
STATE-DEPENDENT FORCES IN COLD QUANTUM GASES

Christopher Billington

Submitted in total fulfilment of the requirements
of the degree of Doctor of Philosophy

Supervisory committee:

Prof Kristian Helmerston

Dr Lincoln Turner

Dr Russell Anderson



School of Physics and Astronomy
Monash University

January, 2017

rev: 31 (422f54686ef6)
author: Chris Billington <chrisjbillington@gmail.com>
date: Wed Dec 21 16:32:45 2016 -0500
summary: Working on split-step methods.

This page intentionally left blank

Contents

Contents	i
3 Quantum mechanics on a computer	1
3.1 From the abstract to the concrete: neglect, discretisation and representation	1
3.2 Solution to the Schrödinger equation by direct exponentiation	3
3.2.1 Matrix exponentiation by diagonalisation	4
3.2.2 Time-ordered exponentials and time-ordered products	5
3.2.3 The operator product/split-step method	7
3.3 Continuous degrees of freedom	11
3.3.1 Spatial discretisation	11
3.4 Discrete degrees of freedom	15
3.4.1 The interaction picture	15
3.4.2 Unitary integration	16
3.5 Continuous degrees of freedom	16
3.5.1 Finite differences	16
3.5.2 The Fourier basis	16
3.5.3 Harmonic oscillator basis functions	17
3.6 Finding ground states	17
3.6.1 Imaginary time evolution	17
3.6.2 Successive over-relaxation	17
3.6.3 Generalisation to excited states via Gram–Schmidt orthonormalisation	17
3.7 The finite-element discrete variable representation	17
References	19

rev: 31 (422f54686ef6)
author: Chris Billington <chrisjbillington@gmail.com>
date: Wed Dec 21 16:32:45 2016 -0500
summary: Working on split-step methods.

This page intentionally left blank

rev: 31 (422f54686ef6)
author: Chris Billington <chrisjbillington@gmail.com>
date: Wed Dec 21 16:32:45 2016 -0500
summary: Working on split-step methods.

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 [1]. Much of the numerical calculations 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 order (6th order or so) finite differences. I think Fourier transforms for spatial derivatives are overrated, and I show that the finite element discrete variable representation (FEDVR) is 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 ?? 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 interaction (neglecting the quantum degrees of freedom in the electromagnetic field).

rev: 31 (422f54686ef6)
 author: Chris Billington <chrisjbillington@gmail.com>
 date: Wed Dec 21 16:32:45 2016 -0500
 summary: Working on split-step methods.

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 [2–4] 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 [5] was designed to measure the effect of Doppler cooling, which does not involve Zeeman sublevels, and it was not until afterwards that theorists determined [6] 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 [7, 8] or using ultrafast (and therefore broad-band) laser pulses [9–11], 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 strictly speaking 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):

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

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

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

other methods of computing matrix exponentials that are more computationally efficient and numerically stable, such as the Padé approximant [CITE]. It should be noted that there is no known ‘best’ matrix exponential algorithm, all make compromises and perform poorly for certain types of matrices [12].

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

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

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

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

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 below:

³Note that the diagonals of D are the eigenvalues of H , and the columns of U 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.

```

1 import numpy as np
2 from numpy.linalg import eig
3
4 def expm(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 = eig(M)
12
13    # Now we compute exp(M) = U exp(D) U^\dagger where U is the matrix of
14    # eigenvectors (as columns) and D is the diagonal matrix of eigenvalues:
15
16    U = evecs
17    U_dagger = U.conj().swapaxes(-1, -2) # Only transpose the matrix dimensions
18    exp_D_diags = np.exp(evals)
19
20    # Compute the 3-term matrix product U*exp_D_diags*U_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', U, exp_D_diags, U_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 $\mathcal{O}(n^3)$ [12]. 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

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 the next section.

3.2.2 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 [CITE SOMETHING]. In this case, once represented in a specific basis, the solution to the Schrödinger equation is again the matrix multiplication

$$\psi(t) = U(t, t_0) \psi(t_0), \quad (3.14)$$

with

$$U(t, t_0) = e^{-\frac{i}{\hbar} \int_{t_0}^t H(t') dt'}, \quad (3.15)$$

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 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* [CITE]:

$$U(t, t_0) = \mathcal{T} \left\{ e^{-\frac{i}{\hbar} \int_{t_0}^t H(t') dt'} \right\}. \quad (3.16)$$

In this expression, \mathcal{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:

$$\mathcal{T} \{ H(4)H(1)H(2)H(5)H(3) \} = H(5)H(4)H(3)H(2)H(1). \quad (3.17)$$

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 (3.16) 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

$$U(t, t_0) = 1 + T \left\{ \sum_{n=1}^{\infty} \frac{1}{n!} \left[-\frac{i}{\hbar} \int_{t_0}^t H(t') dt' \right]^n \right\} \quad (3.18)$$

$$= 1 + \sum_{n=1}^{\infty} \frac{1}{n!} \left(-\frac{i}{\hbar} \right)^n T \left\{ \left[\int_{t_0}^t H(t') dt' \right]^n \right\}. \quad (3.19)$$

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:

$$U(t, t_0) = 1 + \sum_{n=1}^{\infty} \frac{1}{n!} \left(-\frac{i}{\hbar} \right)^n \int_{t_0}^t dt'_1 \int_{t_0}^{t'_1} dt'_2 \cdots \int_{t_0}^{t'_{n-1}} dt'_n \cdots T \{ H(t'_1) H(t'_2) \cdots H(t'_n) \}, \quad (3.20)$$

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

$$U(t, t_0) = 1 + \sum_{n=1}^{\infty} \left(-\frac{i}{\hbar} \right)^n \int_{t_0}^t dt'_n \cdots \int_{t_0}^{t'_{n-1}} dt'_2 \int_{t_0}^{t'_2} dt'_1 H(t'_n) \cdots H(t'_2) H(t'_1). \quad (3.21)$$

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 [CITE]. 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 [CITE] 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 n 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 corresponds to a Feynman diagram with n nodes [CITE].

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

⁶By sampling the integrand on a uniform grid, using a quadrature method, or Monte-Carlo integration [13] which is widely used for high dimensional integrals such as these.

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

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:

$$U(t, t_0) = \mathcal{T} \left\{ e^{-\frac{i}{\hbar} \int_{t_0}^t H(t') dt'} \right\} \quad (3.22)$$

$$= \lim_{N \rightarrow \infty} \prod_{n=N-1}^0 e^{-\frac{i}{\hbar} H(t_n) \Delta t}, \quad (3.23)$$

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 (3.21) is equivalent to (3.23) 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 (3.23) as a Taylor Series and collecting terms of equal powers of H then reveals that the two Taylor series are identical.

In any case, (3.23) 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 intervals. 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

$$\psi(t_{n+1}) = e^{-\frac{i}{\hbar} H(t_n) \Delta t} \psi(t_n) + \mathcal{O}(\Delta t^2) \quad (3.24)$$

$$\Rightarrow \psi(t) = \left(\prod_{n=N-1}^0 e^{-\frac{i}{\hbar} H(t_n) \Delta t} \right) \psi(t_0) + \mathcal{O}(\Delta t). \quad (3.25)$$

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 $\mathcal{O}(n^3)$.

3.2.3 The operator product/split-step method

