

Efficient methods for linear Schrödinger equation in the semiclassical regime with time-dependent potential

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Abstract

We build efficient and unitary (hence stable) methods for the solution of the linear time-dependent Schrödinger equation with explicitly time-dependent potentials in a semiclassical regime. The Magnus–Zassenhaus schemes presented here are based on a combination of the Zassenhaus decomposition (Bader, Iserles, Kropielnicka & Singh 2014) with the Magnus expansion of the time-dependent Hamiltonian. We conclude with numerical experiments.

1 Introduction

Rapid advances in laser technologies over the recent years have led to a significant progress in the control of systems at the molecular level (Shapiro & Brumer 2003). Pioneering work in the control of chemical systems at the quantum level was done in the study of photo-dissociation and bimolecular reactions. Various control techniques such as the pump-dump quantum control scheme (Kosloff, Rice, Gaspard, Tersigni & Tannor 1989) and the coherent control schemes (Shapiro & Brumer 1997) have had numerous experimental validations and applications (Zhu, Kleiman, Li, Lu, Trentelman & Gordon 1995, Vogt, Nuernberger, Brixner & Gerber 2006).

These experimental successes and a dramatic improvement in our ability to shape femtosecond laser pulses over the recent years has led to a great deal of interest in the development of a systematic way of designing controls (shaping laser pulses) and a requirement for rigorous mathematical analysis of issues such as controllability (Le Bris, Maday & Turinici 2002).

In the case of laser induced breakdown (photo-dissociation) of a molecule, for instance, there is a great deal of interest in designing lasers that achieve efficient

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breakdown. The fact that the dissociation timescales are often themselves in femtoseconds means that it cannot generally be assumed that the laser pulse causes near-instantaneous and efficient excitation of a molecule sitting in the ground state, having no other influence thereafter – the correct dynamics require taking into account the time-dependent nature of the electric potential (laser) throughout the evolution of the wavefunction.

To analyse the control exerted by these lasers we need efficient means of computing the Schrödinger equation featuring time-dependent Hamiltonians, existing strategies for which are either low accuracy or become prohibitively expensive with higher orders of accuracy.

Optimal control schemes for designing laser pulses is often posed as an inverse problem that is solved via optimisation schemes requiring repeated solutions of Schrödinger equations with modified time-dependent Hamiltonians. An ability to efficiently solve these Schrödinger equations with moderately large time steps and high accuracy becomes crucial here, creating a need for high-order methods (Kormann, Holmgren & Karlsson 2008).

In this paper, we are interested in the numerical computation of the linear, time-dependent Schrödinger equation in a semiclassical regime,

$$\partial_t u(x, t) = -\frac{i}{\varepsilon} H u(x, t) = i[\varepsilon \Delta - \varepsilon^{-1} V(x, t)] u(x, t), \quad x \in [-1, 1], \quad t \geq 0, \quad (1.1)$$

equipped with an initial condition $u(x, 0) = u_0(x)$, and periodic boundary conditions. We assume, that the potential function $V(\cdot, t) \in C^\infty[-1, 1]$ is periodic, for $t \geq 0$.

Solving the Schrödinger equation on the torus in this manner is a standard practice in state-of-art numerical methods, particularly when it comes to the semiclassical regime, even though the true problem ought to be solved on \mathbb{R} . The rationale for this choice is the absence of efficient methods for spectral discretisation on \mathbb{R} which is necessitated by the highly oscillatory nature of the solutions. In practice, the typical workaround is to ensure that the boundaries are kept sufficiently far away so that the solution is nearly zero at the boundaries. In this work we additionally assume that the spatial domain is normalised to $[-1, 1]$.

In principle, since the splittings presented in this paper are developed without discretising in space first, it should be possible to combine them with other discretisation strategies and boundary conditions – so long as the the spatial oscillations are resolved correctly and the differential operator has a skew-symmetric discretisation.

The equation (1.1) is posed on a Hilbert space $\mathcal{H} = L_2[-1, 1]$, and the squared modulus of the solution is the probability density of finding the particle in state x at time t . For this reason, the initial condition $u_0(x)$ is normalised to one and it is easy to see that the norm of the solution is an invariant,

$$\|u(x, t)\|_{L^2}^2 = \int_{-1}^1 |u(x, t)|^2 dx = \|u(x, 0)\|_{L^2}^2.$$

The wave function undergoes unitary evolution, which we wish to preserve under discretisation – both because of physical significance, and since, as we mention in Section 4, it implies stability of the numerical method.

The regularity that we require from the potential V depends on the order of desired accuracy of the Magnus–Zassenhaus scheme. However, for convenience we have as-

sumed that it is smooth in its domain. The initial condition is usually a high-frequency wave packet, but even if it is non-oscillatory it can be shown, cf. the analysis in (Jin, Markowich & Sparber 2011), that the solution to this Schrödinger equation is highly oscillatory, with frequency of at least $\mathcal{O}(\varepsilon^{-1})$. This, as a matter of fact, is the main reason why finding an effective numerical method for (1.1) is such a challenging task. In this regime finite differences have been found to require a very fine grid resolution (Markowich, Pietra & Pohl 1999) – instead the usual methodology consists of a semidiscretisation in space via spectral methods followed by an exponential splitting (Bao, Jin & Markowich 2002).

The first step in approximating (1.1) usually is spatial discretisation, which yields the following system of ODEs

$$\mathbf{u}'(t) = i(\varepsilon\mathcal{K}^2 - \varepsilon^{-1}\mathcal{D}_{V(\cdot,t)})\mathbf{u}(t), \quad t \geq 0, \quad (1.2)$$

where \mathcal{K}^2 and $\mathcal{D}_{V(\cdot,t)}$ are $M \times M$ matrices representing the discretisation of second derivative and the multiplication by $V(\cdot, t)$, respectively. We understand, that $\mathbf{u}(t) \in \mathbb{C}^M$ is a vector representing an approximation to the solution (1.1) at time t and $\mathbf{u}(0)$ is derived from the initial conditions.

The exponential midpoint rule, a standard second order method, can be obtained by freezing the matrix $\mathcal{D}_{V(\cdot,t)}$ in the middle of interval $[0, t]$ and applying the Strang splitting,

$$\mathbf{u}(t) = e^{\frac{1}{2}it\varepsilon\mathcal{K}^2}e^{-it\varepsilon^{-1}\mathcal{D}_{V(\cdot,t/2)}}e^{\frac{1}{2}it\varepsilon\mathcal{K}^2}\mathbf{u}(0) + \mathcal{O}(t^3).$$

This splitting has the advantage of separating scales (ε and ε^{-1}) as well as easily computable exponentials. Using spectral collocation or spectral spatial discretisation methods, the matrices \mathcal{K} and $\mathcal{D}_{V(\cdot,t/2)}$ are either diagonal (thus exponentiated directly), or circulant (thus approximated by FFT).

The exponential midpoint rule is the lowest order method among a family of methods based on Magnus expansion that have been found to be effective in computational chemistry (Tal Ezer & Kosloff 1984). Higher order methods of this type are obtained by representing the solution to (1.2) through a Magnus expansion,

$$\mathbf{u}(t) = e^{\Theta(t)}\mathbf{u}(0),$$

where $\Theta(t) \in \mathfrak{u}_M(\mathbb{C})$ is a time-dependent $M \times M$ skew-Hermitian matrix obtained as an infinite series $\sum_{k=1}^{\infty} \Theta^{[k]}(t)$ with each $\Theta^{[k]}(t)$ composed of k nested integrals and commutators of the matrices $i\varepsilon\mathcal{K}^2$ and $i\varepsilon^{-1}\mathcal{D}_V$. In practice finite truncations of this series are used in conjunction with small time steps, h , in order to keep truncation errors small while ensuring convergence of the series.

The convergence of Magnus expansion based methods for solving linear Schrödinger equation under $\varepsilon = 1$ was analysed by Hochbruck & Lubich (2003), where authors conclude that Magnus expansion based methods achieve their prescribed orders of accuracy when the time step h is such, that for some constant c , the inequality $h\|\mathcal{K}\| \leq c$ holds. As we note later in the paper, this forces us to use very small time step of order $\mathcal{O}(\varepsilon)$.

Another serious drawback of this approach lies in the costly approximation of the exponential $e^{\Theta(t)}$. As it occurs, the exponent $\Theta(t)$ ends up to be of a large size (both:

spectral and dimensional), and neither diagonal nor circulant. Indeed, observe, that the highly oscillatory nature of the solution to (1.1) requires a large number of degrees of freedom in the spatial discretisation, $M = \mathcal{O}(\varepsilon^{-1})$. Since the differentiation matrix \mathcal{K} scales as $\mathcal{O}(M) = \mathcal{O}(\varepsilon^{-1})$, the operator $\Theta(t)$, as a sum of nested commutators of $i\varepsilon\mathcal{K}^2$ and $i\varepsilon^{-1}\mathcal{D}_{V(\cdot,t)}$, occurs to be a large matrix which does not possess any favourable structure that could allow an effective approximation of the exponential $\exp(\Theta(t))$.

The standard Magnus–Lanczos schemes which rely on exponentiation of the truncated Magnus expansion via Lanczos iterations become prohibitively expensive in the semiclassical regime due to the large spectral radius while Yōsida type splittings feature an exponential growth in costs as we seek higher orders of accuracy.

Powerful tools like Zassenhaus splitting or Baker–Campbell–Hausdorff formula were historically avoided in splitting methods due to the large computational cost of nested commutators. However as it happens, choosing the correct, infinite-dimensional Lie algebra in case of the Schrödinger vector field, these commutators lose their unwelcome features and enable the derivation of effective, asymptotic splittings.

