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Why the quantitative condition fails to reveal quantum adiabaticity

Dafa Li^{1,2} and Man-Hong Yung^{3,4}

¹Department of Mathematical Sciences, Tsinghua University, Beijing 100084, People's Republic of China

²Center for Quantum Information Science and Technology, Tsinghua National Laboratory for Information Science and Technology (TNL), Beijing 100084, People's Republic of China

³Center for Quantum Information, Tsinghua University, Beijing 100084, People's Republic of China

⁴Department of Chemical and Chemical Biology, Harvard University, Cambridge MA, USA
E-mail: dli@ah.si.gha.ed.ac and m.g.singha@gmail.com

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Abstract

The quantitative adiabatic condition (QAC), a quantitative condition, is a necessary (*a priori*) sufficient condition for adiabaticity, but it is not sufficient. However, the quantitative condition for QAC is not sufficient. It has been shown that QAC can be efficient for guaranteeing the validity of the adiabatic approximation, but the quantitative condition for QAC could be necessary, but not sufficient. Furthermore, it is believed that the validity of the QAC condition is a necessary condition for adiabaticity, but it is not sufficient. However, it is clear that a sufficient condition is not sufficient. Hence, the quantitative condition is not sufficient for adiabaticity, but it is necessary. As a posteriori condition for adiabaticity, the quantitative condition is not sufficient for adiabaticity, but it is necessary. The quantitative condition for QAC could be efficient for guaranteeing the validity of the adiabatic approximation, but it is not sufficient. Hence, the quantitative condition is not sufficient for adiabaticity, but it is necessary.

Keywords: adiabatic condition, quantitative adiabatic condition, adiabatic approximation



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1. Introduction

The rapid adiabatic theorem (QAT) [1, 2] suggests that physical systems initialized in an eigenstate $|E_n(t=0)\rangle$ (corresponding to the ground state) of a time-varying Hamiltonian $H(t)$, will remain in the eigenstate $|E_n(t)\rangle$ (corresponding to the ground state) of $H(t)$ at all times provided the Hamiltonian $H(t)$ varies sufficiently slowly. The adiabatic theorem is first proved by Born and Fock in the context of quantum electrodynamics [3], where it is used to justify the adiabatic approximation of the Dirac equation [4]. Born and Fock's original proof of the adiabatic theorem is based on the perturbation theory, e.g. the Rayleigh-Schrödinger perturbation theory; however, it is not applicable to systems with degeneracy, e.g. the hydrogen atom. This is corrected by Kato in 1950 [5], where he introduces the adiabatic theorem for the case of degeneracy. Kato's original proof is based on the perturbation theory, e.g. the Rayleigh-Schrödinger perturbation theory, where he uses the projection operator technique to deal with the degeneracy. Later, the proof of the adiabatic theorem is extended to the case of degeneracy [6].

This rigorous physical proof of the rapid adiabatic theorem is significant in the field of quantum computing [7], quantum control [8], quantum information theory [9], e.g. the quantum state transfer [10, 11], adiabatic quantum computing [12–18], quantum simulation (see e.g. the review [19]), and quantum control [20].

1.1. Quantitative adiabatic condition

Despite its long history, the QAT is still a challenging field for each. Many works have been devoted to improving the accuracy of the adiabatic theorem. In particular, the possible failure of the slow adiabatic theorem has been carefully studied. Traditionally [1, 2, 14, 21] the so-called (e.g. see [22]) rapid adiabatic condition (QAC) is usually used to describe the failure of the QAT:

$$\left| \frac{\langle E_m(t) | \dot{E}_n(t) \rangle}{E_m(t) - E_n(t)} \right| \ll 1, \tag{1}$$

where $|E_m(t)\rangle$ and $|E_n(t)\rangle$ are the eigenstates of $H(t)$ (see also the definition of the Hamiltonian and eigenstates). However, QAC is not a rigorous condition because it is difficult to evaluate the fidelity of the final state [23]. Furthermore, it has been shown that QAC is not sufficient in the QAT [24] and is not sufficient for the adiabatic approximation [22], especially for the case of degeneracy [25]. The accuracy of the adiabatic approximation of the QAC is discussed in [22] for the case of degeneracy.

where A and B are Hermitian operators and $H_a(t)$ is the Hamiltonian

$$H_b(t) = -U_a(t)^\dagger H_a(t) U_a(t) \tag{2}$$

where B is related to the system A through a unitary transformation $U_a(t) = T \exp\left(-i \int_0^t H_a(t') dt'\right)$ where T is the time-ordering operator. It is shown that the system A and B satisfy the QAC, but the system B does not satisfy the QAC, but the system A does.

adiabatic approximation. This classification is consistent with the results referred to in a NMR review [26].

1.2. Related studies in the literature

Major studies (e.g. [27–33]) have been made regarding the adiabatic theorem established by Majorana and Sados [24]. It is argued [29–33] that a series of perturbations can be applied to the system to achieve the adiabatic theorem. A refined adiabatic condition has been found [34], which takes into account the effects of a perturbation on the system.

On the other hand, the validity of the adiabatic theorem is also addressed by the asymptotic approach [35, 36], which provides a diagnostic tool for the adiabatic condition and yields the adiabatic condition (1) as the first-order approximation. It is argued [37] that the adiabatic condition is sufficient for the adiabatic approximation in the Hilbert space and has all the properties. Furthermore, generalizing QAC to the case of systems [27, 38] and a broader case [39] have been achieved. Efforts for finding conditions that can replace QAC are made [40–42].

