is still accurate to the same order in Δt no matter how many pieces we split a banded matrix into,

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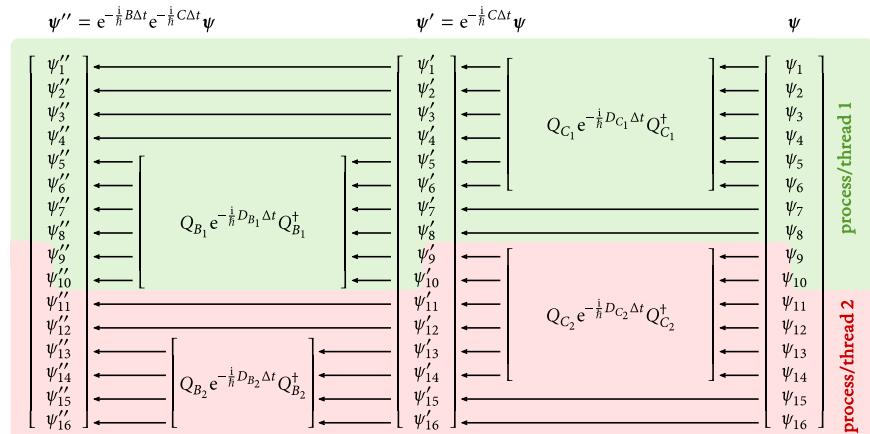


Figure 3.1: Schematic of data flow for parallel split-step method. Computation proceeds right to left. To compute the action of $e^{-\frac{i}{\hbar}B\Delta t} e^{-\frac{i}{\hbar}C\Delta t}$ on a vector ψ , one can treat the submatrices B_1, B_2, C_1 and C_2 , of the block-diagonal matrices B and C separately. First, the exponentiation of C can be applied to ψ by applying the exponentiations of C_1 and C_2 to ψ . If diagonalisations $C_1 = Q_{C_1} D_{C_1} Q_{C_1}^\dagger$ and $C_2 = Q_{C_2} D_{C_2} Q_{C_2}^\dagger$ are known, then the two operations can be applied as a series of matrix-vector multiplications and the exponentiation of diagonal matrices. Because the two submatrices are non-overlapping, they can be applied to the vector completely independently in separate computer processes, CPU or GPU threads, or cluster nodes. The exponentiation of B can then be applied to the resulting intermediate vector ψ' , similarly via two independent operations acting on nonoverlapping elements of ψ' . However, because each submatrix of C overlaps each submatrix of B by $b = 2$ elements on either side (b being the bandwidth of the original matrix $A = B + C$), threads/processes must share these elements (here the ninth and tenth elements), sending them to each other whenever they have been updated. For simplicity this example has two compute threads and two submatrices in each term B and C , but in general there can be any number of either. When a single thread must apply the exponentiation of multiple submatrices to the state vector, it is advantageous for it to first compute the result of any submatrix whose output is required by another thread, then all submatrices that are independent of other threads, and finally any submatrix which requires input from another thread. In this way, data required by other threads can be sent as early as possible, and data needed from other threads called upon as late as possible, minimising the time that threads are waiting for each other whilst there is useful work to be done.

we clearly introduce additional ‘commutation error’ every time we split A into additional submatrices. So one might think that the number of pieces ought to be minimised, and hence the submatrix size maximised. With regard to this, one might decide to split A into $2n_{\text{threads}}$ submatrices (where n_{threads} is the number of independent computational threads available), half of which will reside in B and half in C . This minimises the extra commutation error subject to the constraint that all threads are put to use—splitting A into yet smaller pieces within one computational thread will only yield unnecessary additional error.

Is this the best option? No. Despite the extra commutation error, there is additional benefit to splitting A into more submatrices than required for parallelisation. Let s be the ‘nominal’ size of each submatrix, that is, the size of the corresponding *non-overlapping* submatrices prior to expanding each one along the diagonal (creating overlap) by a number of elements equal to the bandwidth b . The computational cost of the matrix-vector

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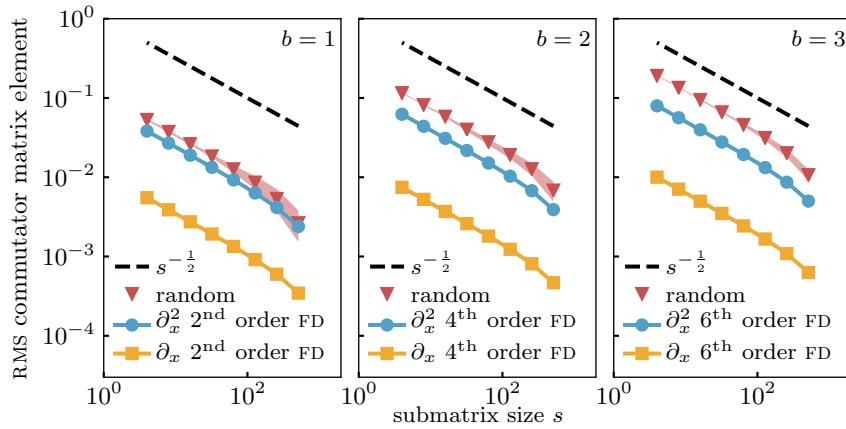


Figure 3.2: Numerical experiment to determine how the commutation error scales with the submatrix size when splitting certain banded matrices into the sum of block-diagonal matrices and using a split-operator method to exponentiate the sum. From left to right, matrices with bandwidth $b = 1$, $b = 1$ and $b = 2$ are considered. For each bandwidth, the matrices for the finite difference schemes of the order resulting in that bandwidth for first and second derivatives are considered, as well as random matrices of the given bandwidth. The random matrices have real and imaginary parts of each element within the band drawn from standard normal distributions. All matrices are 1024×1024 . Calculations with random matrices were performed on 20 independent random matrices, and the mean and standard deviation (shaded) of the results plotted. The error metric is the RMS element of the commutator $[B, C]$, where the orginal matrix A is split into the sum of block-diagonal matrices $B + C$ using a submatrix size of s (described in-text). The RMS error in a vector propagated with the split-operator method is proportional this error metric. The result in all cases is that the commutator error, and therefore the error in the split-operator method scales with the submatrix size as $s^{-\frac{1}{2}}$.

multiplications for computing the action of $e^{-\frac{i}{\hbar}B\Delta t} e^{-\frac{i}{\hbar}C\Delta t}$ on a vector is then $\mathcal{O}((s+b)^2)$ per submatrix, since $s+b$ is the size of the submatrices in terms of their nominal size and the bandwidth. The total number of submatrices required to cover A is ns^{-1} , where n is the size of A , and so the total cost of applying the exponentiations of all submatrices to a vector is $\mathcal{O}(ns^{-1}(s+b)^2)$. The cost *per unit time* of simulation is then $\mathcal{O}(n\Delta t^{-1}s^{-1}(s+b)^2)$. Here we see that splitting into smaller submatrices is desirable from the point of view of speed: because matrix-vector multiplication runs in quadratic time, a larger number of smaller matrices can be multiplied by vectors faster than a smaller number of larger matrices.

But the more submatrices, the more commutation error. Extra error can be made up for by making the timestep smaller, which increases the cost per unit time once more. So is it worth it? The extra commutation error from splitting up A into more and more pieces in general depends on the form of A . I performed a small numerical experiment to see how the commutator $[B, C]$ (which the commutation error is proportional to) scales with s for random banded matrices, as well as those corresponding to 2nd, 4th and 6th order finite differences for first and second derivatives. The result in all cases was that the commutator scaled as $s^{-\frac{1}{2}}$ (see Figure 3.2).

Back to our question—does decreasing the timestep size Δt to compensate for the additional commutation error result in more, or less computational cost per unit time than if we hadn't split A into more pieces than required for parallelisation? Taking into account the nominal submatrix size s and the method's error in terms of Δt , the total

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error of integrating using the split-step method (assuming the $s^{-\frac{1}{2}}$ commutator scaling holds) is $\mathcal{O}(\Delta t^a s^{-\frac{1}{2}})$, where a is the order in Δt of the accuracy of the specific split-step method used (1, 2, or 4 for those I've discussed). Using these two pieces of information: the total error, and the total cost per unit time; we can now answer the question "What value of s minimises the computational cost per unit time, assuming constant error?".

For constant error we set $\Delta t^a s^{-\frac{1}{2}} \propto 1$ and get that the computational cost per unit time at constant error is $\mathcal{O}(ns^{-\frac{1}{2a}-1}(s+b)^2)$. For $a > \frac{1}{2}$, this expression has a minimum at $s = b$, which is the smallest possible submatrix size in any case. For our example banded matrix A , decomposition into the smallest possible submatrices looks like this:

So the conclusion is: use the smallest submatrices possible when decomposing a banded matrix into a sum of two block-diagonal matrices. The decrease in computational costs, even when the increased error is compensated for by a smaller timestep, is worth it.

Limitations and nonlinearity

The split-step method is quite powerful and general. It allows you to approximately decompose the exponentiation of a Hamiltonian into exponentiations of its component terms, in the subspaces that they act on, and avoids exponentiating large banded matrices; saving on computing power immensely compared to exponentiating the full Hamiltonian. It is unitary (when the Hamiltonian is actually Hermitian), and stable—that is, unlike Runge–Kutta methods, the method’s truncation error does not grow without limit as the simulation proceeds but is bounded (disregarding floating point rounding error, which is much smaller than the truncation error of either method) [62]. This is extremely appealing. And, whilst higher order split-step methods quickly become unwieldy [62], fourth order accuracy is quite acceptable for many problems.

Applying all the tricks described in the above sections results in $\mathcal{O}(\Delta t)$, $\mathcal{O}(\Delta t^2)$ and $\mathcal{O}(\Delta t^4)$ accurate timestepping methods for solving the Schrödinger equation with total computational cost scaling as $\mathcal{O}(n \sum_i b_i)$, where $n = \prod_i n_i$ (for a product space of subspaces with dimensionalities $\{n_i\}$) is the dimensionality of the total Hilbert space, and $\{b_i\}$ are the bandwidths of the matrix representations of each term in the Hamiltonian in the chosen basis.¹¹ Furthermore, the method is efficiently parallelisable, provided the

¹¹Now that we are here, we can finally say something about the best basis for simulating in: to minimise computational costs, the best basis is the one that minimises precisely this sum of bandwidths. The exception to this is when fast Fourier transforms are involved, which I discuss later.

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maximum size-to-bandwidth ratio $\max_i(n_i/b_i)$ of the terms in the Hamiltonian is much larger than the number of parallel computing threads available. Although the constant factors that big-O notation neglects may not be optimal, this scaling would seem to be the best one could hope for—for each of the n elements in the state vector one must consider the $\sum_i b_i$ elements (including itself) that the Hamiltonian couples it with, and no more.¹²

The main downside of this otherwise excellent method of exponentiating Hamiltonians is that the evolution modelled must actually be described by a linear system of equations. One cannot add arbitrary terms and nonlinear operators to the Hamiltonian, as the split-step method requires that one can evaluate each time-dependent term in the Hamiltonian at specific times, including times at which the solution for the state vector is not yet available. This would seem to limit the split-step method strictly to modelling *linear* dynamics, that is, terms in the Hamiltonian must not depend explicitly on the state vector they are operating on. Whilst nature might fundamentally be described by linear dynamics, once approximations of various kinds are made in order to make problems tractable, it's common to end up with a nonlinear pseudopotential or nonlinear effective Hamiltonian. Note that non-*Hermitian*, pseudo-Hamiltonians—leading to non-unitary evolution—are fine. The split-step method has made no assumptions that rules them out, it only assumes the differential equation can be expressed in the form

$$\frac{d}{dt}\psi(t) = \sum_n H_n(t)\psi(t), \quad (3.62)$$

where $\{H_n\}$ are a set of linear operators, Hermitian or not.

Fortunately, one specific form of nonlinearity that cold atom physicists are particularly interested in—the nonlinear term in the Gross–Pitaevskii equation—can be incorporated without much difficulty. As mentioned, the problem with nonlinearity is that all but the first-order split step methods require you to evaluate terms in the Hamiltonian at some future time at which the state vector is not yet known, that is the algorithm contains steps akin to $\psi(t_{n+1}) = U(t_{n+1}, t_n; \psi(t_{n+1}))\psi(t_n)$ for some nonlinear unitary matrix $U(t_{n+1}, t_n; \psi(t_{n+1}))$ — U cannot be explicitly constructed because it both requires and is required by $\psi(t_{n+1})$.

For the Gross–Pitaevskii effective Hamiltonian, the second-order split-step method (from which the fourth order method is constructed) for a single step might be naively written

$$\psi(t_{n+1}) = e^{-\frac{i}{\hbar}K\frac{\Delta t}{2}} e^{-\frac{i}{\hbar}(g\rho(t_{n+1})+V(t_{n+1}))\frac{\Delta t}{2}} e^{-\frac{i}{\hbar}(g\rho(t_n)+V(t_n))\frac{\Delta t}{2}} e^{-\frac{i}{\hbar}K\frac{\Delta t}{2}} \psi(t_n), \quad (3.63)$$

where $\psi(t)$ is the state vector¹³ represented in a discrete position basis, g is the nonlinear interaction constant, $V(t)$ and $\rho(t) = \psi(t)\psi^\dagger(t)$ are diagonal matrices for the external potential and the density matrix for $\psi(t)$ in the same position basis, and K is a discrete approximation to the kinetic energy operator in the position basis.

As written, this can't be evaluated because $\psi(t_{n+1})$ —required to evaluate $\rho(t_{n+1})$ —is not yet known. The order we choose to exponentiate the terms in our Hamiltonian is arbitrary, however (so long as we alternately reverse that order each half-step as required by the second order split-step method), and so swapping the order gives

$$\psi(t_{n+1}) = e^{-\frac{i}{\hbar}(g\rho(t_{n+1})+V(t_{n+1}))\frac{\Delta t}{2}} e^{-\frac{i}{\hbar}K\Delta t} e^{-\frac{i}{\hbar}(g\rho(t_n)+V(t_n))\frac{\Delta t}{2}} \psi(t_n), \quad (3.64)$$

which incidentally has the benefit that since K is (ordinarily) time independent, the two adjacent exponentials containing it can be combined into one. Now $\rho(t_{n+1})$ is contained within the leftmost exponential, and so it is the last operator to be applied in the timestep. Note that since $g\rho(t)$ is real and diagonal in the position basis (as is $V(t)$), this leftmost unitary merely changes the phase of the state vector at each point in space,

¹²Assuming the banded matrices are otherwise dense within their band—further improvements would be possible if some elements within the band were always zero.

¹³Usually when modelling the GPE the single-particle state vector is normalised to the number of particles, rather than unity, and so strictly speaking it cannot be called a state vector, though it otherwise can be treated as one in most respects.

having no effect on its density. This means that $\rho(t)$ is, in fact, unaffected by this last unitary evolution operator. Hence, $\rho(t_{n+1})$ can be computed simply as the density matrix of the intermediate state vector that this unitary was to act on:

$$\rho(t_{n+1}) = \psi(t_{n+1})\psi^\dagger(t_{n+1}) = \tilde{\psi}\tilde{\psi}^\dagger, \quad (3.65)$$

where

$$\tilde{\psi} = e^{-\frac{i}{\hbar}K\Delta t} e^{-\frac{i}{\hbar}(g\rho(t_n) + V(t_n))\frac{\Delta t}{2}} \psi(t_n). \quad (3.66)$$

The inclusion of $V(t)$ with $g\rho(t)$ is optional, and if $V(t)$ was not real valued, would not be valid (since in that case the density *would* be affected by the evolution induced by $V(t)$). In such a case the Hamiltonian would have to be split into three terms with $V(t)$ and $g\rho(t)$ treated separately.

So now we can put some conditions on what types of nonlinear operators can be used within the second-order split step method. The first condition is that at most one nonlinear operator can be included, since it must be placed last in the sandwich of exponentials (otherwise its value at the end of the timestep cannot be inferred immediately prior to acting on the state vector). The second condition is that the nonlinear operator must be invariant with respect to the evolution that it itself induces in the state vector. Here we have an operator that depends only on the state vector's absolute value, but for which the corresponding unitary only evolves the state vector's phase. Another example might be an operator that depends only on the state vector's phase gradients, but evolves the state vector's absolute value, and so forth.

Although I find this argument compelling for second-order split-step, it's less obvious that it should hold for fourth-order split-step as well, which, even though it is based on multiple applications of second-order split-step, involves a substep that is backwards in time. 'Evaluate the nonlinear operator based on the state vector at this specific time' becomes ambiguous when that moment in time is traversed in both directions by two different sub-steps. However, [63] have verified using computer symbolic algebra that indeed, up to even higher order split-step methods, putting the nonlinear density term last in the splitting and always evaluating it using the value of the intermediate state vector immediately prior results in the method having the same order accuracy in Δt as for linear operators only. Given this and my argument above, as well as the reasoning that the second-order split-step method has no way of 'knowing' whether it is acting backward in time or not when embedded in a higher-order split step method, I would expect the same to hold for all nonlinear operators meeting the above two conditions, though I haven't shown this explicitly.

3.3 For everything else, there's fourth-order Runge–Kutta

Fourth-order Runge–Kutta (RK4) is the enduring workhorse of numerical integration methods. Of the Runge–Kutta methods, it offers a good balance of accuracy and computational cost. When a problem does not have the properties that allow manifestly unitary or error-bounded methods to be used, or when enough computing power can be deployed so as to make these concerns irrelevant, and the programmer's time more important, fourth order Runge–Kutta is a good choice. Its downsides are that it is not manifestly unitary, and its error is not bounded. Nonetheless for many problems this is not a concern in practice.

The advantages of fourth-order Runge–Kutta are compelling: it has global error that is fourth order in the integration timestep, involves four function evaluations per timestep, requires only linear arithmetic operations outside of the function evaluations,

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and can be applied any problem that can be written in the form

$$\frac{d}{dt} \mathbf{x} = f(\mathbf{x}, t), \quad (3.67)$$

for some (possibly nonlinear) function f , where \mathbf{x} is a vector of dynamical variables, often in our case the components of a state vector, or for classical dynamics the positions and velocities of an ensemble of particles¹⁴. Note that this formulation allows for initial value problems with coupled ODES, discretised PDES, as well as second or higher order differential equations, since an equation of the form

$$\frac{d^2}{dt^2} \mathbf{x} = f(\mathbf{x}, t) \quad (3.68)$$

can be rewritten

$$\frac{d}{dt} (\mathbf{x}, \dot{\mathbf{x}}) = (\dot{\mathbf{x}}, f(\mathbf{x}, \dot{\mathbf{x}}, t)), \quad (3.69)$$

treating the time derivative of each element of \mathbf{x} as simply another coupled dynamical variable.

Not unrelated to its ease of implementation, the method itself can be stated concisely. Propagation of the dynamical variables for one timestep from time t to time $t + \Delta t$ is computed as follows:

$$\begin{aligned} \mathbf{k}_1 &= f(\mathbf{x}(t), t), \\ \mathbf{k}_2 &= f(\mathbf{x}(t) + \frac{1}{2}\mathbf{k}_1\Delta t, t + \frac{1}{2}\Delta t), \\ \mathbf{k}_3 &= f(\mathbf{x}(t) + \frac{1}{2}\mathbf{k}_2\Delta t, t + \frac{1}{2}\Delta t), \\ \mathbf{k}_4 &= f(\mathbf{x}(t) + \mathbf{k}_3\Delta t, t + \Delta t), \\ \mathbf{x}(t + \Delta t) &= \mathbf{x}(t) + \frac{\Delta t}{6}(\mathbf{k}_1 + 2\mathbf{k}_2 + 2\mathbf{k}_3 + \mathbf{k}_4). \end{aligned} \quad (3.70)$$

Each step is self-contained, that is, the algorithm does not contain any state from previous steps. This is appealing as it means that f can change discontinuously between timesteps without giving rise to Runge's phenomenon [65] in the approximate solutions $\mathbf{x}(t)$, as can be the case with multistep methods which do retain some dependency on previous steps. $\mathbf{x}(t)$ can also be modified discontinuously between steps without causing problems, as is required by Monte Carlo wavefunction or quantum jump methods [66,67], or even in the imaginary time evolution method (see section 3.5.1) which may require normalisation of the state vector in between steps. Fourth order Runge–Kutta is therefore quite compatible with stochastic processes and many other models which may not be described by a differential equation alone.

One downside is that the timestep used must be much smaller than the timescale on which \mathbf{x} changes - even if \mathbf{x} 's time variation is highly regular. If \mathbf{x} 's time variation is dominated by the simple accumulation of complex phase at some angular frequency—as is often the case in quantum mechanics—the timesteps used for RK4 must be small enough to resolve these circles about the complex plane. This is in contrast to the exponentiation methods, for which an energy offset (and hence overall change in the angular frequency at which the state vector's elements accumulate phase) is largely irrelevant. Energy offsets or use of an interaction picture can mitigate this problem but requires some foresight, and may not be possible if the required energy offsets change in time or are not known analytically in advance. The method I develop in section 3.6 is a partial remedy for this specific problem, which in my opinion is the biggest weakness of fourth order Runge–Kutta as applied to quantum state evolution, when compared to the unitary methods.

¹⁴For classical particle dynamics, often lower order symplectic methods such as the leapfrog method [64] can be preferable, but nonetheless it is hard to overstate the usefulness of a general purpose algorithm such as RK4 that can be deployed as a first attempt, or as a plan B once the assumptions required by another algorithm are violated

3.3.1 Complexity and parallelisability for the Schrödinger equation

Excluding the evaluation of f , RK4 requires a number of arithmetic operations proportional to the number of elements in \mathbf{x} , and so contributes time-complexity $\mathcal{O}(n)$ to the overall calculation, where n is the number of elements in \mathbf{x} . Since f itself usually doesn't run in linear time, the computational time complexity of RK4 is usually therefore that of evaluating f . Its parallelisability also comes down to that of f , since equations (3.70) above treat each element of \mathbf{x} completely independently, with any couplings computed within f .

For the Schrödinger equation in a concrete basis, f is

$$f(\psi, t) = -\frac{i}{\hbar} H(t)\psi, \quad (3.71)$$

¹⁵For the Gross–Pitaevskii equation this would be a nonlinear $H(\psi, t)$, and everything else in this section still applies.

where ψ is the state vector in the given basis and $H(t)$ is the matrix representation of the Hamiltonian in that basis¹⁵. Fourth order Runge–Kutta is therefore as computationally expensive and parallelisable as computing the matrix–vector product $H(t)\psi$. In the worst case this multiplication is $\mathcal{O}(n^2)$ in the size of the Hilbert space and barely parallelisable at all, if H is dense. But in the common case (as mentioned in section 3.2.4), of H being a sum of terms which act on different subspaces, i.e.

$$H(t) = \sum_{i=1}^N \mathbb{I}_{n_1} \otimes \cdots \otimes \mathbb{I}_{n_{i-1}} \otimes H_i(t) \otimes \mathbb{I}_{n_{i+1}} \otimes \cdots \otimes \mathbb{I}_{n_N} \quad (3.72)$$

where n_i is the dimensionality of the subspace acted on by the i^{th} term and \mathbb{I}_m is the $m \times m$ identity matrix, then things are much better. If each term is dense, then the overall cost of evaluating the product $H(t)\psi$ is $\mathcal{O}(n \sum_i n_i)$, where $n = \prod_i n_i$ is the dimensionality of the total Hilbert space, which can be considerably less than the $\mathcal{O}(n^2)$ of evaluating the single matrix–vector product for the total Hamiltonian. And if each term is a banded matrix with bandwidth b_i , then the cost of applying a single term in the Hamiltonian becomes $\mathcal{O}(nb_i)$ instead of $\mathcal{O}(nn_i)$, and so the cost of evaluating $H(t)\psi$ becomes $\mathcal{O}(n \sum_i b_i)$. This is identical to the earlier result for the split-operator method once the trick of splitting up banded matrices into block-diagonal matrices was applied (section 3.2.4), but with much less work (in the sense of programmer effort rather than computational complexity). Representing $H(t)$ as a sum of terms over different subspaces is no extra work—this is the form we are likely to be writing Hamiltonians in already, and constructing the total $H(t)$ is often not required if one desires only to propagate a state vector in time.¹⁶ No, we are already applying these operators to different subspaces of the Hilbert space and then summing the results without thinking twice about it, and so the above analysis mostly serves as a reminder that what we are already doing most of the time is in fact very efficient.

Parallelisation, provided one or more of the matrices are banded, is also straightforward, since if A is banded with bandwidth b , then

$$(A\psi)_i = \sum_{j=-b}^b A_{i,i+j} \psi_j, \quad (3.73)$$

that is, calculation of an element of $A\psi$ requires only the corresponding element of ψ and its nearest b neighbours on each side. One can therefore divide up the state vector into contiguous regions (in the subspace in which A acts), and compute different elements of the product on different compute threads, requiring an exchange of only b elements at each boundary¹⁷ between neighbouring threads. An example of this is to split two dimensional space into into a number of pieces in each dimension (resulting in a 2D grid) so as to compute the application of (finite difference approximations to) the x and y second derivative operators on a state vector in parallel.

¹⁶A note about minimising the effect of latency: at the start of an RK4 substep, have each thread send the required elements at the edges of its region in each subspace to its neighbouring threads first. Then compute $H(t)\psi$ on the interior elements. If each thread has a sufficiently large workload, then by the time this is complete, the data from neighbouring regions will have arrived and threads will not have spent any time waiting for each other.

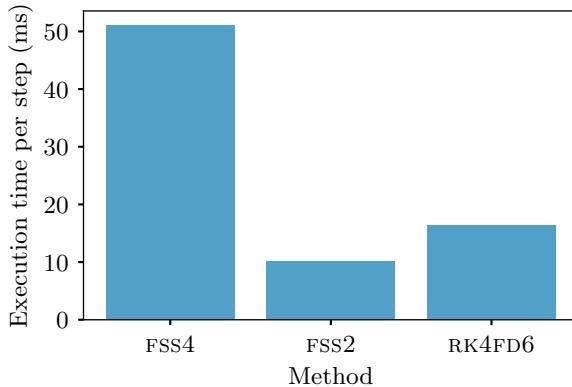


Figure 3.3: The average computation time per step of second and Fourth order Fourier split-step (see Section 3.4.1) and fourth-order Runge–Kutta using sixth order finite differences (see Section 3.4.2) for the simulations detailed in Figure 3.4. It should be noted that the use of finite differences is not significant—as mentioned in Section 3.4.2, it is difficult to observe any difference between the speed of finite differences and Fourier derivatives in practice on present computers. The difference in speed is almost entirely due to the cost of computing exponentials, as can be seen by comparing FSS2 and FSS4, which differ by a factor in five both in speed and in the number of exponentials in their implementations. For fourth order accuracy, RK4 is about three times faster than Fourier split-step.

An additional benefit of parallelised RK4 when compared to split-operator is that the same amount of data needs to be sent between threads regardless of the number of terms in the Hamiltonian. In split-step methods, each term in the Hamiltonian is exponentiated separately (multiple times for the higher order schemes), requiring an exchange of data each time. The amount of data needing to be exchanged per step therefore scales with the number of terms in the Hamiltonian being parallelised, whereas for RK4 it is constant.

The constant factors that big-O notation disregards also favour RK4 when compared to split-operator methods. For one, addition and multiplication are much cheaper than exponentiation, meaning that the exponentials in split-operator methods may add considerable computational cost if the Hamiltonian itself is simple. Furthermore, parallelised or not, each term in fourth order split-operator adds 20 matrix-vector products in the space the term acts, whereas RK4 requires only one additional matrix-vector product per term. In practice due to these properties, RK4 runs considerably faster (see Figure 3.3) than split-operator methods, even for simple systems, and the gap widens as complexity increases. This combined with the relatively minor improvement in accuracy or stability (see Figure 3.4) between the methods make RK4 an enduring choice for cold atom problems.

3.4 Continuous degrees of freedom

The single-particle, non-relativistic, scalar Schrödinger wave equation, as distinct from the general Schrödinger equation (3.1), is:

$$i\hbar \frac{\partial}{\partial t} \psi(\mathbf{r}, t) = \left[-\frac{\hbar^2}{2m} \nabla^2 + V(\mathbf{r}) \right] \psi(\mathbf{r}, t). \quad (3.74)$$

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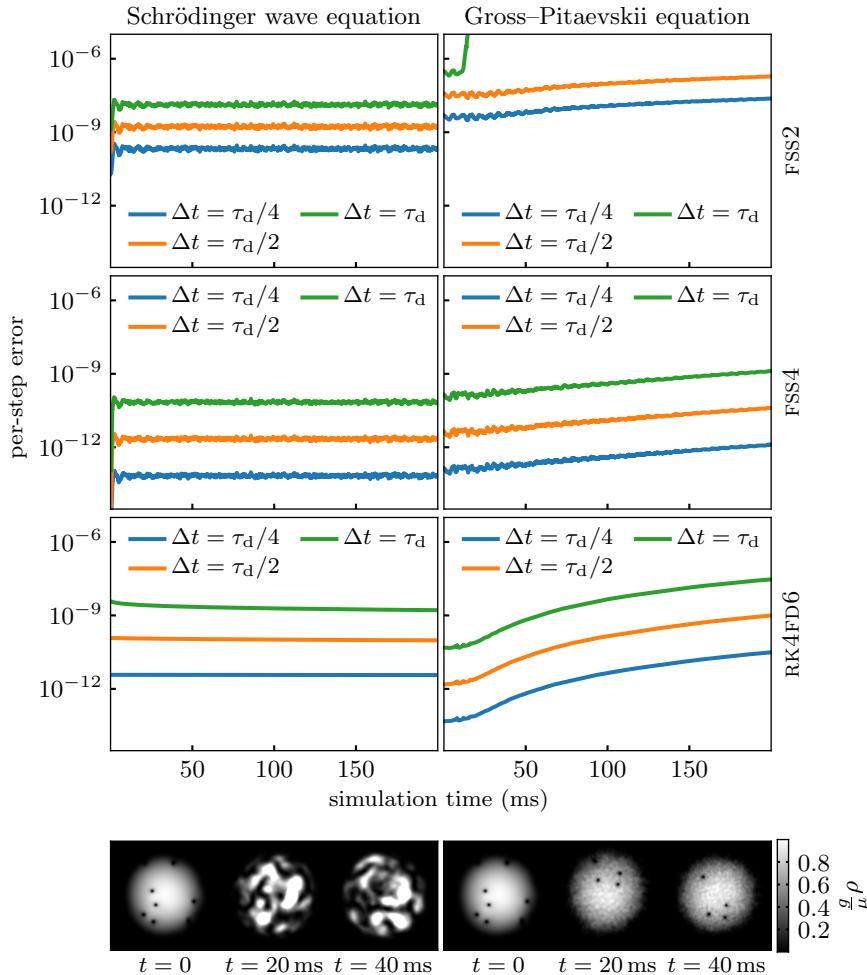


Figure 3.4: A comparison of the RMS per-step accuracy of second and Fourth order Fourier split-step (see Section 3.4.1) and fourth-order Runge–Kutta using sixth order finite differences (see Section 3.4.2). In the left column, the linear Schrödinger equation is solved for a 2D wavefunction and on the right the Gross–Pitaevskii equation for a 2D condensate wavefunction, both for the case of ^{87}Rb . The spatial grid is 256×256 and represents a spatial region of $10 \mu\text{m}$. The same initial conditions are used for each, produced in the following way: first a Gross–Pitaevskii groundstate (for the sake of comparison, even though this is not a groundstate of the Schrödinger equation) in a harmonic trap chosen to give a Thomas–Fermi radius of $7.5 \mu\text{m}$ is produced using successive overrelaxation (Section 3.5.2). Vortices are created by multiplying by a spatially varying complex phase factor to produce the appropriate phase winding for a number of randomly chosen vortices, as described in the example in Section 3.6. This is then relaxed using imaginary time evolution (Section 3.5.1) of the Gross–Pitaevskii equation for a duration equal the chemical potential timescale $\tau_\mu = \frac{2\pi\hbar}{\mu}$ to produce a physically realistic initial condition. This is then evolved in time using the three methods, with three different timesteps, expressed as multiples of the dispersion timescale (described in Section 3.4.4) $\tau_d = \frac{m\Delta x^2}{\pi\hbar}$ where Δx is the grid spacing. Second order Fourier split-step, although unitary and thus having bounded error, reaches that bound for the smallest timestep. The two fourth order methods have comparable accuracy, with Fourier split step being slightly more accurate and the error in RK4 tending to increase faster with time (for the case of the Gross–Pitaevskii equation). Comparing with Figure 3.3, the difference is even more marginal if we compare by equal computational cost per unit simulation time—which allows RK4 to take timesteps about three times smaller before it consumes the same computational resources as Fourier split step.

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Similarly, as mentioned in Section 2.2, the equation for the single-particle wavefunction of an atom in a single-component Bose–Einstein condensate is the Gross–Pitaevskii equation

$$i\hbar \frac{\partial}{\partial t} \psi(\mathbf{r}, t) = \left[-\frac{\hbar^2}{2m} \nabla^2 + V(\mathbf{r}) + g|\psi(\mathbf{r}, t)|^2 \right] \psi(\mathbf{r}, t), \quad (3.75)$$

where $\psi(\mathbf{r}, t) = \sqrt{N} \langle \mathbf{r} | \psi(t) \rangle$ is the single-particle wavefunction scaled by the square root number of atoms N .

Both these equations are partial differential equations involving both spatial and temporal derivatives. But in numerical quantum mechanics all state vectors are mapped to column vectors and all operators to matrices. Spatial wavefunctions are no exception to the former and differential operators such as ∇^2 are no exception to the latter—these objects can be thought of as infinite-dimensional vectors and matrices. So the above two equations are specific instances of the general Schrödinger equation (3.1), given specific Hamiltonians, and represented in a concrete—albeit infinite-dimensional—basis. But we can only perform a finite number of computations, so what do these vectors and operators look like once we reduce them to something finite? That depends on whether we choose to discretise space on a grid, or use a functional basis (and on which functional basis we choose). As we'll see, however, spatial discretisation can be a special case of a functional basis, namely the Fourier basis, plus an additional approximation or two. The resulting matrices are *banded*, justifying the previous two sections' attention to dealing with banded matrices efficiently.

3.4.1 Spatial discretisation on a uniform grid: the Fourier basis

Imagine a two dimensional spatial region within which we are solving the single-component Gross–Pitaevskii equation, evolving an initial condensate wavefunction in time. Having specified which degrees of freedom we want to simulate (two continuous degrees of freedom, one for each spatial dimension), the next step according to the method outlined in section 3.1 is to choose a basis in which to represent this state vector.

Let's say we discretise space in an equally-spaced $n_x \times n_y = 7 \times 7$ rectangular grid,¹⁸ with spacings Δx and Δy , and only represent the wavefunction at those 49 points in space. The state vector can then be represented by a list of 49 complex numbers, each taken to be the wavefunction's value at the spatial position corresponding to one gridpoint. This 49-vector is now a concrete representation of our state vector (Figure 3.5).

¹⁸For the sake of example—256 × 256 is a more realistic minimum.

We can also evaluate the potential (and nonlinear term in the case of the Gross–Pitaevskii equation) at each gridpoint and declare this a diagonal operator (Figure 3.6). Finally, we could use finite differences to compute the Laplacian - equivalent to replacing the Laplacian with a matrix

$$L = L_x \otimes \mathbb{I}_{n_y} + \mathbb{I}_{n_x} \otimes L_y \quad (3.76)$$

where L_x and L_y are (banded) matrices for finite-difference approximations to second derivatives in each direction. We might also use discrete Fourier transforms to evaluate the Laplacian, since

$$\nabla^2 \psi(\mathbf{r}) = \mathcal{F}^{-1} \left\{ -k^2 \mathcal{F}\{\psi(\mathbf{r})\}(\mathbf{k}) \right\}, \quad (3.77)$$

where $k = |\mathbf{k}|$.

This is all well and good, and it works. But at what point did we choose a basis just now—what are the basis vectors? This just looks like discretising space at a certain resolution, rather than the formal process of choosing a basis and projecting the state vector

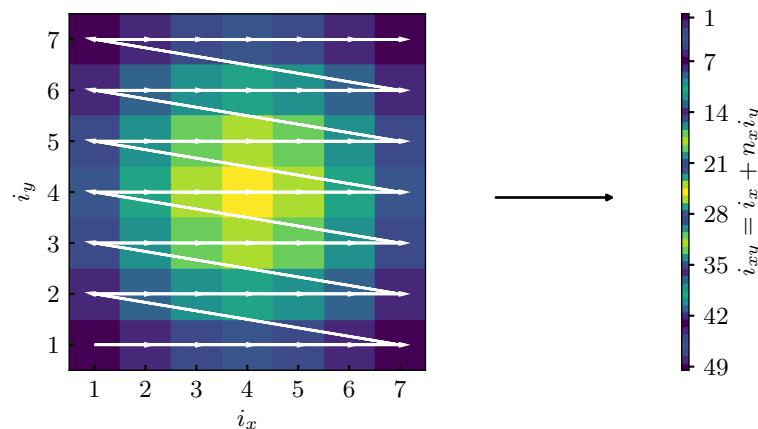


Figure 3.5: Discretising a function over two dimensional space on a grid yields a list of coefficients, one for each gridpoint. These can be arranged as a column vector, and in this way a two dimensional wavefunction approximated by a finite-dimensional state vector. Computationally we don't normally treat this state vector as a column vector—it is more convenient to leave it as a two-dimensional array. But conceptually it is a single vector living in the product space of the discretised x and y spaces.

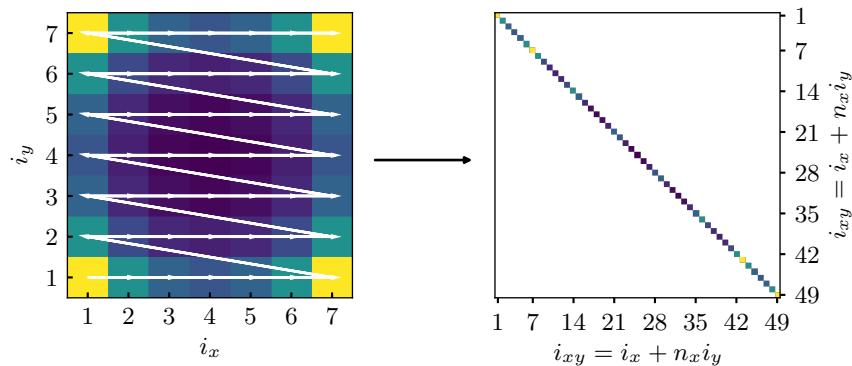


Figure 3.6: A discretised potential energy operator can be formed by evaluating the function for the potential at a set of gridpoints. The resulting matrix is diagonal and square with $n_x \times n_y$ rows and columns. Since multiplying this operator by a state vector entails multiplying each element of the state vector by one diagonal element of the potential operator, both the state vector and the diagonals of the potential operator can be stored in simulation code as 2D arrays and multiplied elementwise, hiding somewhat the fact that the operation is still a matrix-vector product. This way of discretising a potential operator is called the pseudospectral approximation, and is described later in this section

and operators onto each basis vector, as outlined in section 3.1. Assuming what we've done is equivalent to choosing a basis, that basis has a finite number (49) of basis vectors, which means it cannot be complete, since state vectors we're approximately representing with it require an infinite number of complex numbers to be described exactly.¹⁹ So what do the basis functions look like, and what state vectors have we implicitly excluded from simulation by choosing a basis that is incomplete?

Prior to discretisation, the spatial wavefunction $\psi(\mathbf{r}, t) = \langle \mathbf{r} | \psi(t) \rangle$ was already the

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¹⁹One for each position within the two dimensional space we're representing.

representation of the abstract state vector $|\psi\rangle$ in the “spatial basis”—a basis in which the basis vectors $\{|\mathbf{r}\rangle\}$ are Dirac deltas positioned at each point in space. The value of the wavefunction $\psi(\mathbf{r})$ for a specific \mathbf{r} is then simply a coefficient saying how much of the basis vector $|\mathbf{r}\rangle$ to include in the overall state vector. What we have *not* done is chosen a subset of these Dirac delta basis functions as our basis. This would be very strange—our representation of the wavefunction would allow it to be nonzero at the gridpoints, but not in between, like a comb. Spatially separated Dirac deltas do not spatially overlap at all; the matrix elements of the kinetic energy operator:

$$\langle \mathbf{r}_i | \hat{K} | \mathbf{r}_j \rangle = \int \delta(\mathbf{r} - \mathbf{r}_i) \left(-\frac{\hbar^2}{2m} \nabla^2 \right) \delta(\mathbf{r} - \mathbf{r}_j) d\mathbf{r} \quad (3.78)$$

would all be zero for $i \neq j$, disallowing any flow of amplitude from one point in space to another by virtue of it not being able to pass through the intervening points.

Neither have we implicitly chosen a set of two-dimensional boxcar functions centred on each gridpoint with width Δx and Δy in each direction respectively. These cover all space in between gridpoints, but are not twice differentiable everywhere, and hence the kinetic energy operator’s matrix elements cannot be evaluated²⁰. No, neither of these bases makes sense. To interpret our spatial grid as a basis, we need a set of functions $\phi_{ij}(\mathbf{r})$ (where i and j are the indices of the gridpoints in the x and y directions respectively) that are orthonormal, are nonzero only at one gridpoint and are zero at all others, and are twice differentiable everywhere in our spatial region. Infinite choices are available, differing in which subspace of the original Hilbert space they cover. A sensible heuristic for choosing one is that we want to be able to represent the state vectors whose wavefunctions do not change much between adjacent gridpoints, and we are happy for the necessary incompleteness of our basis to exclude wavefunctions with any sort of structure in between gridpoints.

²⁰Delta functions aren’t twice differentiable either, so this itself isn’t a fatal flaw—but even if one defines the boxcar functions as the limit of twice differentiable functions becoming increasingly square with some parameter, the limit for the kinetic energy matrix elements between adjacent boxcars goes to infinity.

The discrete Fourier transform to the rescue

It turns out that discretising space in this way can indeed be equivalent to choosing a sensible basis. This is made clearer by first discretising in Fourier space instead, and seeing how this can imply a discretisation in real space.

One possible basis for representing all possible state vectors is the Fourier basis $\{|\mathbf{k}_{ij}\rangle\}$. With it, any state vector (whose wavefunction is nonzero only within the 2D region) can be represented as the sum of basis vectors whose wavefunctions are 2D plane waves, also localised to the 2D region:

$$\langle \mathbf{r} | \mathbf{k}_{ij} \rangle = \begin{cases} \frac{1}{\sqrt{A}} e^{i\mathbf{k}_{ij} \cdot \mathbf{r}} & (\mathbf{r} \text{ within 2D region}) \\ 0 & (\mathbf{r} \text{ not within 2D region}), \end{cases} \quad (3.79)$$

where A is the area of the 2D region and the wavevector of each plane wave is

$$\mathbf{k}_{ij} = \left[\frac{2\pi i}{L_x}, \frac{2\pi j}{L_y} \right]^T, \quad (3.80)$$

where i and j are (possibly negative) integers. Any state vector whose wavefunction is localised to the 2D region can then be written as the infinite sum:

$$|\psi\rangle = \sum_{i=0}^{\infty} \sum_{j=0}^{\infty} \langle \mathbf{k}_{ij} | \psi \rangle | \mathbf{k}_{ij} \rangle \quad (3.81)$$

$$\Rightarrow \psi(\mathbf{r}) = \langle \mathbf{r} | \psi \rangle = \begin{cases} \sum_{i=0}^{\infty} \sum_{j=0}^{\infty} \langle \mathbf{k}_{ij} | \psi \rangle \frac{1}{\sqrt{A}} e^{i\mathbf{k}_{ij} \cdot \mathbf{r}} & (\mathbf{r} \text{ within 2D region}) \\ 0 & (\mathbf{r} \text{ not within 2D region}). \end{cases} \quad (3.82)$$

So $\{\langle \mathbf{k}_{ij} | \psi \rangle\}$ are simply the coefficients of the 2D Fourier series of $\psi(\mathbf{r})$.