In (Bader et al. 2014), the current authors established a new framework for a numerical approach to the linear time-dependent problem with an autonomous potential,

$$\partial_t u(x, t) = i[\varepsilon\Delta - \varepsilon^{-1}V(x)]u(x, t), \quad x \in [-1, 1], \quad t \geq 0,$$

where the underlying problem is considered to evolve in a certain Lie group, and the splitting of the linear operator on the right is followed by semidiscretisation. Due to the choice of a suitable Lie algebra, the authors were able to derive a new exponential splitting, that is the *asymptotic exponential splitting* of the following form:

$$e^{ih(\varepsilon\mathcal{K} - \varepsilon^{-1}\mathcal{D})} = e^{\frac{1}{2}W^{[0]}} e^{\frac{1}{2}W^{[1]}} \dots e^{\frac{1}{2}W^{[s]}} e^{\mathcal{W}^{[s+1]}} e^{\frac{1}{2}W^{[s]}} \dots e^{\frac{1}{2}W^{[1]}} e^{\frac{1}{2}W^{[0]}} + \mathcal{O}(\varepsilon^{2s+2}), \quad (1.3)$$

where

$$\begin{aligned} W^{[0]} &= W^{[0]}(h, \varepsilon, \mathcal{K}, \mathcal{D}) = \mathcal{O}(\varepsilon^0), \\ W^{[k]} &= W^{[k]}(h, \varepsilon, \mathcal{K}, \mathcal{D}) = \mathcal{O}(\varepsilon^{2k-2}), \quad k = 1, \dots, s, \\ \mathcal{W}^{[k+1]} &= \mathcal{W}^{[k+1]}(h, \varepsilon, \mathcal{K}, \mathcal{D}) = \mathcal{O}(\varepsilon^{2s}). \end{aligned}$$

Here \mathcal{K} and \mathcal{D} are matrices that approximate the differential operator and multiplication by the potential V , respectively. Such asymptotic exponential splittings derived in (Bader et al. 2014) are superior to standard exponential splittings in a number of ways.

First of all, instead of quantifying the errors in terms of the step size, h , which could have been misleading due to large hidden constants, the errors are quantified in terms of the inherent semiclassical parameter ε , taking into account the $\mathcal{O}(\varepsilon^{-1})$ oscillations characteristic of the Schrödinger equation in the semiclassical regime.

Secondly, these require far fewer exponentials than classical splittings to attain a given order. To be precise, the number of exponentials is shown to grow linearly, rather than exponentially, with the order. Moreover, the exponents decay increasingly more rapidly in powers of ε , yielding an *asymptotic splitting*.

Thirdly, each of these exponentials can be computed fairly easily. The exponents $W^{[0]}$ and $W^{[1]}$ are either diagonal or circulant matrices and their exponentials can be

computed either directly or through FFT, respectively. Remaining exponents are very small and their exponentials can be computed cheaply using low-dimensional Lanczos methods.

The overall cost is quadratic in the desired order, in contrast to the exponential costs of Yōsida type splittings which becomes increasingly prohibitive once the Hamiltonian to be split features more than two terms. Moreover, the resulting Magnus–Zassenhaus schemes are effective for much larger time steps such as $h = \mathcal{O}(\sqrt{\varepsilon})$.

The aim of the paper is to derive asymptotic exponential splittings for Schrödinger equations with time-varying potentials. To develop such a splitting, we must first resort to the Magnus expansion. We follow the approach of (Munthe-Kaas & Owren 1999) in Section 3, discretising the integrals in the Magnus expansion using Gauss-Legendre quadratures. However, unlike the traditional Magnus expansion for ODEs, we work with infinite dimensional operators to evaluate the commutators. To arrive at such a commutator-free expression, we work in the *free Lie algebra* of the infinite dimensional operators ∂_x^2 and V discussed in Section 2. Following the framework of (Bader et al. 2014), a symmetric Zassenhaus splitting is carried out on the commutator-free Magnus expansion, to present, eventually, the Magnus–Zassenhaus scheme of the fifth order (3.29). Obviously, following this derivation one can obtain the method of any desired order, see Table 1. Implementation and numerical examples are discussed in Section 4.

Convergence and unitarity of our method follows from exactly the same argument that was presented in (Bader et al. 2014). Namely it can be easily shown that all the exponents appearing in the derived splitting (3.29) are skew-Hermitian, hence the exponentials are unitary, which suffices for the stability of the method. Now, given consistency of our method (indeed, our scheme will be shown to be of local accuracy much higher than the order one required for the method to be consistent), we can use Lax-equivalence theorem and conclude the convergence of the method.

Realistic systems in quantum chemistry could involve time-dependent matrix-valued, highly oscillatory and stochastic potentials, among others. The first of these will require an extension of our Lie algebraic framework and is under active investigation, while extensions of an alternative scheme that was developed in a recent work (Iserles, Kropielnicka & Singh 2015) could prove promising for oscillatory and low regularity potentials. In this approach the integrals appearing in the Magnus expansion are discretised at the very last stage, following a symmetric Zassenhaus splitting.

2 Lie-group setting

Following the established framework in (Bader et al. 2014), we suppress the dependence on x in (1.1) and analyse the following abstract ODE

$$\partial_t u(t) = \mathcal{A}(t)u(t), \quad u(0) = u_0, \quad (2.4)$$

where $\mathcal{A}(t) := i\varepsilon\partial_x^2 - i\varepsilon^{-1}V(t)$. Because the operator $\mathcal{A}(t)$ belongs to $\mathfrak{u}(\mathcal{H})$, the Lie algebra of (infinite-dimensional) skew-Hermitian operators acting on the Hilbert space \mathcal{H} , its flow is unitary and resides in $\mathcal{U}(\mathcal{H})$ – the Lie group corresponding to $\mathfrak{u}(\mathcal{H})$.

The vector field in the Schrödinger equation is a linear combination of the action of two operators, ∂_x^2 and multiplication by the interaction potential V . Since

our main tools, Magnus expansion and exponential splitting methods, entail nested commutation, we consider the free Lie algebra,

$$\mathfrak{F} = \text{FLA}\{\partial_x^2, V\},$$

i.e., the linear-space closure of all nested commutators generated by ∂_x^2 and V . Following (Bader et al. 2014), we describe their action on sufficiently smooth functions, e.g.

$$[V, \partial_x^2]u = V(\partial_x^2 u) - \partial_x^2(Vu) = -(\partial_x^2 V)u - 2(\partial_x V)\partial_x u$$

which means that $[V, \partial_x^2] = -(\partial_x^2 V) - 2(\partial_x V)\partial_x$. In general, we note that all terms in \mathfrak{F} belong to the set

$$\mathfrak{G} = \left\{ \sum_{k=0}^n y_k(x) \partial_x^k : n \in \mathbb{Z}_+, y_0, \dots, y_n \in C_p^\infty[-1, 1] \right\},$$

where the subscript p means periodicity in $[-1, 1]$. It is trivial to observe that \mathfrak{G} is itself a Lie algebra with the commutator

$$\begin{aligned} \left[\sum_{i=0}^n f_i(x) \partial_x^i, \sum_{j=0}^m g_j(x) \partial_x^j \right] &= \sum_{i=0}^n \sum_{j=0}^m \sum_{\ell=0}^i \binom{i}{\ell} f_i(x) (\partial_x^{i-\ell} g_j(x)) \partial_x^{\ell+j} \\ &\quad - \sum_{j=0}^m \sum_{i=0}^n \sum_{\ell=0}^j \binom{j}{\ell} g_j(x) (\partial_x^{j-\ell} f_i(x)) \partial_x^{\ell+i}. \end{aligned} \quad (2.5)$$

In similar vein to (Bader et al. 2014), we proceed in the pursuit of stability to replace all odd powers of ∂_x that are accompanied by i . The identities,

$$\begin{aligned} y \partial_x &= -\frac{1}{2} \left[\int_0^x y(\xi) d\xi \right] \partial_x^2 - \frac{1}{2} \partial_x y + \frac{1}{2} \partial_x^2 \left[\int_0^x y(\xi) d\xi \cdot \right], \\ y \partial_x^3 &= -(\partial_x y) \partial_x^2 - \frac{1}{4} \left[\int_0^x y(\xi) d\xi \right] \partial_x^4 + \frac{1}{4} \partial_x^3 y - \frac{1}{2} \partial_x^2 [(\partial_x y) \cdot] + \frac{1}{4} \partial_x^4 \left[\int_0^x y(\xi) d\xi \cdot \right], \\ y \partial_x^5 &= \frac{4}{3} (\partial_x^3 y) \partial_x^2 - \frac{5}{3} (\partial_x y) \partial_x^4 - \frac{1}{6} \left[\int_0^x y(\xi) d\xi \right] \partial_x^6 - \frac{1}{2} \partial_x^5 y + \frac{7}{6} \partial_x^2 [(\partial_x^3 y) \cdot] \\ &\quad - \frac{5}{6} \partial_x^4 [(\partial_x y) \cdot] + \frac{1}{6} \partial_x^6 \left[\int_0^x y(\xi) d\xi \cdot \right], \end{aligned}$$

where y is a C^1 function, suffice for our presentation. The general form for expressing $y \partial_x^{2s+1}$ as a linear combination of even derivatives is reported in (Bader et al. 2014).

In the Zassenhaus splitting for time-independent potentials (Bader et al. 2014), the commutators arise solely from the *symmetric Baker–Campbell–Hausdorff formula* where each commutator has an odd number of *letters*. In the case of the Schrödinger equation, where our operators ∂_x^2 and V are each multiplied by i , this translates into an odd power of i for each commutator.

The Magnus expansion, however, does not possess such a desirable structure – it has commutators with odd as well as even number of letters. As a consequence, we

have odd and even powers of i accompanying our terms and it is not enough to blindly replace odd powers of ∂_x . Instead, we replace all odd powers of ∂_x when accompanied by an odd power of i and all even powers of ∂_x when accompanied by an even power of i . A general formula for the replacement of even derivatives by odd derivatives can be proven along similar lines as (Bader et al. 2014). For all practical purposes, however, we only require the identities

$$\begin{aligned} y &= - \left[\int_0^x y(\xi) d\xi \right] \partial_x + \partial_x \left[\int_0^x y(\xi) d\xi \cdot \right], \\ y \partial_x^2 &= - \frac{1}{3} \left[\int_0^x y(\xi) d\xi \right] \partial_x^3 - \frac{2}{3} (\partial_x y) \partial_x - \frac{1}{3} \partial_x [(\partial_x y) \cdot] + \frac{1}{3} \partial_x^3 \left[\int_0^x y(\xi) d\xi \cdot \right], \\ y \partial_x^4 &= - \frac{1}{5} \left[\int_0^x y(\xi) d\xi \right] \partial_x^5 - \frac{4}{3} (\partial_x y) \partial_x^3 + \frac{8}{15} (\partial_x^3 y) \partial_x + \frac{7}{15} \partial_x [(\partial_x^3 y) \cdot] \\ &\quad - \frac{2}{3} \partial_x^3 [(\partial_x y) \cdot] + \frac{1}{5} \partial_x^5 \left[\int_0^x y(\xi) d\xi \cdot \right], \end{aligned}$$

which can be easily verified directly.