Here I'll present decompositions similar to (3.24), 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 (3.24) 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. 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 the next section on continuous degrees of freedom, 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 a (matrix representation of a) Hamiltonian which is the sum of two non-commuting terms:

$$H(t) = H_1(t) + H_2(t). \quad (3.26)$$

The unitary for the solution to the Schrödinger equation is then:

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

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

$$U(t, t_0) = \prod_{n=N-1}^0 U(t_{n+1}, t_n), \quad (3.28)$$

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

$$U(t_{n+1}, t_n) = \mathcal{T} \left\{ e^{-\frac{i}{\hbar} \int_{t_n}^{t_{n+1}} H(t') dt'} \right\}. \quad (3.29)$$

Evaluating the integral using a one-point rectangle rule gives:

$$U(t_{n+1}, t_n) = e^{-\frac{i}{\hbar} H(t_n) \Delta t + \mathcal{O}(\Delta t^2)}, \quad (3.30)$$

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:

$$U(t_{n+1}, t_n) = e^{-\frac{i}{\hbar} H(t_n) \Delta t} + \mathcal{O}(\Delta t^2). \quad (3.31)$$

So far all we've done is justify (3.24). Now we'll acknowledge that H is a sum of two terms and write:

$$U(t_{n+1}, t_n) = e^{-\frac{i}{\hbar} (H_1(t_n) + H_2(t_n)) \Delta t} + \mathcal{O}(\Delta t^2). \quad (3.32)$$

Using the Baker–Campbell–Hausdorff formula [CITE], we can expand the exponential as:

$$e^{-\frac{i}{\hbar} (H_1(t_n) + H_2(t_n)) \Delta t} = e^{-\frac{i}{\hbar} H_1(t_n) \Delta t} e^{-\frac{i}{\hbar} H_2(t_n) \Delta t} e^{\frac{1}{2\hbar^2} [H_1(t_n), H_2(t_n)] \Delta t^2 + \mathcal{O}(\Delta t^3)} \quad (3.33)$$

$$= e^{-\frac{i}{\hbar} H_1(t_n) \Delta t} e^{-\frac{i}{\hbar} H_2(t_n) \Delta t} + \mathcal{O}(\Delta t^2) \quad (3.34)$$

$$\Rightarrow U(t_{n+1}, t_n) = e^{-\frac{i}{\hbar} H_1(t_n) \Delta t} e^{-\frac{i}{\hbar} H_2(t_n) \Delta t} + \mathcal{O}(\Delta t^2). \quad (3.35)$$

So we see that the ‘commutation error’ in (3.34) caused by treating the two terms as if they commute is no greater than the ‘integration error’ in (3.31) 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:

$$U(t, t_0) = \prod_{n=N-1}^0 U_1(t_{n+1}, t_n) + \mathcal{O}(\Delta t), \quad (3.36)$$

where

$$U_1(t_{n+1}, t_n) = e^{-\frac{i}{\hbar} H_1(t_n) \Delta t} e^{-\frac{i}{\hbar} H_2(t_n) \Delta t}. \quad (3.37)$$

Using the diagonalisation method of matrix exponentiation discussed in [SECREP], 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:

$$U_1(t_{n+1}, t_n) = \left(e^{-\frac{i}{\hbar} \tilde{H}_1(t_n) \Delta t} \otimes \mathbb{I}_{n_2} \right) \left(\mathbb{I}_{n_1} \otimes e^{-\frac{i}{\hbar} \tilde{H}_2(t_n) \Delta t} \right). \quad (3.38)$$

By applying (3.34) recursively for the case where either H_1 or H_2 is itself the sum of two further terms, (3.37) immediately generalises to the case where H is a sum of an arbitrary number of non-commuting terms:

$$U_1(t_{n+1}, t_n) = \prod_{m=1}^M e^{-\frac{i}{\hbar} H_m(t_n) \Delta t}, \quad (3.39)$$

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 (3.29), 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:

$$U(t_{n+1}, t_n) = e^{-\frac{i}{\hbar} (H(t_n) + H(t_{n+1})) \frac{\Delta t}{2}} + \mathcal{O}(\Delta t^3) \quad (3.40)$$

$$= e^{-\frac{i}{\hbar} (H_1(t_n) + H_2(t_n) + H_1(t_{n+1}) + H_2(t_{n+1})) \frac{\Delta t}{2}} + \mathcal{O}(\Delta t^3). \quad (3.41)$$

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 $H_1(t_n) + H_1'(t_n) \Delta t + \mathcal{O}(\Delta t^2)$, one can show that if the individual exponentials are ordered in the following way, then the remaining commutation error is also $\mathcal{O}(\Delta t^3)$:

$$U(t_{n+1}, t_n) = e^{-\frac{i}{\hbar} H_2(t_{n+1}) \frac{\Delta t}{2}} e^{-\frac{i}{\hbar} H_1(t_{n+1}) \frac{\Delta t}{2}} e^{-\frac{i}{\hbar} H_1(t_n) \frac{\Delta t}{2}} e^{-\frac{i}{\hbar} H_2(t_n) \frac{\Delta t}{2}} + \mathcal{O}(\Delta t^3). \quad (3.42)$$

This gives the second order split-step method:

$$U(t, t_0) = \prod_{n=N-1}^0 U_2(t_{n+1}, t_n) + \mathcal{O}(\Delta t^2), \quad (3.43)$$

where

$$U_2(t_{n+1}, t_n) = e^{-\frac{i}{\hbar} H_2(t_{n+1}) \frac{\Delta t}{2}} e^{-\frac{i}{\hbar} H_1(t_{n+1}) \frac{\Delta t}{2}} e^{-\frac{i}{\hbar} H_1(t_n) \frac{\Delta t}{2}} e^{-\frac{i}{\hbar} H_2(t_n) \frac{\Delta t}{2}}, \quad (3.44)$$