Another line of research related to the adiabatic theorem is the scaling of the fidelity of the final state. Under the general condition of Hilbert space, it is found [43] that the series of perturbations scales as $O(1/T^2)$ for a large T . When the adiabatic condition is satisfied, it is shown [44, 45] that the fidelity of the final state is given by the expression $L \equiv \int_0^1 \|\partial_r \psi(r)\| dr$, where $r(t)$ is a piecewise linear approximation to the Hilbert space.

2. Motivation

Instead of asking the validity of the adiabatic condition as a diagnostic for the adiabaticity, we ask the question: ‘Under what additional conditions would QAC become necessary?’.

The answer to this question has been clear [46–48]. One can be interested in the condition for the QAC to be necessary, the necessary condition for the adiabatic state $|\psi_n^{adi}(t)\rangle$ [8, 46], which is defined by a single eigenvalue of the Hamiltonian $H(t)$ (see also the Berry phase [8]) $e^{i\beta_n(t)}$ the eigenstate $|E_n(t)\rangle$, i.e.,

$$|\psi_n^{adi}(t)\rangle \equiv e^{i\beta_n(t)} |E_n(t)\rangle, \tag{3}$$

where

$$\beta_n(t) \equiv - \int_0^t E_n(x) dx + i \int_0^t \langle E_n(x) | \dot{E}_n(x) \rangle dx. \tag{4}$$

The key result obtained in [46] is that the series of perturbations is a linear combination $c_m(t) = \langle E_m(t) | \psi(t) \rangle$ of the eigenstate $|E_m(t)\rangle$ and it is given by the following series:

Table 1. Summary of definitions and notations.

Terms	Meaning
Quantum adiabatic theorem	This theorem states that for general physical systems initially in an eigenstate (e.g. ground state) of the unperturbed Hamiltonian, the system remains in the corresponding eigenstate of the perturbed Hamiltonian if the Hamiltonian is sufficiently slow.
Quantum adiabatic condition (QAC) (or quantum adiabatic condition)	A condition for the validity of the adiabatic approximation (see eq. (1)), i.e., $\left \frac{\langle E_m(t) \dot{E}_n(t) \rangle}{E_m(t) - E_n(t)} \right \ll 1.$
Adiabatic approximation	An approximation where the exact state $ \psi(t)\rangle$ in the adiabatic case is $ \psi_n^{adi}(t)\rangle$, which leads to $ c_m(t) \ll 1$ for all $m \neq n$.
Adiabatic state $ \psi_n^{adi}(t)\rangle$	Defined by $ \psi_n^{adi}(t)\rangle \equiv e^{i\beta_n(t)} E_n(t)\rangle$, where $ E_n(t)\rangle$ is the n th eigenstate (with eigenvalue $E_n(t)$) of the Hamiltonian $H(t)$, and $\beta_n(t) \equiv -\int_0^t E_n(x) dx + i \int_0^t \langle E_n(x) \dot{E}_n(x) \rangle dx$ (see eq. (3)).
Difference vector $ D(t)\rangle$	Defined by $ D(t)\rangle \equiv \psi(t)\rangle - \psi_n^{adi}(t)\rangle$, the difference between the exact state $ \psi(t)\rangle$ and the adiabatic state $ \psi_n^{adi}(t)\rangle$ (see eq. (8)).

$$c_m(t) \approx ie^{i\beta_n(t)} \frac{\langle E_m(t) | \dot{E}_n(t) \rangle}{E_m(t) - E_n(t)}, \tag{5}$$

which leads to the condition for the adiabatic approximation to be valid, i.e., the probability amplitude c_m for all eigenstates $m \neq n$ are small, $|c_m(t)| \ll 1$, hence the QAC (cf eq. (1)) essentially holds (for a summary of the definitions of terms, see table (1)).

For convenience, as illustrated in (1), we give the notation of Schiff [1] as

$$c_m(t) \approx \frac{\langle E_m(t) | \dot{H} | E_n(t) \rangle}{i(E_m(t) - E_n(t))^2} (e^{i(E_m - E_n)t} - 1), \tag{6}$$

$$= i \frac{\langle E_m(t) | \dot{E}_n(t) \rangle}{E_m(t) - E_n(t)} (e^{i(E_m - E_n)t} - 1). \tag{7}$$

The definition of the fidelity has been provided in section B. These expressions (eqs (5) and (7)) predict the validity of the adiabatic approximation when the quantum adiabatic condition (cf eq. (1)) is satisfied.

Here, the result [46] is crucial [47, 48]. Zha and Wu [47] argued that the condition for the validity of the result [46] is derived. Chua [48] pointed out that the condition of the result [46] leads to a contradiction if the system is adiabatic. This result is a direct consequence of the result [49] established by the authors in the adiabatic case. In fact, the result [46] is not valid.

3. Summary of results

We use the following results:

1. See the definition of the linear [46–49] by the definition of the exactness (cf eq. (9)) for the system $c_m(t)$, which is a special case of the result [46] (cf eq. (5)), and
2. The result of the exactness (cf eq. (22)) for eq. (5), which holds for the exactness of the QAC and the adiabatic condition.