Once appropriate odd and even differential operators are replaced, operators of the form $f \partial_x^k + \partial_x^k [f \cdot]$ start appearing ubiquitously in our analysis. Far from being unique to the Magnus expansion, they are characteristic of the free Lie algebra of ∂_x^2 and V – these algebraic forms also appear in Zassenhaus splittings for time-independent potentials (Bader et al. 2014). We introduce a convenient notation,

$$\langle f \rangle_k := f \bullet \partial_x^k = \frac{1}{2} \{f \circ \partial_x^k + \partial_x^k \circ f\} = \frac{1}{2} \{f \partial_x^k + \partial_x^k [f \cdot]\}, \quad f \in C_p^\infty([-1, 1]; \mathbb{R}),$$

where \bullet is the *Jordan product* on the associative algebra of \circ (operatorial composition). In this notation $\langle 1 \rangle_2 = \partial_x^2$ and $\langle V \rangle_0 = V$.

It is worth noting that there is rich algebraic theory behind these structures which will feature in another publication, but not much is lost here by considering these as merely a notational convenience. For the purpose of this work we make observations which can be verified using the machinery of (2.5) in conjunction with the odd and even derivative replacement rules. We present identities which suffice for simplifying all commutators appearing in this work,

$$\begin{aligned} [\langle f \rangle_4, \langle g \rangle_0] &= 4 \langle f(\partial_x g) \rangle_3 - 2 \langle 3(\partial_x f)(\partial_x^2 g) + f(\partial_x^3 g) \rangle_1, \\ [\langle f \rangle_3, \langle g \rangle_0] &= 3 \langle f(\partial_x g) \rangle_2 - \frac{1}{2} \langle 3(\partial_x f)(\partial_x^2 g) + f(\partial_x^3 g) \rangle_0, \\ [\langle f \rangle_2, \langle g \rangle_2] &= 2 \langle f(\partial_x g) - (\partial_x f)g \rangle_3 + \langle 2(\partial_x^2 f)(\partial_x g) - 2(\partial_x f)(\partial_x^2 g) + (\partial_x^3 f)g - f(\partial_x^3 g) \rangle_1, \\ [\langle f \rangle_2, \langle g \rangle_1] &= \langle 2f(\partial_x g) - (\partial_x f)g \rangle_2 - \frac{1}{2} \langle 2(\partial_x f)(\partial_x^2 g) + f(\partial_x^3 g) \rangle_0, \\ [\langle f \rangle_2, \langle g \rangle_0] &= 2 \langle f(\partial_x g) \rangle_1, \\ [\langle f \rangle_1, \langle g \rangle_1] &= \langle f(\partial_x g) - (\partial_x f)g \rangle_1, \\ [\langle f \rangle_1, \langle g \rangle_0] &= \langle f(\partial_x g) \rangle_0. \end{aligned} \tag{2.6}$$

The terms $i\partial_x^2 = i\langle 1 \rangle_2$ and $iV = i\langle V \rangle_0$ reside in

$$\mathfrak{H} = \{i^{k+1} \langle f \rangle_k : f \in C_p^\infty([-1, 1]; \mathbb{R}), k \geq 0\}$$

and, as evident through a few examples in (2.6), all commutators of elements of \mathfrak{H} also reside in \mathfrak{H} . In other words, \mathfrak{H} is a Lie algebra such that

$$\text{FLA}\{\text{i}\partial_x^2, \text{i}V\} \subseteq \mathfrak{H},$$

and it suffices to work directly in \mathfrak{H} using the rules (2.6) instead of proceeding via (2.5) followed by the odd-even derivative replacement rules.

For a real valued f , $\langle f \rangle_k$ is symmetric if k is even and skew-symmetric otherwise. This property is preserved under discretisation once we use spectral collocation on a uniform grid. In that case ∂_x is discretised as a skew-symmetric matrix \mathcal{K} and V is discretised as a diagonal matrix \mathcal{D}_V . The term $\langle f \rangle_k$ is discretised as $(\mathcal{K}^k \mathcal{D}_f + \mathcal{D}_f \mathcal{K}^k) / 2$ which is clearly symmetric when k is even and skew-symmetric otherwise. Consequently, elements of \mathfrak{H} such as $\text{i}^{k+1} \langle f \rangle_k$, which are skew-Hermitian operators, discretise to skew-Hermitian matrices of the form $\text{i}^{k+1} (\mathcal{K}^k \mathcal{D}_f + \mathcal{D}_f \mathcal{K}^k) / 2$.

This structural property of \mathfrak{H} is responsible for unitary evolution and numerical stability of our schemes since exponentials of skew-Hermitian matrices are unitary.

Definition 1 *The height of a term is defined as*

$$\text{ht} \left(\sum_{i=0}^n \alpha_i \langle f_i \rangle_{k_i} \right) = \max\{k_1, \dots, k_n\}.$$

These terms benefit from a remarkable property of *height reduction* which is stated here without proof,

$$\text{ht}([\langle f \rangle_k, \langle g \rangle_l]) \leq k + l - 1.$$

For the commutators relevant to this work, this property can be verified by a quick inspection of the identities (2.6).

For the largest part, our work will proceed in the language of the undiscretised operators introduced in this section. At the very last stage we will resort to spectral collocation on the uniform grid over $[-1, 1]$ for spatial discretisation. For this purpose we will need at least $M = \mathcal{O}(\varepsilon^{-1})$ points since (regardless of initial conditions) the solution of the Schrödinger equation develops spatial oscillations of order $\mathcal{O}(\varepsilon^{-1})$ (Jin et al. 2011, Bao et al. 2002). Consequently, \mathcal{K} scales like $\mathcal{O}(\varepsilon^{-1})$ and $(\mathcal{K}^k \mathcal{D}_f + \mathcal{D}_f \mathcal{K}^k) / 2 = \mathcal{O}(\varepsilon^{-k})$. Keeping eventual discretisation in mind, we abuse notation and write $\langle f \rangle_k = \mathcal{O}(\varepsilon^{-k})$.

More formally, following (Bao et al. 2002) we assume that the solution $u(t)$, which is known to feature $\mathcal{O}(\varepsilon^{-1})$ oscillations, obeys the bounds,

$$\|\partial_t^m \partial_x^k u(x, t)\| \leq C_{m,k} \varepsilon^{-m-k}, \quad t \in [0, T]. \quad (2.7)$$

In this context

$$\|\langle f \rangle_k u\| = \frac{1}{2} \|(\mathcal{K}^k \mathcal{D}_f + \mathcal{D}_f \mathcal{K}^k) u\| \leq \|V\|_\infty C_k \varepsilon^{-k} = \mathcal{O}(\varepsilon^{-k}).$$

Although it is possible to work in a more rigorous language throughout, the shorthand $\langle f \rangle_k = \mathcal{O}(\varepsilon^{-k})$ is indeed seen to be based on firm theoretical grounds while simplifying exposition greatly. We also remind the reader that the growth of derivatives of the

potential, while certainly effecting error constants in our splittings (and therefore of concern in the context of moderately small values of ε), are irrelevant in the asymptotic limit of $\varepsilon \rightarrow 0$ since they don't scale with ε and don't effect the asymptotic analysis carried out here.

The property of height reduction leads to a systematic decrease in the size of terms with commutation,

$$[\langle f \rangle_k, \langle g \rangle_l] = \mathcal{O}(\varepsilon^{-k-l+1}).$$

Going further, we want to analyse all terms in the common *currency* of the inherent semiclassical parameter ε and assume that our choice of the time-step, h , is governed by $h = \mathcal{O}(\varepsilon^\sigma)$, for some $0 < \sigma \leq 1$. Larger values of σ correspond to very small time steps and are best avoided.

3 The solution

3.1 The Magnus expansion

To look for the solution of (2.4) one needs to take into account some features of the operator $\mathcal{A}(t)$. First of all it depends on time and it cannot be assumed that its values in different points of time commute, i.e. we assume that $[\mathcal{A}(t_1), \mathcal{A}(t_2)] \neq 0$ and give up the hope that the solution is of the simple form $e^{\int_0^t \mathcal{A}(\xi) d\xi} u_0$. Secondly $\mathcal{A}(t)$ evolves in a Lie algebra so the solution of (2.4) resides in a corresponding Lie group. Both properties can be dealt with elegantly using the famous result from (Magnus 1954) by writing the solution as single exponential,

$$u(t) = e^{\Theta(t)} u(0), \quad (3.8)$$

where the infinite series $\Theta(t) = \sum_{k=1}^{\infty} \Theta_k(t)$, also called as *Magnus expansion*, is an element of the underlying Lie algebra. Its convergence has been shown in (Iserles & Nørsett 1999), (Moan & Niesen 2008), (Hochbruck & Lubich 2003) for sufficiently small time-steps. Obviously we truncate this series and advance with adequately small time step h

$$u(t+h) = e^{\Theta(t,t+h)} u(t), \quad (3.9)$$

starting from the initial step,

$$u(h) = e^{\Theta(0,h)} u(0), \quad (3.10)$$

where we understand that the operator $e^{\Theta(t,t+h)}$ is a flow evolving the solution from t to $t+h$. Let us observe now, that the aim of the paper consists in a derivation of the asymptotic exponential splitting for a certain function of type $e^{\Theta(t,t+h)}$. This means, that the algorithm we are going to present will advance in small time steps h (exactly like the method (1.3) does). For the clarity of exposition, however, we will focus on the first step of Magnus expansion, i.e. (3.10), noting that (3.9), when required for any time window $[t, t+h]$, is easily recovered from (3.10) by a straightforward translation of the vector field $\mathcal{A}(\xi)$ to $\mathcal{A}(t+\xi)$. For convenience we shorten the notation, writing $\Theta(h)$ instead of $\Theta(h, 0)$.