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

$$U_2(t_{n+1}, t_n) = \left(\prod_{m=M}^1 e^{-\frac{i}{\hbar} H_m(t_{n+1}) \frac{\Delta t}{2}} \right) \left(\prod_{m=1}^M e^{-\frac{i}{\hbar} H_m(t_n) \frac{\Delta t}{2}} \right). \quad (3.45)$$

U_2 can be concisely written in terms of U_1 as:

$$U_2(t_{n+1}, t_n) = U_1^\dagger(t_n + \frac{\Delta t}{2}, t_{n+1}) U_1(t_n + \frac{\Delta t}{2}, t_n), \quad (3.46)$$

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 (3.44) 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 [CITE]:

$$U(t, t_0) = \prod_{n=N-1}^0 U_4(t_{n+1}, t_n) + \mathcal{O}(\Delta t^4), \quad (3.47)$$

where

$$\begin{aligned} U_4(t_{n+1}, t_n) = & U_2(t_{n+1}, t_{n+1} - p\Delta t) \\ & \times U_2(t_{n+1} - p\Delta t, t_{n+1} - 2p\Delta t) \\ & \times U_2(t_{n+1} - 2p\Delta t, t_n + 2p\Delta t) \\ & \times U_2(t_n + 2p\Delta t, t_n + p\Delta t) \\ & \times U_2(t_n + p\Delta t, t_n), \end{aligned} \quad (3.48)$$

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

Expressions such as (3.44) 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.

rev: 31 (422f54686ef6)
author: Chris Billington <chrisjbillington@gmail.com>
date: Wed Dec 21 16:32:45 2016 -0500
summary: Working on split-step methods.

Dealing with nonlinearity

It is difficult to understate how powerful a method the split-step method is. It allows you to decompose exponentiating a Hamiltonian into exponentiating its component parts, in the subspaces that they act on, saving on computing power immensely compared to exponentiating the full Hamiltonian. It is unitary, and stable—that is, unlike Runge–Kutta methods, the method’s truncation error does not grow exponentially in time but is bounded (floating point rounding error still grows with time, but linearly) [TODO CONFIRM BASED ON FEDVR PAPER’S COMMENTS]. So there is no need to take much smaller timesteps than one otherwise would at earlier times in order to ensure the error remains acceptable at later times—an appropriate timestep at earlier times will be appropriate at later times, *ceteris paribus*. This is very appealing. Whilst higher order split-step methods quickly become unwieldy [CITE], fourth order accuracy is quite acceptable for many problems.

Actual Hamiltonians can be complicated, and are rarely. Otherwise, if the Hamiltonian can be decomposed into a *sum* of terms which are individually diagonalisable, the split-step method can be used (see sec [SECREf]). Failing that, a general purpose integration method such as fourth-order Runge-Kutta may have to be resorted to.⁹ I will elaborate on these comments with examples in the following sections.

⁹Though if I admit it, `RK4` is usually what I reach for first.

- absorbing boundary conditions, reflective boundary conditions, periodic boundary conditions

3.3 Continuous degrees of freedom

The single-particle, non-relativistic, scalar Schrödinger wave equation, as distinct from the general Schrödinger equation [EQREF], is:

[EQUATION].

As mentioned in [SECREf], the equation for the single-particle wavefunction of an atom in a single-component Bose–Einstein condensate is the Gross–Pitaevskii equation [EQUATION],

where $\psi(\mathbf{r}) = \sqrt{N} \langle \mathbf{r} | \psi \rangle$ 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 quantum mechanics all state vectors can be 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. So what do these vectors and operators look like? 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 below, however, spatial discretisation is actually just a particular choice of functional basis, namely the Fourier basis.

3.3.1 Spatial discretisation

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 [SECREf ABOVE] is to choose a basis in which to represent this state vector.

Lets say we discretise space in an equally-spaced $N_x \times N_y = 5 \times 5$ rectangular grid,¹⁰ with spacing Δx , and only represent the wavefunction at those 25 points in space. The state vector can then be represented by a list of 25 complex numbers, each taken to be

¹⁰For the sake of example— 256×256 is a more realistic minimum.