Our approach can be formulated in the following: the definition of $|D(t)\rangle$, which is defined by the difference between the exact state $|\psi(t)\rangle$ and the adiabatic state $|\psi_n^{adi}(t)\rangle$ (cf eq. (3)):

$$|D(t)\rangle \equiv |\psi(t)\rangle - |\psi_n^{adi}(t)\rangle. \tag{8}$$

3.1. Our main result and its consequences

Our main result is an exactness for $c_m(t)$, and the condition is eq. (5)

$$c_m(t) = \underbrace{ie^{i\theta_n(t)} \frac{\langle E_m(t) | \dot{E}_n(t) \rangle}{E_m(t) - E_n(t)}}_{\equiv Q_m(t) \text{ Result in [46]}} - \underbrace{E_n(t) \frac{\langle E_m(t) | D(t) \rangle}{E_m(t) - E_n(t)} + i \frac{\langle E_m(t) | \dot{D}(t) \rangle}{E_m(t) - E_n(t)}}_{\equiv R_m(t) \text{ Correction terms}}, \tag{9}$$

which reduces to the result [46] (cf eq. (5)) when the magnitude of $|R_m(t)|$ for the exactness is small; for example, when $|D(t)\rangle = 0$ and $|\dot{D}(t)\rangle = 0$. These conditions are satisfied by the exactness [46]. For the result [46] (cf eq. (5)) is a special case of the result [46].

The exactness in eq. (9) implies the following results:

- The result [46] is satisfied by the system if $|D(t)\rangle$ and $|\dot{D}(t)\rangle$ can be ignored, i.e., $|D(t)\rangle \approx 0$ and $|\dot{D}(t)\rangle \approx 0$. Here, [46] did not show that all these should be. Our exactness gives the following result:

$$\| |D(t)\rangle \| \ll \left| \frac{E_m(t) - E_n(t)}{E_n(t)} \right| \quad \text{and} \quad \| |\dot{D}(t)\rangle \| \ll |E_m(t) - E_n(t)|. \quad (10)$$

- Finally, for the ϵ -regime, each basis state $|i\rangle$ [46] becomes a linear combination of eigenstates

$$\| |i\dot{D}(t)\rangle - E_n(t) |D(t)\rangle \| \ll |E_m(t) - E_n(t)|. \quad (11)$$

- One of the criticals of the algorithm [46] is the careful choice of the initial state $|i\rangle$ to avoid the adiabatic condition [48]. One can choose the initial state to be the ground state $|0\rangle$ (9).
- The linear combination of the basis states $|i\rangle$ is defined as the initial state $|D(t)\rangle$ (5), i.e.,

$$c_m(t) = ie^{i\beta_n(t)} \frac{\langle E_m(t) | \dot{E}_n(t) \rangle}{E_m(t) - E_n(t)}, \quad (12)$$

where the initial state is

$$|i\dot{D}(t)\rangle = E_n(t) |D(t)\rangle. \quad (13)$$

- We also find that the condition $|\dot{D}(t)\rangle = 0$ implies $c_m(t) = 0$ and $|D(t)\rangle = 0$. In the ϵ -regime, for a given state $|\psi(t)\rangle$, if $|\dot{\psi}(t)\rangle = |\dot{\psi}_n^{adi}(t)\rangle$, the initial state is equal to $|\psi_n^{adi}(t)\rangle$ as well, i.e., $|\psi(t)\rangle = |\psi_n^{adi}(t)\rangle$.

Finally, we have the fact that the algorithm is a *posteriori* algorithm, i.e., it is necessary to check the success of the algorithm, as has been done [47, 50] for the algorithm [24] and [46].

3.2. Organization of the report

Before going into the details, we have a general remark. It is well known that the QAC is a *a priori* condition for the adiabatic algorithm, although the algorithm is clear. Instead, *a posteriori* condition, e.g., the *why the QAC fails to reveal the adiabatic approximation*—a problem has gained considerable interest in the community. Although the theoretical results are clear, the lack of a clear explanation is still a challenge.

The rest of this paper is organized as follows:

- In section ‘Details of the algorithm’: we provide a detailed description of the algorithm for the algorithm (9).

In section ‘Discussion’: we discuss the relevance of the first and the second series in eq (9). The discussion of the adiabatic condition is discussed. The implications for the series are also listed.

In section ‘Illustrative example’: we consider the following based on the Schrodinger equation for a 1/2 harmonic oscillator. This demonstrates, in detail, the effect of the first-order perturbation theory on the adiabatic condition. The second-order perturbation theory is also considered. The adiabatic condition is also discussed. The adiabatic condition is also discussed.