What does this have to do with our discretised space? These basis functions $\{\langle \mathbf{r} | \mathbf{k}_{ij} \rangle\}$ don't have the required properties for a spatial discrete basis. For one, there are an infinite number of them, and we require 49 for our 7×7 example. Secondly, all of them are nonzero everywhere within the 2D region, whereas we require each basis function to be nonzero at exactly one of our 49 gridpoints.

We can solve the first problem by truncating the Fourier series. By only including basis vectors $|\mathbf{k}_{ij}\rangle$ for which:

$$\begin{cases} i \in [-\frac{n_x}{2}, \frac{n_x}{2} - 1] & (n_x \text{ even}) \\ i \in [-\frac{n_x-1}{2}, \frac{n_x-1}{2}] & (n_x \text{ odd}) \end{cases} \quad (3.83)$$

and

$$\begin{cases} j \in [-\frac{n_y}{2}, \frac{n_y}{2} - 1] & (n_y \text{ even}) \\ j \in [-\frac{n_y-1}{2}, \frac{n_y-1}{2}] & (n_y \text{ odd}) \end{cases} \quad (3.84)$$

we include only the n_x and n_y longest wavelengths in each respective spatial dimension. This is a sensible truncation with a physically meaningful interpretation. By making it, we are no longer able to represent state vectors with short wavelength components. Because the kinetic energy operator, when represented in the Fourier basis, is:

$$\langle \mathbf{k}_{ij} | \hat{K} | \mathbf{k}_{i'j'} \rangle = \frac{\hbar^2 k^2}{2m} \delta_{ii'} \delta_{jj'}, \quad (3.85)$$

²¹Because a *square* region in Fourier space is being carved out, by limiting each of k_x and k_y to finite ranges rather than the total wavenumber $k = \sqrt{k_x^2 + k_y^2}$, there is no single kinetic energy cutoff so to speak. Nonetheless there is a maximum wavenumber $k_{\max} = \min(\{|k_x|\} \cup \{|k_y|\})$ defining a kinetic energy cutoff $K_{\max} = \hbar^2 k_{\max}^2 / (2m)$ below which kinetic energies definitely are representable and above which they may not be.

where $k = |\mathbf{k}_{ij}|$, by excluding basis vectors with larger wavevectors, we are excluding state vectors with large kinetic energy. Thus the truncation is a kinetic energy cutoff, and is an accurate approximation whenever a simulation is such that the system is unlikely to obtain kinetic energies above the cutoff.²¹ It also matches our earlier intuition that our basis should represent wavefunctions that don't vary much between gridpoints—here we are discarding short wavelengths and hence limiting wavefunctions we can represent to ones that vary slowly in space compared to the cutoff wavelength.

Now we have a set of basis vectors—a discrete Fourier basis—but their spatial wavefunctions still don't have the property of being nonzero only at a single gridpoint each. On the contrary, each plane wave has a constant amplitude everywhere in space. But consider the following superposition of Fourier basis vectors:

$$|\mathbf{r}_{ij}\rangle = \sum_{i'} \sum_{j'} e^{-i\mathbf{k}_{i'j'} \cdot \mathbf{r}_{ij}} |\mathbf{k}_{i'j'}\rangle \quad (3.86)$$

with $\mathbf{r}_{ij} = (i\Delta x, j\Delta y)^T$. The set of vectors $\{|\mathbf{r}_{ij}\rangle\}$ are also an orthonormal basis, related to the discrete Fourier basis by a unitary transformation with matrix elements:

$$U_{\text{DFT2}, i'j'ij} = \langle \mathbf{k}_{i'j'} | \mathbf{r}_{ij} \rangle = e^{-i\mathbf{k}_{i'j'} \cdot \mathbf{r}_{ij}}. \quad (3.87)$$

This unitary transformation is in fact a two-dimensional discrete Fourier transform (hence the subscript), and the basis vectors $\{|\mathbf{r}_{ij}\rangle\}$ have spatially localised wavefunctions that are nonzero only at one of the spatial gridpoints. Vectors and matrices can be transformed from their discrete Fourier space representation to their discrete real-space representation and back using the unitary U_{DFT2} :

$$\psi_{\text{real}} = U_{\text{DFT2}}^\dagger \psi_{\text{Fourier}} \quad (3.88)$$

$$\psi_{\text{Fourier}} = U_{\text{DFT2}} \psi_{\text{real}} \quad (3.89)$$

$$A_{\text{real}} = U_{\text{DFT2}}^\dagger A_{\text{Fourier}} U_{\text{DFT2}} \quad (3.90)$$

$$A_{\text{Fourier}} = U_{\text{DFT2}} A_{\text{Real}} U_{\text{DFT2}}^\dagger. \quad (3.91)$$

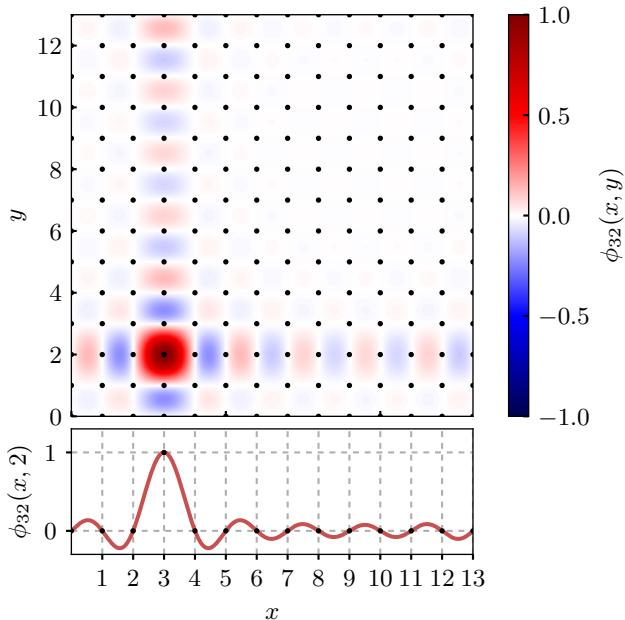


Figure 3.7: An example of the spatial representation of one of the basis vectors obtained by transforming the discrete Fourier basis using the discrete Fourier transform. Since the discrete Fourier transform is a unitary transformation, the set of these basis functions forms an equivalent orthonormal basis for representing state vectors and operators.

where ψ_{real} is the vector of coefficients $\psi_{\text{real},ij} = \langle \mathbf{r}_{ij} | \psi \rangle$ for representing a state vector in the discrete real space basis, ψ_{Fourier} is the vector of coefficients $\psi_{\text{Fourier},ij} = \langle \mathbf{k}_{ij} | \psi \rangle$ for the state vector in the discrete Fourier basis, and A_{real} and A_{Fourier} are the representations of some operator \hat{A} in the discrete real and Fourier bases respectively, with matrix elements $A_{\text{real},iji'j'} = \langle \mathbf{r}_{ij} | \hat{A} | \mathbf{r}_{i'j'} \rangle$ and $A_{\text{Fourier},iji'j'} = \langle \mathbf{k}_{ij} | \hat{A} | \mathbf{k}_{i'j'} \rangle$.

The spatial representation of the basis vectors $\{|\mathbf{r}_{ij}\rangle\}$ can be computed using (3.79) as:

$$\phi_{ij}(\mathbf{r}) = \langle \mathbf{r} | \mathbf{r}_{ij} \rangle = \sum_{i'j'} e^{-i\mathbf{k}_{i'j'} \cdot \mathbf{r}_{ij}} \langle \mathbf{r} | \mathbf{k}_{i'j'} \rangle \quad (3.92)$$

$$\Rightarrow \phi_{ij}(\mathbf{r}) = \begin{cases} \sum_{i'j'} \frac{1}{\sqrt{L_x L_y}} e^{i\mathbf{k}_{i'j'} \cdot (\mathbf{r} - \mathbf{r}_{ij})} & (\mathbf{r} \text{ within 2D region}) \\ 0 & (\mathbf{r} \text{ not within 2D region}), \end{cases} \quad (3.93)$$

where L_x and L_y are the spatial extents of the x and y dimensions. An example of one of these basis vectors in the spatial representation is plotted in Figure 3.7.

These functions are sometimes called *periodic sinc functions*, *band-limited delta functions*, or the *Dirichlet kernel* [68, p. 619; 69]. Each of them is zero at all of the gridpoints except one, and they form an orthonormal basis set. They satisfy all of our requirements to be a basis corresponding to our gridded discretisation of space.

One thing to note is that these functions are periodic. By using the Fourier basis in the way we have to restrict our basis to cover only a finite region of both Fourier space and real space, we have necessarily imposed periodicity on the problem. This periodicity shows itself when we compute matrix elements of operators in this basis - if we compute the kinetic energy operator's matrix elements for example, it will couple basis states across the boundary of the region, resulting in spatial periodicity—a wavepacket moving rightward through the right boundary will emerge moving rightward from the left boundary.

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²²With the possible exception of the region of Fourier space being simulated corresponding to the first Brillouin zone of a lattice potential, in which case these velocity reversals would correspond to Bloch oscillations.

²³The wavefunctions of eigenstates of the 3D harmonic oscillator are products of radial polynomials and spherical harmonics, the combination of which is a good spectral basis for many problems.

Less obviously, the basis is also periodic in Fourier space, and so a wavepacket moving out of the region of Fourier space simulated will also wrap around to the opposite side of Fourier space. In real space, this may appear as a wavepacket undergoing acceleration only to suddenly reverse its velocity as if reflected off a barrier. This effect is unphysical²² and should be taken as a sign that the spatial grid is not fine enough for the dynamics being simulated.

The discrete Fourier basis we've described is one example of a *spectral basis*, and a numerical method which represented the state vector solely in this basis would be called a *spectral method*. Other choices of basis functions, such as polynomials (see section 3.4.3) or spherical harmonics lead to other spectral methods.²³ The discrete spatial basis discussed above, despite being related to the Fourier basis by a unitary transformation, is often called a *pseudospectral basis*, however, and a method representing the state vector in this basis a *pseudospectral method*. Although it is "just another basis", the fact that the basis functions are zero at all spatial points bar one each leads to the possibility of a further approximation when representing some operators, which makes bases with this property especially useful. I describe this in the following section.

Vector and matrix elements in the Fourier and pseudospectral bases

We now have a finite basis $\{|\mathbf{r}_{ij}\rangle\}$ that matches our intuitions somewhat for representing a wavefunction at a set of gridpoints. A state vector can be approximated as a linear sum of these basis vectors, with the coefficient for each one being equal to the projection of state vector's wavefunction onto the basis vector's wavefunction:

$$|\psi\rangle \approx \sum_{ij} \psi_{ij} |\mathbf{r}_{ij}\rangle, \quad (3.94)$$

where

$$\psi_{ij} = \int \phi_{ij}^*(\mathbf{r}) \psi(\mathbf{r}) d\mathbf{r}. \quad (3.95)$$

In practice however this integral is rarely done. Instead, ψ_{ij} is simply taken to be the value of the exact wavefunction $\psi(\mathbf{r}) = \langle \mathbf{r} | \psi \rangle$ at the point $\mathbf{r} = \mathbf{r}_{ij}$:

$$\psi_{ij} \approx \psi(\mathbf{r}_{ij}) \quad (3.96)$$

To see why this is a good approximation, imagine that the approximation (3.94) were exact, that is, $\psi(\mathbf{r})$ were exactly equal to a linear sum of the functions $\{\phi_{ij}(\mathbf{r})\}$. Since all the basis functions are zero at the point \mathbf{r}_{ij} except for ϕ_{ij} , the value of $\psi(\mathbf{r}_{ij})$ must come solely from the $\psi(\mathbf{r})$'s projection onto the basis function $\phi_{ij}(\mathbf{r})$. Since approximating (3.94) underlies the results of any simulation that discretises a state vector in this way, treating it as exact for the initial projection onto the discrete basis is making an assumption no worse than that already being relied upon.

With a basis and initial discrete state vector in hand, we can now proceed to calculate matrix elements of the Hamiltonian, after which we can proceed to solve the differential equation (3.3) to determine how the coefficients $\{\psi_{ij}\}$ evolve in time.

The specific properties of our Fourier/pseudospectral basis make it quite useful for a range of common Hamiltonians. For example, let's take the single particle Schrödinger Hamiltonian:

$$\hat{H}_{\text{Schrö}} = \frac{\hbar^2 \hat{k}^2}{2m} + V(\hat{\mathbf{r}}), \quad (3.97)$$

where $\hat{k} = |\hat{\mathbf{k}}|$.

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The two terms, kinetic and potential, are each diagonal in different bases. The kinetic term is diagonal in the Fourier basis:

$$\langle \mathbf{k}' | \frac{\hbar^2 \hat{k}^2}{2m} | \mathbf{k} \rangle = \frac{\hbar^2 k^2}{2m} \delta(\mathbf{k} - \mathbf{k}'), \quad (3.98)$$

where $k = |\mathbf{k}|$, and the potential term is diagonal in the spatial basis:

$$\langle \mathbf{r}' | V(\hat{\mathbf{r}}) | \mathbf{r} \rangle = V(\mathbf{r}) \delta(\mathbf{r} - \mathbf{r}'). \quad (3.99)$$

The kinetic term is also diagonal our discrete Fourier basis:

$$K_{\text{Fourier}, i'j'ij} = \langle \mathbf{k}_{i'j'} | \frac{\hbar^2 \hat{k}^2}{2m} | \mathbf{k}_{ij} \rangle = \frac{\hbar^2 k_{ij}^2}{2m} \delta_{i'i} \delta_{j'j}, \quad (3.100)$$

however—perhaps surprisingly—the potential term is not diagonal in the pseudospectral basis, only approximately so:

$$V_{\text{real}, i'j'ij} = \langle \mathbf{r}_{i'j'} | V(\hat{\mathbf{r}}) | \mathbf{r}_{ij} \rangle \approx V(\mathbf{r}_{ij}) \delta_{i'i} \delta_{j'j}. \quad (3.101)$$

Nonetheless, equation (3.101) is in practice treated as exact for the purpose of computing matrix elements in the discrete basis of operators that are diagonal in the full spatial basis. As above with projecting a state vector using simply the values of a wavefunction at the gridpoints (rather than doing integrals), treating this approximation as exact is equivalent to making the assumption that the potential $V(\mathbf{r})$ is already accurately representable in the discrete basis as a diagonal operator:

$$V(\mathbf{r}) \approx \sum_{ij} V(\mathbf{r}_{ij}) \phi_{ij}^*(\mathbf{r}) \phi_{ij}(\mathbf{r}), \quad (3.102)$$

and so similarly is going to be a good approximation whenever the discrete basis is well able to represent the potential. So long as your discrete basis can accurately represent the state vectors and spatially-diagonal operators you will be using, it makes little difference whether you project those vectors and operators onto the basis using integrals or using simply their values at the gridpoints.

This alternate method of projection has a name, it is called *collocation* [55, p. 227]. Using collocation instead of vector projection amounts to treating our basis vectors as a scheme for interpolating state vectors and operators in between gridpoints, given their values at the points, rather than as basis vectors to project upon. Another way to interpret collocation is to say that we are evaluating the vector projections using integrals after all, however, we're numerically computing those integrals using a quadrature scheme, only evaluating the integrand at the gridpoints and performing a discrete sum [55, p. 283]. Collocation is what puts the *pseudo* in pseudospectral—if we evaluated all these operators using integrals instead, we would be treating the discrete spatial basis exactly the same as any spectral basis.

Compared to a purely spectral method, pseudospectral methods are of comparable accuracy [70]. This makes intuitive sense—discrete sums instead of integrals is how we are going to do all inner products once we are in the discrete basis—so it can't be much worse to use the same method to approximate operators and initial state vectors. The sole downside of pseudospectral methods, according to [70], is that the error can lead to instability in the presence of certain nonlinearities. Specifically, if long wavelength waves interact in a way that would produce wavelengths shorter than two grid spacings (the Nyquist wavelength), pseudospectral methods will produce longer wavelength waves instead, whereas in a purely spectral method the interaction would not occur, being due

to couplings to a Fourier mode outside the discrete Fourier basis. This *aliasing* can cause instability, but can be circumvented with smoothing techniques [71]. This is no great downside: if you want to simulate short length scales, you need to choose a grid spacing small enough to represent them, whereas if short wavelengths are produced despite your willingness to ignore them, you must smooth them away before they are aliased into long wavelengths that you do care about.

Finally, armed with a kinetic energy operator in Fourier space and a pseudospectral approximation to the potential operator in real space, we can write all the matrix elements of $\hat{H}_{\text{Schrö}}$ in a single basis, and thus our discretised, pseudospectral two-dimensional Schrödinger wave equation:

$$i\hbar \frac{d}{dt} \psi_{i'j'}(t) = \sum_{ij} H_{i'j',ij}(t) \psi_{ij}(t) \quad (3.103)$$

$$\Rightarrow i\hbar \frac{d}{dt} \psi_{i'j'}(t) = \sum_{ij} \left[U_{\text{DFT2}}^\dagger K_{\text{Fourier}} U_{\text{DFT2}} + V_{\text{real}}(t) \right]_{i'j',ij} \psi_{ij}(t), \quad (3.104)$$

where we have used the discrete Fourier transform to transform the kinetic energy operator into the discrete real space basis, and allowed the potential operator to be possibly time dependent.

The right hand side of this expression can now simply be evaluated, yielding the time derivative of each component $\psi_{ij}(t)$ of the discrete state vector $\psi(t)$:

$$\frac{d}{dt} \psi(t) = -\frac{i}{\hbar} \left[U_{\text{DFT2}}^\dagger K_{\text{Fourier}} U_{\text{DFT2}} \psi(t) + V_{\text{real}}(t) \psi(t) \right] \quad (3.105)$$

$$\Rightarrow \frac{d}{dt} \psi(t) = -\frac{i}{\hbar} \left[\text{FFT}_2^{-1} \left\{ \frac{\hbar^2 \tilde{\mathbf{k}}^{\odot 2}}{2m} \odot \text{FFT}_2 \{ \psi(t) \} \right\} + V(\tilde{\mathbf{r}}, t) \odot \psi(t) \right], \quad (3.106)$$

where FFT_2 is the two dimensional fast Fourier transform, an efficient implementation of the discrete Fourier transform (taking time $\mathcal{O}(n \log n)$ in the size of each dimension), $\tilde{\mathbf{r}}$ is a vector (of vectors) containing each discrete position vector (such that $V(\tilde{\mathbf{r}}, t)$ is a vector (of scalars) containing the potential evaluated at each discrete position), $\tilde{\mathbf{k}}$ is a vector (of vectors) containing each discrete k -vector, such that $\tilde{\mathbf{k}}^{\odot 2}$ is a vector (of scalars) containing the squared magnitude of each discrete k -vector, and \odot represents elementwise multiplication (or exponentiation) of vectors. Other than ψ , these vectors are more akin to arrays used in programming languages than to members of a vector space, hence the somewhat clunky notation in (3.106). Comparison with the continuous version of equation (3.106) (i.e. the Schrödinger wave equation (3.74)), since elementwise multiplication of functions is a more common operation in mathematics, might be clarifying:

$$\frac{\partial}{\partial t} \psi(\mathbf{r}, t) = -\frac{i}{\hbar} \left[\mathcal{F}^{-1} \left\{ \frac{\hbar^2 k^2}{2m} \mathcal{F} \{ \psi(\mathbf{r}, t) \} (\mathbf{k}) \right\} (\mathbf{r}) + V(\mathbf{r}, t) \psi(\mathbf{r}, t) \right], \quad (3.107)$$

where \mathcal{F} is the continuous Fourier transform, and k as always is $|\mathbf{k}|$. So we see that the discretised Schrödinger equation for a single particle in a potential really is the same as evaluating the continuous equation at a set of gridpoints, evaluating spatial derivatives in Fourier space, and replacing the continuous Fourier transform with its discrete equivalent.

Here is an example of how one might compute the RHS of (3.106) in Python code:

```

1 import numpy as np
2 from numpy.fft import fft2, ifft2, fftfreq
3
4 pi = np.pi
5 u = 1.660539e-27 # unified atomic mass unit

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```

6 m = 86.909180*u # 87Rb atomic mass
7 omega = 15 # Harmonic trap frequency
8 hbar = 1.054571726e-34 # Reduced Planck's constant
9
10 # Space:
11 nx = ny = 256
12 x_max = y_max = 100e-6
13
14 # Arrays of components of position vectors. The reshaping is to ensure that
15 # when used in arithmetic with each other, these arrays will be treated as if
16 # they are two dimensional with repeated values along the dimensions of size
17 # 1, up to the size of the other array (this is called broadcasting in numpy).
18 x = np.linspace(-x_max, x_max, nx, endpoint=False).reshape(1, nx)
19 y = np.linspace(-y_max, y_max, ny, endpoint=False).reshape(ny, 1)
20
21 # Grid spacing:
22 dx = x[0, 1] - x[0, 0]
23
24 # Arrays of components of k vectors.
25 kx = 2*pi*fftfreq(nx, d=dx).reshape(1, nx)
26 ky = 2*pi*fftfreq(ny, d=dx).reshape(ny, 1)
27
28 # The kinetic energy operator in Fourier space (shape ny, nx).
29 K_fourier = hbar**2 * (kx**2 + ky**2)/(2*m)
30
31 # The potential operator in real space (shape ny, nx)
32 V_real = 0.5 * m * omega**2 * (x**2 + y**2)
33
34 def dpsi_dt(t, psi):
35     """Return a 2D array for the time derivative of the 2D array psi
36     representing a discretised wavefunction obeying the Schrodinger wave
37     equation"""
38     K_real_psi = ifft2(K_fourier * fft2(psi))
39     return -1j/hbar * (K_real_psi + V_real * psi)

```

Where the example is for a time-independent potential. If the potential were time dependent, `V_real` within the function `dpsi_dt(t, psi)` would need to be replaced with a call to a function that returned an array for `V_real` at time `t`.

Starting with some initial discrete wavefunction, this could then be solved with a forward differencing scheme like fourth order Runge–Kutta (section 3.3):

```

1 def rk4(t, t_final, dt, psi, dpsi_dt):
2     """Evolve the initial array psi_initial forward in time from time t to
3     t_final according to the differential equation dpsi_dt using fourth order
4     Runge-Kutta with timestep dt"""
5     while t < t_final:
6         k1 = dpsi_dt(t, psi)
7         k2 = dpsi_dt(t + 0.5 * dt, psi + 0.5 * k1 * dt)
8         k3 = dpsi_dt(t + 0.5 * dt, psi + 0.5 * k2 * dt)
9         k4 = dpsi_dt(t + dt, psi + k3 * dt)
10
11     psi[:] += dt/6.0 * (k1 + 2 * k2 + 2 * k3 + k4)
12
13     t += dt
14
15     return psi

```

The discretised differential equation (3.105) can also be solved using a split-step method (section 3.2.4), since the Hamiltonian matches the requirements of being written as a sum of terms for which individually an eigenbasis is known (the discrete real space basis for the potential term, and the Fourier basis for the kinetic term). For example, here is how one might implement second or fourth-order split-step (only a single timestep shown):

```

1 def split_step2(t, psi, dt):
2     """Evolve psi in time from t to t + dt using a single step of the second
3     order Fourier split-step method with timestep dt"""
4
5
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```