```
rev:      31 (422f54686ef6)
author:   Chris Billington <chrisjbillington@gmail.com>
date:     Wed Dec 21 16:32:45 2016 -0500
summary:  Working on split-step methods.
```


the wavefunction's value at the spatial position corresponding to one gridpoint. This 25-vector is now a concrete representation of our state vector.

[FIGURE SHOWING UNWRAPPING OF 2D REGION INTO COLUMN VECTOR. PERHAPS 2D REGION IS CONTINUOUS WITH A GRID SUPERIMPOSED ON IT, AND LINES JOIN EACH GRIDPOINT TO THE VECTOR AT RIGHT]

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 and operators onto each basis vector, as outlined in [SECREP ABOVE]. Assuming what we've done is equivalent to choosing a basis, that basis has a finite number (25) 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.¹¹ 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?

¹¹One for each position within the two dimensional space we're representing.

As a sidenote, spatial wavefunctions are often described as the representation of wavefunctions in the “spatial basis”—a basis in which the basis vectors $\{|\mathbf{r}\rangle\}$ are Dirac-deltas centred on each point in space [CITE DIRAC'S BOOK MAYBE]. The wavefunction $\psi(\mathbf{r})$ is then simply a coefficient saying how much of the basis vector $|\mathbf{r}\rangle$ (the spatial representation of which is a Dirac delta centred on the position \mathbf{r}) to include in the overall state vector. What we have *not* done is chosen a subset of these Dirac delta basis functions¹² as our basis. This would be very strange—our representation of the wavefunction would allow it to have a value at one gridpoint, and at the next gridpoint, but not in between. Spatially separated Dirac deltas do not spatially overlap at all; the matrix elements of the kinetic energy operator:

¹²Strictly, distributions, not functions, but “basis distributions” just doesn't have the right ring to it.

$$\langle \mathbf{r}_i | \hat{K} | \mathbf{r}_j \rangle = \int \delta(\mathbf{r} - \mathbf{r}_i) \left(-\frac{\hbar^2}{2m} \nabla^2 \right) \delta(\mathbf{r} - \mathbf{r}_j) d\mathbf{r} \quad (3.49)$$

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 chosen a set of two-dimensional boxcar functions centred on each gridpoint with width Δx in each direction. These cover all space in between gridpoints, but are no good because they are not twice differentiable everywhere, and hence the kinetic energy operator's matrix elements cannot be evaluated. No, neither of these bases will do. To interpret our spatial grid as a basis, we need a set of functions $[FUNC_{ij}(\mathbf{r})]$ (where i and j are the indices of the gridpoints in the x and y directions respectively) that have unit norm, are nonzero only at one gridpoint and are zero at all others, and are twice differentiable everywhere in our spatial region. Infinite choices present themselves to us, differing only in their incompleteness—the choice of which state vectors they will and won't be able to represent. A sensible choice 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.

The discrete Fourier transform to the rescue

It turns out that discretising space in this way can indeed be equivalent to choosing an entirely sensible basis. This is made clearer by thinking in terms of what we have done in Fourier space, which I'll quickly go through now.

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:

$$\langle \mathbf{r} | \mathbf{k}_{ij} \rangle = \begin{cases} \frac{1}{\sqrt{A}} e^{i\mathbf{k}_{ij} \cdot \mathbf{r}} & (\mathbf{r} \text{ within 2D region}) \\ 0 & (\mathbf{r} \text{ not within 2D region}), \end{cases} \quad (3.50)$$

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

$$\mathbf{k}_{ij} = \left[\frac{2\pi i}{L_x}, \frac{2\pi j}{L_y} \right]^T, \quad (3.51)$$

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

$$|\psi\rangle = \sum_{i=0}^{\infty} \sum_{j=0}^{\infty} \langle \mathbf{k}_{ij} | \psi \rangle |\mathbf{k}_{ij}\rangle \quad (3.52)$$

$$\Rightarrow \psi(\mathbf{r}) = \begin{cases} \sum_{i=0}^{\infty} \sum_{j=0}^{\infty} \langle \mathbf{k}_{ij} | \psi \rangle \frac{1}{\sqrt{A}} e^{i\mathbf{k}_{ij} \cdot \mathbf{r}} & (\mathbf{r} \text{ within 2D region}) \\ 0 & (\mathbf{r} \text{ not within 2D region}). \end{cases} \quad (3.53)$$

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 discrete basis. For one, there are an infinite number of them, and we require 25. 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 25 gridpoints.

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

$$\begin{cases} i \in \left[-\frac{N_x}{2}, \frac{N_x}{2} - 1\right] & (N_x \text{ even}) \\ i \in \left[-\frac{N_x-1}{2}, \frac{N_x-1}{2}\right] & (N_x \text{ odd}) \end{cases} \quad (3.54)$$