4. Derivation of the main result

We are now ready to derive the adiabatic condition in eq (9). This is done by considering the following:

$$\langle E_m(t) | \left(i \frac{d}{dt} - E_n(t) \right) | D(t) \rangle, \tag{14}$$

which can be seen as a differential equation, i.e.,

$$\underbrace{\langle E_m(t) | \left(i \frac{d}{dt} - E_n(t) \right) | \psi(t) \rangle}_{=(E_m(t)-E_n(t))c_m(t)} - \underbrace{\langle E_m(t) | \left(i \frac{d}{dt} - E_n(t) \right) | \psi_n^{adi}(t) \rangle}_{=ie^{i\beta_n(t)}\langle E_m(t) | \dot{E}_n(t) \rangle}, \tag{15}$$

from the definition (cf eq (8)) of the differential equation $|D(t)\rangle$. These equations can be simplified as follows:

First, the first term is equal

$$(E_m(t) - E_n(t))c_m(t), \tag{16}$$

which comes from the Schrödinger equation which has

$$i \langle E_m(t) | \dot{\psi}(t) \rangle = \langle E_m(t) | H(t) | \psi(t) \rangle, \tag{17}$$

followed by the Hermiticity condition, $H^\dagger(t) = H(t)$, for $H(t)$ has

$$\langle E_m(t) | H(t) = E_m(t) \langle E_m(t) |, \tag{18}$$

and the definition of the coefficient $c_m(t) = \langle E_m(t) | \psi(t) \rangle$.

Second, we have the following general condition

$$\langle E_m(t) | \psi_n^{adi}(t) \rangle = \langle E_m(t) | E_n(t) \rangle = 0 \tag{19}$$

for all eigenstates $m \neq n$. The second term in the first line is $i \langle E_m(t) | \frac{d}{dt} | \psi_n^{adi}(t) \rangle$, which is equal to $ie^{i\beta_n(t)} \langle E_m(t) | \dot{E}_n(t) \rangle$ from the definition (cf eq (3)) of the adiabatic state $|\psi_n^{adi}(t)\rangle$.

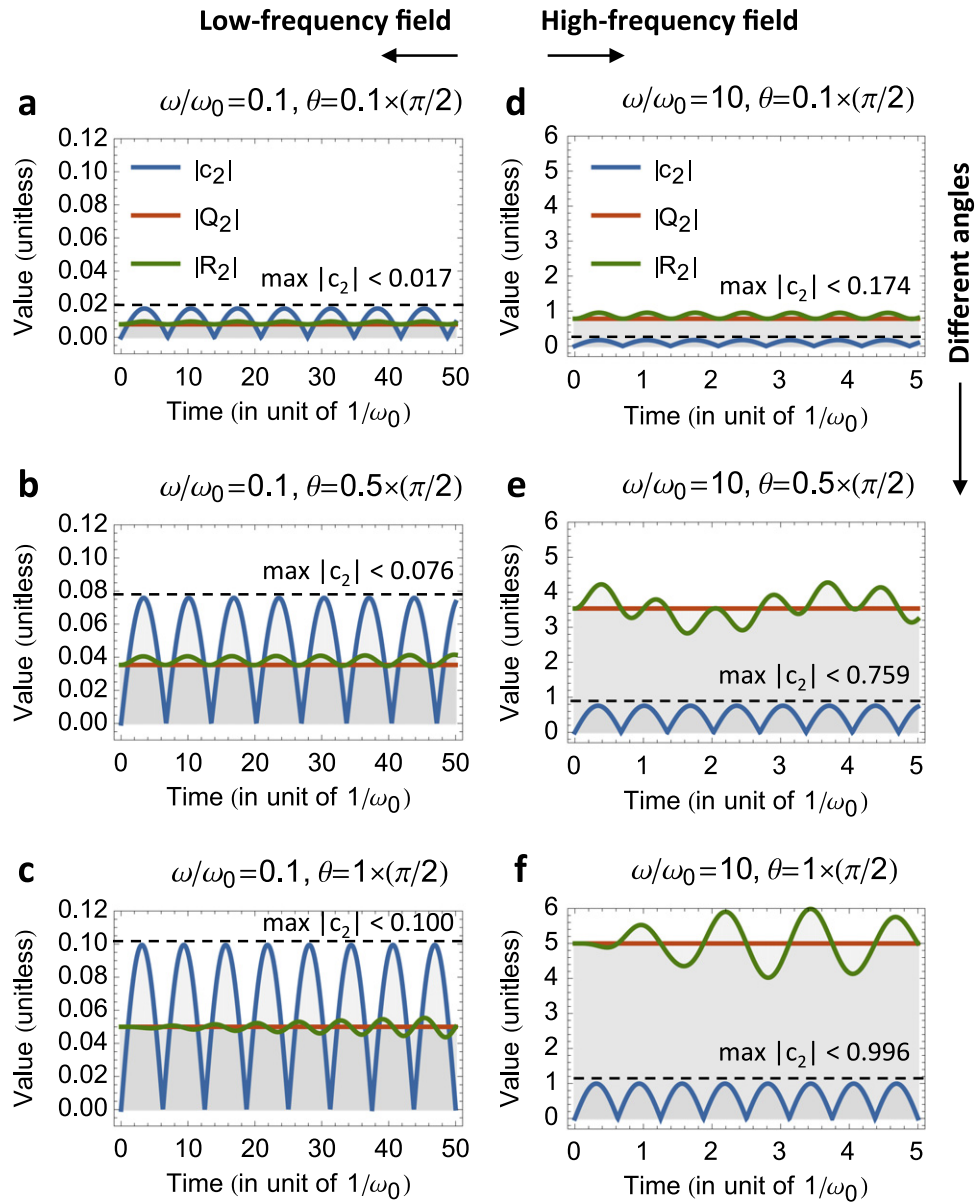


Figure 1. Time evolutions of the magnitudes of $|c_2(t)|$, $|Q_2(t)|$, and $|R_2(t)|$ of the essential state (9). (a)–(c) correspond to the case of the low-frequency field ($\omega/\omega_0 = 0.1$) and (d)–(f) are for the high-frequency field ($\omega/\omega_0 = 10$) and different angles.