Simple differentiation of the *ansatz* in (3.10) together with elementary algebra, see (Iserles & Nørsett 1999) or (Blanes, Casas, Oteo & Ros 2009) for details, lead to the conclusion that the exponent $\Theta(t)$ satisfies the *dexpinv equation*,

$$\dot{\Theta}(h) = \text{dexp}_{\Theta(h)}^{-1} \mathcal{A}(h) = \sum_{k=0}^{\infty} \frac{B_k}{k!} \text{ad}_{\Theta(h)}^k \mathcal{A}(h), \quad \Theta(0) = 0, \quad (3.11)$$

where B_k are Bernoulli numbers ($B_0 = 1$, $B_1 = -\frac{1}{2}$, $B_2 = \frac{1}{6}$, $B_3 = 0$, $B_4 = -\frac{1}{30}$, $B_5 = 0$, $B_6 = \frac{1}{42}$) and the adjoint representation is defined recursively by $\text{ad}_A^0 V = V$ and $\text{ad}_A^{k+1} V = [A, \text{ad}_A^k V]$. The solution of (3.11) is an infinite series and can be obtained using Picard iterations. It was proposed in (Magnus 1954) and widely analysed in (Iserles & Nørsett 1999, Iserles, Munthe-Kaas, Nørsett & Zanna 2000, Blanes et al. 2009).

The first few terms of the Magnus expansion ordered by size in h are

$$\begin{aligned} \Theta(h) = & \int_0^h \mathcal{A}(\xi) d\xi - \frac{1}{2} \int_0^h \int_0^{\xi_1} [\mathcal{A}(\xi_2), \mathcal{A}(\xi_1)] d\xi \\ & + \frac{1}{12} \int_0^h \int_0^{\xi_1} \int_0^{\xi_1} [\mathcal{A}(\xi_2), [\mathcal{A}(\xi_3), \mathcal{A}(\xi_1)]] d\xi \\ & + \frac{1}{4} \int_0^h \int_0^{\xi_1} \int_0^{\xi_2} [[\mathcal{A}(\xi_3), \mathcal{A}(\xi_2)], \mathcal{A}(\xi_1)] d\xi + \dots \end{aligned} \quad (3.12)$$

We say that a multivariate integral of a nested commutator, \mathcal{I} , is of grade m if $\mathcal{I} = \mathcal{O}(h^m)$ for every smooth \mathcal{A} . Truncating the Magnus expansion at grade p to $\Omega_p(h) = \Theta(h) + \mathcal{O}(h^{p+1})$, preserves time symmetry (Iserles et al. 2000), (Iserles, Nørsett & Rasmussen 2001). Time symmetry means that not only the exact flow φ , but also the numerical flow $\Phi = e^{\Omega_p(h)}$, satisfy

$$\varphi(h, 0) \circ \varphi(0, h) = I, \quad \Phi_{h,0} \circ \Phi_{0,h} = I. \quad (3.13)$$

As one can observe, the time symmetry of the numerical flow is equivalent to the fact that

$$\Omega_p(0, h) = -\Omega_p(h, 0). \quad (3.14)$$

Time symmetry is a desirable feature because truncation by power with odd p leads to a gain of an extra unit of order, see (Iserles et al. 2000). This means that if we aim for a numerical method of order six it suffices to consider the truncation of the Magnus expansion only to the terms listed in (3.12).

3.2 Magnus expansion in practice

It turns out that the multivariate integrals can be efficiently computed using simple univariate quadrature rules of Munthe-Kaas & Owren (1999). We will follow their approach and evaluate the potential at the Gauss-Legendre quadrature points ($t_1 = \frac{1}{2} - \frac{\sqrt{15}}{10}$, $t_2 = \frac{1}{2}$, $t_3 = \frac{1}{2} + \frac{\sqrt{15}}{10}$) which is then transformed (Iserles & Nørsett 1999) to

obtain a far less costly quadrature. As a result, to obtain order six approximation, all the effort of approximation of the solution boils down to the following formula

$$\begin{aligned} \Theta(h) = & B_1 + \frac{1}{12}B_3 - \frac{1}{12}[B_1, B_2] + \frac{1}{240}[B_2, B_3] + \frac{1}{360}[B_1, [B_1, B_3]] \\ & - \frac{1}{240}[B_2, [B_1, B_2]] + \frac{1}{720}[B_1, [B_1, [B_1, B_2]]] + \mathcal{O}(h^7), \end{aligned} \quad (3.15)$$

where

$$B_1 = h\mathcal{A}(t_2), \quad B_2 = \frac{\sqrt{15}}{3}h(\mathcal{A}(t_3) - \mathcal{A}(t_1)), \quad B_3 = \frac{10}{3}h(\mathcal{A}(t_3) - 2\mathcal{A}(t_2) + \mathcal{A}(t_1)). \quad (3.16)$$

See (Iserles et al. 2000) and (Blanes et al. 2009) for comprehensive information and ways to approximate the Magnus expansion using different quadrature rules and to higher orders. The former could be relevant if the time-dependent potential is only known at certain grid-points as might be the case in some control setups.

Substituting $\mathcal{A}(t)$ with the given Hamiltonian as $\mathcal{A}(t) = -iH(t)/\varepsilon$ and working in the free Lie algebra \mathfrak{H} , we can derive a commutator free expansion using the identities (2.6). Keeping the notation of the previous section in mind, we approximate the time derivatives of the potential by central differences, cf. (3.16),

$$V_0 = V(t_2), \quad V_1 = \frac{\sqrt{15}}{3h}(V(t_3) - V(t_1)), \quad V_2 = \frac{10}{3h^2}(V(t_3) - 2V(t_2) + V(t_1)),$$

so that

$$B_1 = ih\varepsilon\partial_x^2 - ih\varepsilon^{-1}V_0, \quad B_2 = ih^2\varepsilon^{-1}V_1, \quad B_3 = ih^3\varepsilon^{-1}V_2.$$

Once these are substituted in (3.15), we use the identities (2.6) along with the observation that $\partial_x^2 = \langle 1 \rangle_2$ and $V_j = \langle V_j \rangle_0$ to arrive at a Magnus expansion in the format $\sum_k i^{k+1}c_k \langle f_k \rangle_k$ with $c_k \in \mathbb{Q}$ and $f_k \in C_p^\infty([-1, 1]; \mathbb{R})$.

The grade one commutators of the self-adjoint basis appearing in (3.15), for instance, can be simplified as follows,

$$\begin{aligned} [B_1, B_2] &= [ih\varepsilon\langle 1 \rangle_2 - ih\varepsilon^{-1}\langle V_0 \rangle_0, -ih^2\varepsilon^{-1}\langle V_1 \rangle_0] \\ &= h^3[\langle 1 \rangle_2, \langle V_1 \rangle_0] = 2h^3\langle \partial_x V_1 \rangle_1, \end{aligned} \quad (3.17)$$

$$\begin{aligned} [B_1, B_3] &= [ih\varepsilon\langle 1 \rangle_2 - ih\varepsilon^{-1}\langle V_0 \rangle_0, -ih^3\varepsilon^{-1}\langle V_2 \rangle_0] \\ &= h^4[\langle 1 \rangle_2, \langle V_2 \rangle_0] = 2h^4\langle \partial_x V_2 \rangle_1, \end{aligned} \quad (3.18)$$

$$\begin{aligned} [B_2, B_3] &= [-ih^2\varepsilon^{-1}\langle V_1 \rangle_0, ih^3\varepsilon^{-1}\langle V_2 \rangle_0] \\ &= 0. \end{aligned} \quad (3.19)$$

Consequently, the grade two commutators appearing in (3.15) are,

$$\begin{aligned}
[B_1, [B_1, B_3]] &= [ih\varepsilon \langle 1 \rangle_2 - ih\varepsilon^{-1} \langle V_0 \rangle_0, 2h^4 \langle \partial_x V_2 \rangle_1] \\
&= 2ih^5\varepsilon [\langle 1 \rangle_2, \langle \partial_x V_2 \rangle_1] + 2ih^5\varepsilon^{-1} [\langle \partial_x V_2 \rangle_1, \langle V_0 \rangle_0] \\
&= 2ih^5\varepsilon (2 \langle \partial_x^2 V_2 \rangle_2 - \frac{1}{2} \langle \partial_x^4 V_2 \rangle_0) + 2ih^5\varepsilon^{-1} \langle (\partial_x V_2)(\partial_x V_0) \rangle_0 \\
&= 4ih^5\varepsilon \langle \partial_x^2 V_2 \rangle_2 - ih^5\varepsilon \langle \partial_x^4 V_2 \rangle_0 + 2ih^5\varepsilon^{-1} \langle (\partial_x V_2)(\partial_x V_0) \rangle_0, \quad (3.20)
\end{aligned}$$

$$\begin{aligned}
[B_2, [B_1, B_2]] &= [-ih^2\varepsilon^{-1} \langle V_1 \rangle_0, 2h^3 \langle \partial_x V_1 \rangle_1] \\
&= 2ih^5\varepsilon^{-1} [\langle \partial_x V_1 \rangle_1, \langle V_1 \rangle_0] \\
&= 2ih^5\varepsilon^{-1} \langle (\partial_x V_1)^2 \rangle_0, \quad (3.21)
\end{aligned}$$

$$\begin{aligned}
[B_1, [B_1, B_2]] &= [ih\varepsilon \langle 1 \rangle_2 - ih\varepsilon^{-1} \langle V_0 \rangle_0, 2h^3 \langle \partial_x V_1 \rangle_1] \\
&= 2ih^4\varepsilon [\langle 1 \rangle_2, \langle \partial_x V_1 \rangle_1] + 2ih^4\varepsilon^{-1} [\langle \partial_x V_1 \rangle_1, \langle V_0 \rangle_0] \\
&= 2ih^4\varepsilon (2 \langle \partial_x^2 V_1 \rangle_2 - \frac{1}{2} \langle \partial_x^4 V_1 \rangle_0) + 2ih^4\varepsilon^{-1} \langle (\partial_x V_1)(\partial_x V_0) \rangle_0 \\
&= 4ih^4\varepsilon \langle \partial_x^2 V_1 \rangle_2 - ih^4\varepsilon \langle \partial_x^4 V_1 \rangle_0 + 2ih^4\varepsilon^{-1} \langle (\partial_x V_1)(\partial_x V_0) \rangle_0. \quad (3.22)
\end{aligned}$$