```

5      # First evolve using the potential term for half a timestep:
6      psi *= np.exp(-1j/hbar * V_real * 0.5 * dt)
7
8      # Then evolve using the kinetic term for a whole timestep, transforming to
9      # and from Fourier space where the kinetic term is diagonal:
10     psi = ifft2(np.exp(-1j/hbar * K_fourier * dt) * fft2(psi))
11
12     # Then evolve with the potential term again for half a timestep:
13     psi *= np.exp(-1j/hbar * V_real * 0.5 * dt)
14
15     return psi
16
17 def split_step4(t, psi, dt):
18     """Evolve psi in time from t to t + dt using a single step of the fourth
19     order Fourier split-step method with timestep dt"""
20     p = 1/(4 - 4**((1/3.0)))
21
22     # Five applications of second-order split-step using timesteps
23     # of size p*dt, p*dt, (1 - 4*p)*dt, p*dt, p*dt
24     for subdt in [p*dt, p*dt, (1 - 4*p)*dt, p*dt, p*dt]:
25         psi = split_step2(t, psi, subdt)
26         t += subdt
27
28     return psi

```

In all the above code examples, a nonlinear term as in the case of the Gross-Pitaevskii equation can be included in the potential simply by adding a term $\mathbf{g} * \text{np.abs}(\psi)**2$ to the potential V_{real} wherever it appears. As discussed in section 3.2.4, the nonlinearity poses no problem for the split-step methods so long as the potential term of the Hamiltonian is evaluated as the outermost sandwich of exponentials in the second-order split step method (which comprises the sub-steps of fourth-order split-step).

3.4.2 Finite differences

Fourier split-step, or using discrete Fourier transforms to evaluate the spatial derivatives at each gridpoint in order to time-evolve using Runge–Kutta are effective and versatile numerical methods.

The use of discrete Fourier transforms in the previous section can be seen as replacing the Laplacian operator in the Schrödinger wave equation (3.74) with the equivalent operation in Fourier space:

$$\nabla^2 \psi(\tilde{\mathbf{r}}) \approx \text{FFT}_2^{-1} \left\{ -\tilde{\mathbf{k}}^{\odot 2} \odot \text{FFT}_2 \left\{ \psi(\tilde{\mathbf{r}}) \right\} \right\}, \quad (3.108)$$

where as before $\tilde{\mathbf{r}}$ is a vector (of vectors) containing the discrete positions, $\tilde{\mathbf{k}}$ is a vector (of vectors) containing discrete k -vectors such that $\tilde{\mathbf{k}}^{\odot 2}$ is a vector (of scalars) containing the squared magnitudes of each k -vector, and \odot represents elementwise multiplication or exponentiation of vectors. More generally for any derivative,

$$\frac{\partial}{\partial x} \psi(\tilde{\mathbf{r}}) \approx \text{FFT}_2^{-1} \left\{ i\tilde{\mathbf{k}}_x \odot \text{FFT}_2 \left\{ \psi(\tilde{\mathbf{r}}) \right\} \right\}, \quad (3.109)$$

where $\tilde{\mathbf{k}}_x$ is a vector (of scalars) containing the discrete angular wavenumbers for the x spatial dimension.

Equations (3.108) and (3.109) are exact for any wavefunction $\psi(\mathbf{r})$ which is periodic and band-limited to the discrete Fourier space (and thus exactly representable as a vector ψ of its values at each gridpoint), which is why the Fourier method of computing derivatives this way is sometimes said to be accurate to “infinite order” [72] in the grid spacings Δx and Δy , in contrast to fixed-order approximations to derivatives which are second, fourth, sixth order etc. In practice the Fourier method for derivatives is often used for wavefunctions which are *not* intended to be periodic (the periodicity imposed by using the method is unphysical), and so for these it has merely very high order accuracy, not actually infinite.

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	$k = -3$	$k = -2$	$k = -1$	$k = 0$	$k = 1$	$k = 2$	$k = 3$
$\Delta x \left(\delta_{\Delta x}^{(2)} \right)_{i,i+k}$			$-\frac{1}{2}$	0	$\frac{1}{2}$		
$\Delta x \left(\delta_{\Delta x}^{(4)} \right)_{i,i+k}$		$\frac{1}{12}$	$-\frac{2}{3}$	0	$\frac{2}{3}$	$-\frac{1}{12}$	
$\Delta x \left(\delta_{\Delta x}^{(6)} \right)_{i,i+k}$	$-\frac{1}{60}$	$\frac{3}{20}$	$-\frac{3}{4}$	0	$\frac{3}{4}$	$-\frac{3}{20}$	$\frac{1}{60}$
$\Delta x^2 \left(\delta_{\Delta x}^{(2)} \right)_{i,i+k}$			1	-2	1		
$\Delta x^2 \left(\delta_{\Delta x}^{(4)} \right)_{i,i+k}$		$-\frac{1}{12}$	$\frac{4}{3}$	$-\frac{5}{2}$	$\frac{4}{3}$	$-\frac{1}{12}$	
$\Delta x^2 \left(\delta_{\Delta x}^{(6)} \right)_{i,i+k}$	$\frac{1}{90}$	$-\frac{3}{20}$	$\frac{3}{2}$	$-\frac{49}{18}$	$\frac{3}{2}$	$-\frac{3}{20}$	$\frac{1}{90}$

Table 3.1: Matrix elements [73] for some finite-differencing schemes for first ($\delta_{\Delta x}^{(n)}$) and second ($\delta_{\Delta x}^{2(n)}$) derivatives using central finite differences of various orders n for uniform grid spacing Δx . All finite difference matrices are banded; each column here shows the matrix elements of k^{th} diagonal, which are all identical. Elements outside of each matrix's band are left blank. Each matrix element is shown multiplied by factors of Δx for clarity.

In any case, such high accuracy is not often necessary—if one is using only an $\mathcal{O}(\Delta t^4)$ accurate timestepping scheme, then the timestepping may be the limiting factor in overall accuracy and it might be wise to decrease the accuracy of computing spatial derivatives if there is otherwise a benefit to doing so.

To that end, the Fourier method of derivatives may be replaced with *finite differences* instead. Although finite differences are usually derived as approximations to derivatives directly from the definition of the derivative without reference to discrete Fourier transforms, they can be considered fixed-order approximations to the Fourier method [72]. Thus operators whose form in Fourier space corresponds to a derivative of some order can be approximated with finite differences:

$$U_{\text{DFT2}}^\dagger k_x U_{\text{DFT2}} \psi(\tilde{r}) = -i \delta_{\Delta x}^{(n)} \otimes \mathbb{I}_{n_y} \psi(\tilde{r}) + \mathcal{O}(\Delta x^n) \quad (3.110)$$

where k_x is the diagonal matrix of the discrete angular wavenumbers for the x spatial dimension and $\delta_{\Delta x}^{(n)}$ is the matrix representing n^{th} order finite difference approximation to the first derivative using grid spacing Δx . The matrix elements of this and some other central finite differences are shown in Table??.

The fact that the finite difference matrices are banded allows them to be computed by applying a “stencil” to a discrete state vector, computing an approximation to some linear sum of derivative operators at each point of discrete space by consideration of only that point and a small number of surrounding points. For example, a $\mathcal{O}(\Delta x^2)$ approximation (assuming $\Delta x = \Delta y$) to the kinetic energy operator in two dimensions may be evaluated at each point as:

$$K_{\text{real}} \psi(\tilde{r}) = U_{\text{DFT2}}^\dagger \frac{\hbar^2(k_x^2 + k_y^2)}{2m} U_{\text{DFT2}} \psi(\tilde{r}) \quad (3.111)$$

$$= -\frac{\hbar^2}{2m} \left[\delta_{\Delta x}^{2(n)} \otimes \mathbb{I}_{n_y} + \mathbb{I}_{n_x} \otimes \delta_{\Delta y}^{2(n)} \right] \psi(\tilde{r}) + \mathcal{O}(\Delta x^2) \quad (3.112)$$

$$\Rightarrow (K_{\text{real}} \psi(\tilde{r}))_{ij} = -\frac{\hbar^2}{2m} \left[-4\psi(x_i, y_j) + \psi(x_{i-1}, y_j) + \psi(x_{i+1}, y_j) + \psi(x_i, y_{j-1}) + \psi(x_i, y_{j+1}) \right] + \mathcal{O}(\Delta x^2). \quad (3.113)$$

Thus the kinetic energy operator, when approximated using finite differences, is an example of an operator that can be written as a sum of banded operators acting on

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different subspaces of the total Hilbert space—the identity matrices in (3.112) each leave a part of the Hilbert space untouched. As mentioned in section 3.2.4, this in principle considerably reduces the computational cost of applying the approximate kinetic energy operator to a discretised state vector. Fourier transforms, when computed with the fast Fourier transform algorithm, are already less computationally expensive than a general matrix-vector multiplication, that is, the fast Fourier transform allows one to multiply the matrix U_{DFT_2} in (3.111) by a vector considerably faster than $\mathcal{O}(n_x^2 n_y^2)$, which would be the computational time-complexity for a general $n_x n_y \times n_x n_y$ matrix. Firstly, a two-dimensional discrete Fourier transform can also be written as the sum of two one-dimensional transformations operating on different subspaces:

$$U_{\text{DFT}_2} = U_{\text{DFT},x} \otimes \mathbb{I}_{n_y} + \mathbb{I}_{n_x} \otimes U_{\text{DFT},y}, \quad (3.114)$$

where $U_{\text{DFT},x}$ and $U_{\text{DFT},y}$ are the unitaries for one-dimensional discrete Fourier transforms in the x and y dimensions respectively. So even if $U_{\text{DFT},x}$ and $U_{\text{DFT},y}$ were arbitrary matrices, this already would reduce the cost of multiplying U_{DFT_2} by a vector to $\mathcal{O}(n_y n_x^2 + n_x n_y^2)$. But they are not arbitrary matrices—each of these one-dimensional Fourier transforms has computational cost $\mathcal{O}(n \log n)$ using the FFT algorithm [74, p. 600], where n is the number of points in the relevant dimension, resulting in an overall cost of $\mathcal{O}(n_y n_x \log n_x + n_x n_y \log n_y)$ for applying the two-dimensional Fourier transform U_{DFT_2} to a vector. Finite differences improves on this further. As discussed in section 3.2.4, since our finite-differences approximation to the kinetic energy operator can written as the sum of banded matrices operating on different subspaces, the computational cost of applying it to a vector is $\mathcal{O}(b n_x n_y)$, where b is the bandwidth of the banded matrices, which depends on which order accuracy is used (for example, for second order finite-differences $b = 1$). This is faster than the Fourier method of computing the kinetic energy operator by a factor of $\mathcal{O}(\log n_x + \log n_y)$. Whilst this seems considerable, the difference is hard to observe in practice. On ordinary computers the number of points needs to be increased so much in order to measure any difference in speed between the algorithms that the data no longer fits in CPU cache and copying the data to and from main memory becomes the bottleneck. Although copying data from memory is a linear-time process, the coefficient of that linear time is large enough to make the asymptotic speed of finite differences vs. Fourier transforms not relevant for ordinary computers at the present time.

No, the practical advantages of finite differences compared to Fourier transforms do not come down to single-core speed. Rather, they are:

1. Being banded matrices, the finite-difference approximations to the kinetic energy operator can be multiplied by a vector, or exponentiated and applied to a vector, in parallel on a cluster computer or GPU using the techniques discussed in section 3.2.4, whereas the fast Fourier transform is less efficiently parallelisable [75]. The speed of finite differences can have very nice scaling with the number of CPU cores or cluster nodes used, since only b points need to be exchanged between cores at each step when applying or exponentiating the kinetic energy operator. For large problems, ‘superscaling’ can even be observed, whereby the speedup factor obtained by moving to multiple CPUs or compute nodes on a cluster is larger than the number of CPUs/nodes used. This is counterintuitive, but comes from more effectively using CPU cache—by spreading the data over multiple cores, one minimises the proportion of the state vector that needs to reside in main memory instead of in (much faster) CPU cache at any one time.
2. One can intervene at the boundaries to impose boundary conditions other than periodicity. Strictly speaking, as an approximation to the Fourier method of computing derivatives, the indices for the matrix elements as given in Table?? should

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be read as wrapping around to the other side of the spatial region whenever they would go out of bounds—that is, $(\delta_{\Delta x}^{2(n)})_{i,i+k}$ should be read as $(\delta_{\Delta x}^{2(n)})_{i,(i+k) \bmod n_x}$. However, as mentioned, periodicity is often an undesired consequence of the Fourier method of derivatives. Alternatively one can simply omit these matrix elements that would couple spatial points across the boundary, which has the result of imposing zero boundary conditions instead of periodic. Judicious deletion of matrix elements at other points in space can also be used to impose zero boundary conditions elsewhere, equivalent to an infinite potential barrier which would otherwise be numerically troublesome if done with the potential energy term of the Hamiltonian. Other interventions in the application of the kinetic energy operator can be used to impose other boundary conditions such as constant-value, constant-gradient, etc., whereas the Fourier method is less flexible in this regard.

3. Finite differences are compatible with non-uniform grids, whereas the Fourier method is limited to uniform grids. Non-uniform grids imply different matrix elements [75] for the finite difference operators, but are otherwise treated exactly the same. This allows more dense placement of gridpoints in regions where wavefunctions may have finer spatial structure, without having to waste computational power on regions of space where the wavefunction is known to have only coarser structure. An example is an electron in a Coulomb potential, an accurate simulation of which would need to capture fine details at small radii but less detail at larger radii. A transformation into spherical coordinates with a non-uniform grid for the radial coordinate could be well treated with finite differences.
4. Finite differences are compatible with the use of split-step methods with some operators that are diagonal in neither Fourier nor real space. For example, the real-space representation of the operator for the z component of angular momentum is $L_z = -i\hbar \left(x \frac{\partial}{\partial y} - y \frac{\partial}{\partial x} \right)$. With diagonal matrices X and Y being used for the \hat{x} and \hat{y} operators in accordance with the pseudospectral method, and finite differences being used to approximate the derivatives, the result is a banded matrix representation of L_z , compatible with the techniques from section 3.2.4 for reducing the problem to that of many small matrices instead of one large one. There is a little more complexity; L_z cannot be represented as a sum of operators acting on the x and y subspaces separately—instead, each term in L_z is a *product* of operators that act on different subspaces. Consider the first term of a discretised version of L_z using fourth-order finite differences. It can be diagonalised in the following way:

$$-i\hbar X \otimes \delta_{\Delta y}^{(4)} = -i\hbar \left(\mathbb{I}_{n_x} X \mathbb{I}_{n_x} \right) \otimes \left(Q_{\delta_y}^\dagger D_{\delta_y} Q_{\delta_y} \right), \quad (3.115)$$

where Q_{δ_y} and D_{δ_y} are the unitary and diagonal matrices that diagonalise $\delta_{\Delta y}^{(4)}$ (X is already diagonal, and so is ‘diagonalised’ by the identity matrix for the x subspace). \mathbb{I}_{n_x} and X act on the x subspace, whereas Q_{δ_y} and D_{δ_y} act on the y subspace, and since matrices operating on different subspaces commute, this can be rearranged to:

$$-i\hbar X \otimes \delta_{\Delta y}^{(4)} = \left(\mathbb{I}_{n_x} \otimes Q_{\delta_y}^\dagger \right) \left(-i\hbar X \otimes D_{\delta_y} \right) \left(\mathbb{I}_{n_x} \otimes Q_{\delta_y} \right), \quad (3.116)$$

yielding a diagonalisation of the original matrix that can be used to apply an exponentiation of the original matrix to a vector with matrix-vector multiplications only in the x any y subspaces and not the total Hilbert space. Here, the matrix-vector multiplication in the x subspace is the identity, but more generally the above

idea can be used to exponentiate any operator that can be written as the product of operators that act on different subspaces:

$$e^{A \otimes B} = (Q_A^\dagger \otimes Q_B^\dagger) e^{(D_A \otimes D_B)} (Q_A \otimes Q_B). \quad (3.117)$$

Since X is already diagonal, $\delta_{\Delta_y}^{(4)}$ can be written as the sum of two block-diagonal matrices as described in section 3.2.4, allowing (3.116) to be evaluated as the sum of two terms $-i\hbar X \otimes \delta_{\Delta_y}^{(4)} = -i\hbar(X \otimes \delta_{\text{even}} + X \otimes \delta_{\text{odd}})$, one for each of the block-diagonal matrices δ_{even} and δ_{odd} . The diagonalisation of each term then has matrix-vector products taking place in a space of the size of each block, that is, in the expression

$$-i\hbar X \otimes \delta_{\text{even}} = (\mathbb{I}_{n_x} \otimes Q_{\text{even}}^\dagger) (-i\hbar X \otimes D_{\text{even}}) (\mathbb{I}_{n_x} \otimes Q_{\text{even}}), \quad (3.118)$$

the unitary Q_{even} is also block-diagonal. This splitting allows for parallel application of the exponentiation of L_z to a vector as well as speeding up single-core computation on account of the smaller matrices.

However If a matrix that were not diagonal were present instead of X , such as another derivative operator, then if we wanted to split *both* matrices each into a sum of two block-diagonal matrices, equation (3.117) would become *four* terms rather than two, and the flow of data for a parallel computation would be somewhat more complicated. For a term in the Hamiltonian that is a product of n operators acting on different subspaces, the number of terms obtained by splitting them all in this way grows exponentially with n . But, when all but one operator in the product is diagonal already, as is the case for angular momentum operators, then splitting can be done as normal.

To conclude this section, there are clear advantages to using finite differences as opposed to Fourier transforms when it comes to parallelisability, boundary conditions, and non-uniform grids, but if these are not a concern then both Fourier transforms and finite differences run at approximately equal speeds in practice, meaning one should use whichever is easiest to implement.

Many of these advantages of finite differences over Fourier methods are also enjoyed by the finite element discrete variable representation (FEDVR) method, discussed in the next section. Whilst it's also claimed that FEDVR has a number of advantages over finite-differences [60, 62], I'll argue that most of the comparisons don't stand up to scrutiny, and that finite differences are often still the right choice for the contexts in which FEDVR is argued to be superior.

3.4.3 Stability and the finite element discrete variable representation

Another method of spatial discretisation is the finite element discrete variable representation (FEDVR). As the name suggests, it is a finite element method, using a set of basis functions within each of a number of spatially distributed *elements*, with adjacent elements linked at their boundaries. Here I will not provide a self-contained description of the FEDVR method itself, more detail can be found in [55, p. 285] and specifically in the context of Bose–Einstein condensation, [60, 62]. I'll instead introduce the points that are relevant to the conclusion that I drew regarding the method for the purposes of time-evolving wavefunctions, which is that all practicalities considered, it is less useful than simple central finite differences for these problems.

As a finite element method, FEDVR divides space into a number of ‘elements’, joined to their neighbouring elements at their edges, within each of which the function being solved for is represented as a linear sum of basis functions. Finite element methods such as

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this are useful for problems with irregular boundary conditions, or requiring variable grid sizes, as the elements need not have the same size (area/volume, etc.) as each other, and parameters on which the accuracy of the numerical method depends can be varied from element to element, such that precision is high where it is needed and low where it is not. In addition, the basis functions used within the elements can be chosen so that conserved properties of the differential equation are also inherently conserved by the simulation resulting in a *geometric integrator*. The FEDVR method though does not make use of this possibility, although it can be used with split-step methods for time propagation, which preserve the wavefunction's norm.

Within each element in the FEDVR method, the wavefunction is represented as a linear sum of a set of polynomial basis functions, specially chosen to have some desirable properties. The polynomials are not orthogonal, but at a set of gridpoints, all but one of them is zero, meaning that as with the Fourier pseudospectral method, they can be used as a pseudospectral basis—computing how much of each polynomial to include in the representation of a wavefunction by using only the wavefunction's value at the gridpoint at which each polynomial is zero, and computing the matrix elements of operators using approximate integration based on a discrete sum taking into account only those same gridpoints.

So far (within each element at least) what I have described does not sound too different to the Fourier pseudospectral method. But in the discrete variable representation—which is what the method used within each element of FEDVR is called when used by itself—the gridpoints are not equally spaced. Rather they are more densely packed toward the edges of the element, and least densely packed in the middle of each element. The locations of the gridpoints are chosen according to a *quadrature rule*, which is a method of approximating the definite integral of a function as a weighted sum of the function's values at a specific set of points. There are many quadrature rules, each with a different set of points and weights, but a common feature to them is that the points are more closely spaced at the edges of the integration region. In the context of the pseudospectral method, the chosen quadrature rule is what we are using when we are evaluating integrals based only on the function's values at discrete points. For equally spaced points, the weights are all equal, but for unequally spaced points they are not (and vary depending on the quadrature rule in use).

The use of unequally spaced points increases the accuracy of integrals evaluated this way, to the extent that the result will be exact if the integrand is a polynomial of degree less than a given degree that depends on the quadrature scheme. One would think that equally spaced points would produce the most accurate approximate integrals, but this is not true: approximating functions as polynomials equal to the value of the function at a discrete set of points is vulnerable to Runge's phenomenon—spurious oscillations in the approximation near the edges of the region²⁴. The oscillations are minimised, however, if the density of points increases near the edge, specifically if the density of points approaches $1/\sqrt{1-x^2}$ for the normalised integratiton region $x \in [-1, 1]$ as the number of points goes to infinity [76].

By choosing a quadrature scheme with two points located exactly on each edge of the integration regions, such as *Gauss-Lobatto* quadrature [60], the elements of the FEDVR method can be joined together, with adjacent elements sharing these edge points (see Figure 3.8). This allows one to represent a wavefunction as a list of coefficients for how much of each basis polynomial in each element contributes to it, with the extra condition that some of these coefficients—those corresponding to the edge points—must be equal in order to ensure the wavefunction is continuous across the boundaries between elements.

With some care taken in regard to the points shared between the elements, the FEDVR method provides one with a means of approximating wavefunctions as a finite sum of basis

²⁴Note that although finite differences are also based on polynomial interpolation, *central* finite differences are not vulnerable to Runge's phenomenon because the interpolated polynomials at the 'edge' of the region are never used for anything: a different polynomial is used for each point, with each polynomial and its derivatives only ever being evaluated at its central interpolation point.

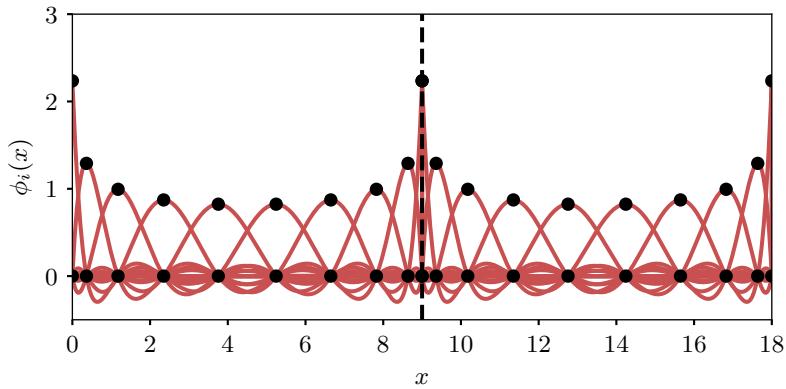


Figure 3.8: An example of the normalised, but non-orthogonal basis polynomials used in the finite-element discrete variable representation, shown here for ten-point Gauss–Lobatto quadrature and two elements. Note that each basis polynomial is nonzero at exactly one point and zero at all others, though it can be nonzero in between points. Outside of its element, a given basis function is defined to be zero, except for the so-called “bridge functions”, which are nonzero in two adjacent elements but zero everywhere else, and surprisingly have a discontinuous derivative at the boundary—necessitating some care in evaluating matrix elements of differential operators.

functions, and of approximately calculating the matrix elements of operators in this basis. After that, one can simply apply the operators to the state vectors in order to compute time derivatives, and propagate in time using fourth order Runge–Kutta (section 3.3) or similar, or one can exponentiate the operators, all at once or one at a time, as part of a split-step method (section 3.2.4). The FEDVR method does not require anything in particular about time propagation—like finite differences or the Fourier pseudospectral method, it is only a way of spatially discretising the differential equation you are trying to solve.

Now we arrive at what makes the FEDVR method exciting for those who want to massively parallelise their simulations. When one calculates the matrix representation of, say, a one-dimensional differential operator in the FEDVR basis, one gets matrices that look like Figure 3.9, with an almost-block-diagonal form.

This is because when the state vector is acted upon by some operator, the result for most points in the resulting vector depends only on the points in the input vector *within the same element*. The exceptions are the points shared between elements—for which the value in the output vector depends on the points in the input vector in *both* adjacent elements, which is why the matrix is not perfectly block diagonal.

Why is this exciting? Well, it means that the operator can be written as a sum of two block diagonal matrices, similar to the splitting of finite-difference matrices in section 3.2.4, and in the same way allows the matrix or its exponentiation to be efficiently applied to a state vector in parallel. The difference between this and the case of finite differences is the number of points of overlap between adjacent blocks, which corresponds to the amount of data that must be transferred between parallel threads or cluster nodes at each step during a parallel computation. In finite differences, the number of points that must be exchanged at boundaries in each step or sub-step of whichever time propagation method is used is equal to the bandwidth of the matrix. In FEDVR, so long as the boundaries of the regions of space assigned to each thread or cluster node align with boundaries between elements, *only one* point need be exchanged. Increasing the number of points within

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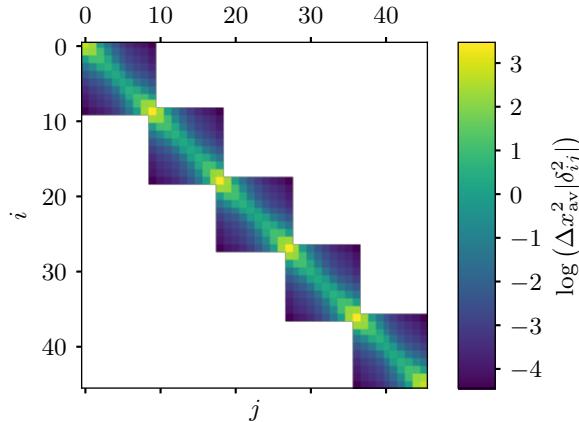


Figure 3.9: The matrix form of the second derivative operator in **FEDVR**, for five elements each with ten DVR points. The colour scale is logarithmic, showing the log of the absolute values of the matrix elements δ_{ij}^2 scaled by Δx_{av}^2 , where Δx_{av} is the average spacing between the unequally spaced grid points. Zero matrix elements are shown as white. One can see that the second derivative at each grid point is computed using only the points within the same element, except for the points shared by adjacent elements, at which the second derivative depends on points in both adjacent elements.

each element—but reducing the number of elements so as to keep the total number of points constant, decreases the discretisation error of the method, but still, only one point needs to be exchanged at boundaries. This contrasts with finite differences, for which increasing the order of the finite difference scheme increases the matrix bandwidth and hence increases the number points required to be exchanged at the boundaries. So it would seem that **FEDVR** ought to scale much better in parallel implementations, which is a large part of its appeal [60, 62].

However, I've noticed that when using both **FEDVR** and finite differences to simulate Bose–Einstein condensates using the Gross–Pitaevskii equation, smaller timesteps are required when using the **FEDVR** method in otherwise comparable setups. By this I mean that, without damping, there is a timestep size at which the GPE simulated using **RK4** (which is a conditionally stable algorithm) is unstable and diverges. Similarly in split-step methods, although the error is bounded, there is a timestep size at which the error rapidly grows to that bound and the wavefunction no longer approximately resembles the true solution. Below this threshold timestep, the error scales as $\mathcal{O}(\Delta t^4)$ for **RK4**, and $\mathcal{O}(\Delta t^4)$, $\mathcal{O}(\Delta t^2)$, $\mathcal{O}(\Delta t)$ for fourth, second, and first order split-step respectively, as expected. But in practice, all these methods are so accurate that one desires to use the largest timestep one can without this blowup (or soft-blowup in the case of the unitary split-step methods) occurring during the time interval one wants to simulate. And my observation was that the threshold timestep is always smaller for **FEDVR** than for finite differences or the Fourier method for determining spatial derivatives, necessitating smaller timesteps for stability or, in the case of the unitary methods, reasonable accuracy.

So why is this? For **RK4**, what is the stability criterion and why might **FEDVR** violate it more easily than finite differences? And how might we understand the sudden decrease in accuracy of the split-step methods at a similar threshold timestep size, despite them being unconditionally stable?

The stability criteron for **RK4** when applied to a linear differential equation of the form:

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$$\frac{d\psi}{dt} = A\psi, \quad (3.119)$$

where A is a matrix with all imaginary eigenvalues (as is the case for Hamiltonian evolution where $A = -\frac{i}{\hbar}H$), is [77]:

$$\Delta t < \frac{2\sqrt{2}}{\rho(A)} \quad (3.120)$$

where $\rho(A)$ is the *spectral radius* of A , defined as the absolute value of the largest (by absolute value) eigenvalue of A .

The Gross–Pitaevskii equation is not linear, but we'll put that to the side for the moment—it will be relevant shortly. In the linear case of the Schrödinger wave equation, an upper bound for the spectral radius of A can be computed from the absolute eigenvalue from the kinetic term, plus the maximum absolute eigenvalue from the potential term. Using the Fourier method to compute spatial derivatives, the largest eigenvalue of the kinetic part of the Hamiltonian is that of the Nyquist mode with $k_{\text{Nyquist}} = \pi/\Delta x$ for grid spacing Δx , leading to:

$$\rho(A) \leq \left| -\frac{i}{\hbar} \frac{\hbar^2 k_{\text{Nyquist}}}{2m} \right| + \left| -\frac{i}{\hbar} V(\mathbf{r}) \right|_{\max} \quad (3.121)$$

$$\Rightarrow \Delta t < \frac{2\sqrt{2}\hbar}{\frac{\hbar^2 \pi^2}{2m\Delta x^2} + |V(\mathbf{r})|_{\max}}. \quad (3.122)$$

In the limit of small Δx , the kinetic term dominates and the stability criterion becomes:

$$\Delta t < \frac{4\sqrt{2}m\Delta x^2}{\pi^2\hbar}, \quad (3.123)$$

whereas in the limit of large potential:

$$\Delta t < \frac{2\sqrt{2}\hbar}{|V(\mathbf{r})|_{\max}}. \quad (3.124)$$

These results match our intuition somewhat in terms of dynamical phase evolution—eigenvalues of the Hamiltonian are energies and determine how quickly the elements of the state vector accumulate dynamical phase. For each circle around the complex plane that a dynamical phase of 2π entails, we expect to require at least a few timesteps to resolve said circle, and the above shows that ‘a few’ means $2\sqrt{2}$. When this dynamical phase evolution is in Fourier space, an accumulated phase of 2π will appear in real space as that component having propagated a distance of one wavelength, so the intuition here is that several timepoints are required in the time interval during which the fastest wave (also the Nyquist mode) propagates a distance of one wavelength at its phase velocity.

As a sidenote, if it is known which term of the Hamiltonian dominates the stability criterion, that term can be removed by use of an *interaction picture* (discussed in section 3.6), essentially treating the dynamical phase evolution due to that term analytically. And if the term is not constant, but is *almost* constant, one can still treat *most* of the dynamical phase evolution analytically, transforming the differential equation onto one with much smaller eigenvalues for that term, in both cases allowing one to take potentially much larger timesteps due to the above reasoning. Fourth order Runge–Kutta in the interaction picture (RK4IP) [78] uses an interaction picture to treat the kinetic term of the

Schrödinger or Gross–Pitaevskii equation analytically, whereas my ‘fourth order Runge–Kutta in an instantaneous, local interaction picture’ method presented in section 3.6 removes (most of) the potential term. Both methods allow larger timesteps to be taken, but in different circumstances depending on which term is dominating the Hamiltonian.

The problem with FEDVR then, is that it requires smaller timesteps than finite differences because its kinetic energy operator has larger eigenvalues than an equally accurate finite difference method. How much larger?

In order to make a fair comparison, we need to know how many DVR basis functions are required per element in order to compute equally accurate second derivatives as a given finite difference scheme. With N points per element, FEDVR represents the state vector within each element in a spectral basis of polynomials up to degree $N - 1$. Therefore a polynomial of degree $N - 1$ or less can be represented exactly, any other function is subject to truncation error²⁵. If one varies the number of points per element, varying the number of elements in order to keep the total number of points constant, the truncation error in representing an arbitrary function in this basis is therefore $\mathcal{O}(\Delta x^N)$, where Δx is either the size of each element, or equivalently the average spacing between grid points (these two differing only by a constant factor if the total number of points is held constant).

In FEDVR the derivative operator, regardless of whether its matrix elements are computed with integrals or the quadrature rule, is exact [60]. The relevant error in a derivative of a wavefunction is therefore determined by this truncation error of representing it in the spectral basis in the first place:

$$\psi_{\text{approx}}(x_i) = \psi_{\text{exact}}(x_i) + \mathcal{O}(\Delta x^N) \quad (3.125)$$

$$\Rightarrow \psi''_{\text{approx}}(x_i) = \psi''_{\text{exact}}(x_i) + \mathcal{O}(\Delta x^{N-2}). \quad (3.126)$$

²⁵Note that this truncation error is distinct from that of evaluating *integrals* using the Gauss–Lobatto quadrature rule, which is exact for integrands that are polynomials of degree $2N - 2$ or less.

Central finite difference approximations to second derivatives on the other hand have error $\mathcal{O}(\Delta x^{N-1})$ where N is the total number of points used to compute the derivative at each point (for example the three-point central finite difference rule for second derivatives has error $\mathcal{O}(\Delta x^2)$, the five-point rule is accurate to $\mathcal{O}(\Delta x^4)$, etc). With this knowledge we can translate our question

Which is less computationally intensive to simulate the GPE or Schrödinger wave equation: m^{th} -order accurate FEDVR or m^{th} -order accurate finite differences?

to

which is less computationally intensive to simulate the GPE or Schrödinger wave equation: $m + 2$ points-per-element FEDVR or $m + 1$ point central finite differences?

The argument in favour of FEDVR is that as m grows, finite differences require an increasing number of points to be exchanged at the boundaries between cluster nodes/threads etc in a parallel implementation, whereas so long as each boundary between spatial regions allocated to different cluster nodes aligns with the boundary between two elements, only one point need be exchanged per timestep in FEDVR, no matter how many points per element there are. Therefore, in the limit of high accuracy, FEDVR wins, it seems.

The problem with this argument is that it compares only the amount of work that needs to be done *per timestep*. But, since the allowed timestep size required for stability (at least for fourth order Runge–Kutta, I’ll argue shortly why I think this generalises to other timestepping schemes as well) depends on the spectral radius of the kinetic energy operator, and the kinetic energy operator is not the same as m grows, more work is needed to show which method is least computationally expensive per unit *simulation time*.

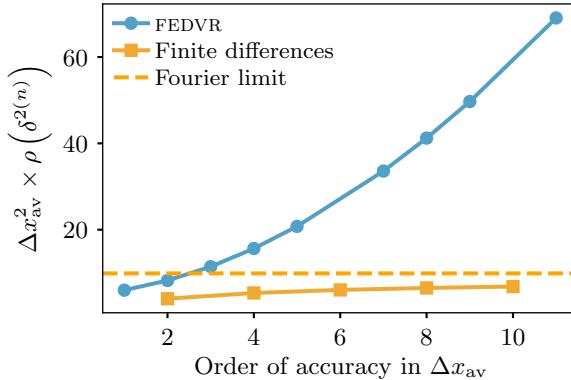


Figure 3.10: Scaling of the spectral radius (maximum absolute eigenvalue) of the second derivative operator with respect to the order of accuracy for finite differences and FEDVR. Results were numerically computed by constructing a 721×721 matrix (chosen to allow various combinations of number of DVR points per element whilst holding total number of points constant) for each order of accuracy and numerically diagonalising. The spectral radius for FEDVR can be seen to scale quadratically with the order of accuracy, whereas for finite differences the maximum eigenvalue approaches that of the Fourier method.

The spectral radius of the kinetic energy operator when approximated using central finite differences does not grow without limit as the number of points used to compute derivatives increases—indeed, its eigenspectrum converges to that of the Fourier method (since the Fourier method can be considered the limit of ‘infinite order’ finite differences [72]), with maximum eigenvalue equal to the kinetic energy of the Nyquist mode. The kinetic energy operator approximated with FEDVR on the other hand does not have a bounded spectral radius as one increases the number of points per element whilst holding the total number of gridpoints constant. Instead its spectral radius increases quadratically with the number of points (equivalently with the order of accuracy of the derivatives), as shown in Figure 3.10.

This result ought to be expected, at least qualitatively. The density of points in FEDVR is higher toward the edges of the elements than in their centres, and if one increases the number of points per element whilst decreasing the number of elements so as to hold the total number of points approximately constant, the smallest spacing between any two adjacent points will be inversely proportional to the number of points per element. We already know that in high order finite differences or the Fourier method, the spectral radius of the kinetic energy operator is simply the kinetic energy of the Nyquist mode, and being the mode described by a wave with half a wavelength spanning one grid spacing, its kinetic energy is proportional to the square of the grid spacing. It is not therefore surprising that in FEDVR when the grid spacing is linearly smaller—albeit locally—that the spectral radius of the kinetic energy operator is quadratically larger. The same result as above could be obtained by asking: “what is the smallest grid spacing, and what is the kinetic energy of the Nyquist mode corresponding to that grid spacing?”, and then declaring the stability criterion to be that one must have at least a few points per period of the frequency corresponding to that energy.

Since the spectral radius of FEDVR operators grows faster with increasing order of accuracy in derivative operators than finite differences, FEDVR requires smaller timesteps for stability when used with RK4. Even at low order accuracy, finite difference operators have smaller spectral radii, and so at all levels of accuracy RK4 is stable with finite differ-

ences for larger timesteps than with **FEDVR**. In the limit of high accuracy then, compared to finite differences **FEDVR** requires *quadratically* more timesteps to be taken for stability, whereas it requires only *linearly* fewer points to be transferred at the boundaries between threads or cluster nodes per timestep. The larger number of timesteps is therefore the dominant effect, more than cancelling out the benefit of having to transfer fewer points at boundaries per timestep than with finite differences. Indeed, even if transferring data between nodes or threads is the bottleneck of a parallel simulation, *more* points are being transferred all up with **FEDVR**—they are just spread out over more timesteps.

That's fourth order Runge–Kutta. What about split-step methods? The split-step methods discussed in section 3.2.4 are unconditionally stable and have bounded error. However, their error can still be large enough to make the results not useful. As shown in (3.46), the leading error term in first-order split-step is $1/2\hbar^2[V, K]\Delta t^2$ where $[V, K]$ is the commutator between the (discretised) kinetic and potential energy operators. The leading term of V is proportional to the discretised \hat{x} operator X , and the kinetic energy operator is proportional to an n^{th} order accurate discretised second derivative operator $\delta^{2(n)}$. Therefore the error per timestep scales with $[X, \delta^{2(n)}]\Delta t^2$. Although this expression is not bounded, the error in first-order split-step is nonetheless bounded because this error appears as the argument of a complex exponential. When expanding the complex exponential as its Taylor series, this error term is the leading term, but when it grows it leads merely to a complex exponential with unbounded phase error, not to unbounded absolute error. Nonetheless unbounded phase error still makes the results of a simulation unlikely to be useful.

The largest (by absolute value) eigenvalue of this commutator sets the rate at which erroneous dynamical phase is accumulated by some eigenstate of the commutator. When this phase becomes comparable to 2π per timestep, the simulation has clearly lost any semblance of accuracy, despite still being technically “stable”. Therefore as before, the spectral radius of the matrix of this commutator can be used to define a *pseudo*-stability criterion, such that when the spectral radius of $[V, K]\Delta t^2/2\hbar^2$ becomes comparable to unity, the error will dominate the simulation results. Although the value of the spectral radius will depend on the details of the potential energy operator V for the particular problem, we can still ask how it scales with increasing order of accuracy of K given that X is the leading term of V . This is done in Figure 3.11, comparing the spectral radius of the commutator when using derivative operators of various orders in both finite differences and **FEDVR** approximations.

The result is that as with **RK4**, the pseudo-stability criterion allows at all orders of accuracy for larger timesteps when using finite differences than when using **FEDVR**, and that the required timestep is bounded from below when increasing the order of accuracy of finite differences, but can become arbitrarily small for increasing accuracy of **FEDVR**. However the situation is not so dire for **FEDVR** when used with split-step as it is with **RK4** timestepping. Because the error term in split-step is this commutator multiplied by Δt^2 , the timestep required for pseudo-stability scales inversely with the *square root* of the spectral radius of the commutator—not with the spectral radius itself as with **RK4**. Furthermore the spectral radius of the **FEDVR** commutator increases only *linearly* with increasing order of accuracy, not quadratically as with **RK4**. Therefore if transferring data between nodes (with time cost proportional to the amount of data transferred) is the bottleneck of one's simulation, then with increasing accuracy, **FEDVR** *does* win out. For high enough order of accuracy n , a factor of \sqrt{n} more timesteps are needed with **FEDVR** compared to finite differences, but a factor of n fewer points are sent per timestep. This results in a factor of $1/\sqrt{n}$ fewer points being sent per unit simulation time with **FEDVR** as compared to finite differences. The question of which method is faster then comes down to which is more costly: extra timesteps, or extra data transfer? Using **FEDVR** will save you a factor of \sqrt{n} in data transfer at a cost of a factor of \sqrt{n} in the number of timesteps. Given InfiniBand interconnects on cluster computers with tens of gigabits per second

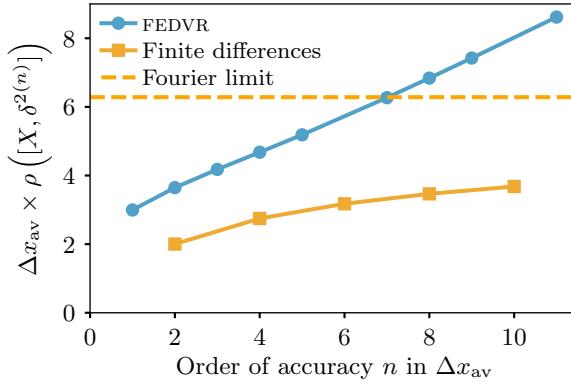


Figure 3.11: Scaling with order of accuracy of the spectral radius of the commutator of position and second derivative operators in the pseudospectral finite difference method and FEDVR. The spectral radius of the commutator grows linearly with increasing order of accuracy for FEDVR, whereas it is bounded for finite differences, approaching that of the Fourier method.

of bandwidth, and latency that does not scale with the amount of data, I suspect that a \sqrt{n} increase in cost of processing within nodes/threads is almost always the larger cost to pay. Therefore I suspect the benefits of FEDVR for parallel simulations of the Schrödinger wave equation or Gross–Pitaevskii equation are limited.

3.4.4 Nonlinearity considerations

The above arguments about stability all disregarded the nonlinearity of the Gross–Pitaevskii equation. Although the stability of a numerical method is very difficult to analyse for nonlinear differential equations, there is a simple argument for heuristically putting an upper bound on the timestep required in order to accurately model the effect of the GPE’s nonlinear term. Essentially, density waves in the condensate propagate not at the phase velocity of the condensate wavefunction, but at its group velocity. The group velocity of the shortest possible wavelength in the system therefore sets an upper bound for the propagation of information due to the nonlinear term. In order to correctly model this term, timesteps must be short enough that one’s numerical method is evaluating the nonlinear term frequently enough that fast moving density waves can’t ‘skip’ gridpoints in between evaluations. If the smallest wavelength is that of the Nyquist mode, or for nonuniform grids the mode whose wavelength is twice the smallest grid spacing, then the maximum group velocity is:

$$v_g(k_{\max}) = \frac{\partial E_K(k_{\max})}{\partial p(k_{\max})} = \frac{\hbar k_{\max}}{m} \frac{\pi \hbar}{m \Delta x_{\min}} \quad (3.127)$$

Asking how long it takes a wave to move a distance of Δx_{\min} at this velocity then results in what I’m calling the *dispersion timescale*:

$$\tau_d = \frac{m \Delta x_{\min}^2}{\pi \hbar}. \quad (3.128)$$

The numerical method being used is now fairly irrelevant: one must evaluate the nonlinear term at least as often as every τ_d in order to be able to model it accurately, otherwise density waves may move past each other faster than they can be resolved by the

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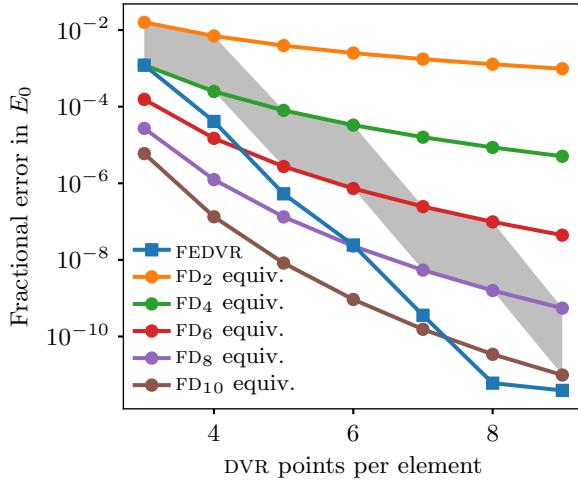


Figure 3.12: Comparison of accuracy of different discretisation methods in representing the groundstate of a harmonic oscillator $V(x) = x^2$ in a spatial region $-10 < x < 10$. For the **FEDVR** results, 20 elements were used with a varying number of DVR points per element, leading to decreasing error in the groundstate energy with increasing number of points. For each number of DVR basis points, finite difference operators are produced using the same total number of points, but on a uniform grid. This is still not a fair comparison however, because it is not equally computationally expensive to apply a **FEDVR** operator or a finite differences operator to a state vector. As mentioned in earlier sections, the cost of applying operators is dependent on their bandwidth—related to how many surrounding points are coupled to each point by the operator. This comparison is made above by computing the *average* bandwidth of each **FEDVR** operator (since it is not the same for all points), and shading in the region between the two finite differencing orders that have bandwidths either side of the result. The resulting shaded region shows **FEDVR** and finite differences are quite comparable when judged by accuracy at fixed cost of applying operators, with the gap small enough to be closed by increasing the bandwidth of a finite differencing scheme only by one.

timestepping, and their interference patterns will be aliased by the too-slow timestepping, resulting in incorrect nonlinear dynamics.

Note that up to constant factors, this criterion for accurate modelling is the same as the stability criterion (3.123) for **RK4** when the kinetic term dominates the Hamiltonian. Therefore although first-order split-step as argued above appears to be more forgiving to **FEDVR** than **RK4** is, this is no longer the case when modelling the GPE as opposed to the linear Schrödinger wave equation, and although I didn't extend the argument in the previous section to second or fourth order split-step (figuring out which commutator is the leading term is much more involved), they too are subject to the same requirement to sample the nonlinear term this frequently, even if their linear pseudo-stability region might be larger.

3.4.5 Conclusion

This leaves us with little remaining benefit to **FEDVR** over finite differences for the Schrödinger wave and Gross–Pitaevskii equations. The argument that **FEDVR** scales better for parallel computation is unconvincing as argued above. The fact that it allows nonuniform grids is also not unique—central finite differences allow for nonuniform

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grids as well [73]. Reference [62] compares the accuracy of finite differences to FEDVR in the context of computing the groundstate energy of a harmonic oscillator using imaginary time evolution (see section 3.5.1), but only uses second-order finite differences in the comparison whereas higher-order FEDVR schemes are used. If this restriction was in order to hold the number of points transferred at boundaries between parallel computing units constant (since second-order finite differences require, like FEDVR, only one per timestep), then this is not an equal comparison as the number of points transferred does not limit computation time as I've shown—had the authors included comparisons with higher order finite differencing schemes, they would have resulted in comparable accuracy (see Figure 3.12) between FEDVR and finite differences, but with finite differences completing in less time owing to the larger timesteps allowed by the increased range of stability of finite differences.

Perhaps for problems with certain unusual boundary conditions or strangely shaped regions, or where one can construct a geometric integrator that conserves some quantities of interest other than merely the wavefunction norm, FEDVR may still make sense. But unless there is a compelling reason, it seems that simple finite differences, as naïve as it might seem at first, remains a practical choice for highly accurate and potentially massively parallelised simulations of the GPE.

3.5 Finding ground states

For systems with continuous degrees of freedom, such as Bose–Einstein condensates, finding groundstates by directly diagonalising the Hamiltonian matrix is impractical due to the size of said matrix, which we normally avoid even constructing for systems with more than one dimension when simply propagating condensate wavefunctions in time. The pseudo-Hamiltonian for the Gross-Pitaevskii equation is not even linear, implying direct diagonalisation would not even return groundstates unless one diagonalised matrices repeatedly, each with an improved estimate of the groundstate until a self-consistent solution was found.

Below are two methods that are typically used instead to find groundstates in these cases, as well as a generalisation (which only applies to linear systems) for finding other excited states of quantum systems.

3.5.1 Imaginary time evolution

Imaginary time evolution [79] is a robust, if somewhat inefficient method of finding groundstates of quantum systems, and can be generalised to finding excited states as well (Section 3.5.3, below). The basic idea is to evolve the following differential equation:

$$\hbar \frac{d\psi}{dt} = -H\psi \quad (3.129)$$

in time instead of the time-dependent Schrödinger equation (3.3), for some wavefunction or state vector ψ and matrix/linear operator H . The name *imaginary time evolution* is due to the fact that this differential equation can be obtained by making the substitution $t \rightarrow -it$ in the time-dependent Schrödinger equation, and hence can be thought of as propagating the Schrödinger equation in (negative) imaginary time. This method is capable of evolving an initial guess for ψ into the groundstate of the given Hamiltonian. The reason behind this can be seen if we rewrite (3.129) in the eigenbasis of H with eigenvalues $\{E_n\}$ and eigenkets $\{|n\rangle\}$:

$$\hbar \frac{d}{dt} \langle n|\psi(t)\rangle = -E_n \langle \phi_n|\psi(t)\rangle, \quad (3.130)$$

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which has solutions:

$$\langle n|\psi(t)\rangle = \exp\left[-\frac{E_n}{\hbar}t\right] \langle n|\psi(0)\rangle, \quad (3.131)$$

that is, the coefficient of each eigenstate exponentially decays with a decay rate proportional its energy. After enough time has elapsed (much larger than the difference in decay timescales between the two lowest lying eigenstates), all coefficients $\langle n|\psi(t)\rangle$ will have decayed, but the coefficient corresponding to the groundstate will have decayed by far the least, leading to any initial superposition evolving into one dominated by the groundstate.

This method is flexible and robust, and can be used to find groundstates subject to imposed conditions such as a phase winding about a point, resulting in a vortex state. It can also be used to ‘smooth’ candidate wavefunctions, making a guessed initial state of a condensate wavefunction more like a physically realistic one by lowering its energy somewhat, but not propagating far enough to obtain the actual groundstate. This smoothing method is used in Section 3.6 to create an approximately correct density profile for a condensate wavefunction after the phase patterns for randomly distributed vortices are artificially imposed, resulting in physically realistic initial conditions for a turbulent condensate.

For practical reasons, exponential decay cannot be carried out indefinitely, as eventually all numbers underflow to zero as they become too small to be represented on a computer. In addition, the wavefunction/state vector obtained after exponential decay will not be normalised. For these reasons, one often normalises the wavefunction/state vector to unit norm in between each step of imaginary time evolution.

Normalisation must be paid attention particularly in the case of the Gross–Pitaevskii equation, since it has a nonlinear pseudo-Hamiltonian. This pseudo-Hamiltonian must always be evaluated with a normalised condensate wavefunction as input, otherwise the imaginary time evolution may, depending on the numerical method employed, produce a final condensate wavefunction which is the groundstate not in the presence of the actual nonlinear term, but in the presence of one that has had one timestep worth of exponential decay applied to it. Since the magnitude of the wavefunction matters in the Gross–Pitaevskii equation, one must normalise every time a nonlinear H is evaluated in order to prevent this problem (which can also be mitigated by taking tiny timesteps to minimise the amount of decay within each timestep, but this is wasteful when larger steps could otherwise be used).

An exception to the above normalisation warning is when a first order explicit method such as the first-order split step method (Section 3.2.4) or the Euler method is used—since these methods only ever evaluate the nonlinear pseudo-Hamiltonian at the start of a timestep, normalising the wavefunction once per timestep is sufficient, whereas when using higher order methods that evaluate the pseudo-Hamiltonian at intermediate times one will have to take care to normalise more often.

The imaginary time evolution method is a *relaxation method*, and as such only the final wavefunction/state vector after a sufficient period of evolution is of interest, not the intermediate states. As such, inaccurate but fast methods such as first order split-step or the Euler method may be readily employed, and used with the largest timesteps that do not lead to numerical instability. Second order Runge–Kutta (the explicit midpoint method) can be a good balance due to being faster per-step than more accurate methods, whilst allowing larger timesteps than the Euler method due to better stability characteristics.

A final note is that energies of quantum systems are arbitrary up to a global additive constant, and so energies can be negative or positive or some mix of both in the same system. In this case, imaginary time evolution can have some or all coefficients of energy eigenstates exponentially *growing* rather than decaying. This does not modify the end result at all, as the groundstate energy, if negative, will be more negative than any other energy and hence its coefficient will exponentially grow faster than all others. If an

²⁶Note that the energy of a BEC using the Gross–Pitaevsii pseudo-Hamiltonian cannot be computed simply as the expectation value of the pseudo-Hamiltonian—this double counts the nonlinear term, which should be halved in such energy calculations. The chemical potential however can be computed as this expectation value.

estimate of the groundstate energy is known ahead of time, it can be advantageous to offset the system Hamiltonian with a constant energy such that the groundstate energy is approximately zero, which minimises the amount of exponential growth/decay and allows for the largest timesteps within the stability range of the numerical method employed. One can even dynamically compute an energy or chemical potential estimate²⁶ to update the energy offset as the computation proceeds. Imaginary time evolution can be slow enough for non-trivial cases that these considerations can be relevant.

Since imaginary time evolution comprises the evolution of a state vector in time, it is amenable to all the parallel processing techniques we have outlined in this chapter depending on the timestepping and spatial discretisation employed.

3.5.2 Successive over-relaxation

Successive over-relaxation (SOR) [80] is a much faster method of finding groundstates than imaginary time evolution, but slightly less flexible. Strictly speaking (that is, before any ad-hoc modifications), SOR is a method for solving linear systems of the form:

$$A\psi = \mathbf{b}, \quad (3.132)$$

for some unknown vector (or discretised function) ψ given a vector (or discretised function) \mathbf{b} for the right hand side and a matrix (or discretised linear operator) A for the left hand side. The method is of most benefit when A is diagonally dominant, lending SOR to use with finite difference methods. As with imaginary time evolution, SOR requires an initial guess for ψ that is repeatedly updated to produce an increasingly accurate approximate solution to the given linear system.

Successive over-relaxation proceeds by updating each element ψ_i of ψ one at a time to be more consistent with the linear system and all other estimated elements of ψ , by solving the i^{th} linear equation of (3.132) for ψ_i in terms of A , b_i , and the other elements of ψ :

$$\psi_i = \frac{b_i - \sum_{j \neq i} A_{ij} \psi_j}{A_{ii}}. \quad (3.133)$$

If one updates every element ψ_i of ψ in this manner, based on the other elements $\psi_{j \neq i}$ of the estimate for ψ at the previous iteration, the method that results is called the Jacobi method, with update rule:

$$\psi_{i \text{ Jacobi}}^{(k+1)} = \frac{b_i - \sum_{j \neq i} A_{ij} \psi_j^{(k)}}{A_{ii}}, \quad (3.134)$$

where $\psi^{(k)}$ is the k^{th} iterative estimate of ψ . If one allows use of already updated elements of ψ when updating subsequent elements of ψ , one obtains the more efficient Gauss–Seidel method:

$$\psi_{i \text{ GS}}^{(k+1)} = \frac{b_i - \sum_{j \neq i} A_{ij} \psi_j^{(\text{latest})}}{A_{ii}}, \quad (3.135)$$

where $\psi_j^{(\text{latest})} = \psi_j^{(k+1)}$ if $\psi_j^{(k+1)}$ has been computed already, and $\psi_j^{(k)}$ if not. Assuming elements of ψ are updated in order of increasing j , one can split the sum into two for these two cases and write:

$$\psi_{i \text{ GS}}^{(k+1)} = \frac{b_i - \sum_{j < i} A_{ij} \psi_j^{(k+1)} - \sum_{j > i} A_{ij} \psi_j^{(k)}}{A_{ii}}, \quad (3.136)$$

which is how the method is often presented; however, since one can update elements in any order (or equivalently, the order of the basis vectors in ψ is arbitrary), I prefer the form (3.135) as it leaves the iteration order undetermined. One reason to impose an update order other than just that of increasing index in the array storing ψ in computer memory is to minimise communication latency in parallel implementations: one may choose to update points near one or more edges of the spatial region assigned to one CPU core or cluster node first, in order that they become available to other cores/nodes sooner, before updating the interior points that other cores/nodes do not depend on. In this case of parallel processing there isn't even a strict order since elements are being updated independently by separate cores/nodes simultaneously. Gauss–Seidel and successive over-relaxation still work well in this case, they just become a little more like the Jacobi method the more independent updates are occurring on different parts of the vector ψ .

The final feature of successive over-relaxation is that elements of ψ are updated not to the values given by the Gauss–Seidel method, but overshoot them by a factor α :

$$\psi_{i \text{SOR}}^{(k+1)} = \psi_i^{(k)} + \alpha \left(\frac{b_i - \sum_{j \neq i} A_{ij} \psi_j^{(\text{latest})}}{A_{ii}} - \psi_i^{(k)} \right). \quad (3.137)$$

where $\alpha < 2$ is the convergence criterion in the case of an infinite system.²⁷

So what are A and b for quantum mechanics problems? One deviation from pure SOR used in cold atom physics is to allow both A and b to have dependence on ψ , which in practice does not prevent the method from finding solutions. For the groundstate of the Gross–Pitaevskii equation, we can write:

$$H(\psi)\psi = \mu\psi \quad (3.138)$$

where μ is the chemical potential of the condensate in its groundstate and $H(\psi)$ is the matrix for some discretisation of the Gross–Pitaevskii pseudo-Hamiltonian,²⁸ to which we can then apply a single iteration of SOR by updating all elements of ψ in some order using (3.137) with $A^{(k)} = H(\psi^{(k)})$ and $b^{(k)} = \mu\psi^{(k)}$ to produce an improved estimate $\psi^{(k+1)}$ of the groundstate based on the previous estimate $\psi^{(k)}$:

$$\psi_{i \text{SOR}}^{(k+1)} = \psi_i^{(k)} + \alpha \left(\frac{\mu\psi_i^{(k)} - \sum_{j \neq i} H_{ij}(\psi^{(k)})\psi_j^{(\text{latest})}}{A_{ii}} - \psi_i^{(k)} \right). \quad (3.139)$$

$A^{(k+1)}$ and $b^{(k+1)}$ can then be computed based on $\psi^{(k+1)}$ and the process repeated to provide increasingly accurate approximate solutions.

A difference between successive over-relaxation and imaginary time evolution is that SOR finds the state with chemical potential μ (replaceable by the energy E for a linear quantum system), as opposed to the groundstate. In this way successive over-relaxation doesn't tell you what the groundstate energy is, it requires that as input. Typically to find a Gross–Pitaevskii groundstate I will analytically estimate the chemical potential for a given number of atoms or peak density, or whatever my requirement is, using the Thomas–Fermi approximation, and then find the state with that chemical potential using SOR. This results in minor differences compared to the peak density or atom number originally targeted, but the requirement is not usually strict enough for this to be of consequence. For the Gross–Pitaevskii equation, the ability of the atom number to vary means a condensate wavefunction with the given chemical potential can always be found.

However, sometimes one desires the groundstate of the Gross–Pitaevskii pseudo-Hamiltonian given a certain atom number, or one may be solving the time-independent linear Schrödinger equation and not know the eigenstate energies. In the former case

²⁷With the deviation from strict SOR by using non-constant A and b mentioned below, and parallel updates as mentioned above, as well as the fact that systems are not infinite, α in the range of 1.6 to 1.8 is realistic to ensure stability in Gross–Pitaevskii equation simulations.

²⁸A subtle consequence of A and b depending on ψ is that we cannot simplify this to $(H(\psi) - \mu)\psi = 0$, as this would admit $\psi = \mathbf{0}$ as a solution. As written, (3.138) also admits $\psi = \mathbf{0}$ as a solution, however, $A^{(k)} = H(\psi^{(k)})$ and $b^{(k)} = \mu\psi^{(k)}$ are treated as constants within a single iteration of SOR, preventing the method from approaching zero as a solution within each step compared to if ψ was present only on the left hand side.

²⁹Because normalisation for the linear Schrödinger equation is arbitrary, one ought to either normalise ψ at every step, or impose a boundary condition at a single point where ψ is nonzero (simply not updating that point in each SOR step), and normalise once at the end in order to prevent ψ growing or shrinking without limit.

one can to normalise ψ to the desired atom number at each step to impose fixed atom number, but then SOR will not converge to a stationary state/eigenstate unless one has by chance guessed the correct chemical potential. SOR will also not converge for the linear Schrödinger equation²⁹ unless the energy given corresponds to an actual eigenstate.

A strategy that works is to dynamically update the targeted μ (or E) as SOR proceeds, by computing the expectation value of H using the current best estimate of ψ after each SOR step:

$$\psi_{i \text{SOR}}^{(k+1)} = \psi_i^{(k)} + \alpha \left(\frac{\mu^{(k)} \psi_i^{(k)} - \sum_{j \neq i} H_{ij}(\psi^{(k)}) \psi_j^{(\text{latest})}}{A_{ii}} - \psi_i^{(k)} \right), \quad (3.140)$$

where

$$\mu^{(k)} = \frac{\psi^{*(k)} \cdot H(\psi^{(k)}) \psi^{(k)}}{\psi^{*(k)} \cdot \psi^{(k)}} \quad (3.141)$$

This process of dynamically updating the targeted chemical potential or energy can find stationary states and eigenstates, but not necessarily the groundstate. Different initial guesses for ψ may result in convergence to different stationary states/eigenstates, and some attention and guidance from the programmer is required in order to craft initial guesses that result in convergence to the desired states.

3.5.3 Generalisation to excited states via Gram–Schmidt orthonormalisation

³⁰Realised independently by me, but not original [81].

Directly diagonalising a Hamiltonian can be costly in a spatial basis. Another approach³⁰ (for the linear Schrödinger equation) is to find the groundstate using one of the above techniques, and then repeat the process with a fresh initial guess, subtracting off the wavefunction's projection onto the already found ground state at every step. This yields the lowest energy state that is orthogonal to the first—i.e. the first excited state. Repeating the process, but subtracting at each step the projection onto *both* eigenstates found so far, then yields the second excited state and so forth. This is simply the Gram–Schmidt process for finding orthonormal vectors, with the additional step of relaxing each vector to the lowest possible energy given the orthogonality requirement—this ensures the eigenstates of the Hamiltonian are produced, rather than a different orthonormal basis. Extra conditions can be imposed on the wavefunction at each relaxation step in order to obtain particular solutions in the case of degenerate eigenstates. For example, a phase winding can be imposed in order to obtain a particular harmonic oscillator state, otherwise this process produces an arbitrary superposition of basis states that have equal energy.

For the case of successive over-relaxation used with the linear Schrödinger equation, as mentioned in Section 3.5.2, one often cannot predict which eigenstate results when the targeted energy is not already known—the groundstate is not necessarily the one produced by SOR. Nonetheless, relaxation to *some* eigenstate or other, subject to the state vector being orthonormal to all eigenstates found so far, will produce an additional orthonormal eigenstate each time, just not necessarily in order from lowest to highest energy.

3.6 Fourth order Runge–Kutta in an instantaneous local interaction picture

Consider the differential equation for the components of a state vector $|\psi(t)\rangle$ in a particular basis with basis vectors $|n\rangle$. This might simply be the Schrödinger equation, or

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perhaps some sort of nonlinear or other approximate, effective or phenomenological equation not corresponding to pure Hamiltonian evolution. Though they may have additional terms, such equations are generally of the form:

$$\frac{d}{dt} \langle n | \psi(t) \rangle = -\frac{i}{\hbar} \sum_m \langle n | \hat{H}(t) | m \rangle \langle m | \psi(t) \rangle, \quad (3.142)$$

where $\langle n | \hat{H}(t) | m \rangle$ are the matrix elements in that basis of the Hamiltonian $\hat{H}(t)$, which in general can be time dependent, or even a function of $|\psi(t)\rangle$, depending on the exact type of equation in use. If $\hat{H}(t)$ is almost diagonal in the $|n\rangle$ basis, then the solution to (3.142) is dominated by simple dynamical phase evolution, that is:

$$|\psi(t)\rangle \approx \sum_m e^{-\frac{i}{\hbar} E_m t} |m\rangle, \quad (3.143)$$

where E_m is the energy eigenvalue corresponding to the eigenstate $|m\rangle$.

A transformation into an interaction picture (IP) [54, p. 317] is commonly used to treat this part of the evolution analytically, before solving the remaining dynamics with further analytics or numerics. For numerical methods, integration in the interaction picture allows one to use larger integration timesteps, as one does not need to resolve the fast oscillations around the complex plane due to this dynamical phase.

Choosing an interaction picture typically involves diagonalising the time-independent part of a Hamiltonian, and then proceeding in the basis in which that time-independent part is diagonal. However, often one has a good reason to perform computations in a different basis, in which the time independent part of the Hamiltonian is only approximately diagonal,³¹ and transforming between bases may be computationally expensive (involving large matrix-vector multiplications). Furthermore, the Hamiltonian may change sufficiently during the time interval being simulated that the original time-independent Hamiltonian no longer dominates the dynamics at later times. In both these cases it would still be useful to factor out the time-local oscillatory dynamics in whichever basis is being used, in order to avoid taking unreasonably small timesteps.

To that end, suppose we decompose $\hat{H}(t)$ into diagonal and non-diagonal (in the $|n\rangle$ basis) parts at each moment in time:

$$\hat{H}(t) = \hat{H}_{\text{diag}}(t) + \hat{H}_{\text{nondiag}}(t), \quad (3.144)$$

and use the diagonal part at a specific time $t = t'$ to define a time-independent Hamiltonian:

$$\hat{H}_0^{t'} = \hat{H}_{\text{diag}}(t'), \quad (3.145)$$

which is diagonal in the $|n\rangle$ basis. We can then use then use $\hat{H}_0^{t'}$ to define an interaction picture state vector:

$$|\psi_1^{t'}(t)\rangle = e^{\frac{i}{\hbar}(t-t')\hat{H}_0^{t'}} |\psi(t)\rangle, \quad (3.146)$$

which obeys the differential equation:

$$\frac{d}{dt} |\psi_1^{t'}(t)\rangle = e^{\frac{i}{\hbar}(t-t')\hat{H}_0^{t'}} \frac{d}{dt} |\psi(t)\rangle + \frac{i}{\hbar} \hat{H}_0^{t'} |\psi_1^{t'}(t)\rangle, \quad (3.147)$$

where:

$$|\psi(t)\rangle = e^{-\frac{i}{\hbar}(t-t')\hat{H}_0^{t'}} |\psi_1^{t'}(t)\rangle \quad (3.148)$$

is the original Schrödinger picture (SP) state vector.

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³¹For example, a spatial basis which allows for partitioning the integration region over multiple nodes on a cluster or cores on a GPU.

This transformation is exact, no approximations or assumptions have been made. If indeed the dynamics of $|\psi(t)\rangle$ in the given basis are dominated by fast oscillating dynamical phases, that is, the diagonals of $\hat{H}_{\text{diag}}(t)$ are much greater than all matrix elements of $\hat{H}_{\text{nondiag}}(t)$ in the $|n\rangle$ basis, then solving the differential equation (3.147) for $|\psi_I^{t'}(t)\rangle$ should allow one to use larger integration timesteps than solving (3.142) directly. And if not, then it should do no harm other than the (small) computational costs of computing some extra scalar exponentials.

Equation (3.146) defines an *instantaneous* interaction picture, in that it depends on the dynamics at a specific time $t = t'$, and can be recomputed repeatedly throughout a computation in order to factor out the fast dynamical phase evolution even as the oscillation rates change over time. It is *local* in that $H_0^{t'}$ is diagonal in the $|n\rangle$ basis, which means that transformations between Schrödinger picture and interaction picture state vectors involves ordinary, elementwise exponentiation of vectors, rather than matrix products. Thus (3.146), (3.147) and (3.148) can be written componentwise as:

$$\langle n|\psi_I^{t'}(t)\rangle = e^{i(t-t')\omega_n^{t'}} \langle n|\psi(t)\rangle, \quad (3.149)$$

$$\frac{d}{dt} \langle n|\psi_I^{t'}(t)\rangle = e^{i(t-t')\omega_n^{t'}} \frac{d}{dt} \langle n|\psi(t)\rangle + i\omega_n^{t'} \langle n|\psi_I^{t'}(t)\rangle, \quad (3.150)$$

and:

$$\langle n|\psi(t)\rangle = e^{-i(t-t')\omega_n^{t'}} \langle n|\psi_I^{t'}(t)\rangle, \quad (3.151)$$

where we have defined:

$$\omega_n^{t'} = \frac{1}{\hbar} \langle n|\hat{H}_0^{t'}|n\rangle \quad (3.152)$$

This is in contrast to fourth order Runge–Kutta in the interaction picture (RK4IP) [78], in which the interaction picture uses the Fourier basis and thus transforming to and from it involves fast Fourier transforms (FFTs). RK4IP was developed to augment computations in which FFTs were already in use for evaluating spatial derivatives, and so its use of FFTs imposes no additional cost. Nonetheless, an interaction picture based on the kinetic term of the Schrödinger equation (which is the term of the Hamiltonian that RK4IP takes as its time-independent part) may not be useful if that term does not dominate the Hamiltonian, as in the case of a Bose–Einstein condensate in the Thomas–Fermi limit. We compare the two methods below.

3.6.1 Algorithm

The *fourth order Runge–Kutta in an instantaneous local interaction picture* RK4ILIP algorithm is now obtained by using (3.146) to define a new interaction picture at the beginning of each fourth-order Runge–Kutta (RK4) integration timestep. The differential equation and initial conditions supplied to the algorithm are in the ordinary Schrödinger picture, and the interaction picture is used only within a timestep, with the Schrödinger picture state vector returned at the end of each timestep. Thus differential equations need not be modified compared to if ordinary RK4 were being used, and the only modification to calling code required is for a function to compute and return $\omega_n^{t'}$.

Being based on fourth order Runge–Kutta integration, this new method enjoys all the benefits of a workhorse method that is time-proven, and—as evidenced by its extremely widespread use—at a sweet-spot of ease of implementation, accuracy, and required computing power [82].

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Below is the resulting algorithm for performing one integration timestep. It takes as input the time t_0 at the start of the timestep, the timestep size Δt , an array ψ_0 containing the components $\{\langle n|\psi(t_0)\rangle\}$ of the state vector at time t_0 , a function $F(t, \psi)$ which takes a time and (the components of) a state vector and returns an array containing the time derivative of each component, and a function $G(t, \psi)$ which takes the same inputs and returns an array containing the interaction picture oscillation frequency ω_n for each component at that time.

For example, for the case of the Gross–Pitaevskii equation [83] in the spatial basis $\psi(r, t) = \langle r|\psi(t)\rangle$, these would be:

$$F(t, \psi(r, t)) = -\frac{i}{\hbar} \left[\underbrace{-\frac{\hbar^2}{2m} \nabla^2}_{\hat{H}_{\text{nondiag}}} + \underbrace{V(r, t) + g|\psi(r, t)|^2}_{\hat{H}_{\text{diag}}} \right] \psi(r, t), \quad (3.153)$$

and

$$G(t, \psi(r, t)) = \frac{1}{\hbar} \underbrace{[V(r, t) + g|\psi(r, t)|^2]}_{\hat{H}_{\text{diag}}}. \quad (3.154)$$

Note that each symbol in bold in the algorithm below denotes an array containing one element for each basis vector $|n\rangle$, subscripts denote the different stages of RK4, and all arithmetic operations between arrays are elementwise³². The only opportunity for non-elementwise operations to occur is within F , which contains the details (via \hat{H}_{nondiag}) of any couplings between basis states for whatever system of equations is being solved, for example, using FFTs or finite differences to evaluate the Laplacian in (3.153).

³²For example, the expression $\mathbf{a} \leftarrow e^{-i\omega\Delta t} \mathbf{b}$ indicates that for all n , $a_n \leftarrow e^{-i\omega_n \Delta t} b_n$, where a_n denotes the n^{th} element of \mathbf{a} etc.

Algorithm 1 RK4ILIP