In addition, we have the following relation:

$$\langle E_m | \left(i \frac{d}{dt} - E_n \right) | D \rangle = (E_m - E_n) c_m - i e^{i\beta_n} \langle E_m | \dot{E}_n \rangle. \quad (20)$$

Nevertheless, although a simple relation between the essential state and the coefficients $c_m(t)$ and \dot{E}_n is established in equation (9).

5. Discussion on the main result

The first one,

$$Q_m(t) \equiv ie^{i\beta_n(t)} \frac{\langle E_m(t) | \dot{E}_n(t) \rangle}{E_m(t) - E_n(t)}, \tag{21}$$

the right hand side of eq. (9) is closely related to the QAC (cf eq. (1)) and is bounded in [46], which assured that the QAC (cf eq. (1)) is necessary subject to the condition that $\langle D(t) \rangle \approx 0$ and $\langle \dot{D}(t) \rangle \approx 0$. Hence, it is sufficient to (9)

is satisfied (cf eq. (23)) here if $\langle D(t) \rangle$ and $\langle \dot{D}(t) \rangle$ are small for QAC (which is needed just if the adiabatic [46]), but also the energy level condition (cf eq. (25)) has to lead to a small value for the adiabaticity of the QAC.

The second one,

$$R_m(t) \equiv -E_n(t) \frac{\langle E_m(t) | D(t) \rangle}{E_m(t) - E_n(t)} + i \frac{\langle E_m(t) | \dot{D}(t) \rangle}{E_m(t) - E_n(t)}, \tag{22}$$

is the second one in the adiabaticity [46]. Remarkably, it is noted that the value of $R_m(t)$ is small compared with 1, i.e., $|R_m(t)| \ll 1$, the QAC in eq. (1) is necessary for the adiabaticity of the adiabatic approximation.

5.1. On the necessity of QAC

Having defined the adiabaticity in eq. (9), we are lead to ask whether the condition of the adiabaticity is necessary. Here we consider the condition that the QAC (cf eq. (1)) becomes necessary for the adiabatic approximation to be valid, i.e., $|c_m| \ll 1$ for all $m \neq n$. We shall ask the following question: 'Does the condition $R_m(t)$ in eq. (9) arise?'

First of all, the condition is closely related to $\langle D(t) \rangle$ and $\langle \dot{D}(t) \rangle$. Clearly, the QAC is necessary for the adiabatic approximation, it is noted that the condition (the necessary condition $\langle E_m(t) \rangle$) is bounded by $\langle D(t) \rangle$ and $\langle \dot{D}(t) \rangle$ are small, compared with $(E_m(t) - E_n(t))/E_n(t)$ and $|E_m(t) - E_n(t)|$ respectively, i.e.,

$$\|\langle D \rangle\| \ll \left| \frac{E_m - E_n}{E_n} \right| \quad \& \quad \|\langle \dot{D} \rangle\| \ll |E_m - E_n|. \tag{23}$$

In fact, in [46] it is explicitly assumed that $\langle D(t) \rangle \approx 0$ and $\langle \dot{D}(t) \rangle \approx 0$ in the adiabaticity [5] (see eqs (7) and (9) in [46]). Hence the condition in eq. (23) is a necessary condition for the adiabaticity $\langle D(t) \rangle \approx 0$ and $\langle \dot{D}(t) \rangle \approx 0$ to be satisfied in [46], and the classification of the cases has been discussed in [47, 48].

5.2. Generalization

Of course, the necessity of QAC (cf eq. (5)) is still valid as long as the correction $R_m(t)$ becomes sufficiently small. Requiring $\langle D(t) | \approx 0$ and $\langle \dot{D}(t) | \approx 0$ as in [46] is just a possibility. Generally, from eq. (22), it is sufficient to require the correction to be linearly small

$$i|\dot{D}(t)\rangle - E_n(t)|D(t)\rangle, \tag{24}$$

besides all conditions in the above, the energy gap $E_m(t) - E_n(t)$, i.e.,

$$\|i|\dot{D}(t)\rangle - E_n(t)|D(t)\rangle\| \ll |E_m(t) - E_n(t)|, \tag{25}$$

which is also a possibility. The latter is equivalent to $\langle D(t) | \approx 0$ and $\langle \dot{D}(t) | \approx 0$. In the above, as long as the condition (25) holds, the QAC (cf eq. (1)) is still the adiabatic approximation if $|c_m(t)| \ll 1$ for all $m \neq n$, and vice versa.