The only grade three commutator that we need is

$$\begin{aligned}
[B_1, [B_1, [B_1, B_2]]] &= [ih\varepsilon \langle 1 \rangle_2 - ih\varepsilon^{-1} \langle V_0 \rangle_0, \quad (3.23) \\
&\quad 4ih^4\varepsilon \langle \partial_x^2 V_1 \rangle_2 - ih^4\varepsilon \langle \partial_x^4 V_1 \rangle_0 + 2ih^4\varepsilon^{-1} \langle (\partial_x V_1)(\partial_x V_0) \rangle_0] \\
&= -4h^5\varepsilon^2 [\langle 1 \rangle_2, \langle \partial_x^2 V_1 \rangle_2] + h^5\varepsilon^2 [\langle 1 \rangle_2, \langle \partial_x^4 V_1 \rangle_0] \\
&\quad - 2h^5 [\langle 1 \rangle_2, \langle (\partial_x V_1)(\partial_x V_0) \rangle_0] - 4h^5 [\langle \partial_x^2 V_1 \rangle_2, \langle V_0 \rangle_0] \\
&= -4h^5\varepsilon^2 (2 \langle \partial_x^3 V_1 \rangle_3 - \langle \partial_x^5 V_1 \rangle_1) + 2h^5\varepsilon^2 \langle \partial_x^5 V_1 \rangle_1 \\
&\quad - 4h^5 \langle (\partial_x^2 V_1)(\partial_x V_0) + (\partial_x V_1)(\partial_x^2 V_0) \rangle_1 - 8h^5 \langle (\partial_x^2 V_1)(\partial_x V_0) \rangle_1 \\
&= -8h^5\varepsilon^2 \langle \partial_x^3 V_1 \rangle_3 + 3h^5\varepsilon^2 \langle \partial_x^5 V_1 \rangle_1 - h^5 \langle 12(\partial_x^2 V_1)(\partial_x V_0) + 4(\partial_x V_1)(\partial_x^2 V_0) \rangle_1.
\end{aligned}$$

Substituting (3.17–3.23) in (3.15) gives us a truncated Magnus expansion for the Schrödinger equation (1.1) in the Lie algebra \mathfrak{H} ,

$$\begin{aligned}
\Omega_5 &= ih\varepsilon \langle 1 \rangle_2 - ih\varepsilon^{-1} \langle V_0 \rangle_0 - \frac{1}{12} ih^3\varepsilon^{-1} \langle V_2 \rangle_0 - \frac{1}{6} h^3 \langle \partial_x V_1 \rangle_1 \\
&\quad + \frac{1}{360} (4ih^5\varepsilon \langle \partial_x^2 V_2 \rangle_2 - ih^5\varepsilon \langle \partial_x^4 V_2 \rangle_0 + 2ih^5\varepsilon^{-1} \langle (\partial_x V_2)(\partial_x V_0) \rangle_0) \\
&\quad - \frac{1}{120} ih^5\varepsilon^{-1} \langle (\partial_x V_1)^2 \rangle_0 + \frac{1}{720} (-8h^5\varepsilon^2 \langle \partial_x^3 V_1 \rangle_3 + 3h^5\varepsilon^2 \langle \partial_x^5 V_1 \rangle_1 \\
&\quad \quad - h^5 \langle 12(\partial_x^2 V_1)(\partial_x V_0) + 4(\partial_x V_1)(\partial_x^2 V_0) \rangle_1) \quad (3.24)
\end{aligned}$$

$$\begin{aligned}
\Omega_5 &= \overbrace{ih\varepsilon\partial_x^2 - ih\varepsilon^{-1}V_0}^{\mathcal{O}(\varepsilon^{\sigma-1})} - \overbrace{\frac{1}{12}ih^3\varepsilon^{-1}V_2 - \frac{1}{6}h^3\langle\partial_x V_1\rangle_1}^{\mathcal{O}(\varepsilon^{3\sigma-1})} \\
&\quad + \overbrace{\frac{1}{360}ih^5\varepsilon^{-1}\left(2(\partial_x V_2)(\partial_x V_0) - 3(\partial_x V_1)^2\right)}^{\mathcal{O}(\varepsilon^{5\sigma-1})} - \overbrace{\frac{1}{180}h^5\langle(\partial_x V_1)(\partial_x^2 V_0) + 3(\partial_x V_0)(\partial_x^2 V_1)\rangle_1}^{\mathcal{O}(\varepsilon^{5\sigma-1})} \\
&\quad + \overbrace{\frac{1}{90}ih^5\varepsilon\langle\partial_x^2 V_2\rangle_2 - \frac{1}{90}h^5\varepsilon^2\langle\partial_x^3 V_1\rangle_3}^{\mathcal{O}(\varepsilon^{5\sigma-1})} - \overbrace{\frac{1}{360}ih^5\varepsilon(\partial_x^4 V_2) + \frac{1}{240}h^5\varepsilon^2\langle\partial_x^5 V_1\rangle_1}^{\mathcal{O}(\varepsilon^{5\sigma+1})} \\
&= \Theta + \mathcal{O}(\varepsilon^{7\sigma-1}).
\end{aligned} \tag{3.25}$$

For $\sigma \leq 1$, the last two terms in Ω_5 , which are $\mathcal{O}(\varepsilon^{5\sigma+1})$, become $\mathcal{O}(\varepsilon^{7\sigma-1})$ and can be discarded. After discarding these terms, the Magnus expansion reduces to

$$\begin{aligned}
\Omega_5 &= \overbrace{ih\varepsilon\partial_x^2 - ih\varepsilon^{-1}V_0}^{\mathcal{O}(\varepsilon^{\sigma-1})} - \overbrace{\frac{1}{12}ih^3\varepsilon^{-1}V_2 - \frac{1}{6}h^3\langle\partial_x V_1\rangle_1}^{\mathcal{O}(\varepsilon^{3\sigma-1})} \\
&\quad + \overbrace{\frac{1}{360}ih^5\varepsilon^{-1}\left(2(\partial_x V_2)(\partial_x V_0) - 3(\partial_x V_1)^2\right)}^{\mathcal{O}(\varepsilon^{5\sigma-1})} \\
&\quad - \overbrace{\frac{1}{180}h^5\langle(\partial_x V_1)(\partial_x^2 V_0) + 3(\partial_x V_0)(\partial_x^2 V_1)\rangle_1}^{\mathcal{O}(\varepsilon^{5\sigma-1})} \\
&\quad + \overbrace{\frac{1}{90}ih^5\varepsilon\langle\partial_x^2 V_2\rangle_2 - \frac{1}{90}h^5\varepsilon^2\langle\partial_x^3 V_1\rangle_3}^{\mathcal{O}(\varepsilon^{5\sigma-1})} = \Theta + \mathcal{O}(\varepsilon^{7\sigma-1}).
\end{aligned} \tag{3.26}$$

We note that, due to the property of height reduction discussed in Section 2, a grade n commutator in the Magnus expansion of $\mathcal{A}(t)$ should be $\mathcal{O}(\varepsilon^{n\sigma-1})$. This can indeed be verified in the above expansion. Asymptotically speaking, in terms of ε , the terms in the expansion are decreasing in size with increasing n for any $\sigma > 0$, so that convergence of the Magnus expansion also occurs for much larger time steps such as $h = \mathcal{O}(\varepsilon^{1/2})$ or $h = \mathcal{O}(\varepsilon^{1/4})$. This is a considerable improvement over existing analysis.

Since Ω_5 includes the term $ih\varepsilon\partial_x^2 - ih\varepsilon^{-1}V_0$, its exponential is, at the very least, as troublesome to approximate as the problem of solving the Schrödinger equation with time-independent potential. Fortunately the Zassenhaus procedure is sufficiently flexible and can tackle such modified Hamiltonians with ease.

3.3 The Zassenhaus splitting algorithm

Let us recall the basic principle for the iterative symmetric Zassenhaus splitting (Bader et al. 2014). Our goal is to compute $e^{\mathcal{W}^{[0]}}$, where $\mathcal{W}^{[0]} = X + Y$ and $X, Y = \mathcal{O}(\varepsilon^p)$. Using the *symmetric Baker-Campbell-Hausdorff* (sBCH) formula (Dynkin 1947, Casas & Murua 2009), we then write

$$e^{\mathcal{W}^{[0]}} = e^{\frac{1}{2}X} e^{\text{sBCH}(-X, \mathcal{W}^{[0]})} e^{\frac{1}{2}X}. \quad (3.27)$$

Grade-three commutators of X and Y are at most $\mathcal{O}(\varepsilon^{3p})$ so that $\text{sBCH}(-X, X + Y) = Y + \mathcal{O}(\varepsilon^{3p})$. Thus we have extracted X from the exponent at the cost of correction terms in form of higher-order commutators. Assuming that the corrections are decreasing in size, it is then enough to identify the largest terms as $W^{[1]}$ in the central exponent $\mathcal{W}^{[1]} = \text{sBCH}(-X, \mathcal{W}^{[0]})$ and to continue the iteration until the desired accuracy is reached,

$$\mathcal{W}^{[k+1]} = \text{sBCH}(-W^{[k]}, \mathcal{W}^{[k]}), \quad \mathcal{W}^{[0]} = X + Y. \quad (3.28)$$

In this notation, the splitting after s steps can be written as

$$\exp(X + Y) = e^{\frac{1}{2}W^{[0]}} e^{\frac{1}{2}W^{[1]}} \dots e^{\frac{1}{2}W^{[s]}} e^{\mathcal{W}^{[s+1]}} e^{\frac{1}{2}W^{[s]}} \dots e^{\frac{1}{2}W^{[1]}} e^{\frac{1}{2}W^{[0]}}.$$

We emphasize that, in principle, we can freely choose the elements $W^{[k]}$ that we want to extract. Except for some special cases, at least one of the exponents in this splitting will feature an infinite series of terms. To construct a finite splitting scheme featuring a certain accuracy we may discard, at each stage, all terms smaller than the desired threshold.