```

1: function RK4ILIP( $t_0, \Delta t, \psi_0, F$ )
2:    $f_1 \leftarrow F(t_0, \psi_0)$                                  $\triangleright$  First evaluation of Schrödinger picture DE
3:    $\omega \leftarrow G(t_0, \psi_0)$                                  $\triangleright$  Oscillation frequencies:  $\hbar\omega_n = \langle n|\hat{H}_{\text{diag}}(t_0)|n\rangle$ 
4:    $k_1 \leftarrow f_1 + i\omega\psi_0$                              $\triangleright$  Evaluate (3.150) with  $t - t' = 0$ 
5:    $\phi_1 \leftarrow \psi_0 + k_1 \frac{\Delta t}{2}$                    $\triangleright$  First RK4 estimate of IP state vector, at  $t = t_0 + \frac{\Delta t}{2}$ 
6:    $\psi_1 \leftarrow e^{-i\omega \frac{\Delta t}{2}} \phi_1$                  $\triangleright$  Convert first estimate back to SP with (3.151)
7:    $f_2 \leftarrow F(t_0 + \frac{\Delta t}{2}, \psi_1)$                $\triangleright$  Second evaluation of Schrödinger picture DE
8:    $k_2 \leftarrow e^{i\omega \frac{\Delta t}{2}} f_2 + i\omega\phi_1$      $\triangleright$  Evaluate (3.150) with  $t - t' = \frac{\Delta t}{2}$ 
9:    $\phi_2 \leftarrow \psi_0 + k_2 \frac{\Delta t}{2}$                    $\triangleright$  Second RK4 estimate of IP state vector, at  $t = t_0 + \frac{\Delta t}{2}$ 
10:   $\psi_2 \leftarrow e^{-i\omega \frac{\Delta t}{2}} \phi_2$                  $\triangleright$  Convert second estimate back to SP with (3.151)
11:   $f_3 \leftarrow F(t_0 + \frac{\Delta t}{2}, \psi_2)$                $\triangleright$  Third evaluation of Schrödinger picture DE
12:   $k_3 \leftarrow e^{i\omega \frac{\Delta t}{2}} f_3 + i\omega\phi_2$      $\triangleright$  Evaluate (3.150) with  $t - t' = \frac{\Delta t}{2}$ 
13:   $\phi_3 \leftarrow \psi_0 + k_3 \Delta t$                            $\triangleright$  Third RK4 estimate of IP state vector, at  $t = t_0 + \Delta t$ 
14:   $\psi_3 \leftarrow e^{-i\omega \Delta t} \phi_3$                      $\triangleright$  Convert third estimate back to SP with (3.151)
15:   $f_4 \leftarrow F(t_0 + \Delta t, \psi_3)$                    $\triangleright$  Fourth evaluation of Schrödinger picture DE
16:   $k_4 \leftarrow e^{i\omega \Delta t} f_4 + i\omega\phi_3$          $\triangleright$  Evaluate (3.150) with  $t - t' = \Delta t$ 
17:   $\phi_4 \leftarrow \psi_0 + \frac{\Delta t}{6} (k_1 + 2k_2 + 2k_3 + k_4)$   $\triangleright$  Fourth RK4 estimate, at  $t = t_0 + \Delta t$ 
18:   $\psi_4 \leftarrow e^{-i\omega \Delta t} \phi_4$                      $\triangleright$  Convert fourth estimate back to SP with (3.151)
19:  return  $\psi_4$                                           $\triangleright$  Return the computed SP state vector at  $t = t_0 + \Delta t$ 
20: end function

```

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Note on imaginary time evolution

When `RK4ILIP` is used for imaginary time evolution (`ITE`) [84], the oscillation frequencies ω may have a large imaginary part. If the initial guess is different enough from the ground state, then the exponentials in (3.149), (3.150) and (3.151) may result in numerical overflow. To prevent this, one can define a clipped copy of ω ,

$$\omega_{\text{clipped}} = \text{Re}(\omega) + i \begin{cases} -\frac{\log X}{\Delta t} & \text{Im}(\omega)\Delta t < -\log X \\ \text{Im}(\omega) & -\log X \leq \text{Im}(\omega)\Delta t \leq \log X \\ \frac{\log X}{\Delta t} & \text{Im}(\omega)\Delta t > \log X \end{cases}, \quad (3.155)$$

where X is very large but less than the largest representable floating-point number, and use ω_{clipped} in the exponents instead. In the below results I used `RK4ILIP` with `ITE` to smooth initial states of a Bose–Einstein condensate after a phase printing, and performed clipping with ³³ $\log X = 400$.

This clipped version of ω should be used in all exponents in the above algorithm, but only in exponents—not in the second term of (3.150). If it is used everywhere then all we have done is chosen a different (less useful) interaction picture, and the algorithm will still overflow. By clipping only the exponents, we produce temporarily “incorrect” evolution³⁴, limiting the change in magnitude of each component of the state vector to a factor of X per step (remembering that X is very large). This continues for the few steps that it takes `ITE` to get all components of the state vector to within a factor of X of the ground state, after which no clipping is necessary and convergence to the ground state proceeds as normal, subject to the ordinary limitations on which timesteps may be used with `ITE`.

3.6.2 Domain of improvement over other methods

For simulations in the spatial basis, `RK4ILIP` treats the spatially local part of the Hamiltonian analytically to first order, and hence can handle larger potentials than ordinary `RK4`. However, since a global energy offset can be applied to any potential with no physically meaningful change in the results, ordinary `RK4` can also handle large potentials—if they are large due a a large constant term which can simply be subtracted off.

So `RK4ILIP` is only of benefit in the case of large *spatial variations* in the potential. Only one constant can be subtracted off potentials without changing the physics—subtracting a spatially varying potential would require modification of the differential equation in the manner of a gauge transformation in order to leave the system physically unchanged³⁵.

However that’s not quite all: large spatial variation in potentials often comes with the prospect of the potential energy turning into kinetic energy, in which case `RK4ILIP` is also of little benefit, since in order to resolve the dynamical phase due to the large kinetic term, it would require timesteps just as small as those which ordinary `RK4` would need to resolve the dynamical phase evolution from the large potential term.

This leaves `RK4ILIP` with an advantage only in the case of large spatial variations in the potential that do not lead to equally large kinetic energies. Hence the examples I show in the next section are ones in which the condensate is trapped in a steep potential well—the trap walls are high and hence involve large potentials compared to the interior, but do not lead to large kinetic energies because the condensate is trapped close to its ground state.

The Fourier split-step (`FSS`) method [85] (see Section 3.4.1) also models dynamical phases due to the potential analytically to low order. As such it is also quite capable of modeling large potentials. However, it requires that all operators be diagonal in either the spatial basis or the Fourier basis [85]. Therefore BECs in rotating frames, due to the

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Method	RK4	RK4IP	RK4ILIP	FSS
Error	$\mathcal{O}(\Delta t^4)$	$\mathcal{O}(\Delta t^4)$	$\mathcal{O}(\Delta t^4)$	$\mathcal{O}(\Delta t^2)$
FFTs per step	4	4	4	2
Large ΔV	No	No	Yes	Yes
Large kinetic term	No	Yes	No	Yes
Arbitrary operators	Yes	Yes [†]	Yes	No
Locally parallelisable	Yes	No	Yes	No
Arbitrary boundary conditions	Yes	No	Yes	No

Table 3.2: Advantages and disadvantages of four timestepping methods for simulating Bose–Einstein condensates. *Large ΔV* refers to whether the method can simulate potentials that vary throughout space by an amount larger than the energy scale $2\pi\hbar/\Delta t$ associated with the simulation timestep Δt . *Arbitrary operators* refers to whether the method permits operators that are not diagonal in either the spatial or Fourier basis, such as angular momentum operators. *Locally parallelisable* means the method can be formulated so as to use only spatially nearby points in evaluating operators, and thus is amenable to parallelisation by splitting the simulation over multiple cores in the spatial basis. [†] Whilst one can include arbitrary operators within the RK4IP method, only operators diagonal in Fourier space can be analytically treated the way RK4IP treats the kinetic term, and so there is no advantage for these terms over ordinary RK4.

Hamiltonian containing an angular momentum operator, are not amenable to simulation with FSS³⁶.

This use of FFTs in both the FSS and RK4IP methods necessarily imposes periodic boundary conditions on a simulation, which may not be desirable. By contrast, if different boundary conditions are desired, finite differences instead of FFTs can be used to evaluate spatial derivatives in the RK4 and RK4ILIP methods, so long as a sufficiently high-order finite difference scheme is used so as not to unacceptably impact accuracy.

Along with the ability to impose arbitrary boundary conditions, finite differences require only local data, that is, only points spatially close to the point being considered need be known in order to evaluate derivatives there. This makes finite differences amenable to simulation on cluster computers [86, p. 100], with only a small number of points (depending on the order of the scheme) needing to be exchanged at node-boundaries each step. By contrast, FFT based derivatives require data from the entire spatial region. Whilst this can still be parallelised on a GPU, where all the data is available, it cannot be done on a cluster without large amounts of data transfer between nodes [75]. Thus, RK4 and RK4ILIP, being implementable with finite difference schemes, are considerably friendlier to cluster computing.

Table 3.2 summarises the capabilities of the four methods considered in the following results section. RK4ILIP is the only method capable of modelling a large spatial variation in the potential term whilst being locally parallelisable, and supporting arbitrary operators and boundary conditions.

³⁶ Split-step with more than these two bases is however possible in other schemes such as the finite element discrete variable representation [62]—each operator can be diagonalised and exponentiated locally in each element and applied as a (relatively small) matrix multiplication rather than using FFTs.

3.6.3 Results

Here I compare four numerical methods: Fourier split-step (FSS), fourth order Runge–Kutta in the interaction picture (RK4IP), ordinary fourth order Runge–Kutta (RK4), and my new method — fourth order Runge–Kutta in an instantaneous local interaction picture (RK4ILIP).

The example chosen is a 2D simulation of a turbulent Bose–Einstein condensate, in both a rotating and nonrotating frame. For the nonrotating frame the differential equation simulated was equation (3.153), and for the rotating frame the same equation

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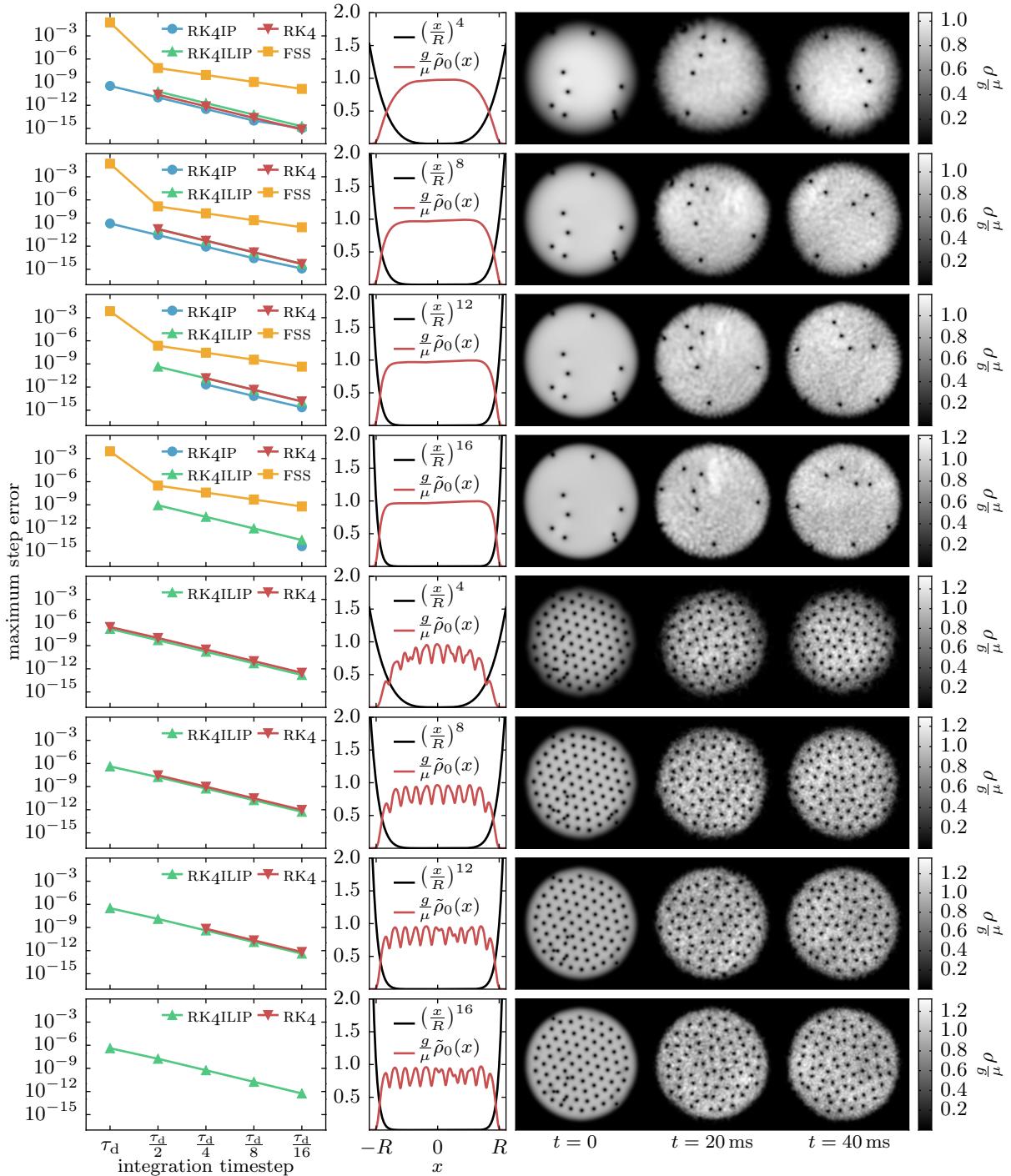


Figure 3.13: Results of simulations to compare RL4ILIP to other timestepping methods. Top four rows: Nonrotating frame simulations with four different radial power-law potentials. Bottom four rows: Rotating frame simulations with same four potentials. Left column: maximum per-step error $\int |\psi - \tilde{\psi}|^2 dr / \int |\tilde{\psi}|^2 dr$ of fourth order Runge-Kutta (RK4), its interaction picture variants (RK4IP and RK4ILIP) and Fourier split-step (FSS) as a function of timestep. Solutions were checked every 100 timesteps against a comparison solution $\tilde{\psi}$ computed using half sized steps for RK4 methods, and quarter sized steps for FSS. Simulations encountering numerical overflow not plotted. Centre column: potential (black) and average density $\tilde{\rho}_0$ of the initial state (red) over a slice of width $R/5$ in the y direction. Right column: Density of solution at initial, intermediate and final times for each configuration simulated (taken from RK4ILIP results). RK4ILIP is the only method usable in rotating frames and not encountering overflow in the steeper traps for the timesteps considered.

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was with an additional two terms added to the Hamiltonian:

$$\hat{H}_{\text{rot}} + \hat{H}_{\text{comp}} = -\boldsymbol{\Omega} \cdot \hat{\mathbf{L}} + \frac{1}{2}\hbar m^2 \Omega^2 r^2 \quad (3.156)$$

$$= i\hbar\Omega \left(x \frac{\partial}{\partial y} - y \frac{\partial}{\partial x} \right) + \frac{1}{2}\hbar m^2 \Omega^2 r^2. \quad (3.157)$$

The addition of the first term transforms the original Hamiltonian into a frame rotating at angular frequency Ω in the (x, y) plane, and is equivalent to the Coriolis and centrifugal forces that appear in rotating frames in classical mechanics [87]. The second term is a harmonic potential that exactly compensates for the centrifugal part of this force. In this way the only potential in the rotating frame is the applied trapping potential, and the only effect of the rotating frame is to add the Coriolis force.

Four trapping potentials were used, all radial power laws with different powers. These examples were chosen to demonstrate the specific situation in which RK4ILIP provides a benefit over the other methods for spatial Schrödinger-like equations, as discussed above.

The results of 120 simulation runs are shown in Figure 3.13. Each simulation was of a ^{87}Rb condensate in the $|F = 2, m_F = 2\rangle$ state, in which the two-body s -wave scattering length is $a = 98.98$ Bohr radii [88]. The simulation region was $20 \mu\text{m}$ in the x and y directions, and the Thomas–Fermi radius of the condensate was $R = 9 \mu\text{m}$. The chemical potential was $\mu = 2\pi\hbar \times 1.91 \text{ kHz}$, which is equivalent to a maximum Thomas–Fermi density $\rho_{\text{max}} = 2.5 \times 10^{14} \text{ cm}^{-3}$ and a healing length $\xi = 1.1 \mu\text{m}$. There were 256 simulation grid points in each spatial dimension, which is 14 points per healing length.

Four different potentials were used, all of the form $V(r) = \mu(r/R)^\alpha$ with $\alpha = 4, 8, 12, 16$. For the rotating frame simulations, the rotation frequency was $\Omega = 2\pi \times 148 \text{ Hz}$. This is 89% of the effective harmonic trap frequency, defined as the frequency of a harmonic trap that would have the same Thomas–Fermi radius given the same chemical potential.

All ground states were determined using successive over-relaxation (See section 3.5.2) with sixth-order finite differences for spatial derivatives. For the nonrotating simulations, convergence was reached with $\Delta\mu/\mu < 1 \times 10^{-13}$, with:

$$\Delta\mu = \sqrt{\frac{\langle \psi | (\hat{H} - \mu)^2 | \psi \rangle}{\langle \psi | \psi \rangle}}, \quad (3.158)$$

where \hat{H} is the nonlinear Hamiltonian and $\langle \psi | \psi \rangle$ is the condensate wavefunction, which does not have unit norm. For the rotating frame simulations the ground states converged to $\Delta\mu/\mu \approx 9 \times 10^{-7}, 2 \times 10^{-6}, 3 \times 10^{-6}$ and 2×10^{-6} for $\alpha = 16, 12, 8$, and 4 respectively.

After each ground state was found, it was multiplied by a spatially varying phase factor corresponding to the phase pattern of a number of randomly positioned vortices:

$$\psi_{\text{vortices}}(x, y) = \psi_{\text{groundstate}}(x, y) \prod_{n=1}^N e^{\pm_n i \arctan 2(y - y_n, x - x_n)} \quad (3.159)$$

where $\arctan 2$ is the two-argument arctan function,³⁷ $N = 30$, \pm_n is a randomly chosen sign, and (x_n, y_n) are vortex positions randomly drawn from a Gaussian distribution centred on $(0, 0)$ with standard deviation equal to the Thomas–Fermi radius R . The same seed was used for the pseudorandom number generator in each simulation run, and so the vortex positions were identical in each simulation run.

After vortex phase imprinting, the wavefunctions were evolved in imaginary time [84]. For the nonrotating frame simulations, imaginary time evolution was performed for a time interval equal to the chemical potential timescale $\tau_\mu = 2\pi\hbar/\mu$, and for the rotating

³⁷Defined as the principle value of the argument of the complex number $x + iy$: $\arctan 2(y, x) = \text{Arg}(x + iy)$.

frame simulations, for $\tau_\mu/10$. This was done to smooth out the condensate density in the vicinity of vortices, producing the correct density profile for vortex cores. However, since imaginary time evolution decreases the energy of the state indiscriminately, it also had the side effect of causing vortices of opposite sign to move closer together and annihilate. This decreased the number of vortices, and is the reason the smoothing step in the rotating frame simulations was cut short to $\tau_\mu/10$, as otherwise all vortices had time to annihilate with one of the lattice vortices. A vortex pair in the process of annihilating is visible in Figure 3.13 as a partially filled hole in the initial density profile near the top of the condensate in the $\alpha = 4, 12$, and 16 rotating frame simulations.³⁸

³⁸The initial states for the four different potentials are not identical, so by chance the corresponding vortex in the $\alpha = 8$ case was not close enough to a lattice vortex to annihilate.

The smoothed, vortex imprinted states were then evolved in time for 40 ms. For each simulation, five different timesteps were used: $\Delta t = \tau_d, \tau_d/2, \tau_d/4, \tau_d/8, \tau_d/16$, where $\tau_d = m\Delta x^2/\pi\hbar \approx 2.68 \mu\text{s}$ is the dispersion timescale associated with the grid spacing Δx , defined as the time taken to move one gridpoint at the group velocity of the Nyquist mode.

For the nonrotating frame simulations, spatial derivatives for the **RK4** and **RK4ILIP** methods were determined using the Fourier method (Section 3.4.1). This was to ensure a fair comparison with the other two methods, which necessarily use Fourier transforms to perform computations pertaining due to the kinetic term.

For the rotating frame simulations, sixth-order finite differences with zero boundary conditions were used instead for the kinetic terms of the **RK4** and **RK4ILIP** methods, which were the only two methods used for those simulations (due to the other methods being incompatible with the angular momentum operator required for a rotating frame). This choice was fairly arbitrary, but did allow the condensate to be closer to the boundary than is otherwise possible with the periodic boundary conditions imposed by use of the Fourier method for spatial derivatives. This is because the rotating frame Hamiltonian is not periodic in space, and so its discontinuity at the boundary can be a problem if the wavefunction is not sufficiently small there.

As shown in Figure 3.13, all methods tested generally worked well until they didn't work at all, with the per-step error of **RK4**-based methods being either small and broadly the same as the other **RK4**-based methods, or growing rapidly to the point of numerical overflow (shown as missing datapoints). The break down of **FSS** was less dramatic, though it too had a clear jump in its per-step error for larger timesteps. Comparing methods therefore came down to mostly whether or not a simulation experienced numerical overflow during the time interval being simulated.

The main result was that **RK4ILIP** and **FSS** remained accurate over the widest range of timesteps and trap steepnesses, with **RK4** and **RK4IP** requiring ever smaller timesteps in order to not overflow as the trap steepness increased.

For the rotating frame simulations, which were only amenable to the **RK4** and **RK4ILIP** methods, the same pattern was observed, with **RK4** only working at smaller timesteps as the trap steepness was increased, and ultimately diverging for all timesteps tested at the maximum trap steepness. By contrast, **RK4ILIP** remained accurate over the entire range of timesteps at the maximum trap steepness.

3.6.4 Discussion

As mentioned, **RK4ILIP** is mostly useful for continuum quantum mechanics only when there are large spatial differences in the potential, which cannot give rise to equally large kinetic energies³⁹. Furthermore, the advantage that **RK4ILIP** has over other methods with that same property is that it does not require a particular form of Hamiltonian or a particular method of evaluating spatial derivatives. The former means it is applicable in rotating frames or to situations with unusual Hamiltonians, and the latter means it can be used with finite differences or **FEDVR** [62] and thus is amenable to parallelisation on a cluster computer.

³⁹This is essentially due to such a situation violating the condition we laid out at the beginning of this section — that the simulation basis must be nearly an eigenbasis of the total Hamiltonian.

The ability to model large spatial variations in the potential provides only a narrow domain of increased usefulness over other methods. If a large kinetic energy results from the large potential, then the method requires just as small timesteps as any other. And if the large potential is supposed to approximate an infinite well, then an actual infinite well may be modelled using zero boundary conditions, negating the need for something like `RK4ILIP`. However, when potential wells are steep, but not infinitely steep, here `RK4ILIP` provides a benefit. The only other model that can handle these large potentials—Fourier split-step—has the disadvantage that it cannot deal with arbitrary operators such as those arising from a rotating frame, and is not parallelisable with local data. The benefits of parallelisability are obvious, and the above results demonstrate `RK4ILIP`'s advantage at simulating BECs in tight traps and rotating frames.

Note that whilst the *Fourier* split-step method can't handle Hamiltonian terms such as $\hat{\mathbf{r}} \cdot \hat{\mathbf{p}}$ that are not diagonal in either real space or Fourier space [55, p. 315], a split-step method based on an approximation to the momentum operator as a banded matrix, such as that obtained with finite differences, can. Using the techniques discussed in section 3.2.4, such a scheme is also parallelisable. The remaining limitations then, when compared to fourth-order Runge–Kutta are the restriction on the types of nonlinearity that can be included, and the complexity of implementation.

For systems with discrete degrees of freedom, `RK4ILIP` may be useful in the case where an approximate diagonalisation of the Hamiltonian is analytically known, and when the Hamiltonian's eigenvalues vary considerably in time (making a single interaction picture insufficient to factor out dynamical phases throughout the entire simulation). In this situation an analytic transformation into the diagonal basis can be performed at each timestep (or the differential equation analytically re-cast in that basis in the first place), and `RK4ILIP` can be used to factor out the time-varying dynamical phase evolution at each timestep. An example may be an atom with a magnetic moment in a time-varying magnetic field which varies over orders of magnitude. The transformation into the spin basis in the direction of the magnetic field can be analytically performed, and if the field varies by orders of magnitude, so do the eigenvalues of the Hamiltonian. Although the eigenvalues in this case and other similar cases can be computed analytically too, unless all time dependence of the Hamiltonian is known in advance of the simulation, it would be difficult to incorporate this into a re-casting of the differential equation in a time-dependent interaction picture. `RK4ILIP` may be useful in these cases to automate this process and evolve the system in the appropriate interaction picture at each timestep.

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Hidden variables for semiclassical models with state-dependent forces

HIDDEN-VARIABLE THEORIES [89,90] are interpretations of quantum mechanics that posit definite states underlying quantum state vectors, such that quantum indeterminacy is an illusion—an emergent phenomenon rather than a fundamental fact. Bell's theorem [91] proves that any such theory must be *nonlocal* in order to explain all the predictions of quantum mechanics, and perhaps in light of this, most physicists surveyed [92] do not believe that hidden variables underlie physical reality.

However, by framing quantum systems in classical terms, hidden-variable theories can provide an excellent computational tool for *approximate* models of quantum systems, when it is reasonable to approximate some degrees of freedom as classical, yet other degrees of freedom need to be modelled quantum mechanically. Just as hidden-variable theories have framed the quantum world in terms that are agreeable to the classical view of the world in the minds of some interpreters of quantum mechanics, so can they bridge the gap between a *simulated* quantum world and a *simulated* classical world coexisting in the same computer simulation.

In this chapter I describe what I call the ‘hidden-variable semiclassical’ (HVSC) method: a method of combining quantum simulations with classical simulations, with hidden variables bridging the gap between the classical and quantum degrees of freedom. In this introduction I describe existing semiclassical methods, the manner in which their quantum and classical parts are typically coupled based on expectation values (which I’m calling the ‘Ehrenfest method’), and in which regimes this can be inaccurate—namely the Stern–Gerlach experiment and similar situations in which considerable entanglement between motional and internal degrees of freedom of atoms can develop.

In Section 6.1 I give an overview of what a semiclassical method is, their most common implementation and in what situations this is insufficiently accurate, motivating the need for an improved method. In Section 6.2 I give the technical definition of a hidden-variable theory, and motivate the use of such a theory for coupling quantum and classical degrees of freedom in such a way that the a semiclassical model can be made to agree more closely than the Ehrenfest method with the underlying fully quantum model it is approximating, and ultimately, with experiment. In Section 6.3 I then go into the implications of welcoming a hidden variable into a semiclassical model, including additional required assumptions and approximations, and in Sections 6.4 and 6.5 I go into detail deriving the equations of motion for the model and presenting some algorithmic and computational details. Finally in Section 6.6, having provided all the background arguments and details, I present the complete algorithm(s), before showing simulation results in Section 6.7 that compare the model to the underlying exact Schrödinger wave

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equation, and concluding in Section 6.8 with further discussion of the method’s benefits and limitations.

During writing this thesis, I discovered that the core idea underpinning this method is (unfortunately) not original, and that the ‘surface hopping’ method enjoys widespread use and continued development in the field of computational chemical physics [93–103]. An earlier version of my model [104] bears a striking resemblance to that presented in a 1991 paper by Tully [93], the pioneer of surface hopping methods. Use of these ideas do not however appear to be common in cold atom physics, but with direct simulation of evaporative cooling and laser cooling schemes becoming increasingly plausible given the increasing availability of computational power as well as efficient molecular dynamics techniques such as Direct Simulation Monte Carlo (DSMC) [105], this sort of convergent evolution of numerical techniques is perhaps not surprising. I did not encounter the surface-hopping literature earlier as I have mostly been thinking of my method as a quantum-trajectories method, and making comparisons to the Monte Carlo wavefunction method [106] and quantum trajectories more generally [107, 108]. I believe surface hopping is another way of doing quantum trajectories, however the two do not seem to overlap much in the literature, and I managed to not notice the former until recently. Throughout this chapter I make comparisons between my own methods and those in the existing surface-hopping literature.

Three long-standing limitations of my method may be resolved in light of what I have now read in the surface-hopping literature. One is that my model was previously limited to one dimension only, as I did not know in what direction atoms should gain or lose momentum upon making a transition. This is a known result in the surface-hopping literature, and discussed in Section 6.4.3, resolving this issue completely to my knowledge. The second limitation is that my model is limited to time-independent Hamiltonians. This is because time-dependent potentials can exchange energy with an atom, whereas part of my model relied on an energy conservation argument to compute the velocity of an atom after a transition. As far as I can tell, the situation is the same in the surface-hopping literature, including an identical argument for velocity correction of atoms based on energy conservation. However, the hidden variable theory primarily used for surface hopping—Tully’s fewest-switches algorithm (discussed in Section 6.4.2)—can be formulated in a way that allows a distinction to be made between transitions caused by spatial variation of the Hamiltonian and those caused by its time variation. I present a way of computing these two sets of transition probabilities, and propose that the energy conservation argument only be applied to the transitions ‘caused’ by the spatial variation. However this is untested, and does not feature in the rest of the presentation of my model. Finally, the fewest-switches algorithm is computationally cheaper than the hidden variable theory I have been using, resolving my concern that the hidden variable theory I have been using—Schrödinger theory, discussed in Section 6.2—is computationally expensive.

On the other hand, I have made contributions that are—to my knowledge—new. I believe the identification of what the surface-hopping literature calls ‘hopping algorithms’ with hidden-variable theories (the latter of which is usually discussed in the context of quantum foundations, philosophy and metaphysics rather than applied computational chemistry) is new. Furthermore, the computationally expensive Schrödinger theory—which is not new, though its use in surface-hopping models is—is more capable in certain circumstances: it can compute transition probabilities for arbitrarily long time intervals, given the unitary describing the corresponding quantum evolution over the time interval. Tully’s fewest-switches, on the other hand, is valid for infinitesimal time-intervals only, and the probabilities need to be integrated over time to obtain transition probabilities over large intervals. It may be the case that in some scenarios Schrödinger theory is computationally cheaper than this integral. Finally, the more sophisticated of my two methods of computing decoherence, discussed in Sections 6.5.4 and 6.5.5, provides ex-

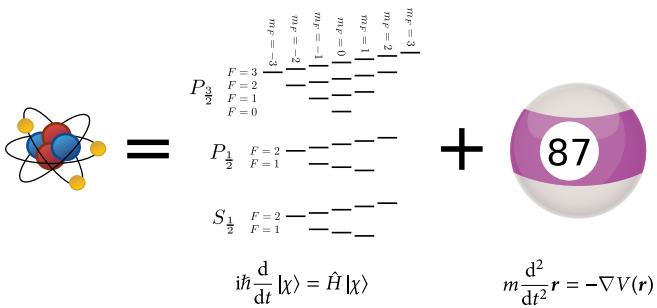


Figure 6.1: Artist's depiction of a semiclassical atom.

cellent agreement with the underlying Schrödinger wave equation it approximates, and I believe it is an improvement over similar methods in the surface hopping literature, which I will discuss at the end of this chapter (Section 6.8).

6.1 Semiclassical models

A semiclassical model is any model in which some degrees of freedom are treated quantum mechanically, and others classically. The most common combination is that of treating an atom's internal electronic state quantum mechanically and its motional degree of freedom classically. This is useful whenever the quantum effects of the atom's motion are not of interest, for example if temperatures are high and thus atomic wavelengths are short—such that quantum effects simply aren't visible in the motion of the particles and so they can accurately be modelled as classical billiard balls. The energy gaps between different electronic states of atoms are so large however that only at very high temperatures (at which atoms ionise anyway) do they start to appear as a continuum compared to thermal energy scales, and the interaction of different spin states of the atom with different optical and magnetic fields does not make them appear as classical continua either. Thus, quantum effects can be ignored for the centre of mass motion of the (relatively heavy) atom, but not for the relative motion of its (much lighter) electrons with respect to the nucleus, or for the nuclear and electronic spin degrees of freedom [93].

In this regime, atoms are often modelled semiclassically, with these internal degrees of freedom modelled using a state vector $|\chi\rangle$ evolving according to a Hamiltonian \hat{H} via the Schrödinger equation, and the centre of mass motion modelled as a position \mathbf{r} and velocity \mathbf{v} evolving according to Newton's second law (Figure 6.1).

Once one has defined a potential function $V(\mathbf{r})$ and a Hamiltonian (also possibly varying with space) $\hat{H}(\mathbf{r})$, and other possible additions,¹ ones job is done and the rest can be left to numerical differential equation solvers to evolve some concrete vector representation χ of $|\chi\rangle$ as well as the state variables \mathbf{r} and \mathbf{v} for motional degree of freedom in time according to the coupled differential equations

$$\frac{d}{dt} |\chi\rangle = -\frac{i}{\hbar} \hat{H}(\mathbf{r}) |\chi\rangle, \quad (6.1)$$

$$\frac{d}{dt} \mathbf{v} = -\frac{1}{m} \nabla V(\mathbf{r}), \quad (6.2)$$

$$\frac{d}{dt} \mathbf{r} = \mathbf{v}. \quad (6.3)$$

¹Such as using a Monte-Carlo wavefunction method [106, 109] to model the effect of spontaneous emission on $|\chi\rangle$, and modifying \mathbf{v} instantaneously by a random-direction recoil velocity upon each photon emission.

In the next subsection I explain why it's not always that simple.

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6.1.1 Stern–Gerlach separation and evaporative cooling

In the Stern–Gerlach experiment [18], particles with quantum spin and a magnetic moment—atoms for example—are fired as a beam through a region of space with a magnetic field gradient. The well known result is that two clusters (if the atoms are spin- $\frac{1}{2}$) of positions are observed once the beam emerges, rather than a continuous smear of positions, indicating that angular momentum—like many quantities in quantum mechanics—is quantised.

This is the case even if one spin-polarises the particles before they are passed through the magnetic field gradient, say putting them in an eigenstate of the \hat{F}_x operator. Then, if the magnetic field is along the z direction, and the gradient is also in the z direction, two clusters of positions are also observed, even though all particles were in the same state when they entered the region in which there was a magnetic field gradient. This is a display of the indeterminacy of quantum mechanics: even though all particles had the same initial state, there were nonetheless different outcomes for each particle.

The outcome of the Stern–Gerlach experiment is a consequence of quantum mechanics, to be sure, but it has little to do with the wave nature of the atoms themselves. If we introduced some double slits for the atoms to pass through in addition to the magnetic field gradient, then we would be seeing the wave nature of the atoms as interference patterns at the detection screen at the end of the experiment. But if we do not, and if the particles have short de-Broglie wavelengths, then quantum mechanics is not apparent in the motion of the particles through space—except via the influence of spin on seemingly choosing one trajectory or the other. The effect is well understood quantum mechanically, but is difficult to model semiclassically because even if we are happy to approximate wavepackets as small, the wavepackets do not take a single trajectory. Rather they split into two wavepackets, with the part of the superposition corresponding to one spin projection state (along the direction of the local magnetic field) moving one way, and the part of the superposition with the other spin projection going the other way. The trajectories can still be quite classical, it's just that there are two of them.

A similar situation exists in RF evaporative cooling (Section 2.1.5) of cold atoms enroute to BEC. Atoms are confined in a magnetic trap, and are spin-polarised so as to be fully spin-down (for ^{87}Rb this is the trapped state) with respect to the local magnetic field at the position of each atom. The magnetic field's direction—not just its magnitude—varies in space, and so the spin vectors of different atoms point in different directions in three-dimensional space, but they are all spin-down with respect to the quantisation axis of their local magnetic field. As the atoms move through space, they move in orbits—punctuated by collisions—about the magnetic field zero at the centre of the trap, since they feel a force $F \propto -\nabla|\mathbf{B}|$ due to the gradient of the Zeeman potential. Provided they are moving slowly (specifically, provided their Larmor precession period is short compared to the time the magnetic field as seen by the atom takes to change by a non-negligible fraction of its current value), the atoms' spins adiabatically follow the local field and remain spin-down, even as the field as seen by each atom fully reverses its direction every half orbital period.

Near the centre of the trap where the atoms are moving faster, the fields are small and therefore have large fractional derivatives and lead to large Larmor periods, adiabaticity no longer holds and the atoms may make spin transitions with respect to their local magnetic field. Once an atom passing close to the field zero has evolved into a superposition of spin-projection states with respect to the local field, it is in a situation identical to the initial condition of the Stern–Gerlach experiment, causing the spin-projection components to spatially separate in the magnetic field gradient. The spin-up component is anti-trapped and repelled from the centre of the trap, and the zero spin-projection component (since the groundstate of ^{87}Rb is spin-1) feels no force and moves in a straight line. The spin-down component continues on an orbit about the field zero which is just as tight as

before, unaffected by the close approach to the field zero other than being reduced in amplitude. Eventually a collision occurs, either with other atoms or with the walls of the vacuum system and the wavefunction collapses to choose one of these options, leading to atoms probabilistically leaving the trap (called Majorana losses [110, 111]) or remaining trapped. Again, the trajectories can still be quite classical, it's just that there are three of them, and which is taken is probabilistic.

How can we model these effects semiclassically? Equations (6.1) to (6.3) are not sufficient, because exists no single classical potential $V(\mathbf{r})$ that can describe the motion of the atoms. Rather, the atoms feel a different force depending on which spin state they are in. Just as the Hamiltonian can be a function of space, so can the potential be a function of the internal state of the atom: $V = V(\mathbf{r}, |\chi\rangle)$. Ehrenfest's theorem [112] states that

$$m \frac{d^2}{dt^2} \langle \hat{\mathbf{r}} \rangle = -\langle \nabla \hat{V} \rangle, \quad (6.4)$$

where the expectation values are over all degrees of freedom, not just motional. If we approximate a small wavepacket centred at the position \mathbf{r} in order to ignore the wave nature of the atoms, this becomes:

$$m \frac{d^2}{dt^2} \mathbf{r} = -\nabla \langle \chi | \hat{V}(\mathbf{r}) | \chi \rangle, \quad (6.5)$$

where the operator $\hat{V}(\mathbf{r})$ now only acts on the subspace of the internal state of the atom, since we have already taken an expectation value over (a small region of) space. Provided all potentials the atom is subjected to are included in the Hamiltonian for its internal state (including any energy offsets that do not depend explicitly on the internal state), this is nothing but

$$m \frac{d^2}{dt^2} \mathbf{r} = -\nabla \langle \chi | \hat{H}(\mathbf{r}) | \chi \rangle, \quad (6.6)$$

where \hat{H} is the Hamiltonian describing the evolution of the atom's internal state. We now can construct the *Ehrenfest semiclassical method* describing how the *expectation value* of a well localised atom's position evolves with time:

$$\frac{d}{dt} |\chi\rangle = -\frac{i}{\hbar} \hat{H}(\mathbf{r}) |\chi\rangle, \quad (6.7)$$

$$\frac{d}{dt} \mathbf{v} = -\frac{1}{m} \langle \chi | \hat{H}(\mathbf{r}) | \chi \rangle, \quad (6.8)$$

$$\frac{d}{dt} \mathbf{r} = \mathbf{v}. \quad (6.9)$$

The Ehrenfest semiclassical method is the same as the simple semiclassical method (6.1) to (6.3), except that it has an answer to the question "What should we use for $V(\mathbf{r})$ when the atom is in a superposition of states that feel different potentials?", which is "use the expectation value".

This is all well and good if the expectation value of position is a good approximation to the situation being modelled. But in the Stern–Gerlach experiment or a Majorana spin flip in a magnetic trap, the expectation value of position is a poor match to reality. In the Stern–Gerlach experiment beginning with spin-polarised atoms, a trajectory subject to the mean force or potential (which are both zero for a 50 : 50 superposition) would land in a single blob in the middle of the screen, rather than two blobs displaced from the centre. In the case of an atom approaching the field zero in a magnetic trap, use of the mean force would result in the atom broadening its orbit somewhat, rather than splitting into

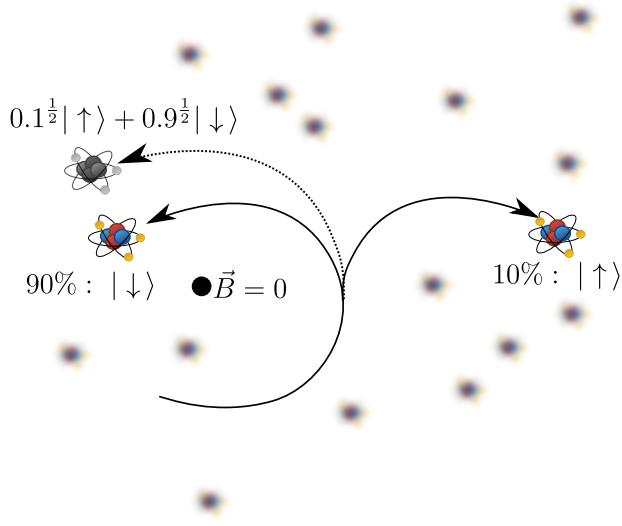


Figure 6.2: When atoms pass near to the field zero in a magnetic trap, their wavepackets diverge in space as multiple trajectories, each corresponding to one local spin projection state. For example, if a spin- $\frac{1}{2}$ atom undergoes a partial spin-flip such that it has a ten percent probability of being spin-up, the end result will be two approximate trajectories. One trajectory describes the motion of a wavepacket that is fully spin-up, and one fully spin-down, with the spin-down wavepacket's squared amplitude equal to 0.1. The Ehrenfest semiclassical method however can only model a single trajectory, and the end result when using this method is a single trajectory being approximately the mean of the two actual trajectories, and retaining a 90 : 10 ratio of spin-down:spin-up state populations.

multiple possible trajectories (Figure 6.2). Semiclassical simulations of evaporative cooling performed by Christopher Watkins (unpublished) displayed an unphysical heating of the atom cloud that I believe is due to the Ehrenfest method's inability to model Stern–Gerlach separation. In a real magnetic trap during evaporative cooling to Bose–Einstein condensation, the mean free path is large enough that the part of the wavepackets that are no longer trapped will usually leave the trap without colliding with any other atoms. This means that the energy the untrapped and anti-trapped components have gained (relative to the trapped component) moving away from the field zero is not usually shared with other atoms upon collision—the extra energy leaves with the atoms. However, if a close approach to the magnetic field zero merely means a broadening of the atoms orbit, then the extra energy does not leave as fast, if at all, and can be shared with other atoms via collisions, turning what would have been an atom loss effect into an overall gradual heating of the cloud.