5.3. Properties of the correction term

In the following, we shall show that the condition $R_m(t) = 0$ is equivalent to the following: for each $m \neq n$, the probability amplitude $c_m(t) = \langle E_m(t) | \psi(t) \rangle$ is given by the following equation (cf eq. (5)):

$$c_m(t) = ie^{i\beta_n(t)} \frac{\langle E_m(t) | \dot{E}_n(t) \rangle}{E_m(t) - E_n(t)}, \tag{26}$$

if and only if

$$|\dot{D}(t)\rangle = -iE_n|D(t)\rangle. \tag{27}$$

In the above, the probability amplitude $c_m(t)$ is given explicitly by eq. (5), in the approximation of the energy gap (26). Furthermore, from eq. (9), it is equivalent to the following relationship:

$$R_m(t) = 0 \Leftrightarrow i|\dot{D}(t)\rangle = E_n(t)|D(t)\rangle. \tag{28}$$

Proof. The following is the backward direction, i.e.,

$$i|\dot{D}(t)\rangle = E_n(t)|D(t)\rangle \Rightarrow R_m(t) = 0, \tag{29}$$

is trivial from the definition of $R_m(t)$ (cf eq. (22)). The other direction is the forward direction, i.e.,

$$R_m(t) = 0 \Rightarrow i|\dot{D}(t)\rangle = E_n(t)|D(t)\rangle, \tag{30}$$

follows.

See 1: For the definition of $R_m(t)$ (cf eq. (22)), for each $m \neq n$, we have

$$\langle E_m(t) | (E_n(t) | D(t) \rangle - i | \dot{D}(t) \rangle) = 0, \quad (31)$$

which implies that the vector $E_n(t) | D(t) \rangle - i | \dot{D}(t) \rangle$ is orthogonal to all the basis vectors $| E_m(t) \rangle$. In other words, this vector belongs to the subspace spanned by $| E_n(t) \rangle$.

See 2: Conversely, we can write

$$E_n(t) | D(t) \rangle - i | \dot{D}(t) \rangle = \lambda | E_n(t) \rangle, \quad (32)$$

for some scalar λ . Since the eigenspaces are assumed to be orthogonal $\langle E_n(t) | E_n(t) \rangle = 1$, we can take

$$E_n(t) \langle E_n(t) | D(t) \rangle - i \langle E_n(t) | \dot{D}(t) \rangle = \lambda. \quad (33)$$

Next, we shall show that λ can be zero, i.e., $\lambda = 0$.

See 3: Let us consider the definition of the difference vector $| D(t) \rangle$ (cf eq. (8)), which gives

$$\langle E_n(t) | \dot{D}(t) \rangle = \underbrace{\langle E_n(t) | \dot{\psi}(t) \rangle}_{=-iE_n(t) \langle E_n(t) | \psi(t) \rangle} - \underbrace{\langle E_n(t) | \dot{\psi}_n^{adi}(t) \rangle}_{=-ie^{i\beta_n(t)} E_n(t)}_{=-iE_n(t) \langle E_n(t) | D(t) \rangle}. \quad (34)$$

For the Schrödinger equation,

$$\langle E_n(t) | \dot{\psi}(t) \rangle = -i \langle E_n(t) | H(t) | \psi(t) \rangle, \quad (35)$$

and for the Hermiticity of $H(t)$, the first term in the right hand side (34) becomes

$$\langle E_n(t) | \dot{\psi}(t) \rangle = -i E_n(t) \langle E_n(t) | \psi(t) \rangle. \quad (36)$$

On the other hand, for the definition of the adiabatic state $|\psi_n^{adi}(t)\rangle$ in eq. (3), we have

$$|\dot{\psi}_n^{adi}(t)\rangle = e^{i\beta_n(t)} |\dot{E}_n(t)\rangle - e^{i\beta_n(t)} (iE_n(t) + \langle E_n(t) | \dot{E}_n(t) \rangle) | E_n(t) \rangle, \quad (37)$$

which implies that the second term in the right hand side (34) becomes

$$\langle E_n(t) | \dot{\psi}_n^{adi}(t) \rangle = -ie^{i\beta_n(t)} E_n(t). \quad (38)$$

Combining eqs. (36) and (38), we finally have

$$\langle E_n(t) | \dot{D}(t) \rangle = -i E_n(t) \langle E_n(t) | D(t) \rangle \quad (39)$$

(we have $\langle E_n(t) | D(t) \rangle = \langle E_n(t) | \psi(t) \rangle - e^{i\beta_n(t)}$ from the definition of $| D(t) \rangle$). This means that λ is a real number, $\lambda = 0$, which further implies that

$$i | \dot{D}(t) \rangle = E_n(t) | D(t) \rangle, \quad (40)$$

and concludes the proof. \square

5.4. Consequences of $|\dot{D}(t)\rangle = 0$

We have shown that the eigenstates $|D(t)\rangle = 0$ and $|\dot{D}(t)\rangle = 0$ (which are also stationary states, cf. eq. (27)), hence can be seen as stationary [46] (cf. eq. (5) and eq. (26)). Hence we shall assume, as usual, that the condition $|\dot{D}(t)\rangle = 0$ implies that the states can be written as eigenstates (generally) $|E_n(t)\rangle$ i.e., all $c_m = 0$ for $m \neq n$.