Assuming that a grade k commutator of X and Y scales as $\mathcal{O}(\varepsilon^{kp})$, convergence of the series requires $p > 0$ at the very least. In the case of the Schrödinger equation, we choose $X = i\hbar\varepsilon\partial_x^2$, $p = \sigma - 1$, and this naively translates to a very stringent time step restriction: $\sigma > 1$. However, the remarkable feature of height reduction means that a grade k commutator in this context scales as $\mathcal{O}(\varepsilon^{kp+(k-1)})$ and convergence requirements become significantly milder: we need $p+1 > 0$ which translates to $\sigma > 0$.

3.4 Applying the Zassenhaus algorithm to Magnus expansions

We perform a Zassenhaus splitting on Ω_5 , choosing to extract the largest terms – analysed in powers of ε – first. We commence the splitting with $W^{[0]} = i\hbar\varepsilon\partial_x^2$, although we could equally well choose $W^{[0]} = -i\hbar\varepsilon^{-1}V_0$, for instance, and arrive at a variant of the splitting presented here. The exponent to be split is $\mathcal{W}^{[0]} = \Omega_5$ and the first step involves computing the sBCH formula. Here, once again, the rules of the free Lie

algebra \mathfrak{H} , (2.6), suffice for arriving at a commutator free expression,

$$\begin{aligned}
\mathcal{W}^{[1]} &= \text{sBCH}(-W^{[0]}, \mathcal{W}^{[0]}) \\
&= -\overbrace{ih\varepsilon^{-1}V_0}^{\mathcal{O}(\varepsilon^{\sigma-1})} + \overbrace{\frac{1}{12}ih^3\varepsilon^{-1}\left(2(\partial_x V_0)^2 - V_2\right) - \frac{1}{6}h^3\langle\partial_x V_1\rangle_1 + \frac{1}{6}ih^3\varepsilon\langle\partial_x^2 V_0\rangle_2}^{\mathcal{O}(\varepsilon^{3\sigma-1})} \\
&\quad - \overbrace{\frac{1}{24}ih^3\varepsilon(\partial_x^4 V_0)}^{\mathcal{O}(\varepsilon^{3\sigma+1})} - \overbrace{\frac{1}{360}ih^5\varepsilon^{-1}\left(8(\partial_x V_0)^2(\partial_x^2 V_0) + 3(\partial_x V_1)^2 - 12(\partial_x V_2)(\partial_x V_0)\right)}^{\mathcal{O}(\varepsilon^{5\sigma-1})} \\
&\quad + \overbrace{\frac{1}{30}h^5\langle 2(\partial_x V_0)(\partial_x^2 V_1) - (\partial_x V_1)(\partial_x^2 V_0)\rangle_1}^{\mathcal{O}(\varepsilon^{5\sigma-1})} \\
&\quad - \overbrace{\frac{1}{720}ih^5\varepsilon\langle 127(\partial_x V_0)(\partial_x^3 V_0) + 130(\partial_x^2 V_0)^2 - 18(\partial_x^2 V_2)\rangle_2}^{\mathcal{O}(\varepsilon^{5\sigma-1})} \\
&\quad + \overbrace{\frac{1}{60}h^5\varepsilon^2\langle\partial_x^3 V_1\rangle_3 - \frac{13}{90}ih^5\varepsilon^3\langle\partial_x^4 V_0\rangle_4}^{\mathcal{O}(\varepsilon^{5\sigma-1})} + \mathcal{O}(\varepsilon^{7\sigma-1}).
\end{aligned}$$

At the second stage we select the largest remaining element $W^{[1]} = -ih\varepsilon^{-1}V_0$, whereby

$$\begin{aligned}
\mathcal{W}^{[2]} &= \text{sBCH}(-W^{[1]}, \mathcal{W}^{[1]}) \\
&= \overbrace{\frac{1}{12}ih^3\varepsilon^{-1}\left(2(\partial_x V_0)^2 - V_2\right) - \frac{1}{6}h^3\langle\partial_x V_1\rangle_1 + \frac{1}{6}ih^3\varepsilon\langle\partial_x^2 V_0\rangle_2}^{\mathcal{O}(\varepsilon^{3\sigma-1})} \\
&\quad - \overbrace{\frac{1}{24}ih^3\varepsilon(\partial_x^4 V_0)}^{\mathcal{O}(\varepsilon^{3\sigma+1})} - \overbrace{\frac{1}{360}ih^5\varepsilon^{-1}\left(13(\partial_x V_0)^2(\partial_x^2 V_0) + 3(\partial_x V_1)^2 - 12(\partial_x V_2)(\partial_x V_0)\right)}^{\mathcal{O}(\varepsilon^{5\sigma-1})} \\
&\quad + \overbrace{\frac{1}{30}h^5\langle 2(\partial_x V_0)(\partial_x^2 V_1) - (\partial_x V_1)(\partial_x^2 V_0)\rangle_1}^{\mathcal{O}(\varepsilon^{5\sigma-1})} \\
&\quad - \overbrace{\frac{1}{720}ih^5\varepsilon\langle 127(\partial_x V_0)(\partial_x^3 V_0) + 130(\partial_x^2 V_0)^2 - 18(\partial_x^2 V_2)\rangle_2}^{\mathcal{O}(\varepsilon^{5\sigma-1})} \\
&\quad + \overbrace{\frac{1}{60}h^5\varepsilon^2\langle\partial_x^3 V_1\rangle_3 - \frac{13}{90}ih^5\varepsilon^3\langle\partial_x^4 V_0\rangle_4}^{\mathcal{O}(\varepsilon^{5\sigma-1})} + \mathcal{O}(\varepsilon^{7\sigma-1}).
\end{aligned}$$

We terminate the procedure by letting $W^{[2]}$ consist of the $\mathcal{O}(\varepsilon^{3\sigma-1})$ terms in $\mathcal{W}^{[2]}$ and are left with $\mathcal{O}(\varepsilon^{5\sigma-1})$ and $\mathcal{O}(\varepsilon^{3\sigma+1})$ terms in $\mathcal{W}^{[3]} = \mathcal{W}^{[2]} - W^{[2]}$ once we ignore $\mathcal{O}(\varepsilon^{7\sigma-1})$ terms. Since $\mathcal{O}(\varepsilon^{3\sigma+1})$ terms can be subsumed into the $\mathcal{O}(\varepsilon^{5\sigma-1})$ terms

for $\sigma \leq 1$, combining them in this way is not a cause for concern. The outcome is the Magnus–Zassenhaus splitting,

$$e^{\Omega_5} = e^{\frac{1}{2}W^{[0]}} e^{\frac{1}{2}W^{[1]}} e^{\frac{1}{2}W^{[2]}} e^{\mathcal{W}^{[3]}} e^{\frac{1}{2}W^{[2]}} e^{\frac{1}{2}W^{[1]}} e^{\frac{1}{2}W^{[0]}} + \mathcal{O}(\varepsilon^{7\sigma-1}), \quad (3.29)$$

with

$$\begin{aligned} W^{[0]} &= i\varepsilon h \partial_x^2 = \mathcal{O}(\varepsilon^{\sigma-1}), \\ W^{[1]} &= -i\varepsilon^{-1} h V_0 = \mathcal{O}(\varepsilon^{\sigma-1}), \\ W^{[2]} &= \frac{1}{12} i\varepsilon^{-1} h^3 (2(\partial_x V_0)^2 - V_2) - \frac{1}{6} h^3 \langle \partial_x V_1 \rangle_1 + \frac{1}{6} i\varepsilon h^3 \langle \partial_x^2 V_0 \rangle_2 = \mathcal{O}(\varepsilon^{3\sigma-1}), \\ \mathcal{W}^{[3]} &= -\frac{1}{24} i\varepsilon h^3 (\partial_x^4 V_0) - \frac{1}{360} i\varepsilon^{-1} h^5 (13(\partial_x V_0)^2 (\partial_x^2 V_0) + 3(\partial_x V_1)^2 - 12(\partial_x V_2)(\partial_x V_0)) \\ &\quad + \frac{1}{30} h^5 \langle 2(\partial_x V)(\partial_x^2 V_1) - (\partial_x V_1)(\partial_x^2 V) \rangle_1 \\ &\quad - \frac{1}{720} i\varepsilon h^5 \langle 127(\partial_x V_0)(\partial_x^3 V_0) + 130(\partial_x^2 V_0)^2 - 18(\partial_x^2 V_2) \rangle_2 \\ &\quad + \frac{1}{60} \varepsilon^2 h^5 \langle \partial_x^3 V_1 \rangle_3 - \frac{13}{90} i\varepsilon^3 h^5 \langle \partial_x^4 V_0 \rangle_4 = \mathcal{O}(\varepsilon^{5\sigma-1}). \end{aligned}$$

4 A numerical scheme

As we have seen, the derivation of the method has two components. First, we choose the desired order of accuracy in the small parameter ε and compute the Magnus expansion up to this order Ω_p . This will lead to an effective exponent of the form (3.15), detailed steps for which can be found in (Munthe-Kaas & Owren 1999, Blanes et al. 2009). Commencing from these expansions we compute the commutator-free Magnus expansion using the rules (2.6) of the Lie algebra \mathfrak{H} . Once we have computed this effective Hamiltonian, we start the Zassenhaus algorithm, detailed in Table 1.