So how can we modify a semiclassical method to choose only one trajectory? Firstly, since each trajectory corresponds to one of the internal states (though some might be degenerate), our model must choose an internal state and use the classical trajectory corresponding to that internal state only. Secondly, since all atoms begin in identical states and yet some take one trajectory and some another, this choice must be probabilistic. Finally, the probabilities must be consistent with those from quantum mechanics, i.e. the Born rule: the probability of an atom taking each trajectory must be proportional to the squared amplitude of the internal state of the atom, projected onto the eigenstate corresponding to that trajectory.

There exists a category of theories dealing with precisely this question of how to choose a specific state of a quantum system in a stochastic way, such that the probability of having chosen a state is equal to that given by quantum mechanics. Such theories are called *hidden-variable* theories, and any parameter, variable or label specifying which state has been chosen is a *hidden variable*.

6.2 Hidden-variable theories

In his paper *Quantum computing and dynamical quantum models*, Aaronson defines a *dynamical quantum model* [113]:

A *dynamical quantum model* assigns an eigenstate to a specified observable even when no measurement is made, and gives a stochastic evolution rule for that eigenstate. Such a model yields a distribution of classical histories of a quantum state.

In a later paper, *Quantum computing and hidden variables*, superseding the first, Aaronson renames dynamical quantum models to *hidden-variable theories* and instead defines [90]:

For us, a hidden-variable theory is simply a way to convert a unitary matrix that maps one quantum state to another, into a stochastic matrix that maps the initial probability distribution to the final one in some fixed basis.

This is what a hidden-variable theory is for the purposes of this chapter as well. Though the first paper is superseded by the second, the definition it gives is more tangible for us—we wish to *assign an eigenstate* to our atoms (choose one of the internal states in order to decide which trajectory to follow), and have a *stochastic evolution rule* for which eigenstate is chosen at any one time (allow the atom to begin taking a different trajectory if it makes transitions between states), probabilistically depending on the change in quantum populations of the states.

But the second definition is more specific. The stochastic evolution rule is in the form of a stochastic matrix, with elements equal to transition probabilities for some time interval. And the rule should be in some way based on the unitary evolution that the quantum system evolves according to in the same interval of time, such that the initial and final probabilities of the stochastic matrix and unitary evolution agree. That is, if a quantum state $|\chi\rangle$ evolves in a certain basis $\{|\chi_i\rangle\}$ according to a unitary $\hat{U}(t', t)$:

$$\begin{bmatrix} c_1(t') \\ c_2(t') \\ c_3(t') \\ \vdots \end{bmatrix} = \begin{bmatrix} U_{11}(t', t) & U_{12}(t', t) & U_{13}(t', t) & \dots \\ U_{21}(t', t) & U_{22}(t', t) & U_{23}(t', t) & \dots \\ U_{31}(t', t) & U_{32}(t', t) & U_{33}(t', t) & \dots \\ \vdots & \vdots & \vdots & \ddots \end{bmatrix} \begin{bmatrix} c_1(t) \\ c_2(t) \\ c_3(t) \\ \vdots \end{bmatrix}, \quad (6.10)$$

where $c_i = \langle \chi_i | \chi \rangle$ and $U_{ij}(t', t) = \langle \chi_i | \hat{U}(t', t) | \chi_j \rangle$, then a hidden-variable theory is a matrix-valued function $S(U(t', t), \chi(t))$, where $\chi(t)$ and $U(t', t)$ are the vector and matrix representations of the state vector $|\chi(t)\rangle$ and unitary $\hat{U}(t', t)$ in the $\{|\chi_i\rangle\}$ basis, that satisfies²

$$\begin{bmatrix} |c_1(t')|^2 \\ |c_2(t')|^2 \\ |c_3(t')|^2 \\ \vdots \end{bmatrix} = \begin{bmatrix} S_{11}(t', t) & S_{12}(t', t) & S_{13}(t', t) & \dots \\ S_{21}(t', t) & S_{22}(t', t) & S_{23}(t', t) & \dots \\ S_{31}(t', t) & S_{32}(t', t) & S_{33}(t', t) & \dots \\ \vdots & \vdots & \vdots & \ddots \end{bmatrix} \begin{bmatrix} |c_1(t)|^2 \\ |c_2(t)|^2 \\ |c_3(t)|^2 \\ \vdots \end{bmatrix}. \quad (6.11)$$

In essence, if quantum unitary evolution takes amplitudes to amplitudes, a hidden-variable theory takes probabilities to probabilities. Thus the elements of S are conditional

²In order to more clearly compare to quantum evolution, we are using *left stochastic* matrices, which multiply column vectors of probabilities from the left, contrary to the most common convention for stochastic matrices which is to multiply row matrices from the right. Thus S has unit column sums, whereas the corresponding right stochastic matrix (its transpose) would have unit row sums.

probabilities—or transition probabilities—for a hidden variable $\eta(t)$, giving the chance that the state $|\chi_\eta\rangle \in \{|\chi_i\rangle\}$ assigned by the hidden variable will change from one to another in the given time interval:

$$\Pr(\eta(t')=i|\eta(t)=j) = S_{ij}(U(t', t), \chi(t)). \quad (6.12)$$

There are many ways to define functions S that satisfy this condition. The simplest is to ignore the unitary completely and set $S_{ij} = |c_i(t')|^2$, which yields:

$$\Pr(\eta(t')=i|\eta(t)=j) = |c_i(t')|^2, \quad (6.13)$$

that is that the hidden variable η is equally likely to transition to a given value regardless of its previous value, and regardless of the unitary, and that the conditional transition probability from all input states is just the squared amplitude of the final state. This theory represents a hidden variable that jumps between states randomly based on their population, with no regard for its history or whether there were actually amplitude flows between the states between which it is transitioning. Nonetheless, this theory, called *product theory* [90], matches the definition of a hidden-variable theory— S is a stochastic matrix, and the hidden variable on average will spend an amount of time in each state consistent with the Born rule.

Aaronson outlines some additional axioms [90] that hidden variables ought to satisfy, if we are to take them seriously, including reasonable statements about symmetries, and insensitivity to small perturbations. He goes on to prove that the properties cannot all be satisfied simultaneously, but some theories satisfy more of them than others.

It is convenient to introduce the matrix P of absolute transition probabilities. Summing (6.12) over all possible initial values of the hidden variable yields the absolute probability of it having a particular final value after the given time interval:

$$\Pr(\eta(t')=i) = \sum_j \Pr(\eta(t')=i|\eta(t)=j) \Pr(\eta(t)=j), \quad (6.14)$$

which, recognising that the final and initial probabilities must be the final and initial squared amplitudes of the state vector, can be rewritten:

$$|c_i(t')|^2 = \sum_j S_{ij}(U(t', t), \chi(t)) |c_j(t)|^2. \quad (6.15)$$

We now define the matrix of absolute transition probabilities P :

$$P_{ij} = S_{ij}(U(t', t), \chi(t)) |c_j(t)|^2, \quad (6.16)$$

such that:

$$|c_i(t')|^2 = \sum_j P_{ij}. \quad (6.17)$$

So we have that P has row sums equal to the final squared amplitudes. And, because $P_{ij} = S_{ij} |c_j(t)|^2$ and S is a stochastic matrix with column sums equal to one, we have that P must have row sums equal to the initial squared amplitudes.

Now that we have introduced P and shown what its row and column sums must be, we come to a particularly simply defined hidden-variable theory called the Schrödinger theory, discussed in Aaronson's hidden variables paper [90]. The idea is to form P by starting with the matrix of absolute values of U , and simply scaling its rows and columns to have the correct values:³

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³One might expect that the absolute value squared might be a better choice, since this matches Fermi's golden rule. Aaronson discusses in his paper, when presenting his own hidden-variable theory *Flow theory*, how absolute values of elements of the unitary have some appealing properties.

$$P = \begin{bmatrix} a_1 b_1 |U_{11}(t', t)| & a_1 b_2 |U_{12}(t', t)| & a_1 b_3 |U_{13}(t', t)| & \dots \\ a_2 b_1 |U_{21}(t', t)| & a_2 b_2 |U_{22}(t', t)| & a_2 b_3 |U_{23}(t', t)| & \dots \\ a_3 b_1 |U_{31}(t', t)| & a_3 b_2 |U_{32}(t', t)| & a_3 b_3 |U_{33}(t', t)| & \dots \\ \vdots & \vdots & \vdots & \ddots \end{bmatrix}, \quad (6.18)$$

that is,

$$P_{ij} = a_i b_j |U_{ij}(t', t)|, \quad (6.19)$$

where the row scalings $\{a_i\}$ and column scalings $\{b_j\}$ satisfy:

$$\sum_j a_i b_j |U_{ij}(t', t)| = |c_i(t')|^2, \quad (6.20)$$

$$\sum_i a_i b_j |U_{ij}(t', t)| = |c_j(t)|^2, \quad (6.21)$$

which can be solved numerically, and then the Schrödinger theory stochastic matrix then able to be extracted by inverting (6.16):

$$S_{ij}(U(t', t), \chi(t)) = a_i b_j \frac{|U_{ij}(t', t)|}{|c_j(t)|^2}. \quad (6.22)$$

Schrödinger theory is the hidden-variable theory I have used in the simulations in the results section of this chapter (Section 6.7). In Section 6.4.1 I detail my efforts to most efficiently find the row and column scalings in order to actually compute the Schrödinger theory S matrix.

Having now discovered Tully's fewest-switches algorithm in the surface hopping literature [93, 95, 100] (also detailed in Section 6.4.2), I suspect that in most circumstances it is the better choice, over Schrödinger theory, for the hidden-variable theory in a hidden-variable semiclassical/surface hopping simulation. Tully's fewest switches has three advantages over Schrödinger theory, two of which are important in my view, and one less important. The first important one, as detailed in Section 6.4.1, is that Schrödinger theory is considerably more computationally expensive to evaluate (except in the case of a two-state system). The second, which is only relevant for time-dependent Hamiltonians, is that fewest-switches can be formulated in a way that allows the distinction to be drawn between motion of a particle through an inhomogeneous field as the cause of non-adiabatic transitions, versus non-adiabatic transitions being caused by explicit time dependence of said fields. The usefulness of this distinction will hopefully be made clear in Sections 6.4.2 and 6.4.3, though since I have not been able to make the distinction in the past using only Schrödinger theory, these improvements are untested and hence speculative. The less important (but still appealing) advantage is in the name: fewest switches makes the *fewest number of switches* consistent with the quantum probabilities in each infinitesimal time interval, whereas Schrödinger theory makes more switches. More switches means larger statistical variation in the outcomes: even though the expectation values for the number of atoms in each state are consistent with the Born rule, the variance is higher, and a larger ensemble will be required to see agreement with the Born rule to within some tolerance. More switches also implies larger statistical variation in the classical trajectories, as atoms are more likely to have spent time following multiple classical trajectories in a region of non-adiabatic coupling rather than just one, perhaps jumping to a different trajectory momentarily before jumping back, rather than just remaining in the first one the whole time. As mentioned in the introduction to this chapter, Schrödinger theory has one advantage over fewest switches: it is formulated in a way that allows one to compute transition probabilities corresponding to quantum evolution over an arbitrary time interval, so long as one has the unitary that describes that

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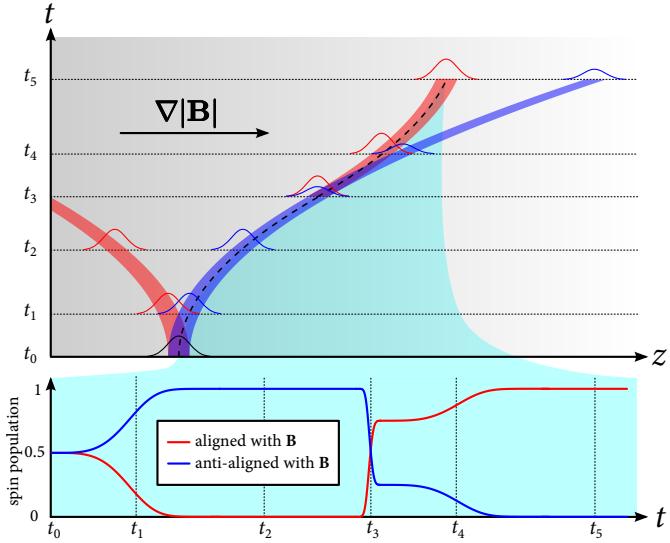


Figure 6.3: Schematic of method. At t_0 an atom is in a 50 : 50 superposition of spin-up:spin-down population, and the two spin components accelerate apart in the magnetic field gradient. If the hidden variable dictates that we follow the spin-down component (the initial trajectory given by the dotted line), then we see a reduced spin-up population at times t_1 and t_2 . At t_3 the field changes direction suddenly, and partly flips the spins (in the local basis with quantisation axis given by the direction of the magnetic field). The stochastic hidden variable transitions to instead select the spin-up component, which we follow thereafter. Following the spin-up component we then see the a reduced spin-down population at t_5 due to the wavepackets separating once more.

evolution. Fewest switches, on the other hand, produces differential probabilities that need to be integrated over time to obtain transition probabilities over time intervals that are not small compared to the dynamical timescales of the problem. This is usually not a problem since one is generally numerically integrating the state vector with timesteps that also need to be smaller than dynamical timescales of the problem. However Schrödinger theory could still be useful for problems where the quantum evolution can be performed analytically, with the computation of transition probabilities at repeated intervals may be the limiting factor in the size of simulation timesteps.

6.3 Overview of method

To motivate the method and highlight some necessary properties that any sensible semiclassical model capable of simulating Stern–Gerlach separation must have, consider, as an example, a spin- $\frac{1}{2}$ particle undergoing Stern–Gerlach separation in a magnetic field gradient. The two spin components accelerate away from each other, with the relative acceleration vector pointing in the same direction as the gradient in field strength. We can ask the question: “What spin-state populations would I see, if I followed the spin-down wavepacket only?” The meaning of “follow” will be made more precise in Section 6.5, but for now we will simply look at the spin populations in the vicinity of the region of space occupied by the spin-down wavepacket.

As shown in Figure 6.3, one sees the spin-up population decreasing over time until only spin-down population remains. The rate at which spin-up population ‘leaves’ the region of space we are watching depends on how quickly the wavepackets are accelerating away from each other, as well as the shape of the wavepackets. Likewise, if one follows

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the spin-up component instead, one sees the spin-down population decrease to zero.

Motivated by this observation, the hidden-variable semiclassical method is a phenomenological model comprised of the following bolted-together pieces:

- One internal state of the atom is chosen at any moment in time, stored along with the state vector as a stochastic hidden variable that makes transitions at each timestep, according to a hidden-variable theory. The hidden-variable theory takes as input the state vector and Hamiltonian or unitary evolution matrix describing the evolution of the internal state of the atom at each timestep, and thus the probability of the hidden variable selecting out a particular internal state is equal to that state's population at each moment in time.
- The internal state of the atom evolves according to the Schrödinger equation, but with the addition of back-action caused by the continuous projection of the atomic wavefunction onto a specific ‘classical’ motional state under assumptions about the form and evolution of all motional states. This is a tricky matter to which I devote several sections later in this chapter.
- Upon a change in the state selected by the hidden variable, the velocity of the atom is modified instantaneously as required by energy conservation, or if this would result in a negative kinetic energy, the transition is disallowed.

Written in the same way as we previously described semiclassical models, the equations of motion for the model are the following coupled differential equations and stochastic transition rule:

$$\frac{d}{dt} |\tilde{\chi}\rangle = -\frac{i}{\hbar} \hat{H}(\mathbf{r}, t) |\tilde{\chi}\rangle - \hat{\Gamma}_\eta(\mathbf{r}, t) |\tilde{\chi}\rangle, \quad (6.23)$$

$$\frac{d}{dt} \mathbf{v} = -\frac{1}{m} \nabla V_\eta(\mathbf{r}, t), \quad (6.24)$$

$$\frac{d}{dt} \mathbf{r} = \mathbf{v} \quad (6.25)$$

$$\Pr(\eta(t')=i|\eta(t)=j) = S_{ij}(U_{\text{eff}}(t', t), \chi(t)), \quad (6.26)$$

where $|\tilde{\chi}\rangle$ is the (non-normalised) internal state vector including the effect of back-action, $\hat{\Gamma}_\eta(\mathbf{r}, t)$ is a non-Hermitian operator that implements this back-action by decaying the amplitude of states not being followed (there are many ways to do this, all approximate, discussed in Section 6.5), η is the hidden variable, V_η is the adiabatic potential experienced by the eigenstate selected by the hidden variable, χ is the vector representation of the (normalised) state vector in the local eigenbasis, and $U_{\text{eff}}(t', t)$ is the unitary matrix describing the evolution of the state vector from time t to t' in the local eigenbasis under the action of the effective Hamiltonian \hat{H}_{eff} , defined in Section 6.4.2. In addition to these evolution rules, the state vector must be normalised at each timestep of simulation,⁴ and the velocity of the atom modified instantaneously whenever a transition is made in order to conserve energy. This latter requirement is detailed in Section 6.4.3.

6.4 Hidden variables: implementation details

In this section I go into the gritty details of numerically evaluating hidden-variable theories, conserving energy when a transition occurs, and some other concerns.

6.4.1 Numerically evaluating Schrödinger theory

As mentioned, Schrödinger theory has a number of disadvantages compared to Tully’s fewest-switches as a hidden-variable theory for a hidden-variable semiclassical model,

⁴It would be relatively simple to include a term in the differential equation to preserve normalisation, but the non-normalisation-preserving differential equation is simpler to write, compute, and understand.

primarily its high computational cost. Nonetheless here I present the results of my investigation into how to minimise said cost.

There are many ways to solve numerically for the row and column scalings required to compute the S matrix of Schrödinger theory (6.22), but there is one unique solution for the resulting scaled matrix.⁵ The simplest method is to simply alternate between scaling the rows to get the right row sums, then scaling the columns to get the right column sums, and repeating, that is, alternating between solving equation (6.20) for all a_i , and solving (6.21) for all b_i , until the result converges. This is called the Sinkhorn–Knopp method of r-c (row-column) scaling [114], but is computationally intensive, with slow convergence [115]. An alternative is the method by Linial et al. [115] which converges much faster. Both are iterative methods, and so in practice one can save the resulting row and column scalings at each integration step of a simulation and use them as the initial guesses for the same computation at the next integration step,⁶ providing a considerable speedup.

For the case of a two-state system, the row and column scalings can be found analytically, with the result

$$a_1 = 1 \quad (6.27)$$

$$a_2 : a_2^2 + \left(\frac{1}{AB} - \frac{A}{B} |c_1(t)|^2 - \frac{B}{A} |c_2(t)|^2 \right) a_2 + \frac{|c_2(t')|^2}{|c_1(t')|^2} = 0 \quad (6.28)$$

$$b_1 = \frac{|c_1(t')|^2}{A + Ba_2} \quad (6.29)$$

$$b_2 = \frac{|c_2(t')|^2}{B + Aa_2}, \quad (6.30)$$

where a_2 is the positive solution to the given quadratic, $A = |U_{11}| = |U_{22}|$ and $B = |U_{12}| = |U_{21}|$. This result only holds in the case of Schrödinger theory for a state vector with two components subject to unitary evolution—not for row-column scaling in general—since it makes use of symmetries of unitary matrices and the fact that the initial and final probabilities given by the squared amplitudes of the state vector components must sum to unity.

Some further notes on numerics: the above expressions for Schrödinger theory and its analytic expression for a two-state system involve dividing by elements of the unitary, and state populations, both of which may be zero or very close to zero. Whilst the relevant limits may exist, we cannot easily compute them numerically, and so I have taken to simply replacing small values of $|U_{ij}|$, $|c_i(t')|^2$ and $|c_j(t)|^2$ with a tiny nonzero constant (I use $\varepsilon = 10^{-100}$), ensuring that the convergence criterion (which represents a tolerance for the sum squared error in the column sums) I pass to Linial’s method is larger than the square of this by some margin, so as to allow convergence even though modifying the matrix elements may make the matrix no longer row-column scalable to higher precisions. I use a convergence criterion of 10^{-16} for Linial’s algorithm, implying the root sum squared error in column sums will be at most 10^{-8} , which is small compared to unity—the sum of all column sums for a perfectly scaled matrix given that the column sums are probabilities that must add to unity. Smaller tolerances imply more iterations before convergence, so tolerances should be as large as is acceptable for the problem if computational cost is a limiting factor.

Potentially faster algorithms exist for row-column scaling of matrices,⁷ for example, ones that treat the problem as a root finding problem or optimisation problem aimed at solving the simultaneous equations (6.20) and (6.21) or minimising the residuals [116]. I had initially great success with Newton’s method for finding a root to this set of equations (after fixing $a_1 = 1$ to make them fully determined), which required considerably fewer iterations than Linial’s method for small (3×3) random unitary matrices and random state

⁵The values of $\{a_i\}$ and $\{b_j\}$ are only determined up to an overall multiplication of each $\{a_i\}$ by a constant and division of each $\{b_j\}$ by the same constant, since only products $a_i b_j$ appear in the resulting scaled matrix

⁶With the caveat that since (as mentioned in the previous footnote) the row and column scalings are only determined up to an overall multiplication/division, occasional multiplication of all $\{a_i\}$ and division of all $\{b_j\}$ by a constant may be necessary to prevent the values numerically overflowing or underflowing in the middle of a simulation.

vectors. However, the unitaries and state vectors in quantum mechanics are not random, and the fact that most elements of the unitary and state vector are zero when there is no evolution and the atom is in an eigenstate resulted in numerical difficulties with Newton's method that Linial's method does not seem to encounter. Similar to many of the methods in reference [116], one could construct a hybrid method that takes a Newton step, and then checks the row sum and column sum residuals, and if they increased compared to the previous step, ignores that step and takes a step of Linial's method instead. I have not attempted this, and for the moment use Linial's method.

A final note is that my hidden-variable semiclassical method is not just agnostic to which matrix scaling algorithm is used, but that it is also not married to any particular hidden-variable theory. An early version of the method [104] was limited to two-component systems, and the probability of transition was computed as:

$$\Pr(\eta(t')=2|\eta(t)=1) = \max(0, |c_2(t')|^2 - |c_2(t)|^2), \quad (6.31)$$

$$\Pr(\eta(t')=1|\eta(t)=2) = \max(0, |c_1(t')|^2 - |c_1(t)|^2), \quad (6.32)$$

that is, I simply inspected the populations each step and declared any positive change in population of a state as a probability of transition from the other state. Since there were only two states, the originating state of the transition was unambiguous, but the method did not generalise to systems with three or more states.⁸ However, it resulted in final populations in simulations on average in agreement with the underlying Schrödinger wave equation, leading me to suspect that the exact hidden-variable theory used is not crucial, so long as it satisfies the most obvious of Aaronson's axioms so as not to behave like the product theory. I chose Schrödinger theory fairly arbitrarily, it being the one that seemed most easily computable out of the two presented in Aaronson's paper [90], but one might try using Aaronson's flow theory, or inventing another altogether. The surface hopping literature uses Tully's fewest-switches algorithm with much success, as well as approximations to it [118] that are even cheaper, computationally speaking.

Interestingly, the main conclusion of Aaronson's paper [90] is that if we could know the entire history of a hidden variable, we could use it to make a computer more powerful than a quantum computer. It is therefore perhaps not surprising that hidden-variable theories ought to be computationally expensive to simulate on a classical computer. Prior to discovering Tully's fewest-switches algorithm, I was therefore somewhat resigned to the fact that any hidden-variable theory would likely be computationally expensive to compute. In light of this it was pleasantly surprising to discover that Tully's fewest-switches (discussed in the next subsection) is computationally cheap. The seeming conflict could be reconciled however if the computational power of hidden variables is crucially dependent on some feature we are not interested in and which fewest-switches does not capture, such as agreement with arbitrary unitaries rather than only those due to evolution over small time intervals, as is assumed by fewest-switches.

⁸This was before I coincidentally came across the definition of a stochastic hidden-variable theory in Aaronson's book *Quantum computing since Democritus* [117] and realised that what I had made was a hidden-variable theory, allowing me to choose a more general one from his paper (and before later still, finding Tully's fewest-switches algorithm).

6.4.2 Time-dependent formulation of Tully's fewest-switches algorithm

In this subsection I derive Tully's fewest-switches algorithm [93, 95] with the distinction that the Hamiltonian may have arbitrary time-dependence. As such, the non-adiabatic transitions that can occur can be due to either the spatial variation of the Hamiltonian, or its time variation. It is important to be able to distinguish which of the two non-adiabatic effects is responsible for a given transition of the hidden variable, as energy conservation only applies in the case of a nonadiabatic transition due to spatial variation of the Hamiltonian, in which case the atom is paying/receiving the energy cost of a transition using its kinetic energy. However for a transition due to temporal variation of the Hamiltonian, energy can be added and removed from the atom by the driving field without conserving its total energy. Thus, when a time dependent potential is used in

a hidden-variable semiclassical simulation, I propose that velocity corrections ought to only be performed following transition of the hidden variable if that transition was due to spatial variation in the Hamiltonian, and not temporal variation.

As flagged earlier, this proposed improvement is speculative and untested. In the results section of this chapter I have only tested the model for the case of time-independent fields. Though accepting this proposed improvement would clearly yield the correct behaviour in both the limit of a time-independent inhomogeneous field (energy is always conserved via velocity jumps) and a time-dependent homogeneous field (energy is never required to be conserved and there are no velocity jumps); for a field with both time and space dependence it is not clear that giving some atoms velocity jumps and some not will yield acceptable results, though it is certainly an improvement over giving velocity jumps to them all. Perhaps upon comparison with the Schrödinger wave equation it will become clear that say, giving all atoms velocity jumps of a smaller magnitude yields good results. How much smaller, however, will likely depend on the same matrix elements computed below. This subsection also serves to present Tully's fewest-switches algorithm and the concepts and notation related to it that I refer to in later sections.

We begin with the time dependent Schrödinger equation for the internal state $|\chi(t)\rangle$ of an atom at position \mathbf{r} :

$$i\hbar \frac{d}{dt} |\chi(t)\rangle = \hat{H}(\mathbf{r}, t) |\chi(t)\rangle \quad (6.33)$$

Now we take the unitary $\hat{U}_H(\mathbf{r}, t)$ that transforms state vectors into the eigenbasis of \hat{H} such that

$$\hat{H}(\mathbf{r}, t) = \hat{U}_H^\dagger(\mathbf{r}, t) \hat{V}(\mathbf{r}, t) \hat{U}_H(\mathbf{r}, t), \quad (6.34)$$

where $\hat{V}(\mathbf{r}, t)$ is a diagonal operator with diagonals equal to the adiabatic potentials that each eigenstate of \hat{H} is subject to in the adiabatic approximation. We can therefore define the state vector in the eigenbasis of \hat{H} as⁹

$$|\chi_H(t)\rangle = \hat{U}_H(\mathbf{r}, t) |\chi(t)\rangle \quad (6.35)$$

$$\Rightarrow |\chi(t)\rangle = \hat{U}_H^\dagger(\mathbf{r}, t) |\chi_H(t)\rangle. \quad (6.36)$$

Substituting (6.35) into (6.33) and premultiplying by $\hat{U}_H(\mathbf{r}, t)$ yields

$$i\hbar \hat{U}_H(\mathbf{r}, t) \frac{d}{dt} \hat{U}_H^\dagger(\mathbf{r}, t) |\chi_H(t)\rangle = \hat{U}_H(\mathbf{r}, t) \hat{H}(\mathbf{r}, t) \hat{U}_H^\dagger(\mathbf{r}, t) |\chi_H(t)\rangle, \quad (6.37)$$

which via the product rule and our definition of $\hat{V}(\mathbf{r}, t)$ simplifies to the differential equation obeyed by the transformed state vector $|\chi_H(t)\rangle$:

$$i\hbar \frac{d}{dt} |\chi_H(t)\rangle = \left[\hat{V}(\mathbf{r}, t) - i\hbar \hat{U}_H(\mathbf{r}, t) \frac{d}{dt} \hat{U}_H^\dagger(\mathbf{r}, t) \right] |\chi_H(t)\rangle \quad (6.38)$$

$$\equiv \hat{H}_{\text{eff}} |\chi_H(t)\rangle. \quad (6.39)$$

This equation has the same form as the Schrödinger equation, with the contents of the brackets comprising an effective Hamiltonian dictating the dynamics of the state vector in the *adiabatic basis* (the basis in which \hat{H} is diagonal). Like a non-inertial reference frame in classical mechanics, use of this transformed basis has resulted in the appearance of an extra term in the Hamiltonian, the non-adiabatic coupling term depending on the time derivative of the transformation \hat{U}_H^\dagger .

Here we differ from Tully by proceeding without assuming that \hat{U}_H has no explicit time dependence. The total time derivative of \hat{U}_H^\dagger includes both its direct time dependence and the effect of motion through space; the latter obtainable via the chain rule:

$$\frac{d}{dt} \hat{U}_H^\dagger(\mathbf{r}, t) = \frac{\partial}{\partial t} \hat{U}_H^\dagger(\mathbf{r}, t) + \mathbf{v} \cdot \nabla \hat{U}_H^\dagger(\mathbf{r}, t), \quad (6.40)$$

where $\boldsymbol{\nu} = \frac{d\mathbf{r}}{dt}$. Thus (6.38) becomes:

$$i\hbar \frac{d}{dt} |\chi_H(t)\rangle = \left[\hat{V}(\mathbf{r}, t) - i\hbar \hat{U}_H(\mathbf{r}, t) \frac{\partial}{\partial t} \hat{U}_H^\dagger(\mathbf{r}, t) - i\hbar \boldsymbol{\nu} \cdot \hat{U}_H(\mathbf{r}, t) \nabla \hat{U}_H^\dagger(\mathbf{r}, t) \right] |\chi_H(t)\rangle. \quad (6.41)$$

The final term is identical to the non-adiabatic coupling term in the equation of motion as usually written in the surface-hopping literature [95] being a matrix with elements (in the eigenbasis):

$$\left(-i\hbar \boldsymbol{\nu} \cdot U_H(\mathbf{r}, t) \nabla U_H^\dagger(\mathbf{r}, t) \right)_{ij} = -i\hbar \boldsymbol{\nu} \cdot \langle \chi_i(\mathbf{r}, t) | \nabla | \chi_j(\mathbf{r}, t) \rangle, \quad (6.42)$$

where $U_H(\mathbf{r}, t)$ is the matrix representation of \hat{U} in any basis that does not vary spatially (i.e, not the eigenbasis), $|\chi_i(\mathbf{r}, t)\rangle$ is the i^{th} eigenvector of $\hat{H}(\mathbf{r}, t)$, and $\langle \chi_i(\mathbf{r}, t) | \nabla | \chi_j(\mathbf{r}, t) \rangle$ is the *non-adiabatic coupling vector* between the i^{th} and j^{th} states referred to in the literature [95]. The second to last term in brackets in (6.41) is the additional contribution due to the time-dependence of the Hamiltonian (more specifically: the time-dependence of its eigenbasis).

We now proceed identically to Tully, computing the rate of change of an eigenstates population $|c_i(t)|^2$ as:

$$\frac{d}{dt} |c_i(t)|^2 = c_i(t) \frac{d}{dt} c_i^*(t) + c_i^*(t) \frac{d}{dt} c_i(t), \quad (6.43)$$

where via (6.41) we have:

$$\frac{d}{dt} c_i(t) = -\frac{i}{\hbar} \sum_j (H_{\text{eff}}(\mathbf{r}, t))_{ij} c_j(t) \quad (6.44)$$

$$\Rightarrow \frac{d}{dt} c_i(t) = -\frac{i}{\hbar} \sum_j \left[V_{ij}(\mathbf{r}, t) - i\hbar \langle \chi_i(\mathbf{r}, t) | (\partial_t + \boldsymbol{\nu} \cdot \nabla) | \chi_j(\mathbf{r}, t) \rangle \right] c_j(t). \quad (6.45)$$

This gives for the time rate of change of the population $|c_i(t)|^2$:

$$\frac{d}{dt} |c_i(t)|^2 = \left[-\frac{i}{\hbar} \sum_j c_i^*(t) (H_{\text{eff}}(\mathbf{r}, t))_{ij} c_j(t) \right] + \text{c.c.} \quad (6.46)$$

$$= -\frac{2}{\hbar} \sum_j \text{Im} \left(c_i^*(t) (H_{\text{eff}}(\mathbf{r}, t))_{ij} c_j(t) \right) \quad (6.47)$$

$$= -\frac{2}{\hbar} \sum_j \text{Im} \left(c_i^*(t) \left[V_{ij}(\mathbf{r}, t) - i\hbar \langle \chi_i(\mathbf{r}, t) | (\partial_t + \boldsymbol{\nu} \cdot \nabla) | \chi_j(\mathbf{r}, t) \rangle \right] c_j(t) \right). \quad (6.48)$$

Since $V(\mathbf{r}, t)$ is diagonal¹⁰ and real, $c_i^*(t) V_{ij}(\mathbf{r}, t) c_j(t)$ is zero when $i \neq j$, and has no imaginary part when $i = j$, leaving us with

$$\frac{d}{dt} |c_i(t)|^2 = 2 \sum_j \text{Re} \left(c_i^*(t) \langle \chi_i(\mathbf{r}, t) | (\partial_t + \boldsymbol{\nu} \cdot \nabla) | \chi_j(\mathbf{r}, t) \rangle c_j(t) \right). \quad (6.49)$$

The change in $|c_i(t)|^2$ in a small interval dt is then:

$$|c_i(t + dt)|^2 - |c_i(t)|^2 = 2dt \sum_j \text{Re} \left(c_i^*(t) \langle \chi_i(\mathbf{r}, t) | (\partial_t + \boldsymbol{\nu} \cdot \nabla) | \chi_j(\mathbf{r}, t) \rangle c_j(t) \right). \quad (6.50)$$

¹⁰This is not always assumed in the surface hopping literature, since additional couplings are sometimes included in V which have not been removed by diagonalisation.

This is the change in the probability of the atom being in the i^{th} state during that time interval. Tully identifies each term in the sum as a probability flow between a pair of states, and if non-negative, equates each term with the (unconditional) probability of a transition from the j^{th} state to the i^{th} state. We do the same, except that we identify two transition probabilities for each originating state, one due to the spatial variation in the eigenbasis, and one due to the temporal variation. To ensure we don't violate the criterion that on a two-state basis only the minimum number of hops consistent with the total probability flow occur, we clip each probability from above to the probability of any transition occurring at all. This gives us transition probability matrix elements

$$P_{ij}^{\text{space}} = \min \left\{ q_{ij}^{\text{total}}, q_{ij}^{\text{space}} \right\}, \quad (6.51)$$

$$P_{ij}^{\text{time}} = \min \left\{ q_{ij}^{\text{total}}, q_{ij}^{\text{time}} \right\}, \quad (6.52)$$

where

$$q_{ij}^{\text{space}} = 2dt \operatorname{Re} \left(c_i^*(t) \langle \chi_i(\mathbf{r}, t) | \mathbf{v} \cdot \nabla | \chi_j(\mathbf{r}, t) \rangle c_j(t) \right), \quad (6.53)$$

$$q_{ij}^{\text{time}} = 2dt \operatorname{Re} \left(c_i^*(t) \langle \chi_i(\mathbf{r}, t) | \partial_t | \chi_j(\mathbf{r}, t) \rangle, c_j(t) \right), \quad (6.54)$$

$$q_{ij}^{\text{total}} = \max \left\{ 0, (q_{ij}^{\text{space}} + q_{ij}^{\text{time}}) \right\}, \quad (6.55)$$

for transitions of the hidden variable from the j^{th} to the i^{th} eigenstate of $\hat{H}(\mathbf{r}, t)$ during the time interval dt due to non-adiabatic spatial and temporal variations in $\hat{H}(\mathbf{r}, t)$ respectively. These expressions can also be numerically integrated with respect to time to obtain transition probabilities over finite time intervals—numerically integrating over a short time interval Δt using a midpoint or higher order method can produce transition probabilities more accurate than $\mathcal{O}(\Delta t)$, which would be the accuracy if Δt were simply used in place of dt in the above expressions.

These matrices have zeros along their diagonals, since the above derivation takes into account only probability changes, and does not count probability remaining in the same state as a transition.¹¹ To be able to construct properly stochastic matrices, we can take into account the probability mass that remains in the same state simply by imposing conservation of overall probability, defining a diagonal matrix P^{stay} for the unconditional probabilities of remaining in a state:

$$P_{ii}^{\text{stay}} = |c_i(t)|^2 - \sum_{j \neq i} (P_{ij}^{\text{space}} + P_{ij}^{\text{time}}). \quad (6.56)$$

The sum of all three of these matrices now satisfy the row sum and column sum requirements in order to be the unconditional transition probabilities for a hidden-variable theory in the eigenbasis of \hat{H} :

$$P = P^{\text{space}} + P^{\text{time}} + P^{\text{stay}}, \quad (6.57)$$

$$\Rightarrow \sum_j P_{ij} = |c_i(t)|^2, \quad (6.58)$$

$$\sum_i P_{ij} = |c_j(t + dt)|^2. \quad (6.59)$$

The corresponding conditional probabilities of a transition to the i^{th} state occurring—given that the hidden variable was already in the j^{th} state—can be obtained via (6.16)

as:

$$S_{ij}^{\text{space}} = \frac{1}{|c_j(t)|^2} P_{ij}^{\text{space}}, \quad (6.60)$$

$$S_{ij}^{\text{time}} = \frac{1}{|c_j(t)|^2} P_{ij}^{\text{time}} \quad (6.61)$$

$$S_{ij}^{\text{stay}} = \frac{1}{|c_j(t)|^2} P_{ij}^{\text{stay}}, \quad (6.62)$$

and the sum of these three matrices of conditional probabilities is the overall (left) stochastic matrix for the fewest-switches hidden-variable theory:

$$S = S^{\text{space}} + S^{\text{time}} + S^{\text{stay}}. \quad (6.63)$$

When using a time-dependent Hamiltonian, under my proposal one would not use this stochastic matrix to make transitions.¹² Rather one would use the individual matrices S^{space} , S^{time} , and S^{stay} in order to distinguish between the different types of transitions, and make the energy conservation part of the surface hopping algorithm conditional on the transition being attributed in this way to the spatial variation of the Hamiltonian rather than its temporal variation.

During a simulation, to choose whether a transition occurs due to spatial or temporal variations in the Hamiltonian, one should not make independent random choices based on the three above matrices of conditional probabilities. Rather one should assemble the probabilities of possible events—transitions from the current state to all others via both spatial and temporal non-adiabatic transitions—into a single list of probabilities, and then take the cumulative sum, resulting in a list of numbers between zero and one. A randomly generated number between zero and one can then be used to determine which event occurs, with the correct probability (Figure 6.4).

¹²Though it is instructive to know that the sum of the other conditional probability matrices is indeed a stochastic matrix, such that Tully's fewest-switches does satisfy this requirement of being a hidden-variable theory as defined by Aaronson.

Framing fewest-switches as a hidden-variable theory

Tully's fewest-switches algorithm allows one to compute transition probabilities for a hidden variable given the Hamiltonian, the state amplitudes, and a small interval of time. Does this satisfy Aaronson's definition of a hidden-variable theory? As written, not quite, since it requires the Hamiltonian rather than the unitary that describes state vector evolution over a particular time interval. However it is simple to reconcile the two sets of requirements. Given that the interval of time is small, the unitary describing evolution in the local basis can be linked to the effective Hamiltonian in (6.39) via

$$\hat{U}(t + dt, t) = e^{-\frac{i}{\hbar} \hat{H}_{\text{eff}}(\mathbf{r}, t) dt}, \quad (6.64)$$

where

$$\hat{H}_{\text{eff}} = \hat{V}(\mathbf{r}, t) - i\hbar \hat{U}_H(\mathbf{r}, t) \frac{d}{dt} U_H^\dagger(\mathbf{r}, t) \quad (6.65)$$

is the effective Hamiltonian in the eigenbasis of \hat{H} , the matrix representation of which can be extracted from the matrix representation of $\hat{U}(t + dt, t)$ as:

$$H_{\text{eff}}(\mathbf{r}, t) dt = i\hbar \text{Log } U(t + dt, t) \quad (6.66)$$

$$\approx i\hbar (\mathbb{I} - U(t + dt, t)) \quad (6.67)$$

where Log is the principal value of the complex matrix logarithm.

Since only the matrix $H_{\text{eff}}(\mathbf{r}, t) dt$ is required to compute transition probabilities according to (6.47), and not its component terms,¹³ specifying the initial state vector

¹³Though we do need its component terms if we wish to distinguish between spatially vs. temporally induced transitions as in Section 6.4.2.

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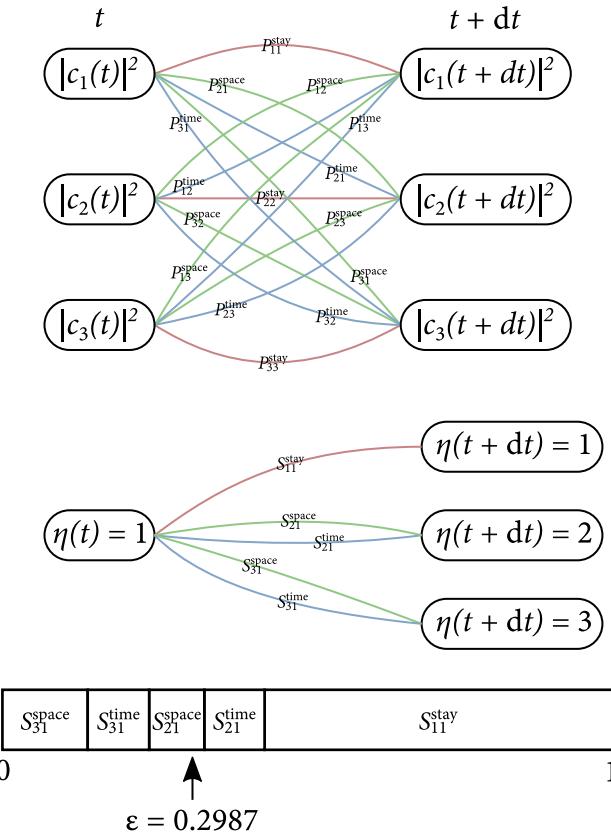


Figure 6.4: Probabilistically choosing a transition. Top: Probability flows between states in a time interval according to the three matrices of unconditional probabilities P^{stay} , P^{space} , and P^{time} in such a way that one total unit of probability is routed from states at the initial time to states at the final time consistent with the populations resulting from quantum mechanical evolution in that timestep. Centre: The hidden variable transitions according to the corresponding conditional probabilities which are elements of the matrices S^{stay} , S^{space} , and S^{time} . Here the hidden variable is in state 1 and we are choosing whether it will remain in that state or if it will transition to state 2 or 3, and whether it transitions via a spatial or temporal non-adiabatic transition. Bottom: All the probabilities for what may happen to the hidden variable in a timestep sum to one, so an array of the cumulative probabilities can be constructed, and a random number drawn from the range $[0, 1]$. The index of the smallest element of the array of cumulative sums that the random number is smaller than corresponds to the event to occur. In the above diagrammatic example, the result of the random draw is that the hidden variable is to transition to state 2 via a spatially induced non-adiabatic transition.

and unitary for an interval of time evolution (both in the local basis) is a sufficient input to be able to compute transition probabilities using Tully's fewest-switches algorithm. Writing the resulting matrix of probabilities in terms of U gives:

$$P_{ij} = \begin{cases} \max \left\{ 0, 2 \operatorname{Re} \left(c_i^*(t) U_{ij}(t + dt, t) c_j(t) \right) \right\} & i \neq j \\ |c_i(t)|^2 - \sum_{j \neq i} P_{ij} & i = j \end{cases} \quad (6.68)$$

The corresponding stochastic matrix S can then be obtained by scaling the columns of

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P by the state populations as in (6.16). Tully's fewest-switches algorithm thus satisfies Aaronson's definition of a hidden-variable theory provided that the time interval is small. What is somewhat remarkable is the low computational complexity of computing probabilities via fewest-switches. There is no matrix scaling, no matrix permanents, or any other large computational expense.

Expressed in Python, Tully's fewest-switches can be computed very simply as shown below, though this calculation is only first-order accurate in the time interval corresponding to the unitaries.