and

$$\begin{cases} j \in \left[-\frac{N_y}{2}, \frac{N_y}{2} - 1\right] & (N_y \text{ even}) \\ j \in \left[-\frac{N_y-1}{2}, \frac{N_y-1}{2}\right] & (N_y \text{ odd}) \end{cases} \quad (3.55)$$

we include only the N_x and N_y (both equal to 5 in our example) 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:

$$\langle \mathbf{k}_{ij} | \hat{K} | \mathbf{k}_{i'j'} \rangle = \frac{\hbar^2 k^2}{2m} \delta_{ii'} \delta_{jj'}, \quad (3.56)$$

where $k = |\mathbf{k}_i|$, 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.¹³

Now we have 25 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:

$$|\mathbf{r}_{ij}\rangle = \sum_{i'j'} e^{-i\mathbf{k}_{i'j'} \cdot \mathbf{r}_{ij}} |\mathbf{k}_{i'j'}\rangle \quad (3.57)$$

¹³Because 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.

with $\mathbf{r}_{i,j} = (i\Delta x, j\Delta y)^T$ and $i, j \in [0, 4]$. This is simply a unitary transformation of the truncated Fourier basis, with the unitary transformation matrix elements:

$$\hat{U}_{\text{DFT2};i',j';i,j} = \langle \mathbf{k}_{i',j'} | \mathbf{r}_{i,j} \rangle = e^{-i\mathbf{k}_{i',j'} \cdot \mathbf{r}_{i,j}}. \quad (3.58)$$

This transformation is in fact a discrete Fourier transform (hence the subscript), and the basis vectors $\{|\mathbf{r}_{i,j}\rangle\}$ have spatially localised wavefunctions that are nonzero only at one of the spatial gridpoints, so we refer to them as a discrete real-space basis. State vectors and operators can be transformed from their discrete Fourier space representation to their discrete real-space representation and back using the unitary \hat{U}_{DFT2} :

$$\langle \mathbf{r}_{i,j} | \psi \rangle = \sum_{i',j'} \hat{U}_{\text{DFT2};i,j;i',j'}^\dagger \langle \mathbf{k}_{i',j'} | \psi \rangle \quad (3.59)$$

$$\langle \mathbf{k}_{i,j} | \psi \rangle = \sum_{i',j'} \hat{U}_{\text{DFT2};i',j';i,j} \langle \mathbf{r}_{i',j'} | \psi \rangle \quad (3.60)$$

[TODO OPERATOR TRANSFORMATIONS ONCE INDEXING IMPROVED AND VECTOR REPRESENTATIONS INTRODUCED]

The spatial representation of the basis vectors $\{|\mathbf{r}_{i,j}\rangle\}$ can be computed using (3.50) as:

$$\phi_{i,j}(\mathbf{r}) = \langle \mathbf{r} | \mathbf{r}_{i,j} \rangle = \sum_{i',j'} e^{-i\mathbf{k}_{i',j'} \cdot \mathbf{r}_{i,j}} \langle \mathbf{r} | \mathbf{k}_{i',j'} \rangle \quad (3.61)$$

$$\Rightarrow \phi_{i,j}(\mathbf{r}) = \begin{cases} \sum_{i',j'} \frac{1}{\sqrt{A}} e^{i\mathbf{k}_{i',j'} \cdot (\mathbf{r} - \mathbf{r}_{i,j})} & (\mathbf{r} \text{ within 2D region}) \\ 0 & (\mathbf{r} \text{ not within 2D region}), \end{cases} \quad (3.62)$$

and are plotted in [TODO MAKE FIGURE].

These functions are sometimes called *periodic sinc functions*, or *band-limited delta functions* [CITE]. Each of them is zero at all of our 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. Thus, the approach of discretising space on a grid is indeed equivalent to choosing a (necessarily incomplete) orthonormal basis in which to work.

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.

Perhaps 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 reflecting off a barrier [SEE FIG TODO FOR EXAMPLE]. This effect is entirely unphysical¹⁴ and should be taken as a sign that the spatial grid is not fine enough for the dynamics being simulated.

Matrix elements using DFT basis

We now have a finite basis $\{|\mathbf{r}_{i,j}\rangle\}$ that matches our intuitions somewhat for representing a wavefunction at a set of gridpoints. A state vector¹⁵ can be represented as a linear sum of these basis vectors, with the coefficient for each one simply being equal to the value of

¹⁴With 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 are Bloch oscillations.

¹⁵Or rather, a state vector's projection onto the subspace being simulated.

that state vector's wavefunction at the gridpoint where that basis vector is nonzero:

$$|\psi\rangle = \sum_{i,j} c_{i,j} |\mathbf{r}_{i,j}\rangle, \quad (3.63)$$

where

$$c_{i,j} = \langle \mathbf{r}_{i,j} | \psi \rangle = \psi(\mathbf{r}_{i,j}). \quad (3.64)$$

Armed with a basis, we can now proceed to calculate matrix elements of the Hamiltonian, and then proceed to solve the differential equation [TODO EQREF 3.3] to determine how the coefficients $\{c_{i,j}\}$ evolve in time.

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

$$\hat{H}_{\text{Schrö}} = -\frac{\hbar^2 \hat{k}^2}{2m} + V(\hat{\mathbf{r}}), \quad (3.65)$$

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

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