Table 1: Zassenhaus algorithm

Initialise:	
$\mathcal{W}^{[0]} := \Omega_{2p+1}$, exploiting time-symmetry for a local error of $\mathcal{O}(h^{2p+3})$	
$W^{[0]} := i\varepsilon h \partial_x^2$, $k = 0$	
<hr/>	
while $k \leq p$	
$\mathcal{W}^{[k+1]}$	$:= \text{sBCH}(-W^{[k]}, \mathcal{W}^{[k]})$
$\mathcal{W}^{[k+1]} - W^{[k+1]}$	$:= \mathcal{O}(\varepsilon^{(2k+1)\sigma-1})$ implicitly defines $W^{[k+1]}$
k	$:= k + 1$
end while	
Final method:	
$e^{\Omega_{2p+1}} = e^{\frac{1}{2}W^{[0]}} e^{\frac{1}{2}W^{[1]}} \dots e^{\frac{1}{2}W^{[p]}} e^{W^{[p+1]}} e^{\frac{1}{2}W^{[p]}} \dots e^{\frac{1}{2}W^{[1]}} e^{\frac{1}{2}W^{[0]}} + \mathcal{O}(\varepsilon^{(2p+3)\sigma-1})$	

For numerical realisation of these splittings schemes it is typical to impose periodic boundary conditions in order to resolve spatial oscillations with spectral accuracy.

Recall that we restrict the domain to $[-1, 1]$, imposing periodic boundaries at $x = \pm 1$. We discretise using spectral collocation on the equispaced grid $\{x_n\}_{n=-N}^N$, $x_n = n/(N + \frac{1}{2})$, $|n| \leq N$, where $M = 2N + 1$ is the number of grid points. The unknowns are $u_n \approx u(x_n)$, $|n| \leq N$. The differential operator ∂_x is discretised as a circulant matrix \mathcal{K} and V as a diagonal \mathcal{D}_V .

All exponents in our splitting (3.29) are of the form $i^{k+1} \langle f \rangle_k$ and are discretised as skew-Hermitian matrices,

$$i^{k+1} \langle f \rangle_k \rightsquigarrow i^{k+1} (\mathcal{D}_f \mathcal{K}^k + \mathcal{K}^k \mathcal{D}_f) / 2.$$

Since the exponential of a skew-Hermitian matrix is unitary, unitary evolution and (consequently) unconditional stability of the method are guaranteed.

The outermost exponentials $W^{[0]}$ and $W^{[1]}$ are replaced by the circulant $i\varepsilon h \mathcal{K}^2$ and the diagonal matrix $-i\varepsilon^{-1} h \mathcal{D}_{V_0}$, respectively. The lowest order scheme of the type (3.29) can be obtained by ignoring the exponents $W^{[2]}$ and $\mathcal{W}^{[3]}$ from (3.29),

$$\mathbf{u}^{n+1} = e^{\frac{1}{2} i h \varepsilon \mathcal{K}^2} e^{-i h \varepsilon \mathcal{D}_{V_0}} e^{\frac{1}{2} i h \varepsilon \mathcal{K}^2} \mathbf{u}^n,$$

which features an $\mathcal{O}(\varepsilon^{3\sigma-1})$ error since the largest term ignored is $W^{[2]}$. Clearly, this is the exponential midpoint rule which resorts to a Strang splitting after freezing the potential in the middle of the interval.

The exponential of the circulant matrix $\frac{1}{2} i \varepsilon h \mathcal{K}^2$ is evaluated to machine precision using Fast Fourier Transform (FFT) in $\mathcal{O}(M \log M)$ operations while the diagonal matrix is exponentiated directly in $\mathcal{O}(M)$ operations.

The first non-trivial splitting is obtained upon including $W^{[2]}$ once more,

$$\mathbf{u}^{n+1} = e^{\frac{1}{2} i h \varepsilon \mathcal{K}^2} e^{-\frac{1}{2} i h \varepsilon \mathcal{D}_{V_0}} e^{\tilde{W}^{[2]}} e^{-\frac{1}{2} i h \varepsilon \mathcal{D}_{V_0}} e^{\frac{1}{2} i h \varepsilon \mathcal{K}^2} \mathbf{u}^n,$$

where

$$\begin{aligned} \tilde{W}^{[2]} = & \frac{1}{12} i \varepsilon^{-1} h^3 \mathcal{D}_{2(\partial_x V_0)^2 - V_2} - \frac{1}{12} h^3 (\mathcal{D}_{\partial_x V_1} \mathcal{K} + \mathcal{K} \mathcal{D}_{\partial_x V_1}) \\ & + \frac{1}{12} i \varepsilon h^3 (\mathcal{D}_{\partial_x V_1} \mathcal{K}^2 + \mathcal{K}^2 \mathcal{D}_{\partial_x^2 V_0}), \end{aligned}$$

is the discretised version of $W^{[2]}$. This splitting commits an error of $\mathcal{O}(\varepsilon^{5\sigma-1})$. We remind the reader that if the pursuit of a splitting with an $\mathcal{O}(\varepsilon^{5\sigma-1})$ error was the objective, it would suffice to start with Ω_3 , which is a lower order Magnus truncation that is easier to obtain and uses merely two Gauss–Legendre quadrature knots, while to obtain splittings that are higher order than the $\mathcal{O}(\varepsilon^{7\sigma-1})$ splitting given in (3.29) we would need to commence with a higher order Magnus expansion than those discussed here.

The exponents $W^{[2]}$ and $\mathcal{W}^{[3]}$ appearing in the non-trivial splittings do not possess a structure amenable to exact exponentiation. However, they are very small – $\mathcal{O}(\varepsilon^{3\sigma-1})$ and $\mathcal{O}(\varepsilon^{5\sigma-1})$, respectively. For $\sigma = 1$, the most costly case we consider, the exponentials of these terms can be evaluated to $\mathcal{O}(\varepsilon^6)$ accuracy using merely three and two Lanczos iterations, respectively (Bader et al. 2014). For $\sigma = \frac{1}{2}$, their exponentiation to an accuracy of $\mathcal{O}(\varepsilon^{5/2})$ requires five and two Lanczos iterations, respectively.

These iterations involve the computation of $\tilde{W}^{[2]}v$, which can be achieved using a few FFTs (remember that \mathcal{K}^k in $\mathcal{D}_{\partial_x V_1} \mathcal{K} + \mathcal{K} \mathcal{D}_{\partial_x V_1}$ and $\mathcal{D}_{\partial_x V_1} \mathcal{K}^2 + \mathcal{K}^2 \mathcal{D}_{\partial_x V_1}$ is a circulant).

We refer the curious reader to (Bader et al. 2014, Iserles et al. 2015) where semi-discretisation strategies, stability analysis and exponentiation methods are addressed in greater detail.

4.1 A numerical example

Consider the evolution of the wave-packet

$$u_0(x) = (\delta\pi)^{-1/4} \exp\left(ik_0 \frac{(x-x_0)}{\delta} - \frac{(x-x_0)^2}{2\delta}\right)$$

with $x_0 = -0.3$ and $k_0 = 0.1$ specifying the initial mean position and momentum, respectively, heading towards the lattice potential

$$V_0 = \rho(4x) \sin(20\pi x),$$

where

$$\rho(x) = \begin{cases} \exp(-1/(1-x^2)) & \text{for } |x| < 1, \\ 0 & \text{otherwise,} \end{cases}$$

is a bump function.

In Figure 4.1, we show the evolution of the wavepacket from $u(0)$ to $u(T)$ at $T = 0.75$ under the influence of the time-independent potential V_0 alone. In these examples, we use $\delta_1 = 10^{-3}$ and $\delta_2 = 10^{-4}$ (specifying the localisation of the wavepacket) while the semiclassical parameter is chosen to be $\varepsilon_1 = 2^{-8} \approx 0.0039$ and $\varepsilon_2 = 5 \times 10^{-4}$ – the same order of magnitude as the choices of δ . The corresponding wavepackets are labeled u^1 and u^2 , respectively.

When we excite the wavepacket using an additional time-varying potential,

$$E(x, t) = \rho(3t - 1)\rho(\sin(2\pi(x - t))),$$

so that the wave packet evolves under $V_E(x, t) = V_0(x) + E(x, t)$, a significantly larger part of the wave packet is able to make it across the lattice to the right hand side (see $u_E^1(T)$ and $u_{2E}^2(T)$ in Figure 4.1).

The excitation pulse is not active for the entire duration since $\rho(3t - 1)$ acts as a smooth envelope simulating the switching on and off of the excitation. The excited potential is evident at $t = T/2$ in Figure 4.1 (*bottom left*).

In Figure 4.2 (*left*) we present the global error at time $T = 0.75$ in the propagation of u_0^1 to u_E^1 under the influence of V_E using the scheme (3.29). Under the scaling $\sigma = 1$, we commit a local L_2 error of $\mathcal{O}(\varepsilon^6)$ per time step in the splitting scheme (3.29). Since the number of time steps is $\mathcal{O}(\varepsilon^{-\sigma})$, the global error is $\mathcal{O}(\varepsilon^5)$. The precise scaling used in our experiments for $\sigma = 1$ is $M \sim 5\varepsilon^{-1}$ and $h \sim \varepsilon/2$.