```

1 import numpy as np
2
3 def S_matrix_fewest_switches(psi, U):
4     """Compute the fewest-switches S matrix for an array psi, shape (M, N),
5     containing N state vectors for an system with M states, and the array of
6     corresponding unitaries unitaries U, shape (M, M, N), describing the evolution
7     of the state vectors over a short time interval."""
8     P = 2 * np.einsum('in,ijn,jn->ijn', psi.conj(), U, psi).real
9     P[P < 0] = 0
10    S = P / np.abs(psi)**2
11    S_diags = np.einsum('in->in', S)
12    S_diags[...] = 0
13    S_diags[...] = 1 - np.sum(S, axis=0)
14
15    return S

```

It is not clear how many of Aaronson's axioms are satisfied by fewest-switches, but since it does depend on the actual coupling strengths between states via the non-adiabatic Hamiltonian, one would expect it to satisfy a few of them. Furthermore, the form of fewest-switches when framed in terms of the unitary for the given time interval is quite similar to that of Schrödinger theory. In Schrödinger theory one takes the absolute value of elements of U and then applies row and column scalings. In fewest-switches one applies row and column scalings (explicitly given as $c_i(t)^*$ and $c_j(t)$, along with a factor of 2), then takes the real part of the result and clips it to zero from below. The two methods are not identical, but the similarity is striking.

6.4.3 Velocity correction and classically disallowed transitions

When a transition of the hidden variable occurs due to a spatial non-adiabatic transition, the kinetic energy of the atom must be adjusted to conserve overall energy. The force on an atom during a transition from the j^{th} state to the i^{th} state due to the spatial non-adiabatic coupling term in the effective Hamiltonian is in the direction of the non-adiabatic coupling vector [95]

$$\mathbf{d}_{ij}(\mathbf{r}, t) = \langle \chi_i(\mathbf{r}, t) | \nabla | \chi_j(\mathbf{r}, t) \rangle \quad (6.69)$$

$$= \left(U_H(\mathbf{r}, t) \nabla U_H^\dagger(\mathbf{r}, t) \right)_{ij} \quad (6.70)$$

where $U_H(\mathbf{r}, t)$ is the matrix representation, in any (spatially and temporally fixed) basis, of the unitary that takes state vectors into the basis in which $\hat{H}(\mathbf{r}, t)$ is diagonal. To conserve energy, the squared component of an atoms velocity in this direction must change by an amount:

$$\Delta V_{ij}^2 = \frac{2}{m} \left(V_j(\mathbf{r}, t) - V_i(\mathbf{r}, t) \right), \quad (6.71)$$

where $V_i(\mathbf{r}, t)$ is the adiabatic potential comprising the i^{th} spatially and/or temporally varying eigenvalue of $\hat{H}(\mathbf{r}, t)$. However, it is possible that a change of this size can leave an atom with a negative kinetic energy. Whilst this is quantum-mechanically permissible, it is forbidden classically, and so we simply disallow such transitions, defining a modified

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matrix of unconditional transition probabilities \tilde{P}^{space} that sets the probability of the disallowed transitions to zero:

$$\tilde{P}_{ij}^{\text{space}} = \begin{cases} P_{ij}^{\text{space}} & \Delta v_{ij}^2 + |\mathbf{v} \cdot \hat{\mathbf{d}}_{ij}|^2 \geq 0 \\ 0 & \Delta v_{ij}^2 + |\mathbf{v} \cdot \hat{\mathbf{d}}_{ij}|^2 < 0 \end{cases} \quad (6.72)$$

¹⁴The surface hopping literature calls these classically disallowed transitions *frustrated transitions*. Prior to encountering the surface-hopping literature, my method of conserving energy was identical to this (for the case of time-independent potentials), with the exception that I did not know what direction the velocity kick ought to be in, limiting my method to one spatial dimension only. There is much interest in alternate methods of energy conservation in chemical physics, however so far this has been outside of my interest in these methods for simulating cold atoms.

where $\hat{\mathbf{d}}_{ij} = \hat{\mathbf{d}}_{ij}(\mathbf{r}, t)$ is the unit vector in the direction of $\mathbf{d}_{ij}(\mathbf{r}, t)$.¹⁴

The diagonal matrix for the probabilities of remaining in each state must be adjusted similarly to absorb the probability discarded this way:

$$\tilde{P}_{ii}^{\text{stay}} = |c_i(t)|^2 - \sum_{j \neq i} (\tilde{P}_{ij}^{\text{space}} + P_{ij}^{\text{time}}) \quad (6.73)$$

The corresponding matrices of conditional probabilities for the hidden variable transitioning, given that it is already in the j^{th} state, are

$$\tilde{S}_{ij}^{\text{space}} = \frac{1}{|c_j(t)|^2} \tilde{P}_{ij}^{\text{space}}, \quad (6.74)$$

$$\tilde{S}_{ij}^{\text{stay}} = \frac{1}{|c_j(t)|^2} \tilde{P}_{ij}^{\text{stay}}, \quad (6.75)$$

Of course, when making probabilistic transitions of the hidden variable, the originating state is known and so only one column of any of the above matrices is used at a time, so the full matrices \tilde{P}^{space} , \tilde{P}^{time} , \tilde{S}^{space} , and \tilde{S}^{time} do not need to be computed at each timestep.

In the case that a transition of the hidden variable does occur to the i^{th} from the j^{th} state due to a spatial non-adiabatic coupling, the velocity kick to be applied to the classical velocity vector in order to conserve energy is:

$$\Delta v_{ij} = \left[\text{sgn}(\mathbf{v} \cdot \hat{\mathbf{d}}_{ij}) \sqrt{(\mathbf{v} \cdot \hat{\mathbf{d}}_{ij})^2 + \frac{2}{m} (V_i(\mathbf{r}, t) - V_j(\mathbf{r}, t))} - (\mathbf{v} \cdot \hat{\mathbf{d}}_{ij}) \cdot \right] \hat{\mathbf{d}}_{ij} \quad (6.76)$$

6.5 Decoherence

In the context of the hidden-variable semiclassical/surface hopping method, *decoherence* refers to the fact that, due to positional separation of different internal states of the atom, those internal states transition from being a coherent superposition into a statistical mixture. For example, in the Stern–Gerlach experiment, a spin may be initially a coherent superposition of spin-up and spin-down, but by the time the two components separate, the superposition is no longer a coherent one. The spin is either up (and the atom off to one side of the screen) or down (and the atom over the other side of the screen). At this point, no interference between the two internal states is observable, as they are not co-located in space. The motional degree of freedom has played the role of an environment, and, having become entangled with the spin of the atom, decohered the spin state. Recoherence can occur if the two wavepackets are brought back together again to have well overlapping positions and velocities, but this is difficult to achieve even intentionally,¹⁵ let alone by accident,

Decoherence was not part of Tully's original surface hopping model [93], such that if used to simulate the Stern–Gerlach experiment, two spots would result on the screen, but both spots would comprise a coherent 50 : 50 superposition of spin-up and spin-down, rather than one being up and one being down. Some decoherence is necessary to include in a hidden-variable semiclassical/surface hopping model in order to avoid this unphysical outcome, and furthermore, to obtain correct transition probabilities in

¹⁵The problem of getting the wavepackets back together again has been coined the “Humpty-Dumpty problem” [93, 119] and has made magnetic separation impractical for large-momentum atom interferometry [120], and methods such as the use of Bragg pulses [121, 122] are needed to create large momentum differences between wavepackets without having them more slowly traverse the intermediate regions of momentum space where they might accumulate coherence-destroying phase noise.

the case that the atom undergoes multiple periods of local spin transitions (as being in a 50 : 50 superposition represents entirely different initial conditions for a Majorana spin flip than being in a single spin state).

Our approach to modelling decoherence requires some assumptions about the motional states of the atom and how they evolve in time. We therefore construct some wavepackets that are as classical as we can make them, being as localised in position and velocity space as they can be. For this we use Gaussian wavepackets of width equal to the thermal De Broglie wavelength $\sigma = \lambda_{\text{th}} = h/\sqrt{2\pi m k_B T}$, evolving without dispersion, with their centre-of-mass position and velocity evolving classically according to the adiabatic potential experienced by the eigenstate to which each motional state corresponds.

In Section 6.5.1 I show the general mechanism by which the divergence of trajectories decoheres internal states of an atom, and specifically how this appears as an apparent decrease in the amplitudes of each other state when one imposes the adiabatic trajectory corresponding to one specific internal state. I then show in Section 6.5.2 why naïvely projecting the wavefunction onto a single trajectory at all times yields nonsensical results, and then in Subsections 6.5.3 and 6.6.2 I present two remedies for how to approximately include decoherence despite this.

The development of the second of these remedies brought to light a possibility I hadn't considered before regarding the form of the 'classical' state that the total wavefunction is projected onto at each timestep, namely that is that it makes some amount of sense for this to be a Dirac delta instead of a Gaussian wavepacket. In light of this, after all that work, in Section 6.5.5 we change our mind about what we mean when we say we are "following" a wavepacket, and produce alternate expressions for the decoherence rates of the two methods under this alternate assumption. I argue why this is appealing and why it resolves some issues that are present when projecting onto a Gaussian wavepacket.

6.5.1 Back action of position measurement on internal state

Consider an atom in a state $|\Psi(t)\rangle$, which is an arbitrary superposition of internal basis states $|\chi_i\rangle$ and motional basis states $|\mathbf{r}\rangle$:

$$|\Psi(t)\rangle = \int \sum_i \psi_i(\mathbf{r}, t) |\chi_i\rangle \otimes |\mathbf{r}\rangle d\mathbf{r}, \quad (6.77)$$

where normalisation requires that

$$\int \sum_i |\psi_i(\mathbf{r}, t)|^2 d\mathbf{r} = 1. \quad (6.78)$$

Recognising that $\psi_i(\mathbf{r}, t)$ is the i^{th} component of a multi-component wavefunction that has one component for each internal state, we define (up to an arbitrary phase factor) a normalised wavefunction $\phi_i(\mathbf{r}, t)$ and corresponding motional state vector $|\phi_i(t)\rangle$ and its coefficient $c_i(t)$ for each component:

$$c_i(t)\phi_i(\mathbf{r}, t) \equiv \psi_i(\mathbf{r}, t) \quad (6.79)$$

$$|\phi_i(t)\rangle \equiv \int \phi_i(\mathbf{r}, t) |\mathbf{r}\rangle d\mathbf{r} \quad (6.80)$$

such that $\int |\phi_i(\mathbf{r}, t)|^2 d\mathbf{r} = 1$, allowing us to write our arbitrary state vector as a sum over the internal basis only:

$$|\Psi(t)\rangle = \sum_i c_i(t) |\chi_i\rangle \otimes |\phi_i(t)\rangle, \quad (6.81)$$

where the spatial state vectors $\{|\phi_i(t)\rangle\}$ are not necessarily orthogonal. To see that spatial separation of the different components leads to decoherence of the internal states, consider the pure density operator corresponding to $|\Psi(t)\rangle$:

$$\hat{\rho}(t) = |\Psi(t)\rangle\langle\Psi(t)| = \sum_{ij} c_i(t)c_j^*(t) |\chi_i(t)\rangle\langle\chi_j(t)| \otimes |\phi_i(t)\rangle\langle\phi_j(t)|, \quad (6.82)$$

which we can write in the $\{|\chi_i\rangle \otimes |\mathbf{r}\rangle\} \equiv \{|\chi_i, \mathbf{r}\rangle\}$ basis, resulting in matrix elements

$$\rho_{ij}(t, \mathbf{r}, \mathbf{r}') = \langle\chi_i | \rho(t) | \chi_j \rangle, \quad (6.83)$$

$$= \psi_i(\mathbf{r}, t) \psi_j^*(\mathbf{r}', t), \quad (6.84)$$

$$= c_i(t)c_j^*(t) \phi_i(\mathbf{r}, t) \phi_j^*(\mathbf{r}', t). \quad (6.85)$$

A partial trace [123] over the motional degree of freedom results in a reduced density operator describing the measurement statistics of the internal degree of freedom only, assuming ignorance of the position degree of freedom:

$$\hat{\rho}^{\text{red}}(t) = \int \langle \mathbf{r} | \hat{\rho}(t) | \mathbf{r} \rangle d\mathbf{r} \quad (6.86)$$

$$\Rightarrow \rho_{ij}^{\text{red}}(t) = \int \rho_{ij}(t, \mathbf{r}, \mathbf{r}') \delta(\mathbf{r} - \mathbf{r}') d\mathbf{r} \quad (6.87)$$

$$= c_i(t)c_j^*(t) \int \phi_i(\mathbf{r}, t) \phi_j^*(\mathbf{r}, t) d\mathbf{r} \quad (6.88)$$

$$= c_i(t)c_j^*(t) \langle \phi_j(t) | \phi_i(t) \rangle. \quad (6.89)$$

So we see that the off diagonals of the density matrix, representing the coherences of the internal states, are reduced by a factor $\langle \phi_j(t) | \phi_i(t) \rangle$, which will be unity only for pairs of motional state vectors that are identical. We therefore see that separation in space of the wavefunctions corresponding to different internal states leads to decoherence of the internal states, with decoherence factor $R_{ji} = \langle \phi_j(t) | \phi_i(t) \rangle$. This same decoherence factor appears when—instead of integrating over all positions—we project the total state vector onto a single motional state corresponding to a classical trajectory, which (along with assumptions about the form and evolution of the motional states), is how we impose classicality on the motional degree of freedom in our model. The effect of decoherence when “following” a specific motional state through space in this way is to reduce the amplitudes of all other states not being followed by the decoherence factor between that state and the one being followed, as shown schematically in Figure 6.3.

To obtain an explicit form for the decoherence factor $R_{ij}(t)$, we need to start imposing our assumptions about the motional states $\{|\phi_i(t)\rangle\}$. As mentioned earlier, we assume that each motional state $|\phi_i(t)\rangle$ is a Gaussian wavepacket propagating—without dispersion—with centre-of-mass motion evolving classically according to the adiabatic potential $V_i(\mathbf{r}, t)$ experienced by the local eigenstate $|\chi_i\rangle$. Thus the wavefunctions of these motional states are

$$\langle \mathbf{r} | \phi_i(t) \rangle = \phi_i(\mathbf{r}, t) = A \exp \left[-\frac{|\mathbf{r} - \mathbf{r}_i(t)|^2}{4\sigma^2} + i \frac{m}{\hbar} \mathbf{v}_i(t) \cdot (\mathbf{r} - \mathbf{r}_i(t)) \right], \quad (6.90)$$

Where A is a real normalisation constant,¹⁶ $\mathbf{r}_{ij}(t) = \mathbf{r}_i(t) - \mathbf{r}_j(t)$ and $\mathbf{k}_{ij}(t) = \frac{m}{\hbar} (\mathbf{v}_i(t) - \mathbf{v}_j(t))$ are the relative mean position and wavevector of the pair of wavepackets, and where the centre-of-mass position and velocity of each wavepacket evolve classically according to

$$\frac{d}{dt} \mathbf{r}_i(t) = \mathbf{v}_i(t) \quad (6.91)$$

$$\frac{d}{dt} \mathbf{v}_i(t) = -\frac{1}{m} \nabla V_i(\mathbf{r}_i, t). \quad (6.92)$$

¹⁶The overall phase is determined by the offset $\mathbf{r} - \mathbf{r}_i(t)$ in the second term in the exponent, and is chosen such that $\phi_i(\mathbf{r}, t)$ is real at $\mathbf{r} = \mathbf{r}_i$, which ensures that $\langle \phi_i(t) | \phi_j(t) \rangle$ has a phase depending only on the relative positions of the two wavepackets rather than their distance from some arbitrary origin. Acquiring an additional arbitrary phase at each timestep when we later apply our reduction operator \hat{R} would not affect the system dynamics, but is unappealing.

This gives for the decoherence factor:

$$R_{ij}(t) = \exp \left[- \left(\frac{1}{8\sigma^2} |\mathbf{r}_{ij}|^2 + \frac{i}{2} \mathbf{r}_{ij}(t) \cdot \mathbf{k}_{ij}(t) + \frac{\sigma^2}{2} |\mathbf{k}_{ij}(t)|^2 \right) \right]. \quad (6.93)$$

As expected, $R_{ij}(t)$ is equal to unity when the two wavepackets have identical positions and velocities, and decays to zero for increasing relative position and velocity.

If the two motional states being considered are identical at $t = 0$ and have constant relative acceleration \mathbf{a}_{ij} , then we have

$$\mathbf{r}_{ij}(t) = \frac{1}{2} \mathbf{a}_{ij} t^2 \quad (6.94)$$

and

$$\mathbf{k}_{ij}(t) = \frac{m}{\hbar} \mathbf{a}_{ij} t, \quad (6.95)$$

which gives for the decoherence factor:

$$R_{ij}(t) = \exp \left[- \frac{|\mathbf{a}_{ij}|^2}{2} \left(\frac{1}{16\sigma^2} t^4 + i \frac{m}{2\hbar} t^3 + \frac{m^2 \sigma^2}{\hbar^2} t^2 \right) \right]. \quad (6.96)$$

We will return later to this explicit form of the decoherence factor under these assumptions.

Returning to our arbitrary state vector (6.79), we find (ignoring normalisation) the projected state vector $|\tilde{\Psi}(t)\rangle = \hat{R}(t)|\Psi(t)\rangle$, where $\hat{R}(t) = \hat{1} \otimes |\phi_\eta(t)\rangle\langle\phi_\eta(t)|$, resulting from projection of the total state vector $|\Psi(t)\rangle$ onto the specific motional state vector $|\phi_\eta(t)\rangle$ corresponding to the state selected by the hidden variable η . This gives the (non-normalised) state vector one would observe conditional on the particle being in that specific motional state at that specific time.

$$|\tilde{\Psi}(t)\rangle = \hat{R}(t)|\Psi(t)\rangle \quad (6.97)$$

$$= \sum_i \langle\phi_\eta(t)|\phi_i(t)\rangle c_i(t) |\chi_i\rangle \otimes |\phi_\eta(t)\rangle \quad (6.98)$$

$$= \sum_i \tilde{c}_i(t) |\chi_i\rangle \otimes |\phi_\eta(t)\rangle \quad (6.99)$$

$$= |\tilde{\chi}(t)\rangle \otimes |\phi_\eta(t)\rangle, \quad (6.100)$$

where we have defined $\tilde{c}_i(t) = \langle\phi_\eta(t)|\phi_i(t)\rangle c_i(t)$ and $|\tilde{\chi}(t)\rangle = \sum_i \tilde{c}_i(t) |\chi_i\rangle$. We can immediately see that the effect of this projection is to reduce the amplitude of all other states ($i \neq \eta$) by the decoherence factor $R_{\eta i}(t)$. We now proceed by considering this a projective measurement to have occurred at time t , and evolving the resulting projected state vector to time $t + dt$ before projecting it again:

$$|\tilde{\Psi}(t+dt)\rangle = \hat{R}(t+dt) \sum_i \left[\tilde{c}_i(t) - \frac{i}{\hbar} \sum_j H_{ij}(t) \tilde{c}_j(t) dt \right] |\chi_i\rangle \otimes |\phi_i(t+dt)\rangle, \quad (6.101)$$

where $H_{ij}(t) = \langle\chi_i|\hat{H}_\eta(t)|\chi_j\rangle = \langle\chi_i|\phi_\eta(t)|\hat{H}(t)|\chi_j|\phi_\eta(t)\rangle$ are the elements of the projected Hamiltonian $\hat{H}_\eta(t)$ dictating the dynamics of the internal state, given the imposed motional state $|\phi_\eta\rangle$ and the total Hamiltonian $\hat{H}(t)$ of the system. Note that because of the previous projection already performed, all motional states $|\phi_i(t)\rangle$ were “reset” to be equal to $|\phi_\eta(t)\rangle$ at time t , and therefore the evolved motional state $|\phi_i(t+dt)\rangle$

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represents the evolution of the motional state corresponding to the i^{th} internal state over an interval dt , starting with the initial condition $|\phi_\eta(t)\rangle$. Accordingly, both motional states will still be approximately equal after this short evolution, allowing us to write

$$\langle \phi_\eta(t+dt)|\phi_i(t+dt)\rangle = \langle \phi_\eta(t)|\phi_i(t)\rangle + \frac{d}{dt} \langle \phi_\eta(t)|\phi_i(t)\rangle dt \quad (6.102)$$

$$= 1 + \frac{dR_{\eta i}(t)}{dt}. \quad (6.103)$$

Using this fact in applying the projection in (6.101) yields a product state once more:

$$|\tilde{\Psi}(t+dt)\rangle = \sum_i \left(1 + \frac{dr_{\eta i}(t)}{dt} \right) \left[\tilde{c}_i(t) - \frac{i}{\hbar} \sum_j H_{ij}(t) \tilde{c}_j(t) dt \right] |\chi_i\rangle \otimes |\phi_\eta(t+dt)\rangle \quad (6.104)$$

$$\Rightarrow |\tilde{\chi}(t+dt)\rangle = \sum_i \left(1 + \frac{dr_{\eta i}(t)}{dt} \right) \left[\tilde{c}_i(t) - \frac{i}{\hbar} \sum_j H_{ij}(t) \tilde{c}_j(t) dt \right] |\chi_i\rangle \quad (6.105)$$

which we can solve for $\frac{1}{dt} (|\tilde{\chi}(t+dt)\rangle - |\tilde{\chi}(t)\rangle)$ to obtain a differential equation for $|\tilde{\chi}(t)\rangle$:

$$\frac{d}{dt} |\tilde{\chi}(t)\rangle = \left[-\frac{i}{\hbar} \hat{H}_\eta(t) - \hat{\Gamma}_\eta(t) \right] |\tilde{\chi}(t)\rangle, \quad (6.106)$$

where $\hat{\Gamma}_\eta(t)$ is a diagonal operator in the $\{|\chi_i\rangle\}$ basis with diagonals

$$(\Gamma_\eta(t))_{ii} = \gamma_{\eta i}(t) = -\frac{dR_{\eta i}(t)}{dt} = -\frac{d}{dt} \langle \phi_\eta(t)|\phi_i(t)\rangle, \quad (6.107)$$

where we have defined $\gamma_{ij}(t) = -\frac{dR_{ij}(t)}{dt}$. The diagonals of $\hat{\Gamma}_\eta(t)$ in the adiabatic basis are *decoherence rates*, and cause exponential damping of all internal states other than the $|\chi_\eta\rangle$ state, with faster damping the faster the motional state corresponding to a given internal state diverges with $|\phi_\eta(t)\rangle$.

6.5.2 The quantum Zeno effect

But now we arrive at a problem. Because all the motional states are reset to be equal to $|\phi_\eta(t)\rangle$ at each timestep, the relative position and wavevector of the two wavepackets are zero at all times, and so our decoherence rates are:

$$\gamma_{ij} = - \left[\frac{dR_{ij}(t)}{dt} \right]_{\substack{k_{ij}=0 \\ v_{ij}=0}} \quad (6.108)$$

$$= - \left[\frac{\partial R_{ij}(t)}{\partial r_{ij}} \cdot \frac{dr_{ij}(t)}{dt} + \frac{\partial R_{ij}(t)}{\partial \mathbf{k}_{ij}} \cdot \frac{d\mathbf{k}_{ij}(t)}{dt} \right]_{\substack{k_{ij}=0 \\ v_{ij}=0}} \quad (6.109)$$

$$= \left[\left(\frac{r_{ij}}{4\sigma^2} + \frac{i}{2} \mathbf{k}_{ij} \right) R_{ij}(t) \cdot \frac{dr_{ij}(t)}{dt} + \left(\sigma^2 \mathbf{k}_{ij} + \frac{i}{2} \mathbf{r}_{ij} \right) R_{ij}(t) \cdot \frac{d\mathbf{k}_{ij}(t)}{dt} \right]_{\substack{k_{ij}=0 \\ v_{ij}=0}} \quad (6.110)$$

$$= 0. \quad (6.111)$$

Every decoherence rate is zero. What is going on? The answer is the quantum Zeno effect [124, 125], which is the name given to the fact that, in the limit of infinitely frequent measurements of whether quantum evolution has occurred, the back action of

the measurement has the effect of preventing the evolution from occurring at all. In our case, the fact that we are projecting constantly onto our assumed motional state causes the amplitude flows to the other motional states to be exactly zero. A sufficiently closely watched quantum pot never boils, and a sufficiently closely watched Schrödinger's cat never may never be poisoned [125]:

In view of the Zeno's paradox formulated above, should we conclude that the particle will never decay? Will the cat escape the cruel death awaiting it, against which it has no defense, provided its vital signs are constantly watched with loving care?

The appearance of the quantum Zeno effect ought to be a reminder that the assumption of infinitely frequent projective measurements is unphysical. In our case it certainly is—we have merely imagined a hypothetical measurement device collapsing our position states because we don't want to simulate them, not because any such measurement device actually exists. Nonetheless the internal states of the atom *do* decohere even in the absence of measurement (and there is eventual measurement when the atoms collide with other atoms or otherwise interact with anything in a position-dependent way), and we wish to model the approximate effect of this, if possible with a differential equation that does not require us to simulate all the quantum details of the motional degree of freedom.

Note that if the decoherence factor had the form of exponential decay:

$$R_{ij}^{\text{Markov}}(t) = e^{-\gamma_{ij} t}, \quad (6.112)$$

then the decoherence rate would be the constant γ_{ij} , and not zero. This is the case for Markovian decoherence [123], which is when the environment has no memory of its past interaction with the system. The memory in our case is due to the wavepackets accelerating away from each other, starting from zero relative position and velocity at each timestep—the relative position and velocity comprise a memory of the past interaction, which we are erasing each time we reproject.

The lack of an environmental memory is only ever approximately true, and no decoherence factors in nature have the form of a decaying exponential at all times. At small enough times the overlap between two states can only move away from unity quadratically owing to unitary evolution on account of the Schrödinger equation, guaranteeing an initial time derivative of zero for all physical decoherence factors. However in many systems of interest, the specifics of the interaction with the system are quickly forgotten by the environmental states, and the decoherence factor does approach a decaying exponential [126]. This is the case, for example, for spontaneous photon emission by atoms, which one can consider a measurement effect in which the electromagnetic field is being regularly measured by the environment in the photon number basis [106, 107, 109]. In the various quantum trajectories methods used to simulate atoms in the presence of spontaneous emission, if one assumes that the measurements are projective, one must simply assume that the frequent measurements take place at large enough timescales that the decoherence factor is in the exponential decay regime, provided that this is much smaller than other dynamical timescales, which is true for spontaneous emission [109]. The measurement interval assumed “should be large enough to allow the photons to get away from the atom” [108]. This can be recast as a continuous *weak* measurement, rather than infrequent projective measurements [109, 127], which is more physically realistic than the assumption of projective measurements at somewhat arbitrarily chosen intervals, but results in the same differential equations.

A decoherence factor that has the form of a decaying exponential implies that any chosen measurement interval results in the same fractional reduction in state amplitudes, since the differential equation $\frac{d\tilde{c}_i(t)}{dt} = -\gamma_i \tilde{c}_i(t)$ has the solution $\tilde{c}(t) = e^{-\gamma_i t} \tilde{c}(0)$, that is,

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repeated consideration of only the first part of the decoherence factor ends up tracing out the whole decoherence curve over time.

How can we replicate something like this for our separating wavepackets? I have two approaches. The first, described in Section 6.5.3, is to gloss over as many of the details of the wavepacket separation as possible and replace the decoherence factor with an exponential one, describing the wavepackets separating on approximately the correct timescale. This is crude, but better than no decoherence at all. The second, described in Section 6.5.4, is to introduce a minimal memory of the separation of wavepackets. In this approach, one (but only one) trajectory is simulated for the motional state corresponding to each internal state. Whenever there is population transfer from the main trajectory to another state, the trajectory corresponding to the other state is replaced with a weighted average of its existing position and velocity with the position and (adjusted for energy conservation) velocity of the main trajectory. These *auxiliary trajectories* are used to compute a dynamic decoherence rate that traces out a more correct decoherence curve as the wavepackets separate.

6.5.3 Approximate Markovian decoherence

A crude way to include decoherence is just to approximate $R_{\eta i}(t)$ as a decaying exponential with roughly the right decay constant. This is crude because $R_{\eta i}(t)$ does not look much like a decaying exponential (see Figure 6.5). In the limit of large t , for the case constant acceleration as in (6.96) its functional form is e^{-t^4} , not the exponential decay required to treat the decoherence as Markovian at any timescale.

To nonetheless find an approximate Markovian decoherence rate, we first construct a “time ignorant” version $\tilde{R}_{ij}(t)$ of the decoherence factor $R_{ij}(t)$ given in (6.96) that answers the question “What is the expected value of $R_{ij}(t)$ at all times $t > 0$ if I don’t know how long before $t = 0$ the two wavepackets began separating?”

We can then take the derivative of this time-ignorant decoherence factor at $t = 0$ to use as an approximate Markovian decoherence rate $\gamma_{ij}^{\text{Markov}}$. Whilst extremely approximate, this method of including decoherence is nonetheless an improvement over the Ehrenfest method (which has no decoherence), and over Tully’s original fewest-hops surface hopping method [93], which also lacked any decoherence. Others [128–130] have since developed various methods to induce approximate damping of states to achieve a similar outcome, including the approximation of an exponential decoherence factor [128], though my development of this method was independent.

Proceeding, we define the time-ignorant decoherence factor $\tilde{R}_{ij}(t)$ as the average of all decoherence factors one would obtain if the two wavepackets began separating at some point in time before $t = 0$:

$$\tilde{R}_{ij}(t) = A \int_{-\infty}^0 \langle \phi_i(t-t') | \phi_j(t-t') \rangle dt' \quad (6.113)$$

$$= A \int_{-\infty}^0 R_{ij}(t-t') dt' \quad (6.114)$$

where A is a normalisation constant such that $\tilde{R}_{ij}(0) = 0$, and where we take that $|\phi_i(0)\rangle = |\phi_j(0)\rangle$ with each thereafter evolving according to the classical motion of their centre of mass with constant relative acceleration as in (6.91) and (6.92) such that the $R_{ij}(t)$ above has the form (6.96).

Our time-ignorant decoherence rate is then the (negative of the) derivative of $\tilde{R}_{ij}(t)$

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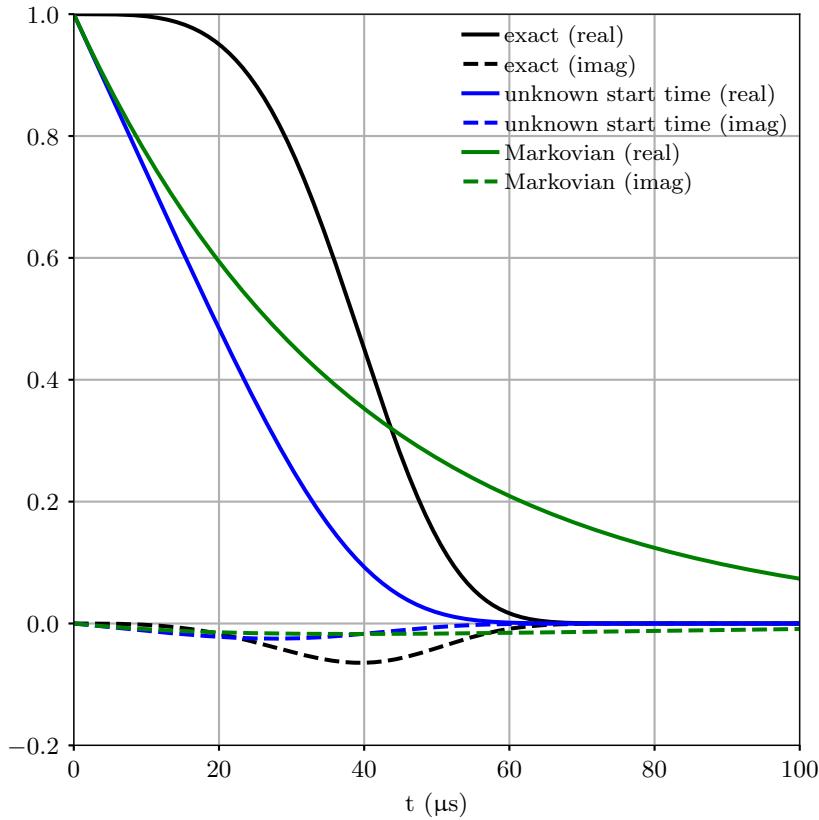


Figure 6.5: Example decoherence factor as a function of time for an experimentally realistic parameters. Plotted in black is the exact decoherence factor $R_{ij}(t)$ (equation (6.96)) between two adjacent ($\Delta m_F = \pm 1$) Zeeman sublevels of the $F = 1$ groundstate of ^{87}Rb , assuming constant relative acceleration due to a magnetic field gradient of 250 G cm^{-1} and a Gaussian wavepacket width $\sigma = \lambda_{\text{th}} = 2.6 \mu\text{m}$, corresponding to the thermal wavelength $\lambda_{\text{th}} = h/\sqrt{2\pi mk_B T}$ at $T = 53 \mu\text{K}$. Shown in blue is the ‘time-ignorant’ decoherence factor described in text. In green is the Markovian approximation to the exact decoherence factor, obtained by extracting a decay constant from the gradient of the time-ignorant decoherence factor at $t = 0$. As can be seen, whilst no decaying exponential is a particularly good fit to the exact decoherence factor, our Markovian approximation is as good as can be expected, decaying to zero over approximately the correct timescale.

at $t = 0$:

$$\gamma_{ij}^{\text{Markov}} = - \frac{\left[\frac{d}{dt} \int_{-\infty}^0 R_{ij}(t-t') dt' \right]_{t=0}}{\left[\int_{-\infty}^0 R_{ij}(t-t') dt' \right]_{t=0}}. \quad (6.115)$$

Moving the derivative inside the integral, noting that $\frac{dR(t-t')}{dt} = \frac{dR(t-t')}{d(t-t')}$ and setting

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$t = 0$, we get:

$$\gamma_{ij}^{\text{Markov}} = -\frac{\int_{-\infty}^0 R'_{ij}(-t') dt'}{\int_{-\infty}^0 R_{ij}(-t') dt'} \quad (6.116)$$

$$= -\frac{\int_0^\infty R'_{ij}(t') dt'}{\int_0^\infty R_{ij}(t') dt'} \quad (6.117)$$

where R'_{ij} is the derivative of R_{ij} with respect to its argument. By the fundamental theorem of calculus the numerator is -1 , since $R_{ij}(t)$ decreases from unity at $t = 0$ to zero as t goes to infinity, leaving us with:

$$\frac{1}{\gamma_{ij}^{\text{Markov}}} = \int_0^\infty R_{ij}(t') dt' \quad (6.118)$$

$$= \int_0^\infty \exp \left[-\left(\frac{1}{8\sigma^2} |\mathbf{r}_{ij}|^2 + \frac{i}{2} \mathbf{r}_{ij}(t') \cdot \mathbf{k}_{ij}(t') + \frac{\sigma^2}{2} |\mathbf{k}_{ij}(t')|^2 \right) \right] dt' \quad (6.119)$$

The expression (6.96) requires a relative acceleration, for which we use the relative acceleration between the pair of adiabatic potentials at the current moment in time and current position of the atom during a simulation:

$$\mathbf{a}_{ij}(\mathbf{r}, t) \approx -\frac{1}{m} (\nabla V_i(\mathbf{r}, t) - \nabla V_j(\mathbf{r}, t)). \quad (6.120)$$

This now gives a decoherence factor depending on position and time:

$$\frac{1}{\gamma_{ij}^{\text{Markov}}(\mathbf{r}, t)} = \int_0^\infty \exp \left[-\frac{|\mathbf{a}_{ij}(\mathbf{r}, t)|^2}{2} \left(\frac{1}{16\sigma^2} t'^4 + i \frac{m}{2\hbar} t'^3 + \frac{m^2\sigma^2}{\hbar^2} t'^2 \right) \right] dt' \quad (6.121)$$