$$\langle \mathbf{k}' | -\frac{\hbar^2 \hat{k}^2}{2m} | \mathbf{k} \rangle = -\frac{\hbar^2 k^2}{2m} \delta(\mathbf{k} - \mathbf{k}'), \quad (3.66)$$

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

$$\langle \mathbf{r}' | V(\hat{\mathbf{r}}) | \mathbf{r} \rangle = V(\mathbf{r}) \delta(\mathbf{r} - \mathbf{r}'). \quad (3.67)$$

Equivalent relations hold for our discrete Fourier and spatial basis:

$$\langle \mathbf{k}_{i',j'} | -\frac{\hbar^2 \hat{k}^2}{2m} | \mathbf{k}_{i,j} \rangle = -\frac{\hbar^2 k_{i,j}^2}{2m} \delta_{i'i} \delta_{j'j}, \quad (3.68)$$

$$\langle \mathbf{r}_{i',j'} | V(\hat{\mathbf{r}}) | \mathbf{r}_{i,j} \rangle = V(\mathbf{r}_{i,j}) \delta_{i'i} \delta_{j'j}. \quad (3.69)$$

But to obtain our differential equation, we need all the matrix elements of $\hat{H}_{\text{Schrö}}$ in a single basis.

3.4 Discrete degrees of freedom

3.4.1 The interaction picture

A common situation in atomic physics is to be simulating the internal state of an atom, armed with the knowledge that only a small number of atomic states are able to become occupied.

- Sometimes called a "rotating frame"
- Is equivalent to basis change where new basis functions differ by a time-dependent phase factor
- Is defined by a time-independent Hamiltonian
- This has the effect of moving some time dependence into the operators (demonstrate, by writing some operators with the unitary in front of them. As you can see it is simply a change of basis - but a time-dependent one.)
- No need to remain in the same interaction picture - can be redefined arbitrarily often throughout a simulation and state vectors transformed into new basis.

3.4.2 Unitary integration

Direct exponentiation via diagonalisation of Hamiltonian

- Unitary - doesn't mean it's accurate but means it won't explode. Great for the bits of your simulation that are explosion-prone but don't matter (like regions of space where the wavefunction is near zero but the potential is large or steep)
- error is order [whatever it is] per timestep, not great compared to RK4
- can be combined with RK4 to improve accuracy (see later subsection)

Approximate exponentiation by operator product

[Comment in this section how the approximate total unitary can be used at each timestep to define an interaction picture, and the remaining dynamics simulated with RK4 like RKILIP does in the spatial basis. Interaction pictures are really useful!]

3.5 Continuous degrees of freedom

Every symbol on the paper has a representation in a computer. State vectors are arrays of complex numbers, operators are matrices - differential operators are no exception. Operators must have a concrete representation, their matrix elements can be computed and then things solved with linear algebra. For discrete degrees of freedom, the matrix representation of the operators may be known, for continuous ones you can find the matrix elements once you define what basis functions you will use [show how]. Or, for the DVR it is a little more subtle (because it's not a basis) but still basically the same process.

- Have to be discretised in some way to simulate on a computer - need basis functions. Often a spatial basis is used. Any spatial basis must be combined with assumptions about what the wavefunction is doing at points in between the basis points, in order to define differential operators. Finite differences approximates wavefunction as low-order polynomial in between points (is this equivalent to a polynomial *basis*? Probably not.). Fourier method assume the Fourier series of the wavefunction at the given points can be used to interpolate between points (or the wavefunction can be Fourier transformed and calculations can be done directly in the Fourier basis). DVR is not actually a spatial basis despite appearances. It assumes the wavefunction is a sum of polynomial 'shape functions', but these shape functions are not basis functions as they are not orthonormal. This is why it is called a representation rather than a basis. Regardless, the shape functions can be used to define an interpolation of the wavefunction between points and thus define differential operators.

3.5.1 Finite differences

Show a matrix representation of a few different finite differences, to show that differential operators really are just matrices. They approximate the function as low-order polynomials about each point. You can take them to arbitrarily high order.

3.5.2 The Fourier basis

Because of properties of Fourier transforms, derivatives can be taken in Fourier space as simple multiplication. This is essentially because differential operators are diagonal in the Fourier basis. So you can use this fact to define a differential operator in the spatial basis