We remind the reader that the Magnus–Zassenhaus methods developed in this paper are highly effective in the semiclassical regime where ε becomes very small. Moreover, the accuracy of these methods is analysed asymptotically. Where ε is moderate and considerations of spatial accuracy are governed by initial conditions

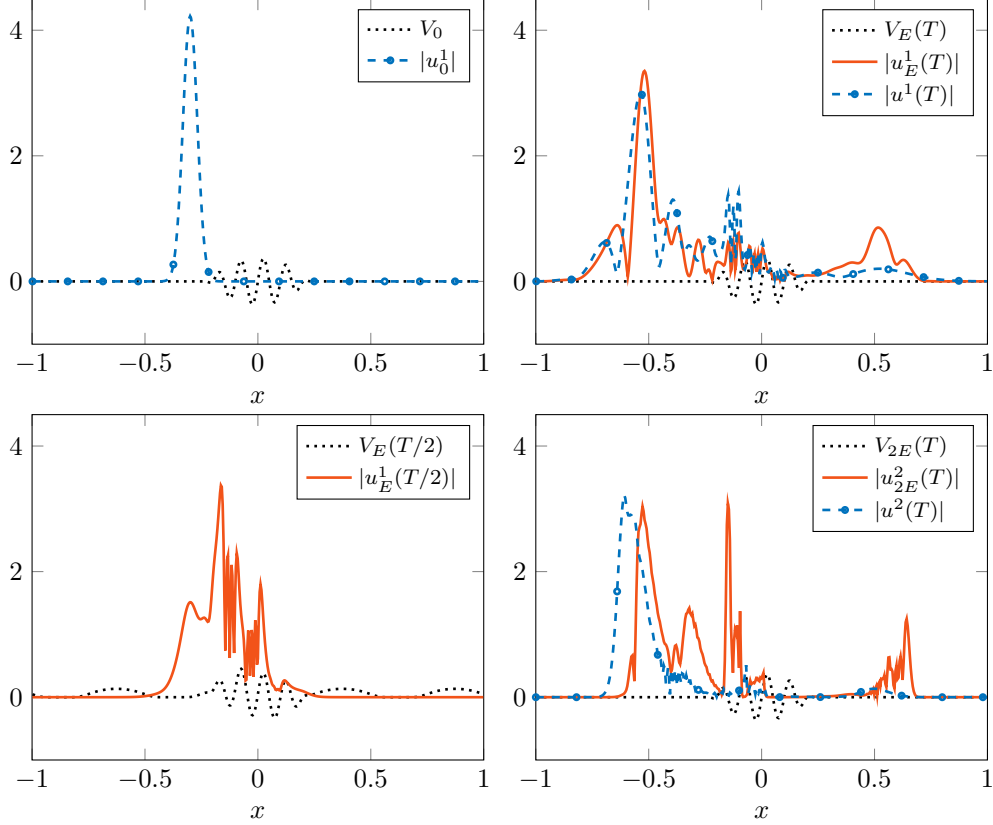


Figure 4.1: Initial wave-packet u_0^1 with localisation $\delta_1 = 10^{-3}$ (top left); final wave-packets at time $T = 0.75$ under $\varepsilon_1 = 2^{-8}$ (top right) – $u^1(T)$ under the influence of V_0 alone and $u_E^1(T)$ under the influence of $V_E(x, t) = V_0 + E(x, t)$. Intermediate values of u_E^1 and V_E at $T/2$ (bottom left); final wave-packets at time $T = 0.75$ under $\varepsilon_2 = 5 \times 10^{-4}$ starting from u_0^2 with localisation $\delta_2 = 10^{-4}$ (top right) – $u^2(T)$ under the influence of V_0 alone and $u_{2E}^2(T)$ under the influence of $V_{2E}(x, t) = V_0 + 2E(x, t)$.

rather than by the high oscillations induced by a small ε , the number of grid points M may end being much larger than ε^{-1} . In this case different scaling laws for the time step h might be found to be optimal and the number of Lanczos iterations required will need to be re-analysed. While this analysis is beyond the scope of this paper, in principle it should be possible to develop effective Magnus–Zassenhaus schemes for moderate ε by utilising the techniques discussed here.

In Figure 4.2 (right) we depict the global L_2 errors in u_E^1 at $T = 0.75$ under different choices of ε , letting the time step vary independently of ε . The number of grid points for this experiment was fixed at $M = 1001$. As expected, the accuracy of these methods is found to be higher when ε is of the same order as $1/M$ (or $M = \mathcal{O}(\varepsilon^{-1})$). Disregarding the semiclassical parameter and the spectral size of the differentiation matrices, the local accuracy of these methods should be $\mathcal{O}(h^5)$

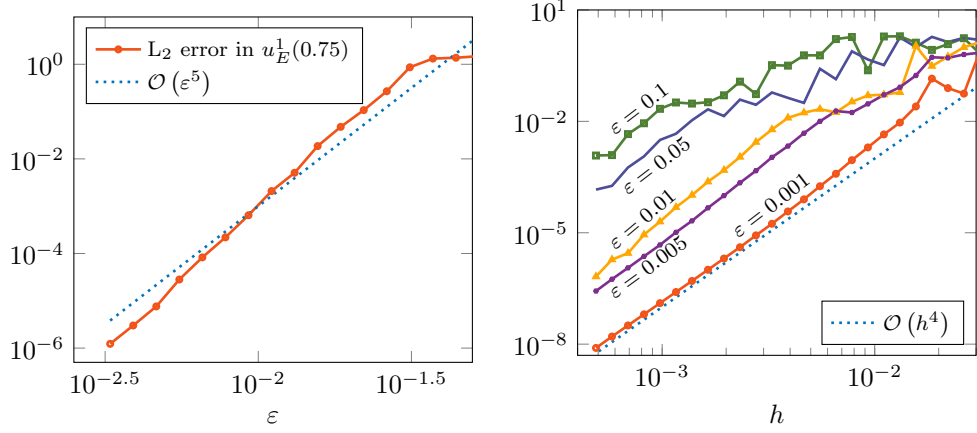


Figure 4.2: Global L_2 errors in u_E^1 at $T = 0.75$ under $\sigma = 1$ (left) and for different choices of ε with time step h chosen independently of ε (right).

while the global accuracy should be $\mathcal{O}(h^4)$. We also note that the three and two Lanczos iterations utilised here for exponentiating $W^{[2]}$ and $\mathcal{W}^{[3]}$, respectively, are wholly inadequate under the choice $\varepsilon = 0.1$ and $M = 1001$, for instance.

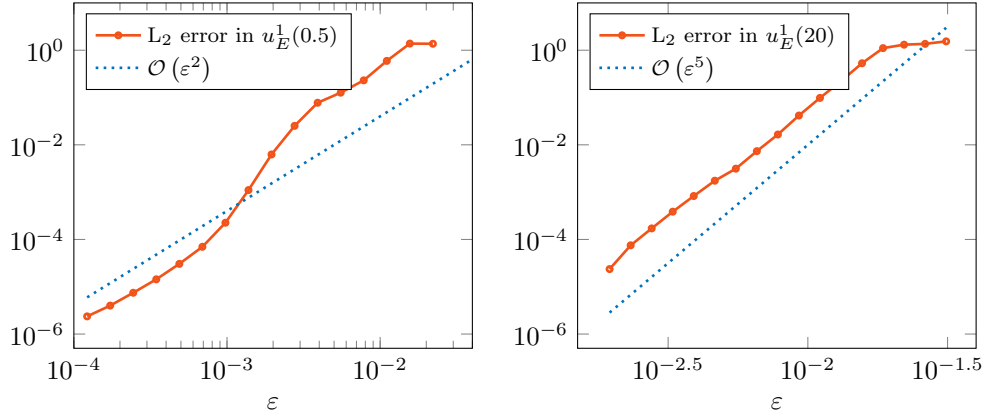


Figure 4.3: Global L_2 errors in u_E^1 at $T = 0.5$ under $\sigma = \frac{1}{2}$ (left) and at $T = 20$ under $\sigma = 1$ (right).

In Figure 4.3 (left), we present the errors under the scaling $\sigma = \frac{1}{2}$. Here, the local L_2 error is $\mathcal{O}(\varepsilon^{5/2})$ while the global error is $\mathcal{O}(\varepsilon^2)$. For experiments under $\sigma = \frac{1}{2}$ we use a scaling $M \sim \varepsilon$ and $h \sim \sqrt{\varepsilon}/4$, keeping in mind that, at the very least, the potential function needs to be correctly resolved. This experiment demonstrates that the Magnus–Zassenhaus methods can achieve very high accuracies under the semiclassical regime despite using relatively larger time steps of order $h = \mathcal{O}(\sqrt{\varepsilon})$.

4.2 Finding a reference solution

Since no analytic solution of (1.1) is available, reference solutions must also be obtained via a numerical approach. We obtain the reference solution \mathbf{u}_R for our numerical experiments by resorting to the exponential midpoint rule,

$$\mathbf{u}_R(t + h_R) = e^{\frac{1}{2}ih_R\varepsilon\mathcal{K}^2} e^{-ih_R\varepsilon^{-1}\mathcal{D}_V(t+h_R/2)} e^{\frac{1}{2}ih_R\varepsilon\mathcal{K}^2} \mathbf{u}_R(t),$$

where $\mathbf{u}_R \in \mathbb{C}^{M_R}$ lives on a much finer grid than the solution of (3.29). In each of the T/h_R time steps required for finding the solution $\mathbf{u}_R(T)$, the potential is frozen in the middle of the interval $[t, t + h_R]$ and a Strang splitting is used.

Since the exponential midpoint rule is also the lowest order method in the Magnus–Zassenhaus family of schemes, we require very small time steps for convergence – certainly $h_R \ll h$ is required for the reference solution to possess an error smaller than the scheme (3.29) whose error we are attempting to quantify.

We rely on this method for producing reliable reference solutions since it is simple and its error is easily analysed. Directly exponentiating a Hamiltonian (via MATLAB’s `expm`, for instance) with potential frozen at the middle of the interval is more expensive but no more accurate than the Strang splitting – this is because freezing the potential is akin to disregarding the nested integrals and commutators in the Magnus expansion which are of the same size (in powers of ε) as the error committed in the Strang splitting.

Another factor we must take into account is the growth of spatial oscillations with decreasing ε . To capture this, starting from $M_R = 3M = 15\varepsilon^{-1}$, we iteratively increase the grid resolution for the reference solution till no high frequencies are clipped and convergence is achieved. In the end, the spatial resolution used for obtaining a reference solution is much greater than that used for (3.29), $M_R \gg M \sim 5\varepsilon^{-1}$.

Using such a low order method for generating reference solutions to a high degree of accuracy in a brute force manner means generating reference solutions is orders of magnitude slower than the splitting method (3.29) requiring validation. The exorbitant cost of reference solutions is what restricts rigorous experimental study of numerical errors to moderate values of ε and T .

Having provided experimental evidence for the $\mathcal{O}(\varepsilon^5)$ global accuracy of the splitting (3.29) under $\sigma = 1$, we use (3.29) with very small time steps and fine grid resolution (finer than prescribed by $\sigma = 1$) for generating reference solutions while analysing the accuracy of our splitting under $\sigma = \frac{1}{2}$ (Figure 4.3 (*left*)).

In Figure 4.3 (*right*) we explore the convergence rate for larger values of T by resorting to the splitting (3.29) for generating reference solutions by using higher spatio-temporal resolutions than prescribed by our scaling laws.

5 Competing interests

We have no competing interests.

6 Authors' contributions

All the authors have contributed equally to this work.

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