In order to obtain an approximate analytic expression for this integral, we consider two limiting cases and then stitch them together in the intermediate regime. In the limit of small wavepackets, σ is small and thus the first term in the exponent in (6.118) is largest, and the third term is smallest. In this regime, which describes when positional separation (as opposed to separation in k -space) dominates the decoherence, we'll neglect the third term in the exponent and treat the second term as small relative to the first. This gives us:

$$\frac{1}{\gamma_{ij}^{\text{Markov}}(\mathbf{r}, t)} \approx \int_0^\infty \exp \left[-\frac{|\mathbf{a}_{ij}(\mathbf{r}, t)|^2}{2} \left(\frac{1}{16\sigma^2} t'^4 + i \frac{m}{2\hbar} t'^3 \right) \right] dt' \quad (6.122)$$

$$\approx \int_0^\infty \exp \left[-\frac{|\mathbf{a}_{ij}(\mathbf{r}, t)|^2}{32\sigma^2} t'^4 \right] \left(1 - i \frac{m|\mathbf{a}_{ij}(\mathbf{r}, t)|^2}{4\hbar} t'^3 \right) dt' \quad (6.123)$$

$$= 2^{\frac{5}{4}} \Gamma(\frac{5}{4}) \sqrt{\frac{\sigma}{|\mathbf{a}_{ij}(\mathbf{r}, t)|}} - 2i \frac{m\sigma^2}{\hbar}, \quad (6.124)$$

where we used a first-order Taylor expansion of an exponential in (6.123). We similarly use a first-order expansion $(x + \varepsilon)^{-1} \approx x^{-1} - \varepsilon x^{-2}$ to take the reciprocal of (6.124) (since the second term is much smaller than the first¹⁷), and arrive at

$$\gamma_{ij}^{\text{Markov}}(\mathbf{r}, t) \approx \frac{1}{2^{\frac{5}{4}} \Gamma(\frac{5}{4})} \sqrt{\frac{|\mathbf{a}_{ij}(\mathbf{r}, t)|}{\sigma}} + \frac{i}{2^{\frac{3}{2}} \Gamma(\frac{5}{4})^2} \frac{m\sigma|\mathbf{a}_{ij}(\mathbf{r}, t)|}{\hbar}. \quad (6.125)$$

¹⁷This isn't necessary in order to obtain a simple expression for $\gamma_{ij}(\text{pos})$ —the reciprocal without this approximation is equally simple—but it leaves us with power laws for the real and imaginary parts of $\gamma_{ij}(\text{pos})$, which are easier to stitch together with those from the large σ regime. It also ensures we don't divide by zero when $\mathbf{a}_{ij}(\mathbf{r}, t)$ is zero.

Similarly for the large σ regime, we neglect the first term in the exponent of (6.118) and consider the second term small relative to the third. This is the regime in which the decrease in overlap of the two wavepackets is dominated by their separation in velocity space. Following the same process as above gives:

$$\frac{1}{\gamma_{ij}^{\text{Markov}}(\mathbf{r}, t)} \approx \int_0^\infty \exp \left[-\frac{|\mathbf{a}_{ij}(\mathbf{r}, t)|^2}{2} \left(i \frac{m}{2\hbar} t'^3 + \frac{m^2 \sigma^2}{\hbar^2} t'^2 \right) \right] dt' \quad (6.126)$$

$$\approx \int_0^\infty \left(1 - i \frac{m |\mathbf{a}_{ij}(\mathbf{r}, t)|^2}{4\hbar} t'^3 \right) \exp \left[-\frac{m^2 \sigma^2 |\mathbf{a}_{ij}(\mathbf{r}, t)|^2}{2\hbar^2} t'^2 \right] dt \quad (6.127)$$

$$= \sqrt{\frac{\pi}{2}} \frac{\hbar}{m\sigma |\mathbf{a}_{ij}(\mathbf{r}, t)|} - i \frac{\hbar^3}{2m^3 \sigma^4 |\mathbf{a}_{ij}(\mathbf{r}, t)|^2} \quad (6.128)$$

$$\Rightarrow \gamma_{ij}^{\text{Markov}}(\mathbf{r}, t) \approx \sqrt{\frac{2}{\pi}} \frac{m\sigma |\mathbf{a}_{ij}(\mathbf{r}, t)|}{\hbar} + i \frac{\hbar}{\pi m \sigma^2} \quad (6.129)$$

Equations (6.125) and (6.129) are our final expressions for the Markovian decoherence rate in the limit of small and large wavepackets respectively. Adding their real parts in quadrature and adding the reciprocals of their imaginary parts then provides a reasonable approximation for $\gamma_{ij}^{\text{Markov}}(\mathbf{r}, t)$ over all wavepacket sizes:

$$\begin{aligned} \gamma_{ij}^{\text{Markov}}(\mathbf{r}, t) &\approx \left[\text{Re}(\gamma_{ij}^{\text{Markov}}(\mathbf{r}, t))^2 + \text{Re}(\gamma_{ij}^{\text{Markov}}(\mathbf{r}, t))^2 \right]^{\frac{1}{2}} \\ &+ i \left[\text{Im}(\gamma_{ij}^{\text{Markov}}(\mathbf{r}, t))^{-1} + \text{Im}(\gamma_{ij}^{\text{Markov}}(\mathbf{r}, t))^{-1} \right]^{-1}. \end{aligned} \quad (6.130)$$

We now have an approximate analytic expression for a Markovian decoherence rate between two Gaussians that is computationally inexpensive to evaluate for each atom in an ensemble at every timestep of a differential equation. An example showing the accuracy of (6.130), compared to the exact expression (6.121) for $\gamma_{ij}^{\text{Markov}}$ over a range of wavepacket sizes is shown in Figure 6.6.

6.5.4 Decoherence with mean auxiliary trajectories

A more accurate approach to including decoherence is to compute a decoherence factor at each timestep of the simulation by modelling the wavepackets as they accelerate away from each other, retaining in the model the position and velocity of each wavepacket. This inclusion of additional positions and velocities in the model comprises a minimal memory of the past interaction between the system and environment—that is, the internal and motional degrees of freedom of the atom.

In quantum mechanics there is not, in general, a single Gaussian wavepacket for each state, and even if the wavefunction corresponding to the originating state is a Gaussian wavepacket, transitions to other states at different times can produce an arbitrary superposition of Gaussian wavepackets in each other state. If we were to simulate this arbitrary superposition for all components of the multi-component wavefunction, we would not be saving any computational power at all, as we would have reverted to a fully quantum description of the motional degree of freedom, which is not at all our intention. Therefore we draw the line at modelling a single trajectory for each state, in order to maintain simplicity whilst including at least some dynamics of the system–environment interaction.

Below I will talk about how to convert a decoherence factor at each point in time into a decoherence *rate* at each point in time, that can be included in a differential equation for the state vector of the atom’s internal state. Then I will introduce the averaging scheme I

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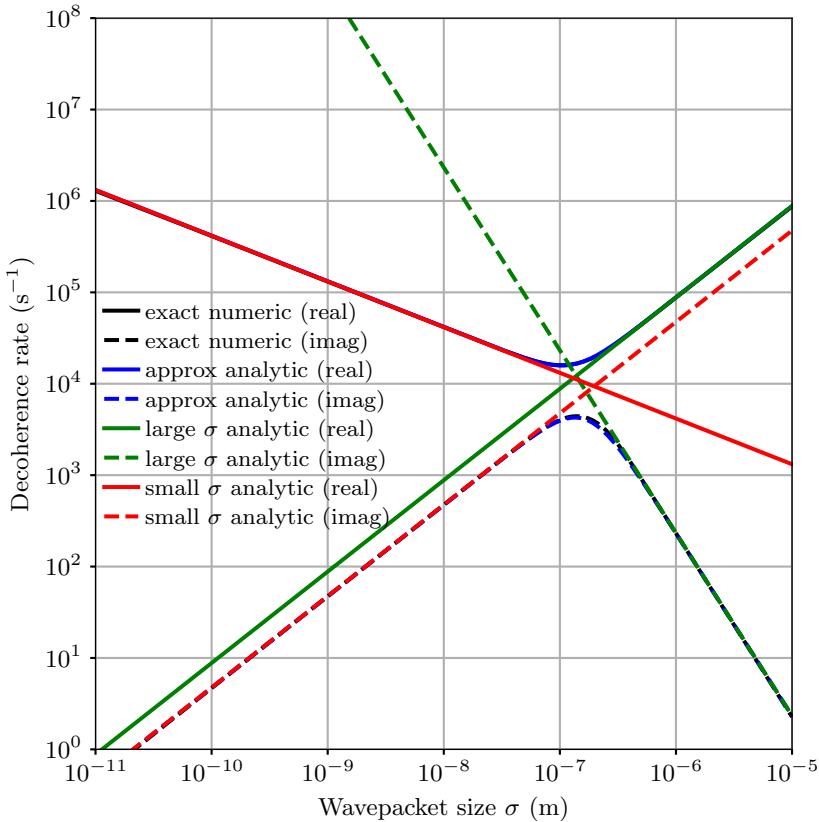


Figure 6.6: Comparison of approximate analytic Markovian decoherence rate with exact Markovian decoherence rate between adjacent ($\Delta m_F = \pm 1$) Zeeman sublevels of the $F = 1$ groundstate of a ^{87}Rb atom in a 250 G cm^{-1} magnetic field gradient. Real (solid) and imaginary (dashed) parts of the small wavepacket (red) and large wavepacket (green) limits are shown, their stitching together in the intermediate wavepacket regime (blue) and the exact expression they approximate (black, not easily visible due to being similar to the blue curves). There is good agreement over all wavepacket sizes.

am using in order to turn what would be many trajectories for each state into one average trajectory.

This use of ‘auxiliary trajectories’ for computing decoherence is common in the surface-hopping literature. There is a range of methods called ‘multiple spawning’ methods, in which additional auxiliary trajectories regularly branch off the main trajectory (the one corresponding to the hidden variable), and are all considered in order to compute the overlap with the main wavepacket and hence the decoherence factor [131–133]. These methods generally contain a parameter controlling the likelihood of new trajectories branching off, and trajectories can be discarded once they recede far enough from the main trajectory. If the parameter is tuned far enough, there can be so many auxiliary trajectories that an extremely accurate decoherence rate can be computed, though the computational cost in this case approaches that of a fully quantum simulation. A method by Shenvi et al. [97], like mine, tracks only one auxiliary trajectory per internal atomic state, probabilistically replacing them with newly spawned trajectories rather than simulating multiple auxiliary trajectories per state. However, the trajectories are spawned at specific moments—when the non-adiabatic coupling strength reaches a local maximum—leading to different be-

haviour in different locations in space. For example, if an atom were orbiting a magnetic field zero in a magnetic trap, this method would only spawn new trajectories at the orbit's point of closest approach, and decoherence rates would be inaccurate at other points in the orbit, even if the non-adiabatic coupling were still high throughout the orbit. My method has the potential to improve upon this due to the fact that it is in a sense spawning new trajectories all the time—rather than only at specific points in space as the method by Shenvi et al.—and avoids the problem of an exponential proliferation of trajectories by only retaining a single, more-or-less representative trajectory for each state rather than many.

Decoherence rate in the absence of projective measurement

Here we generalise our conception of the decoherence rate by removing the assumption that motional states are “reset” at each timestep. This resetting, which we assumed in Section 6.5.1, set each motional state to be equal to $|\phi_\eta(t)\rangle$ at each timestep, and this led to all decoherence rates being zero due to the quantum Zeno effect (Section 6.5.2). What is the reduction in the state amplitudes at each timestep if the wavepackets are not reset, but are left as arbitrary states for the time being? We know that if we were to perform a projective measurement at time t , the i^{th} state amplitude would be reduced by a factor of $R_{\eta i}(t)$ compared to the original state amplitude $c_i(t)$:

$$\tilde{c}_i(t) = R_{\eta i}(t)c_i(t). \quad (6.131)$$

This implies:

$$\frac{d\tilde{c}_i(t)}{dt} = \frac{dR_{\eta i}(t)}{dt}c_i(t) \quad (6.132)$$

$$= \frac{1}{R_{\eta i}(t)} \frac{dR_{\eta i}(t)}{dt} \tilde{c}_i(t). \quad (6.133)$$

and so we see that a decoherence rate that does not reset the environment states is the (negative of the) logarithmic derivative of the decoherence factor at any time, rather than just its derivative at $t = 0$. This case encompasses the earlier case in which $R_{\eta i}(t) = 1$ at all times, in which case the logarithmic and ordinary derivatives are the same. Thus, given the decoherence factor between two states, along with its time derivative, we can compute a decoherence rate

$$\gamma_{ij}^{\text{MM}}(t) = -\frac{1}{R_{ij}(t)} \frac{dR_{ij}(t)}{dt}, \quad (6.134)$$

such that in the absence of Hamiltonian evolution

$$\frac{d\tilde{c}_i(t)}{dt} = -\gamma_{\eta i}^{\text{MM}}(t) \tilde{c}_i(t), \quad (6.135)$$

where “MM” stands for “minimal-memory”.

The (negative) logarithmic derivative of the decoherence factor for our fixed-width Gaussian wavepackets (6.93) gives the decoherence rates

$$\gamma_{ij}^{\text{MM}}(t) = \left(\frac{\mathbf{r}_{ij}}{4\sigma^2} + \frac{i}{2}\mathbf{k}_{ij} \right) \cdot \frac{d\mathbf{r}_{ij}(t)}{dt} + \left(\sigma^2 \mathbf{k}_{ij} + \frac{i}{2} \mathbf{r}_{ij} \right) \cdot \frac{d\mathbf{k}_{ij}(t)}{dt}, \quad (6.136)$$

which for a given relative acceleration \mathbf{a}_{ij} between the states, based on the potential gradient at each modelled position, and written in terms of relative velocity instead of

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wavenumber, is

$$\gamma_{ij}^{\text{MM}}(t) = \frac{1}{4\sigma^2} \mathbf{r}_{ij}(t) \cdot \mathbf{v}_{ij}(t) + \frac{\sigma^2 m^2}{\hbar^2} \mathbf{v}_{ij}(t) \cdot \mathbf{a}_{ij}(t) + i \frac{m}{2\hbar} (|\mathbf{v}_{ij}(t)|^2 + \mathbf{r}_{ij}(t) \cdot \mathbf{a}_{ij}(t)), \quad (6.137)$$

where $\mathbf{r}_{ij}(t) = \mathbf{r}_i(t) - \mathbf{r}_j(t)$ and $\mathbf{v}_{ij}(t) = \mathbf{v}_i(t) - \mathbf{v}_j(t)$ are the relative position and velocity of the two states, and $\mathbf{a}_{ij}(t)$ is their relative acceleration:

$$\mathbf{a}_{ij}(t) = -\frac{1}{m} (\nabla V_i(\mathbf{r}_i, t) - \nabla V_j(\mathbf{r}_j, t)). \quad (6.138)$$

Note this is slightly different to the relative acceleration used in the Markovian decoherence calculation in Section 6.5.3. Here, we actually have (approximate) positions $\mathbf{r}_i(t)$ and $\mathbf{r}_j(t)$ for both states' wavepackets, and so we can evaluate the potential gradient separately at the two positions rather than having to use the position of the state currently selected by the hidden variable for both.

Equation (6.5.4) is our minimal-memory decoherence rate, and requires as input the centre-of-mass position and velocity of each Gaussian wavepacket as input. Let's now move on to how these trajectories are computed.

From many trajectories, one

As mentioned, this method tracks only one trajectory per state. However, it continuously considers the spawning of new trajectories, and rather than actually tracking multiple trajectories per state or probabilistically replacing the existing trajectories with the new ones, it simply averages the trajectories together, weighted by their quantum probabilities. At the end of each timestep, positions and velocities (other than those of the main trajectory) are updated according to a weighted sum of their present values and those of the main trajectory:

$$\mathbf{r}_{i \neq \eta}(t') \leftarrow Q_{i\eta}(t', t) \mathbf{r}_\eta(t') + (1 - Q_{i\eta}(t', t)) \mathbf{r}_i(t'), \quad (6.139)$$

$$\begin{aligned} \mathbf{v}_{i \neq \eta}(t') &\leftarrow Q_{i\eta}^{\text{time}}(t', t) \mathbf{v}_\eta(t') + \tilde{Q}_{i\eta}^{\text{space}}(t', t) (\mathbf{v}_\eta(t') + \Delta \mathbf{v}_{i\eta}(t')) \\ &+ (1 - Q_{i\eta}^{\text{time}}(t', t) - \tilde{Q}_{i\eta}^{\text{space}}(t', t)) \mathbf{v}_i(t'), \end{aligned} \quad (6.140)$$

¹⁸Note that $|c_i(t')|^2$ is the i^{th} population at time t' due to unitary evolution only over the time interval, without taking into account decoherence, since all population flows are computed from unitary evolution only and decoherence is treated separately.

where $Q_{i\eta}(t', t) = P_{i\eta}(t', t)/|c_i(t')|^2$ is the fraction of the population¹⁸ of state i at time t' that flowed to it from state η in the interval t to t' , with $Q_{i\eta}^{\text{time}}(t', t)$ and $\tilde{Q}_{i\eta}^{\text{space}}(t', t)$ defined similarly in terms of $P_{i\eta}^{\text{time}}$ and $\tilde{P}_{i\eta}^{\text{time}}$ as defined in Section 6.4.3, and where $\Delta \mathbf{v}_{i\eta}(t')$ is the required velocity correction for a transition from state η to state i state at time t' , as discussed in Section 6.4.3;

$$\Delta \mathbf{v}_{i\eta}(\mathbf{r}, t') = \left[\text{sgn}(\mathbf{v} \cdot \hat{\mathbf{d}}_{i\eta}) \sqrt{(\mathbf{v} \cdot \hat{\mathbf{d}}_{i\eta})^2 + \frac{2}{m} (V_i(\mathbf{r}, t') - V_\eta(\mathbf{r}, t'))} - (\mathbf{v} \cdot \hat{\mathbf{d}}_{i\eta}), \right] \hat{\mathbf{d}}_{i\eta}, \quad (6.141)$$

where $\mathbf{r} = \mathbf{r}_\eta(t')$ and $\mathbf{v} = \mathbf{v}_\eta(t')$ are the position and velocity of the main trajectory at time t' , and the non-adiabatic coupling unit vector $\hat{\mathbf{d}}_{i\eta} = \hat{\mathbf{d}}_{i\eta}(\mathbf{r}_\eta(t'))$ is evaluated at the position along the main trajectory at time t' .

The above scheme is constructed to track a mean trajectory for each state, where ‘mean’ is defined as a weighted sum of positions and velocities with those of the main trajectory whenever transitions from the latter occur, taking into account velocity jumps and frustrated transitions, with the energy conservation calculation for the velocity jumps based on the potentials at the location of the main trajectory.

6.5.5 What are we “following” exactly?

Now that we’ve worked out a way to approximate the result of projecting the multi-component wavefunction onto a specific motional state (a Gaussian wavepacket) without resetting all motional states (also Gaussian wavepackets) to be equal to the one we are projecting onto, why not take the process further? We can use the same method to project onto any wavefunction we like, in order to inspect what the multi-component wavefunction looks like projected onto the wavefunction in question. This does not have to imply anything about the form of the motional states, which remain Gaussians at all time. It just defines what we ‘see’ at each timestep of our simulation.

For example, if we projected onto a Dirac delta, that would tell us what the multi-component wavefunction looks like at a specific point in space, say, the centre-of-mass position of the wavepacket corresponding to the hidden variable.

There are reasons why this is appealing. Firstly, consider the projected Hamiltonian $\hat{H}_\eta(t)$ first introduced in Section 6.5.1. Since we do not want to have to actually construct Gaussian wavepackets in order to compute the integral which defines the projected Hamiltonian, our only recourse is to approximate the projected Hamiltonian as the value of the full Hamiltonian at the centre-of-mass position of the main trajectory’s Gaussian wavepacket:

$$\langle \chi_i | \hat{H}_\eta(t) | \chi_j \rangle = \langle \chi_i \phi_\eta(t) | \hat{H}(t) | \chi_j \phi_\eta(t) \rangle \quad (6.142)$$

$$\approx \langle \chi_i \mathbf{r}_\eta(t) | \hat{H}(t) | \chi_j \mathbf{r}_\eta(t) \rangle \quad (6.143)$$

$$\equiv H_{ij}(\mathbf{r}_\eta, t), \quad (6.144)$$

which can be considered a small-wavepacket approximation, as mentioned in Section 6.1 in the context of the Ehrenfest semiclassical model. If an arbitrary position is used in place of \mathbf{r}_η , this results in an operator valued function of space $\hat{H}(\mathbf{r}, t)$, which operates only on the internal degrees of freedom of the atom.

This is the operator we actually use in the algorithm, and with it, there are no longer any Gaussian wavepackets to compute integrals over or anything else—all details of the wavepackets are encapsulated and parametrised by the approximations and analytics in this and the preceding sections, leaving us to focus on the classical dynamics of them atoms’ centre-of-mass motion and the quantum evolution of their internal states according to $\hat{H}(\mathbf{r}, t)$ evaluated at the classical positions (representing the centre-of-mass position of the wavepackets) of the atoms.

Although this can be considered a small wavepacket approximation, we cannot use arbitrarily small wavepacket sizes in our decoherence calculations, as the decoherence rates (under both the Markovian and auxiliary trajectories approach) do not converge to a constant as the wavepacket sizes decrease—rather they become arbitrarily large. Very high spin decoherence rates have the potential to prevent spin flips altogether, also a result of the quantum Zeno effect, since high decoherence rates imply strong measurement, and strong measurement of an observable prevents evolution of that variable.¹⁹ So there is a possible inconsistency here: can we make the wavepackets small or not?

On the other hand, note what happens in Figure 6.6 when wavepackets get large: the decoherence rate becomes large as well. Our expression for the non-Markovian decoherence rate (6.137). In both cases this is decoherence due to separation in velocity space. While this is the correct result, given our assumptions, it causes problems when combined with the above way of approximating the projected Hamiltonian. Essentially, rapid velocity-induced decoherence occurs when the wavepackets are large because spatially large Gaussians are small in velocity space, and hence do not need to accelerate much to be non-overlapping any more in velocity space. But do we really expect our wavepackets to be minimal uncertainty wavepackets? Certainly not when they are large compared to the structure of a spatially varying Hamiltonian. A real wavepacket, whilst

¹⁹I have sometimes wondered if there is a way to harness this to prevent Majorana losses in evaporation to BEC, by measuring very strongly whether they have occurred.

undergoing a transition to another state, does not do so everywhere in space the same, as in a single-mode approximation, but our formulation so far assumes it does. In the limit of a very large spin- $\frac{1}{2}$ wavepacket undergoing a complete majorana spin flip due to a magnetic field zero for example, would appear, as its centre-of-mass passed over the zero, as approximately two half-Gaussians, one in each state. I don't know what the half-Gaussians' velocity distributions would look like, but I doubt they are narrow enough to lead to decoherence rates as high as the expressions we've derived so far suggest. Narrow velocity distributions are therefore inconsistent with the other assumptions of our model and any consequences of rapid velocity separation in the model should be viewed with scepticism.

Another argument is that, if two wavepackets are co-located in space but not in velocity space, the amplitude of one of them as seen from the centre-of-mass position of the other is not zero. It's the fact that their relative phase varies so much over the extent of the wavepackets that causes their projection onto each other to be small when their velocities are very different.²⁰ So it's only after this averaging process—integrating over all space—that the wavepackets look like they are small from each other's perspective. One of the common themes of this method as a whole is that we are often happy with *representative* results rather than average ones, drawing results from approximately correct statistical distributions, and only observing the complete distribution when we analyse a large sample of results. Continuing in this vein, why not consider the value of the multi-component wavefunction along a specific trajectory in space, rather than its projection onto a Gaussian? The projection of a multi-component wavefunction onto a Gaussian isn't an any more useful or intuitive quantity than tracing a curve through spacetime and asking "what is the wavefunction here?" This way instead of a simulation telling us "There was, on average, no population transfer to such-and-such state", it would give a (more likely to be useful in my opinion) answer "there was a lot of population transfer to this state, but the phase is essentially random".

The final argument is empirical: I've tried simulating the model with the Gaussian projections, and it doesn't work well unless wavepackets are small. The decoherence rates are too high, presumably for the speculative reasons discussed above. This damps the population of states other than that selected by the hidden variable too quickly, such that the simulation says that the population of a state is zero even though according to the Schrödinger wave equation it should not be (it merely has an inhomogeneous phase). Simply deleting the terms caused by velocity separation from the decoherence rate expressions improves the results somewhat. But if project onto a Dirac delta then there are no velocity-separation terms to argue for the deletion of, and I obtain modestly better results in any case (as shown in Section ??).

Shenvi et al. [97] also seem to find the velocity-separation term of their decoherence factor troublesome, and opt to remove it by gauge transforming the wavepackets to have the same velocity. They argue that because at least a small initial velocity difference between wavepackets is required by energy conservation (i.e. the velocity jumps from Section 6.4.3), that it is 'unfair' to have an initial sudden drop in coherence because of this. I agree that it is strange, but suspect that it is because we are insisting on projecting onto Gaussians. After all, velocity jumps or not, unless there is a mistake in our maths these are the actual decoherence factors between two Gaussians. If we don't like the answer, perhaps the projection of the total wavefunction onto a Gaussian isn't the question we meant to ask. Perhaps we should be asking simply what value a multi-component wavefunction has at a specific point in space. Shenvi et al. gauge transform away the entire velocity difference between the wavepackets, not only the part of it mandated by energy conservation, leading me to suspect that some of the same concerns I outlined above occurred to them too, or that they simply observed, like me, that the velocity separation was responsible for what appeared to be unphysically large decoherence rates in their results.

In light of all this, I now redefine “follow” to mean “look at the wavefunction’s value at the centre-of-mass position of the classical trajectory corresponding to the state being followed”, and project onto a Dirac delta instead, resulting in a decoherence factor

$$R_{ij}^{\text{DD}}(t) = \langle \mathbf{r}_i(t) | \phi_j(t) \rangle \quad (6.145)$$

where $|\phi_j(t)\rangle$ is one of our usual Gaussian motional states and DD stands for “Dirac delta”. Following identical arguments to the previous two sections, we obtain decoherence rates for the Markovian case as

$$\gamma_{ij}^{\text{Markov DD}}(\mathbf{r}, t) \approx \frac{1}{2\Gamma(\frac{5}{4})} \sqrt{\frac{|\mathbf{a}_{ij}(\mathbf{r}, t)|}{\sigma}} + i \frac{1}{2\Gamma(\frac{5}{4})^2} \frac{m\sigma |\mathbf{a}_{ij}(\mathbf{r}, t)|}{\hbar}, \quad (6.146)$$

which is identical to the position-only approximation (6.125) up to a factor of $2^{\frac{1}{4}}$, and for minimal-memory non-Markovian case as

$$\gamma_{ij}^{\text{MMDD}}(t) = \frac{1}{2\sigma^2} \mathbf{r}_{ij}(t) \cdot \mathbf{v}_{ij}(t) - i \frac{m}{\hbar} \mathbf{r}_{ij}(t) \cdot \mathbf{a}_{ij}(t), \quad (6.147)$$

with all symbols as defined in Section 6.5.3 and Section 6.5.4, respectively.

Now that we are projecting onto a Dirac delta instead of a Gaussian, calculation of the projected Hamiltonian as simply the value of the full Hamiltonian at a specific point in space is no longer a contradiction, as it is exactly what we would get if we projected the full Hamiltonian onto a Dirac delta instead of a Gaussian.

6.6 Algorithms

Now we have presented all the pieces from which the two versions of the hidden-variable semiclassical method are constructed. In this section I present step-by-step instructions for implementing the two versions of the method.

6.6.1 Markovian hidden-variable semiclassical method

This is the simplest version of my method, using crude Markovian decoherence as discussed in Section 6.5.3, and simulating only one trajectory corresponding to the centre-of-mass motion of the state corresponding to the hidden variable at each moment in time.

State variables

The Markovian variant of the hidden-variable semiclassical method has the following state variables describing the state of each simulated atom at each moment of time:

- The internal state vector of the atom $|\tilde{\chi}(t)\rangle$, in any chosen basis (though the local eigenbasis is likely to be computationally simplest to perform calculations in).
- The position $\mathbf{r}(t)$ and velocity $\mathbf{v}(t)$ of the atom.
- The hidden variable $\eta(t)$, equal to an integer index specifying one of the eigenstates $|\chi_\eta(\mathbf{r}, t)\rangle$ of the Hamiltonian $\hat{H}(\mathbf{r}, t)$ atom in the local energy eigenbasis.

Initial conditions

Strictly, the hidden-variable semiclassical method is intended to simulate the classical centre-of-mass motion of thermal wavelength sized Gaussian wavepackets. Therefore the initial conditions for the positions and velocities of the atoms in an ensemble are arbitrary, classical initial conditions, and so one might draw them, along with the internal states, randomly from a Boltzmann distribution for a given confining potential and temperature, such that the initial internal state of each atom is an eigenstate $|\chi_i(\mathbf{r}, t = 0)\rangle$. In this case the initial value $\eta(t = 0)$ of the hidden variable for each atom should correspond to its chosen eigenstate.

However, one might want to convert an initial *wavefunction* into an ensemble of positions and velocities, in which case one can draw positions and momenta from the quantum probability distributions $\langle\psi(t = 0)|\hat{\mathbf{r}}|\psi(t = 0)\rangle$ and $\langle\psi(t = 0)|\hat{\mathbf{p}}|\psi(t = 0)\rangle$ for an initial motional state vector $|\psi\rangle$. Similarly, one can set the initial internal state vector of each atom to an arbitrary desired internal state vector appropriate for the problem at hand, in which case the initial values of the hidden variable for each atom in an ensemble being simulated should be chosen randomly with probability equal to the internal state populations $|\langle\chi_i(\mathbf{r}, t = 0)|\tilde{\chi}(\mathbf{r}, t = 0)\rangle|^2$ in the local basis at the position of each atom. More complex initial conditions are possible, there is no reason why one can't draw positions, velocities and internal states from arbitrary joint probability distributions, or use hand-crafted initial states, provided that the hidden variables are chosen randomly at the end with probabilities equal to the internal state populations in the local basis of each atom so as to be consistent with the Born rule.

Evolution

Evolving the Markovian hidden-variable semiclassical method involves three coupled differential equations, a stochastic evolution rule, and some additional steps to be performed in between numerical integration steps. Because the modifications in between steps can be discontinuous, multi-step integration schemes that retain information from previous integration steps should be avoided. One step of the Markovian hidden variable semiclassical method comprises the following steps to evolve the state variables from time t to $t' = t + \Delta t$, where Δt is small compared to dynamical timescales of the problem.

1. Evolve the internal state vector and motional state variables of each atom from time t to t' according to the following coupled differential equations, using one or more steps of your favourite non-multi-step numerical integration method:

$$\frac{d}{dt} |\tilde{\chi}(t)\rangle = -\frac{i}{\hbar} \hat{H}(\mathbf{r}, t) |\tilde{\chi}(t)\rangle - \hat{f}_\eta^{\text{Markov DD}}(\mathbf{r}, t) |\tilde{\chi}(t)\rangle, \quad (6.148)$$

$$\frac{d}{dt} \mathbf{v}(t) = -\frac{1}{m} \nabla V_\eta(\mathbf{r}, t), \quad (6.149)$$

$$\frac{d}{dt} \mathbf{r}(t) = \mathbf{v}(t), \quad (6.150)$$

where $V_\eta(\mathbf{r}, t)$ is the (space- and/or time-dependent) eigenvalue of $\hat{H}(\mathbf{r}, t)$ corresponding to the eigenstate $|\chi_\eta(\mathbf{r}, t)\rangle$, and the Markovian decoherence operator is

$$\hat{f}_\eta^{\text{Markov DD}}(\mathbf{r}, t) = \sum_i |\chi_i(\mathbf{r}, t)\rangle \langle \chi_i(\mathbf{r}, t)| \gamma_{\eta i}^{\text{Markov DD}}(\mathbf{r}, t), \quad (6.151)$$

where $\gamma_{ij}^{\text{Markov DD}}(\mathbf{r}, t)$ are the approximate Markovian decoherence rates given by (6.146). The integration of the internal state can be performed in the local eigenbasis using the transformed Hamiltonian defined in (6.38), in which case the decoherence operator is diagonal with the i^{th} diagonal equal to $\gamma_{\eta i}^{\text{Markov DD}}(\mathbf{r}, t)$.

2. Normalise the internal state vector to unit norm:

$$|\tilde{\chi}(t')\rangle \leftarrow \frac{|\tilde{\chi}(t')\rangle}{\sqrt{\langle \tilde{\chi}(t') | \tilde{\chi}(t') \rangle}} \quad (6.152)$$

3. Evaluate the S^{time} , S^{space} and S^{stay} matrices over this time interval for Tully's minimal-switches algorithm as described in Section 6.4.2, equations (6.60–6.62). This may require numerical differentiation of the basis vectors of the Hamiltonian with respect to space and time, and possibly numerical diagonalisation of the Hamiltonian to obtain the basis vectors as functions of space and time, if analytic expressions are not known. If Δt is not small enough to be considered infinitesimal for the purposes of evaluating (6.53) and (6.54), and one numerically integrates them for a more accurate result, note that these expressions require the state populations $\{c_i(t)\}$ in the local basis in the *absence* of decoherence; one therefore cannot do this integral in tandem with integrating the above differential equation for the internal state.²¹ Alternately, compute the overall S matrix of your favourite hidden-variable theory for evolution over the given time interval in the local eigenbasis (similarly keeping in mind that the effective unitary given to the hidden variable theory must be due to the evolution of $\hat{H}_{\text{eff}}(t)$ only, and not the decoherence term). If it saves computational power to do so, only compute column η of S , since only transitions from state η to other states need be considered.