Moreover, for all $m \neq n$ and $E_i(t) \neq 0$,

$$|\dot{D}(t)\rangle = 0 \Rightarrow |D(t)\rangle = 0 \text{ \& } c_m(t) = 0. \quad (41)$$

Proof. For all, for eq. (8), we can write

$$c_n(t) = e^{i\beta_n(t)} + \frac{i \langle E_n(t) | \dot{D}(t) \rangle}{E_n(t)}. \quad (42)$$

Now since $|\dot{D}(t)\rangle = 0$ implies

$$c_n(t) = \langle E_n(t) | \psi(t) \rangle = e^{i\beta_n(t)}. \quad (43)$$

Since the eigenstates are orthonormal, i.e., $\langle \psi(t) | \psi(t) \rangle = 1$, it means that all $c_m = 0$ for $m \neq n$, and

$$|\psi(t)\rangle = e^{i\beta_n(t)} |E_n(t)\rangle \equiv |\psi_n^{adi}(t)\rangle \quad (44)$$

(i.e., $|D(t)\rangle = 0$). □

For eq. (9), hence we shall assume that $\langle E_m(t) | \dot{E}_n(t) \rangle = 0$ for $m \neq n$ hence $|\dot{D}(t)\rangle = 0$.

6. Illustrative example

Here we will describe the behavior of a system (cf. eq. (9)), which is a bilinear system, and the Schrodinger equation is half-Hermitian [51],

$$H(t) = \vec{\sigma} \cdot \vec{B}(t) \equiv \frac{\hbar\omega_0}{2} (\sigma_x \sin\theta \cos\omega t + \sigma_y \sin\theta \sin\omega t + \sigma_z \cos\theta), \quad (45)$$

with the ansatz:

$$H_S(t) = \frac{\hbar\omega_0}{2} \begin{pmatrix} \cos\theta & \sin\theta e^{-i\omega t} \\ \sin\theta e^{i\omega t} & -\cos\theta \end{pmatrix}. \quad (46)$$

This Hamiltonian describes a time-dependent field (with frequency ω) at an angle θ , where the fields are given by ω_0 . The exact solution can be found analytically (e.g. see [32, 46, 51]), which is also available [2].

Table 2. Schrodinger's half Hamiltonian.

Terms	Expressions
Hamiltonian:	$H_S(t) = \frac{\omega_0}{2} \left((\sigma_x \cos \omega t + \sigma_y \sin \omega t) \sin \theta + \sigma_z \cos \theta \right)$
Eigenvalues:	$E_1(t) = -\omega_0/2$ and $E_2(t) = \omega_0/2$
Eigenstates:	$ E_1(t)\rangle = \left(e^{-i\omega t/2} \sin(\theta/2), -e^{i\omega t/2} \cos(\theta/2) \right)^T$, $ E_2(t)\rangle = \left(e^{-i\omega t/2} \cos(\theta/2), e^{i\omega t/2} \sin(\theta/2) \right)^T$.
Initial state:	$ \psi(t=0)\rangle = E_1(t=0)\rangle$
Time evolution:	$ \psi(t)\rangle = c_1(t) E_1(t)\rangle + c_2(t) E_2(t)\rangle$, hence $c_1(t) = \cos(\tilde{\omega}t/2) + i \sin(\tilde{\omega}t/2)(\omega_0 - \omega \cos \theta)/\tilde{\omega}$, $c_2(t) = i(\omega/\tilde{\omega}) \sin \theta \sin(\tilde{\omega}t/2)$, $\tilde{\omega} = \sqrt{\omega_0^2 + \omega^2 - 2\omega_0\omega \cos \theta}$.

6.1. Calculations of $Q_2(t)$ and $R_2(t)$

The imaginary part $Q_2(t)$,

$$Q_2(t) \equiv ie^{i\beta_1(t)} \frac{\langle E_2(t) | \dot{E}_1(t) \rangle}{E_2(t) - E_1(t)}, \quad (47)$$

in the relevant case (cf. eq. (21)) of the first energy level has the side feature (9). For all, the sign is listed in table 2, we have

$$\langle E_1(t) | \dot{E}_1(t) \rangle = i \frac{\omega}{2} \cos \theta. \quad (48)$$

This gives the expression for $\beta_1(t)$:

$$\beta_1(t) = \frac{\omega_0 t}{2} - \frac{\omega t}{2} \cos \theta. \quad (49)$$

Similarly, the corresponding

$$\langle E_2(t) | \dot{E}_1(t) \rangle = \frac{-i\omega}{2} \sin \theta. \quad (50)$$

Therefore, we have a nice expression for $Q_2(t)$:

$$Q_2(t) = e^{i\beta_1(t)} (\omega/2\omega_0) \sin \theta. \quad (51)$$

On the other hand, the imaginary part $R_2(t)$ is the relevant case of the corresponding eigenvalue listed in [46]. It can be calculated in the knowledge of $c_2(t)$ and $Q_2(t)$, i.e., $R_2(t) = c_2(t) - Q_2(t)$, which is

$$R_2(t) = \omega \sin \theta \left[i \sin \left(\frac{\tilde{\omega}t}{2} \right) / \tilde{\omega} - e^{i\beta_1(t)} / 2\omega_0 \right]. \quad (52)$$

6.2. Numerical results

The time evolution of the amplitudes $|c_2(t)|$, $|Q_2(t)|$, and $|R_2(t)|$ are shown in figure 1 for various values of $\omega/\omega_0 = 0.1$ and $\omega/\omega_0 = 10$ driving fields. With small driving fields (figure 1(a)–(c)), the system is always close ($|c_2| \ll 1$) to the state $|E_1(t)\rangle$ of the Hamiltonian $H_S(t)$ as expected, independent of the value of θ . For large driving fields (figure 1(d)–(f)), the system can also be close to the state $|E_1(t)\rangle$ for all values of θ . Particularly, the case in figure 1(d) is related to the debate [47–49] on the validity of [46], where it is suggested [48] that the case of the adiabatic evolution ($|c_2| \ll 1$ for all times) is not QAC (i.e., $|Q_2(t)|$ is small). Our numerical calculations indicate that in this case, the $R_2(t)$ excitation of the $Q_2(t)$ excitation makes the $|c_2(t)|$ small, as expected from equation (9). In the end, we have identified a case where the adiabatic evolution holds, but $Q_2(t)$ is small, which means that QAC is necessary.