```

rev:      31 (422f54686ef6)
author:   Chris Billington <chrisjbillington@gmail.com>
date:     Wed Dec 21 16:32:45 2016 -0500
summary:  Working on split-step methods.

```

[show matrix] ...or, you could just implement it with Fourier transforms, since FFTs are faster than matrix-vector multiplication ($O(n \log(n))$ rather than $O(n^2)$)

Split operator method

- Equivalent to approximate exponentiation via operator product with the discrete case

3.5.3 Harmonic oscillator basis functions

3.6 Finding ground states

3.6.1 Imaginary time evolution

3.6.2 Successive over-relaxation

3.6.3 Generalisation to excited states via Gram–Schmidt orthonormalisation

Directly diagonalising a Hamiltonian can be costly in a spatial basis. Another approach is to find the ground state using one of the above techniques, and then repeat the process, 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 off *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 for each one - this ensures the eigenstates of the Hamiltonian are produced, rather than a different orthogonal 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.

3.7 The finite-element discrete variable representation

- Plots of representation of sine wave as function of number of points between FEDVR and FD. RMS error of a derivative operator perhaps.

[explanation of how it works, comparison of implementation with RSP4 vs something like RK4. RK4 is more general purpose, method does not need to be modified for different Hamiltonians. Main limitation is inability to factor out fast dynamical phases, see RK4ILIP for solution to this. MPI implementation and scaling properties with increasing cores. Superscaling at low number of cores. Mention how my implementation can tolerate high network latency due to early sending of data before all local computations are complete. Mention that it is ripe for GPU processing. Limitations: vulnerable to Runge's phenomenon for sharp potentials. Can't increase the order of the polynomials much because small spacing at the edges requires tiny timesteps. Possible solution: preconditioning the potential to be an approximation better representable in the DVR basis.]

rev: 31 (422f54686ef6)
 author: Chris Billington <chrisjbillington@gmail.com>
 date: Wed Dec 21 16:32:45 2016 -0500
 summary: Working on split-step methods.

This page intentionally left blank

References

- [1] L. M. Bennie, P. B. Wigley, S. S. Szigeti, M. Jasperse, J. J. Hope, L. D. Turner, and R. P. Anderson. *Precise wavefunction engineering with magnetic resonance*. arXiv:1412.6854 [cond-mat, physics:quant-ph] (2014). [p 1]
- [2] A. Görlitz, J. M. Vogels, A. E. Leanhardt, C. Raman, T. L. Gustavson, J. R. Abo-Shaeer, A. P. Chikkatur, S. Gupta, S. Inouye, T. Rosenband, and W. Ketterle. *Realization of Bose-Einstein Condensates in Lower Dimensions*. Physical Review Letters **87**, 130402 (2001). DOI: [10.1103/PhysRevLett.87.130402](https://doi.org/10.1103/PhysRevLett.87.130402). [p 2]
- [3] S. Hofferberth, I. Lesanovsky, B. Fischer, T. Schumm, and J. Schmiedmayer. *Non-equilibrium coherence dynamics in one-dimensional Bose gases*. Nature **449**, 324 (2007). DOI: [10.1038/nature06149](https://doi.org/10.1038/nature06149). [p 2]
- [4] B. Rauer, P. Grišins, I. E. Mazets, T. Schweigler, W. Rohringer, R. Geiger, T. Langen, and J. Schmiedmayer. *Cooling of a One-Dimensional Bose Gas*. Physical Review Letters **116**, 030402 (2016). DOI: [10.1103/PhysRevLett.116.030402](https://doi.org/10.1103/PhysRevLett.116.030402). [p 2]
- [5] Paul D. Lett, Richard N. Watts, Christoph I. Westbrook, William D. Phillips, Phillip L. Gould, and Harold J. Metcalf. *Observation of Atoms Laser Cooled below the Doppler Limit*. Physical Review Letters **61**, 169 (1988). DOI: [10.1103/PhysRevLett.61.169](https://doi.org/10.1103/PhysRevLett.61.169). [p 2]
- [6] J. Dalibard and C. Cohen-Tannoudji. *Laser cooling below the Doppler limit by polarization gradients: simple theoretical models*. JOSA B **6**, 2023 (1989). DOI: [10.1364/JOSAB.6.002023](https://doi.org/10.1364/JOSAB.6.002023). [p 2]
- [7] M. Saffman, T. G. Walker, and K. Mølmer. *Quantum information with Rydberg atoms*. Reviews of Modern Physics **82**, 2313 (2010). DOI: [10.1103/RevModPhys.82.2313](https://doi.org/10.1103/RevModPhys.82.2313). [p 2]
- [8] E. Urban, T. A. Johnson, T. Henage, L. Isenhower, D. D. Yavuz, T. G. Walker, and M. Saffman. *Observation of Rydberg blockade between two atoms*. Nature Physics **5**, 110 (2009). DOI: [10.1038/nphys1178](https://doi.org/10.1038/nphys1178). [p 2]
- [9] B. B. Blinov, Jr R. N. Kohn, M. J. Madsen, P. Maunz, D. L. Moehring, and C. Monroe. *Broadband laser cooling of trapped atoms with ultrafast pulses*. JOSA B **23**, 1170 (2006). DOI: [10.1364/JOSAB.23.001170](https://doi.org/10.1364/JOSAB.23.001170). [p 2]
- [10] A. J. McCulloch, D. V. Sheludko, M. Junker, and R. E. Scholten. *High-coherence picosecond electron bunches from cold atoms*. Nature Communications **4**, 1692 (2013). DOI: [10.1038/ncomms2699](https://doi.org/10.1038/ncomms2699). [p 2]

rev: 31 (422f54686ef6)
author: Chris Billington <chrisbillington@gmail.com>
date: Wed Dec 21 16:32:45 2016 -0500
summary: Working on split-step methods.

- [11] Thomas Brabec and Ferenc Krausz. *Intense few-cycle laser fields: Frontiers of nonlinear optics*. Reviews of Modern Physics **72**, 545 (2000). DOI: [10.1103/RevModPhys.72.545](https://doi.org/10.1103/RevModPhys.72.545). [p 2]
- [12] C. Moler and C. Van Loan. *Nineteen Dubious Ways to Compute the Exponential of a Matrix, Twenty-Five Years Later*. SIAM Review **45**, 3 (2003). DOI: [10.1137/S00361445024180](https://doi.org/10.1137/S00361445024180). [p 4]
- [13] Stefan Weinzierl. *Introduction to Monte Carlo methods*. arXiv:hep-ph/0006269 (2000). [p 6]

rev: 31 (422f54686ef6)
author: Chris Billington <chrisjbillington@gmail.com>
date: Wed Dec 21 16:32:45 2016 -0500
summary: Working on split-step methods.

Word count

Total

Words in text: 23300

Words in headers: 287

Words outside text (captions, etc.): 3280

Number of headers: 75

Number of floats/tables/figures: 26

Number of math inlines: 481

Number of math displayed: 101

Files: 9

Subcounts:

text+headers+captions (#headers/#floats/#inlines/#displayed)

1753+28+325 (9/2/26/8) File(s) total: atomic_physics.tex

725+17+385 (4/4/7/0) File(s) total: experiment.tex

50+0+0 (0/0/0/0) File(s) total: front_matter.tex

1370+19+65 (3/2/39/13) File(s) total: hidden_variables.tex

2043+7+145 (3/1/4/0) File(s) total: introduction.tex

10171+148+910 (35/2/327/76) File(s) total: numerics.tex

4435+33+760 (11/9/1/0) File(s) total: software.tex

2717+22+690 (7/6/77/4) File(s) total: velocimetry.tex

36+13+0 (3/0/0/0) File(s) total: wave_mixing.tex

rev: 31 (422f54686ef6)
author: Chris Billington <chrisjbillington@gmail.com>
date: Wed Dec 21 16:32:45 2016 -0500
summary: Working on split-step methods.