4. Compute \tilde{P}^{space} (6.72) and hence \tilde{S}^{space} (6.74) to set the probability of classically disallowed transitions to zero as described in Section 6.4.3.

5. Modify the hidden variable with probability

$$\Pr(\eta \xrightarrow{\text{space}} i) = \tilde{S}_{i\eta}^{\text{space}} \quad (6.153)$$

that it transitions to the state ($i \neq \eta$) via a spatially induced transition, and probability

$$\Pr(\eta \xrightarrow{\text{time}} i) = S_{i\eta}^{\text{time}} \quad (6.154)$$

that it transitions to a state ($i \neq \eta$) via a temporally induced probability, and probability

$$\Pr(\eta \rightarrow \eta) = 1 - \sum_i (\tilde{S}_{i\eta}^{\text{space}} + S_{i\eta}^{\text{time}}) \quad (6.155)$$

that it keeps its current value. These choices can be randomly chosen between with the correct probabilities as described in Figure 6.4.

6. If the hidden variable did transition to a different state $i \neq \eta$, and if it was a spatially induced transition, instantaneously adjust the atom's velocity as required by energy conservation:

$$\mathbf{v}(t') \leftarrow \mathbf{v}(t') + \Delta\mathbf{v}_{\eta i}(\mathbf{r}, t), \quad (6.156)$$

where $\Delta\mathbf{v}_{\eta i}(\mathbf{r}, t)$ is given by (6.76).

This completes the Markovian hidden-variable semiclassical method.

6.6.2 Mean auxiliary trajectories hidden-variable semiclassical method

This is the more complex version of my method. Additional classical trajectories are propagated—one for each internal state of each atom—but other than classically evolving these additional trajectories, the increase in computational cost is minimal.

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²¹Though with split-operator methods and some care, the decoherence term can be treated separately such that the two integrals can be evaluated without repeating calculations unnecessarily.

State variables

The non-Markovian/Mean auxiliary trajectories variant of the hidden-variable semiclassical method has the following state variables describing the state of each simulated atom at each moment of time:

- The internal state vector of the atom $|\tilde{\chi}(t)\rangle$, in any chosen basis (though the local eigenbasis is likely to be computationally simplest to perform calculations in).
- A set of positions $\{\mathbf{r}_i(t)\}$ and velocities $\{\mathbf{v}_i(t)\}$, one position and velocity for each internal state of the atom.
- The hidden variable $\eta(t)$, equal to an integer index specifying one of the eigenstates $|\chi_\eta(\mathbf{r}_\eta, t)\rangle$ of the Hamiltonian $\hat{H}(\mathbf{r}_\eta, t)$ atom in the local energy eigenbasis at the location \mathbf{r}_η of the main trajectory.

Initial conditions

The manner of generating initial conditions of the mean auxiliary trajectories variant of the hidden-variable semiclassical method is identical to that of the Markovian variant, with the exception that multiple initial positions and velocities are required instead of just one. The initial positions and velocities of all auxiliary trajectories can however simply be set equal to those of the main trajectory for each atom.

Evolution

The evolution of the state variables for each atom from time t to $t' = t + \Delta t$ in the non-Markovian hidden-variable semiclassical method proceeds similarly to the Markovian variant.

1. Evolve the internal state vector and set of motional state variables of each atom from time t to t' according to the following coupled differential equations, using one or more steps of your favourite non-multi-step numerical integration method:

$$\frac{d}{dt} |\tilde{\chi}(t)\rangle = -\frac{i}{\hbar} \hat{H}(\mathbf{r}_\eta, t) |\tilde{\chi}(t)\rangle - \hat{F}_\eta^{\text{MMDD}}(t) |\tilde{\chi}(t)\rangle, \quad (6.157)$$

$$\frac{d}{dt} \mathbf{v}_i(t) = -\frac{1}{m} \nabla V_i(\mathbf{r}_i, t), \quad (6.158)$$

$$\frac{d}{dt} \mathbf{r}_i(t) = \mathbf{v}_i(t), \quad (6.159)$$

where $V_i(\mathbf{r}_i, t)$ is the (space- and/or time-dependent) eigenvalue of $\hat{H}(\mathbf{r}_i, t)$ corresponding to the eigenstate $|\chi_i(\mathbf{r}_i, t)\rangle$ at the position of the i^{th} trajectory, which may be the main ($i = \eta$) or an auxiliary ($i \neq \eta$) trajectory; and the non-Markovian decoherence operator is

$$\hat{F}_\eta^{\text{MMDD}}(t) = \sum_i |\chi_i(\mathbf{r}_\eta, t)\rangle \langle \chi_i(\mathbf{r}_\eta, t)| \gamma_{\eta i}^{\text{MMDD}}(t), \quad (6.160)$$

where $\gamma_{ij}^{\text{MMDD}}(t)$ are the minimal-memory non-Markovian decoherence rates given by (6.147); $\gamma_{\eta j}^{\text{MMDD}}(t)$ comprising the diagonals of $\hat{F}_\eta^{\text{MMDD}}(t)$ if working in the local eigenbasis along the main trajectory. When computing $\gamma_\eta^{\text{MMDD}}(t)$, clip its real part from below to zero. Although such recoherence can in principle be physically meaningful, it can arise out of tiny amplitudes of states with auxiliary trajectories far from the main trajectory, and is hence subject to large floating point error.²² I have not observed this problem with the Dirac-delta based decoherence rates, but it does seem to appear if the Gaussian-projection based decoherence rates are used.

²²Large error in *positive* (real parts of) decoherence rates in the same context is not a problem, as the state amplitudes are already very close to zero, so the accuracy of the rate at which they more closely approach zero is of no consequence.

2. Normalise the internal state vector to unit norm:

$$|\tilde{\chi}(t')\rangle \leftarrow \frac{|\tilde{\chi}(t')\rangle}{\sqrt{\langle\tilde{\chi}(t')|\tilde{\chi}(t')\rangle}} \quad (6.161)$$

3. Evaluate the S^{time} , S^{space} and S^{stay} matrices as in Step 3 for the Markovian case
4. Compute \tilde{S}^{space} as in Step 4 for the Markovian case. We treat transitions to have occurred at the location of the originating trajectory, therefore for the purposes of computing \tilde{S}^{space} , the expressions for $\Delta v_{ij}^2(\mathbf{r}, t)$ (6.71) and $\hat{d}_{ij}(\mathbf{r}, t)$ (6.69) must be evaluated using $\mathbf{r} = \mathbf{r}_i(t')$, and the conditional in \tilde{P}^{space} (6.72) evaluated with $\mathbf{v} = \mathbf{v}_i(t')$. Since only transitions from the current state specified by the hidden variable to other states need to be considered, one only needs to compute column η of \tilde{S}^{space} , in which case only \mathbf{r}_η and \mathbf{v}_η are necessary. This is then identical to the non-Markovian case, in which the main trajectory is the only trajectory simulated.
5. Compute the energy-conservation mandated velocity difference $\Delta v_{\eta i}$ given by (6.76) for all $i \neq \eta$. As with Step 4, since we treat transitions to have occurred at the location of the main trajectory, the velocity difference expression should be evaluated using the position and velocity of the main trajectory, as emphasised in (6.141).
6. Instantaneously modify the positions and velocities of all auxiliary trajectories ($i \neq \eta$) according to (6.139) and (6.140). This in essence spawns new auxiliary trajectories at the location of the main trajectory, but does not keep them: the averaging expressions replace each auxiliary trajectory with a weighted sum of its previous position and velocity and the position and velocity of the new trajectory. If the population of the state corresponding to a given auxiliary trajectory falls below some threshold (I use $|c_{i \neq \eta}(t)|^2 < 10^{-8}$), then discard the trajectory and replace it with the position and velocity of the main trajectory.
7. Probabilistically make a transition of the hidden variable as in Step 5 for the Markovian case.
8. In the case that the hidden variable transitioned, replace the trajectory of the new state with the position and velocity of the previous main trajectory. In the case that it was a spatially-induced non-adiabatic transition, adjust the velocity of the new state's trajectory with the previously computed $\Delta v_{\eta \xi}$, where η is the previous value of the hidden variable and ξ is the new value.

This completes the mean auxiliary trajectories hidden-variable semiclassical method.

6.7 Results

Here I demonstrate the two variants of the method, each for two different spatially dependent Hamiltonians in one dimension, and compare them to the numerical result from the full Schrödinger wave equation. All simulations are for ^{87}Rb atoms in the $F = 1$ groundstate subject to the Zeeman Hamiltonian for a given magnetic field profile. In scenario I, atoms in the trapped state (for 87 , the locally spin-down state is magnetically trappable) pass over a field minimum, which induces Majorana transitions, allowing the spin-flipped atoms to escape. In scenario II, the magnetic field profile is modified to create two closely spaced field minima. This scenario demonstrates the importance of accurate decoherence in surface-hopping/hidden-variable semiclassical methods, as the two regions of non-adiabatic coupling are encountered in a time comparable to the decoherence timescale where the exact form of the decoherence model can strongly

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influence the outcome of transitions. I use Schrödinger theory for the hidden variable theory throughout. Since the Hamiltonians are time-independent, as mentioned these results do not test my proposed method of conditionally applying velocity jumps to transitioned atoms—velocity jumps are always applied.

I simulated both scenarios first using the Schrödinger wave equation, as the gold-standard against which to benchmark the other methods. Numerically I solved the for the three-component wavefunction using fourth order Runge–Kutta (Section 3.3) and the Fourier method of evaluating spatial derivatives (Section 3.4.1). I added imaginary potentials near the boundaries of the region I wished to simulate to absorb untrapped wavepackets rather than have them reflect or wrap around due to periodic boundary conditions inherent in the Fourier method of derivatives. The initial condition was a Gaussian wavepacket in the locally spin-down eigenstate, $m_F = -1$, though I simulated in the eigenbasis of \hat{F}_z , transforming as necessary to create the initial condition and later to compare state populations in the local eigenbasis.

I then simulated both scenarios with 1×10^4 semiclassical atoms using the Ehrenfest semiclassical method, the Markovian hidden-variable semiclassical method, and the mean auxiliary trajectories hidden-variable semiclassical model. The initial conditions were set by drawing positions and velocities randomly for each semiclassical particle from the position and velocity probability distributions implied by the Gaussian wavepacket.

For both semiclassical methods I used simple first-order split-step evolution for the quantum evolution in the local eigenbasis, and evolved the classical trajectories analytically using constant acceleration over each timestep:

$$|\tilde{\chi}_H(t')\rangle \approx e^{-\hat{F}_\eta(t)\Delta t} \hat{U}_H(z', t') \hat{U}_H^\dagger(z, t) e^{-\frac{i}{\hbar}\hat{V}(z, t)\Delta t} |\tilde{\chi}_H(t)\rangle, \quad (6.162)$$

$$z' \approx r - \frac{\Delta t^2}{2m} \nabla V_\eta(z) \quad (6.163)$$

$$v' \approx v - \frac{\Delta t}{m} \nabla V_\eta(z), \quad (6.164)$$

where z , v and t are the position, velocity and time at the start of a timestep, z' , v' and t' at the end of the timestep, and $\Delta t = t' - t$, is the timestep, which was $\Delta t = 500$ ns. These methods are not particularly accurate, but were sufficient for these demonstrations and allowed the implementation to be simple. In particular, the ordering of the individual operators for evolving the state vector over a timestep was chosen to allow everything to the right of the decoherence factor be the unitary part of the evolution for the timestep, which could be input to Schrödinger theory as the unitary describing evolution in the local basis for that timestep.

Scenario I

In this scenario, a ^{87}Rb atom begins as an initially stationary Gaussian wavepacket of width 600 nm centred at $z_0 = -20$ μm in the Zeeman potential experienced by the $F = 1$ groundstate manifold (Section 2.3.3) with the following magnetic field profile:

$$\mathbf{B}_{\text{trap}}(z) = (B_x, B_y, B_z) = (B_\perp, 0, B'_z z) \quad (6.165)$$

where the transverse magnetic field is $B_\perp = 162.5$ mG and the z -gradient is $B'_z = 250$ G cm^{-1} . This potential provides a region close to $z = 0$ where the magnetic field rapidly reverses direction too quickly for the atom's spin to adiabatically follow, inducing Majorana transitions. The atom accelerates under the initial potential gradient toward this region. The adiabatic potentials resulting from the Zeeman Hamiltonian with this magnetic field are shown in Figure 6.7. Results are shown in Figure 6.9 for the Markovian model and Figure 6.10 for the auxiliary trajectories model.

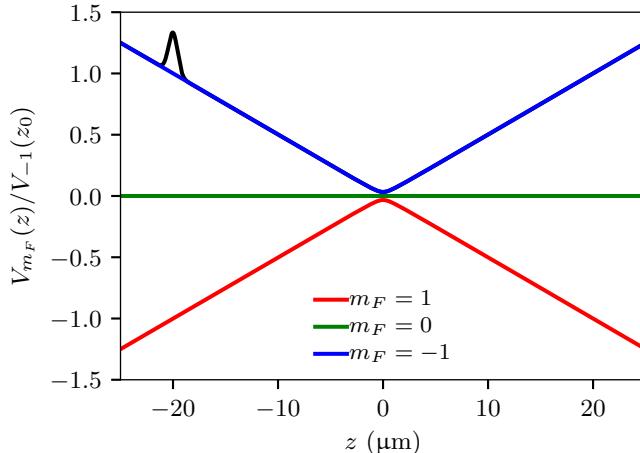


Figure 6.7: Adiabatic potentials for scenario I with a single region of coupling. The initial probability distribution of the atoms is depicted (in arbitrary units), offset for clarity by the adiabatic potential of the initial state.

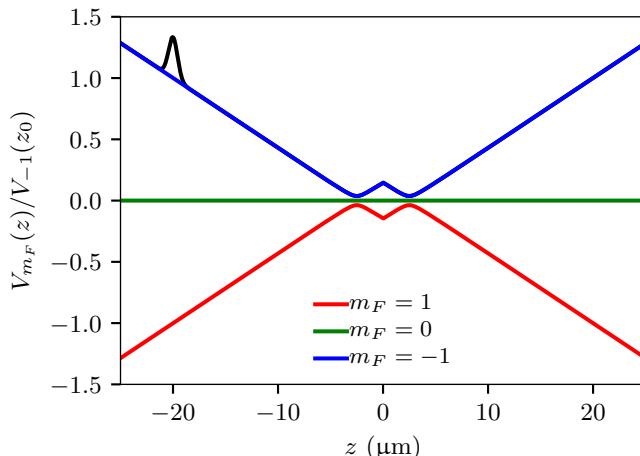


Figure 6.8: Adiabatic potentials for scenario II, with two closely spaced coupling regions. The initial probability distribution of the atoms is depicted (in arbitrary units), offset for clarity by the adiabatic potential of the initial state.

Scenario II

Scenario II is identical to scenario I, except that the magnetic field profile is instead²³

$$\mathbf{B}_{\text{trap}}(z) = (B_x, B_y, B_z) = \begin{cases} (B_{\perp}, 0, B'_z(z + z_{\text{fieldmin}})) & z < 0 \\ (B_{\perp}, 0, -B'_z(z + z_{\text{fieldmin}})) & z \geq 0 \end{cases} \quad (6.166)$$

²³This field profile is unphysical, but serves as a good test that highlights the improvement of the mean auxiliary trajectories method over the Markovian method.

where $z_{\text{fieldmin}} = 2.5 \mu \text{m}$ is the distance of each field minimum to the origin, such that the two minima are separated by $5 \mu \text{m}$, and with the other variables having the same values as in scenario I. The adiabatic potentials resulting from the Zeeman Hamiltonian with this magnetic field are shown in Figure ???. Results are shown in Figure 6.11 for the Markovian model and Figure 6.12 for the auxiliary trajectories model.

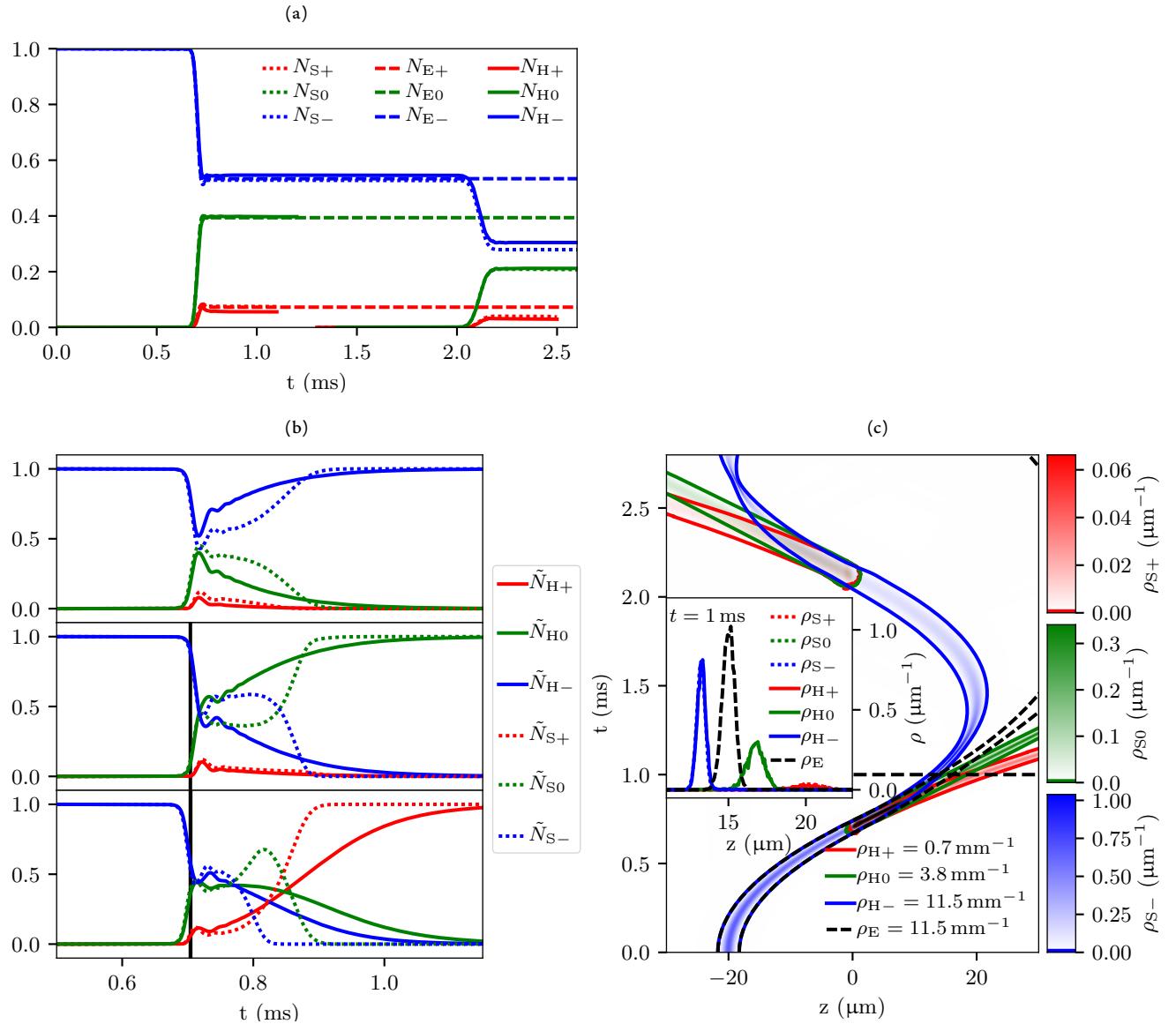


Figure 6.9: Scenario I results for the approximate Markovian decoherence hidden variable semiclassical method, compared to the Schrödinger wave equation and Ehrenfest semiclassical method. (a) The populations N in each spin projection state (denoted N_{\pm} , N_0 for $m_F = \pm 1$ and $m_F = 0$ respectively) over time, obtained for the Schrödinger wave equation (N_S) by integration in the local basis over all space, for the Ehrenfest semiclassical model (N_E) as an expectation value over all semiclassical atoms' internal state vectors, and for the hidden-variable semiclassical method (N_H) as the proportion of atoms with the hidden variable specifying each state at each moment of time. The lines ending just after 1 ms is due to atoms/amplitude leaving the integration region. There is good agreement between all methods for the first approach of the field minimum, though the Ehrenfest trajectories do not make a second approach within the simulation time. The remaining deviations between the Schrödinger and hidden-variable curves are within the random variation of the stochastic method. (b) The state populations along specific trajectories corresponding to classical motion in $m_F = -1$ (top) for all time, switching to $m_F = 0$ at the first field minimum (middle), and to $m_F = 1$ (bottom). Obtained by interpolation and normalisation of the Schrödinger wave equation results and by selecting the semiclassical atoms making hidden-variable transitions closest to the chosen points (black lines). An Ehrenfest atom is chosen for comparison that was close to the given trajectory until the field minimum, after which it diverged. This plot shows the shape of the decoherence curves and how our exponential approximation is a poor fit, but does have the correct timescale and is better than the Ehrenfest method's absence of any decoherence. (c) Probability density of trajectories. Schrödinger wave equation results are shaded, and Ehrenfest and hidden-variable semiclassical results shown with a single contour line for a chosen density. Inset: The probability distribution at $t = 1$ ms. There is good agreement between the hidden-variable semiclassical and Schrödinger wave equation results.

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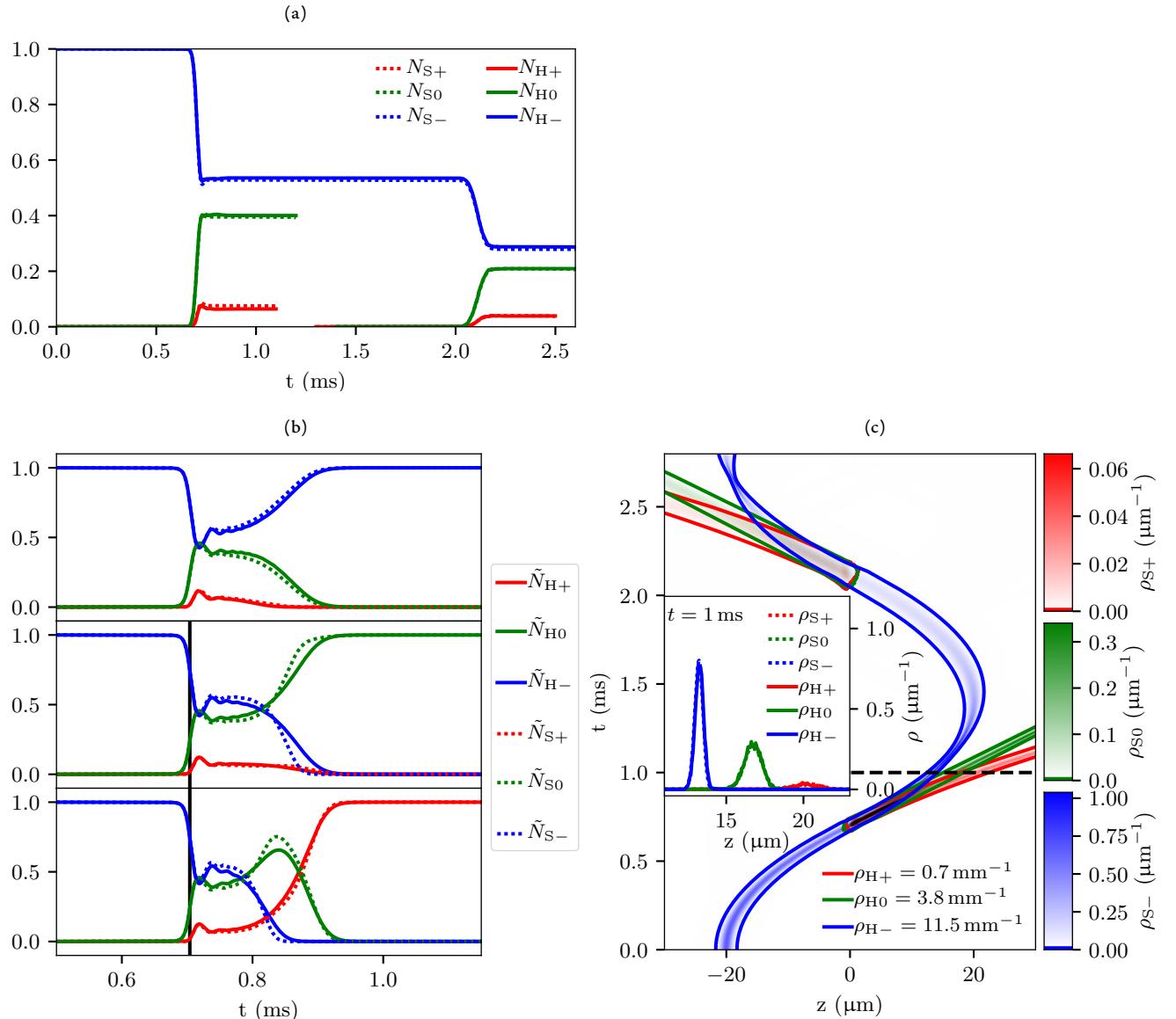


Figure 6.10: Scenario I results for the mean-auxiliary-trajectories hidden variable semiclassical method compared to the Schrödinger wave equation (Ehrenfest results excluded from this and subsequent results plots for clarity). (a) as described in Figure 6.9. Once again, there is good agreement between all methods for the first approach to the field minimum, and good agreement between the hidden-variable semiclassical method for the second approach, though the Ehrenfest trajectories do not approach a second time. Again, the remaining deviation is within the range expected given the stochastic algorithm and could be reduced by increasing the sample size. (b) Populations after projecting onto specific classical trajectories, as and comparing with semiclassical atoms closest to those trajectories, as also described in Figure 6.9. Here we see that the decoherence curves traced out by the auxiliary-trajectories method are in much closer agreement with the results of the Schrödinger wave equation than those of the Markovian model. Again, the discrepancies are within the statistical variation of the model; shown in each subplot is the atom that was closest to the given trajectory but none are exactly aligned with the trajectories. (c) Probability density of trajectories of the three methods, as described in Figure 6.9. Once again there is excellent agreement between the Schrödinger wave equation and hidden-variable semiclassical methods, with the Ehrenfest trajectories showing the expected unphysical behaviour.

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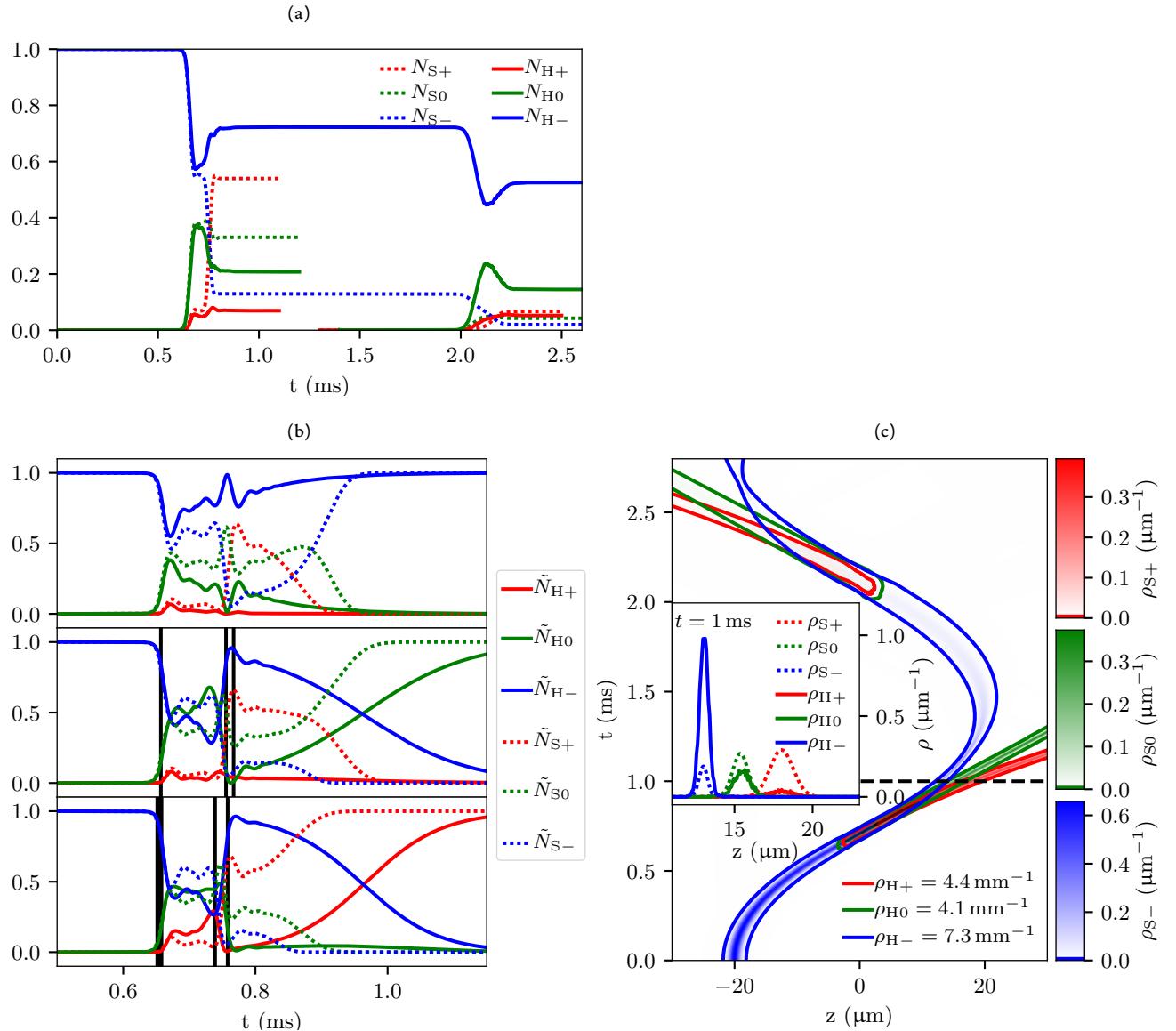


Figure 6.11: Scenario II results for the approximate Markovian decoherence hidden variable semiclassical method compared to the Schrödinger wave equation. All subfigures as described in Figure 6.9. There are two minima in the magnetic field, closely spaced near $z = 0$. Here we see the breakdown of the Markovian model in the presence of a second region of non-adiabatic coupling that the atoms encounter before they have fully decohered after the first. (a) The inaccurate modelling of decoherence has resulted in the wrong relative populations and phases between the states when the atom encounters the second region of coupling, leading to completely different populations after it emerges from the second field minimum. (b) Populations along specific classical trajectories. Incorrect decoherence curves were inconsequential when there were no closely spaced field minima, but here the wrong curves mean the wrong initial conditions as the atom begins to make transitions again near second field minimum. If the minima were farther separated, the exponential decoherence curve would decay the system back to a single eigenstate, which would be correct. At shorter times, it matters whether the decoherence curves actually have the right shape. Because our semiclassical atoms near the centre of the wavepacket have very small population in $m_F = 0$ and $m_F = -1$ states after the second region of coupling, my algorithm for finding the best atom to plot was unable to find any atoms whose hidden variable remained in these two states after the first region of coupling. I therefore allowed it to include atoms that made additional spin flips, so long as they quickly transitioned back to the desired state. This can be seen as groups of multiple black lines indicating a number of closely spaced spin flips. As mentioned earlier, this is one disadvantage of Schrödinger theory: it often makes multiple transitions in a region of coupling in cases where Tully's fewest-switches algorithm would make only one, making it harder to pick out which atoms made certain transitions. (c) Probability distributions of trajectories. The trajectories themselves are in excellent agreement with the Schrödinger wave equation—but there are the wrong proportion of atoms in each cluster of trajectories.

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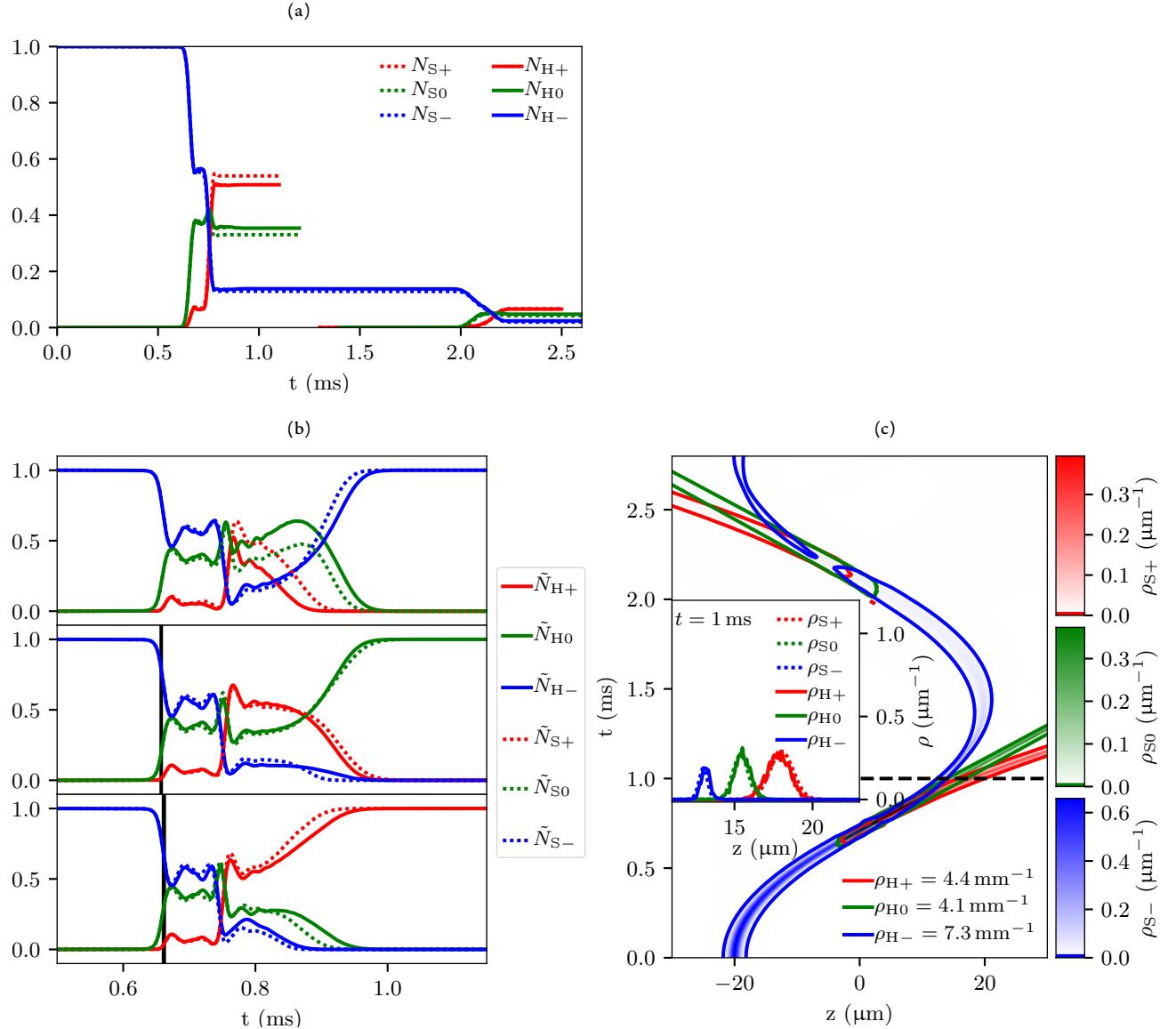


Figure 6.12: Scenario II results for the mean-auxiliary-trajectories hidden variable semiclassical method compared to the Schrödinger wave equation. All subfigures as described in Figure 6.9. (a) Here we see that the auxiliary-trajectories method of modelling decoherence leads to correct populations (again, they are within statistical error), even when there are two field minima closely spaced. (b) Here we see why: the decoherence curves match those of the Schrödinger wave equation in between the two field minima (identifiable by the two large vertical drops in the $m_F = -1$ population), providing the correct initial conditions as the atoms approach the second field minimum. (c) The trajectories and probability density functions agree well.

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6.7.1 Gaussian projection results

Here I demonstrate the difference between the above results, which use the Dirac-delta decoherence rates, with what one obtains if projecting onto a Gaussian instead as we originally proposed before changing our minds in Section 6.5.5. Here I only use the mean-auxiliary-trajectories method; though as argued in Section 6.5.5, the Markovian decoherence rate has its problems too when derived assuming Gaussian projections.

For these results I simulated scenario II, in which there are two closely spaced magnetic field minima, using two different decoherence rates. In the first, I take the decoherence rate $\gamma_{ij}^{\text{MM}}(t)$ given by (6.137) at face value, velocity-space separation included. This leads to poor results as the decoherence is clearly too fast, as shown in Figure 6.13.

In the second, I make a similar change to that of Shenvi et al. [97], setting $k_{ij} = 0$ for the relative wavenumber of the two Gaussian wavepackets in the decoherence factor (6.93), and otherwise proceeding as normal. This gives a decoherence rate

$$\gamma_{ij}^{\text{MM}}|_{k_{ij}=0}(t) = \frac{1}{4\sigma^2} \mathbf{r}_{ij}(t) \cdot \mathbf{v}_{ij}(t). \quad (6.167)$$

This leads to much better results, shown in Figure 6.14 but still not as good as the Dirac-delta projection based decoherence rate.

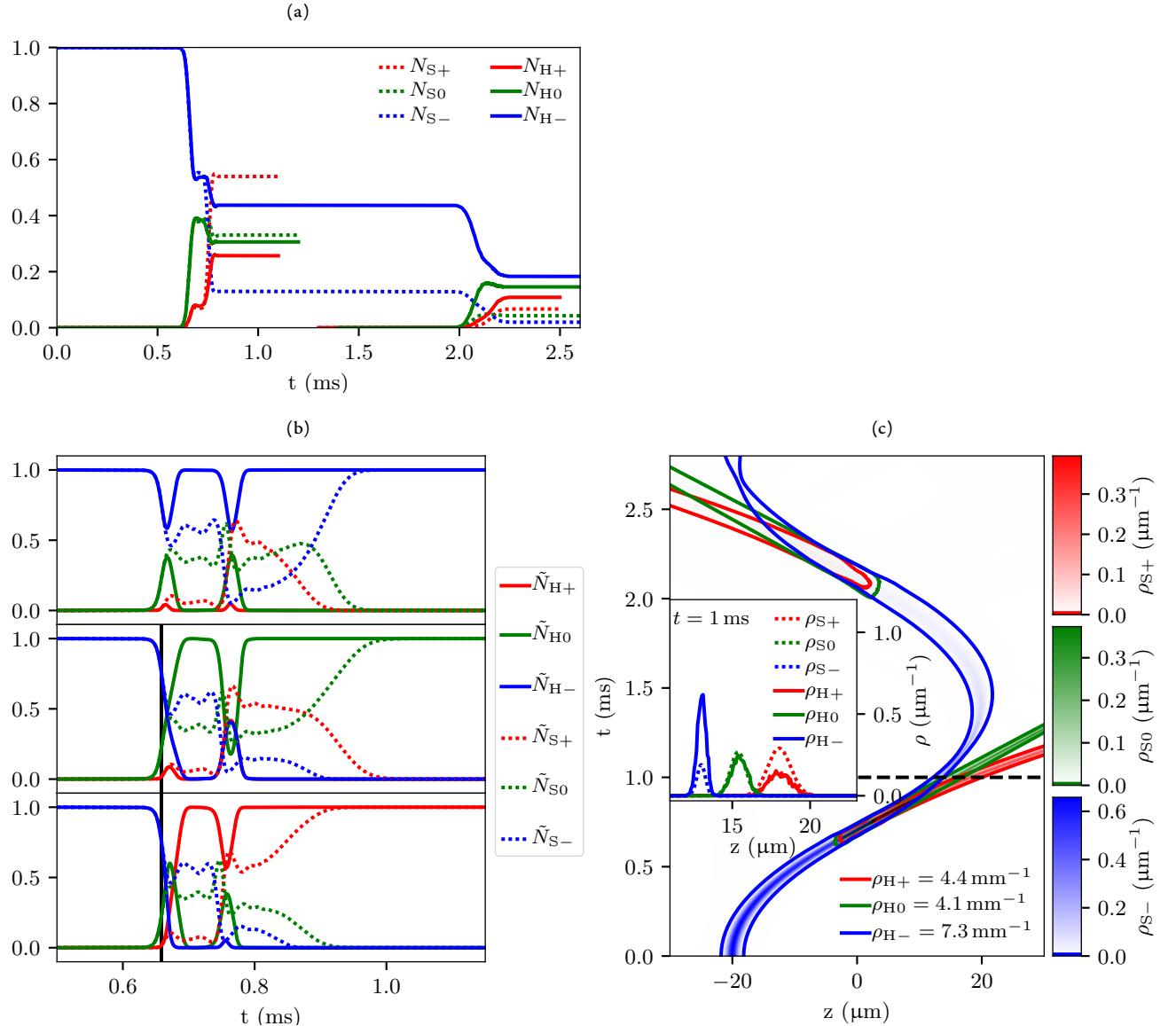


Figure 6.13: Scenario II results with Gaussian-projection based decoherence with the mean-auxiliary-trajectories hidden-variable semiclassical model, compared to the Schrödinger wave equation. All subplots as described in Figure 6.9. (a) Populations after the first region of coupling are correct, but are clearly not correct after the second region of coupling. (b) We see that the cause of this is decoherence that is much too fast decoherence, with population of states other than that selected by the hidden variable decaying away at a rate as fast or faster than that at which it appears, leading to symmetrical profiles in population around the coupling regions. Of course, since this plot is comparing the state vector's populations to the exact wavefunction along specific trajectories, and the decoherence model for this simulation is one based on projecting onto Gaussians rather than Dirac deltas centred exactly on the trajectory, we should not expect them to exactly agree (for that we would need to compare to projections of the exact wavefunctions onto Gaussians). However the disagreement is stark enough to make our point despite this. (c) As with the other results with poor modelling of decoherence, the trajectories are correct but the populations are not.

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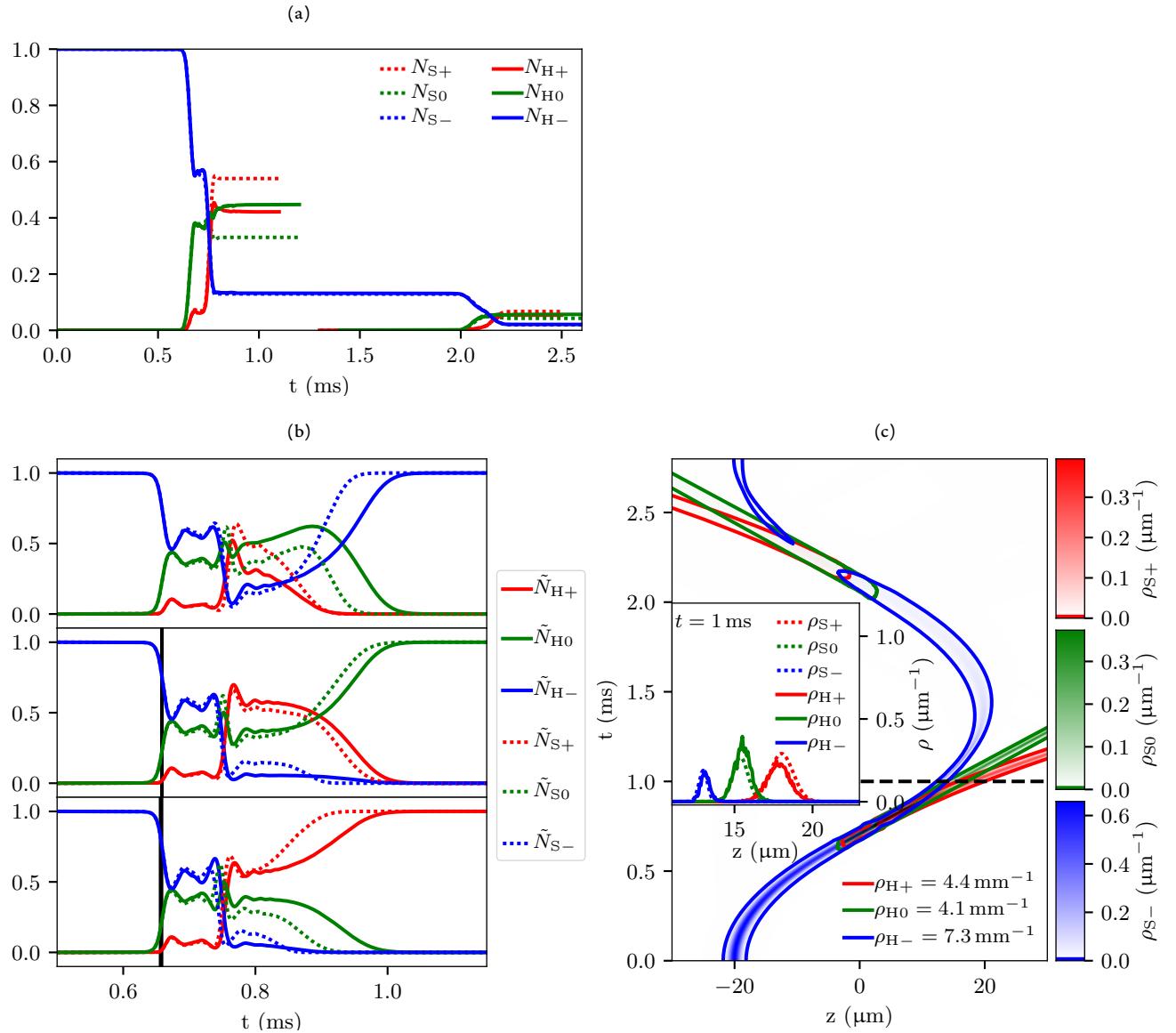


Figure 6.14: Scenario II results with Gaussian-projection based decoherence under the assumption of zero relative wavenumber between the Gaussians being projected onto each other when computing the decoherence rate. This decoherence rate $\gamma_{ij}^{\text{MM}}|_{k_{ij}=0}(t)$ is used with the mean-auxiliary-trajectories hidden-variable semiclassical model, and the results compared to the Schrödinger wave equation. All subplots as described in Figure 6.9. (a) State populations are in much closer agreement with the Schrödinger wave equation results than when we did not set the relative wavenumber of the Gaussians to zero (Figure 6.13). However, they are still not as accurate as our Dirac-delta based decoherence model (Figure 6.12). (b) As with Figure 6.13, we should not expect the Schrödinger wave equation results along the trajectories of the semiclassical atoms to agree with the state populations. However, for the three semiclassical atoms plotted (chosen as their trajectories were closest to the centre of the Gaussians and their transitions closest to the centre of the first region of coupling), they appear to agree quite well, including immediately after the second region of coupling. This may be a coincidence, or perhaps population transfers for atoms near the centre of the Gaussian are computed accurately with this decoherence model, but less accurately for atoms on other trajectories. (c) As usual, the trajectories are correct. We can see the minor error in populations in the inset, and it does not appear that the populations are more accurate in the centre as we just speculated. So any region of high accuracy in population transfer must be small enough to be hidden within the bins of the histogram, and is therefore not particularly of interest.

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6.8 Discussion and conclusion

The methods presented in this chapter give acceptable results in different scenarios. The Markovian decoherence model (Section 6.5.3 and modifications in Section 6.5.5) is able to obtain accurate trajectories and state populations for semiclassical atoms in the case of non-adiabatic coupling regions that are encountered at intervals that are large compared to the decoherence time, but produces inaccurate populations if multiple coupling regions are encountered at smaller time intervals. The mean auxiliary trajectories method (Section 6.6.2 and modifications in Section 6.5.5) solves this problem by more accurately modelling dynamic decoherence curves, and is therefore able to obtain accurate trajectories and state populations even in the presence of multiple regions of coupling encountered at intervals comparable to the decoherence timescale. Both methods are improvements over the Ehrenfest semiclassical method, which has no decoherence and no separation of trajectories, and over Tully's original model, which had separation of trajectories but no decoherence.

The problem of velocity separation of wavepackets causing unexpected and undesired additional decoherence appears to be resolvable by instead modelling a projected state vector that is a projection of the multi-component wavefunction onto a Dirac delta—rather than a Gaussian wavepacket—equivalent to evaluating the wavefunction at a specific point in space following a classical trajectory over time (Section 6.5.5). The total multi-component wavefunction is still modelled as a collection of Gaussian wavepackets following classical trajectories. I believe that this is a more natural way to obtain state vectors that are statistically representative of the underlying multi-component wavefunctions we are approximating. By not being based on integrating wavepacket overlaps over space, population in a state with a different phase gradient is not treated as having vanished merely because it has an inhomogeneous phase—instead we obtain some population, and some (perhaps rapidly varying) phase.

The above solution to the velocity separation problem, along with my method of averaging auxiliary trajectories in order to model only one trajectory per internal state of the atom, are modest improvements over similar methods such as that by Shenvi et al. [97]. The decoherence rate that results when Shenvi et al. remove velocity separation from their decoherence calculation are still based on the overlap of two Gaussians, and is therefore still a kind of average over the entire wavepacket. The real part of the resulting decoherence rate is different from mine, based on Dirac deltas, by a factor of two, and lacks an imaginary part (velocity separation in my model still causes a phase shift). As I've argued, using the value at the centre of the wavepacket is more natural than averaging, as this gives you something closer to a value the wavefunction actually has at some point in space, and this is borne out by the simulation results in the previous section.

Schrödinger theory is new to surface hopping as far as I know, and I've shown in the results section of this chapter that it gives good results. However, as discussed in Section 6.4.1, it is considerably computationally expensive to compute transition probabilities from, though it can take as input arbitrary unitaries, allowing one to compute transition probabilities consistent with quantum evolution over arbitrary time intervals, rather than integrating infinitesimal transition probabilities over a finite time interval as with Tully's fewest switches.

The fact that Tully's fewest-switches algorithm gives smaller transition probabilities than Schrödinger theory (even though both are consistent with the quantum probabilities) makes it more desirable for use in a surface-hopping/hidden-variable semiclassical model. This is because every transition the hidden variable of a semiclassical atom makes is an opportunity for the atom to spend some time subject to a different potential. Even if the probability of the atom ending up in the right state is correct, the more potential surfaces it visits on its way there, the larger the variance in the integrated force it experi-

enced will be, and hence the larger the variance in its final trajectory. A larger number of transitions also implies a larger variance in the number of atoms in each state (as given by their hidden variable), even though the expectation value might agree with the quantum populations. A ‘fewest-switches’ algorithm is therefore likely to be the most accurate in this sense of reduced statistical variation. One final advantage of fewest hops in the context of my mean-auxiliary-trajectories method is that it means the algorithm will not discard as many auxiliary trajectories upon switching to them, resulting in possibly more accurate decoherence.

On that note, one could also imagine a variation of the method where an auxiliary trajectory is simply considered the main trajectory upon a transition, and not replaced by the previous main trajectory. In that case the number of hops would not influence whether the decoherence history was destroyed more than necessary. I have not thought of a way to make this retention of auxiliary trajectories workable in light of conditional velocity jumps, though of course since the conditional velocity jump idea is untested, it may not turn out to be a problem.

Although the core idea of using a stochastically evolving variable to choose which classical force to subject a semiclassical atom to wasn’t original, in rediscovering it I have identified that hidden-variable theories and hopping algorithms are in fact the same thing. With the ideas from the surface-hopping literature that my method previously lacked, and with my own improvements, I look forward to evangelising the method’s use in the cold atom physics communities, as I think it has its place among the Monte Carlo wavefunction methods and stochastic Schrödinger equations in simulating evaporative cooling, laser cooling, and other phenomena in cold atom physics whenever thermal atoms are subject to state-dependent forces.

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