7. Conclusion

In summary, we have considered a one-dimensional fermionic system (cf. equation (9)) and identified the significance of the eigenstate $|E_1(t)\rangle$ [46]. For this system, a reasonable definition of the adiabatic QAC becomes a valid criterion for the adiabatic evolution. As a illustration, we take the example of the Schrieffer-Hamiltonian as defined in section 2. For the case of the fermionic system, the adiabatic evolution is identified as a consequence of the definition of the adiabatic QAC. In addition, the fact that the adiabatic evolution is not always a consequence of the definition of the adiabatic QAC [46–49], which is the reason why we have added the condition of QAC becomes necessary? Our numerical results show that a case where the adiabatic evolution holds, but $Q_2(t)$ is small (cf. equation (10)), which means that the evolution is not QAC [46].

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Appendix A. Hamiltonian and eigenvectors

We consider a one-dimensional Hamiltonian $H(t)$ which consists of the following N -dimensional basis states. For a given $i \in \{1, 2, 3, \dots, N\}$, let the eigenstate $|E_i(t)\rangle$ and

the eigenstate $|E_i(t)\rangle$ is a stationary eigenstate and the corresponding eigenvalue of the Hamiltonian $H(t)$ is real, i.e.,

$$H(t)|E_i(t)\rangle = E_i(t)|E_i(t)\rangle. \quad (\text{A.1})$$

As usual, the evolution of the state $|\psi(t)\rangle$ is governed by the Schrödinger equation,

$$i\dot{|\psi(t)\rangle} = H(t)|\psi(t)\rangle. \quad (\text{A.2})$$

Here we assume the system is initialized at $t=0$ in the eigenstate, i.e.,

$$|\psi(0)\rangle = |E_n(0)\rangle, \quad (\text{A.3})$$

indicating the ground state. Furthermore, the eigenstate $|\psi(t)\rangle$ can be expanded by the complete set of eigenstates $\{|E_i(t)\rangle\}$ of the eigenbasis

$$|\psi(t)\rangle = \sum_i c_i(t)|E_i(t)\rangle, \quad (\text{A.4})$$

where $c_i(t) = \langle E_i(t)|\psi(t)\rangle$ is the corresponding coefficient in the expansion, i.e., $|c_i(t)| \leq 1$, since the Hamiltonian is Hermitian and bounded, i.e., $\| |\psi(t)\rangle \| = 1$.

Appendix B. Transformation of Schiff's expression

In the book of Schiff [1], the following expression (in our notation) is given:

$$c_m(t) \approx \frac{\langle E_m(t)|\dot{H}|E_n(t)\rangle}{i(E_m(t) - E_n(t))^2} (e^{i(E_m - E_n)t} - 1). \quad (\text{B.1})$$

We are going to show that this is not correct.

First, we have

$$\langle E_m(t)|E_n(t)\rangle = 0 \quad (\text{B.2})$$

for all $m \neq n$. Therefore, the following

$$\frac{d}{dt} \langle E_m(t)|E_n(t)\rangle = 0 \quad (\text{B.3})$$

is not correct.

$$\langle \dot{E}_m(t)|E_n(t)\rangle = -\langle E_m(t)|\dot{E}_n(t)\rangle. \quad (\text{B.4})$$

Second, since it is known that we have

$$\langle E_m(t)|H|E_n(t)\rangle = 0 \quad (\text{B.5})$$

for all $m \neq n$, we have

$$\frac{d}{dt} \langle E_m(t) | H(t) | E_n(t) \rangle = 0, \quad (\text{B.6})$$

which implies that

$$\langle \dot{E}_m(t) | H(t) | E_n(t) \rangle + \langle E_m(t) | H | \dot{E}_n(t) \rangle + \langle E_m(t) | \dot{H}(t) | E_n(t) \rangle = 0, \quad (\text{B.7})$$

and hence

$$E_n(t) \langle \dot{E}_m(t) | E_n(t) \rangle + \langle E_m(t) | \dot{H} | E_n(t) \rangle + E_m(t) \langle E_m(t) | \dot{E}_n(t) \rangle = 0. \quad (\text{B.8})$$

Combining these results, we have

$$\frac{\langle E_m(t) | \dot{H} | E_n(t) \rangle}{E_m(t) - E_n(t)} = - \langle E_m(t) | \dot{E}_n(t) \rangle, \quad (\text{B.9})$$

which changes Schiff's expression to

$$c_m(t) \approx i \frac{\langle E_m(t) | \dot{E}_n(t) \rangle}{E_m(t) - E_n(t)} \left(e^{i(E_m - E_n)t} - 1 \right). \quad (\text{B.10})$$

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