


RESEARCH

Open Access



Hermitian and non-Hermitian shortcuts to adiabaticity for fast creation of maximum coherence and beam splitting

Kai Tang^{1,2}, Zhengfeng Hu^{3,4*}, Xi Chen^{1,5*}  and Chengpu Liu^{2*}

Abstract

We theoretically exploit the shortcuts to adiabaticity (STA) technique in Hermitian and non-Hermitian quantum systems to realize the maximum coherence and beam splitting by eliminating the nonadiabatic coupling. Compared with the conventional adiabatic passage (AP) technique with the Gaussian and Allen-Eberly schemes, the operation time can be significantly shortened by three order using STA technique. This STA-based fast creation of maximum coherence or beam splitting are in use ranging from quantum sensing and metrology in a noisy environment to optical gain/loss coupled waveguides in an analogous fashion.

Keywords: Shortcuts to adiabaticity, Non-Hermitian, Beam splitting

Introduction

Preparation and transfer of maximum coherent quantum states are requisite in the applications of quantum measurements and sensing. Traditionally, adiabatic passage (AP), in order to transfer the population coherently and adiabatically in a two-level quantum system, is one of the most popular protocols used for initializing and manipulating quantum systems [1–4], with many applications ranging from large-area atom interferometry [5], atomic frequency standards [6], quantum information processing [7], nuclear magnetic resonance [8], coherent superposition of molecular vibration states [9], to quantum dots [10], etc.

AP approach and its variants including (fractional) stimulated Raman adiabatic passage (STIRAP) depend on the off-resonant coupling between discrete states and laser pulse to achieve a steady-state superposition. In the pseudo polarization representation [11], a two-level

state superposition is considered as a state vector, and the phase, intensity, and detuning of the optical field is defined as a field vector. The state vector precesses about this adjustable field vector. Provided that this field vector varies sufficiently slowly, the population can be transferred from one state to the other with an arbitrary fidelity, that is, the state can be transferred perfectly. If this adiabatic condition is not met, the nonadiabatic coupling between eigenstates results in decoherence and the target state can not be reached perfectly. Therefore, in an AP process, the price to get a perfect target state is at expense of long time, which is the limitation for the application of AP technique in the case where a fast state preparation is demanding under the existence of decoherence.

Subsequently, a technique called shortcuts to adiabaticity (STA) [12–23] was proposed theoretically and demonstrated experimentally to speed up slow AP, whose concepts, methods, and applications are well reviewed by Guery et al. [23]. One of STA is counterdiabatic driving [14, 16, 17], in which an additional counterdiabatic term is introduced to eliminate the nonadiabatic coupling, and thus make it possible for quantum state preparation to be rapidly realized. However, the cost for this

*Correspondence: zfhu@siom.ac.cn; xchen@shu.edu.cn; chpliu@siom.ac.cn

¹International Center of Quantum Artificial Intelligence for Science and Technology (QuArtist) and Department of Physics, Shanghai University, 200444 Shanghai, China

⁵Department of Physical Chemistry, University of the Basque Country, 48080, Bilbao, Spain

Full list of author information is available at the end of the article

fast state preparation is the increase of energy consumption due to the introduction of counterdiabatic driving. The existing studies about STA have focused on a Hermitian system, but recently, an increasing interest is blooming in a non-Hermitian system [24–26]. For example, as for a parity-time (PT) symmetric non-Hermitian system [27, 28], it can produce a faster-than-Hermitian evolution, while keeping the eigenenergy difference fixed [27, 29]. Some proposals for the realization of PT-symmetric Hamiltonian have been put forth [24] as well as some extensions related to the Landau-Zener (LZ) model [25, 30, 31]. In addition, the approximated adiabatic condition for a non-Hermitian Hamiltonian system has also been obtained [32]. And dissipation-assisted adiabatic passages was applied to prepare entangled-state [33]. However, as far as we know, no one consider STA-based creation of maximum coherence or beam splitting in a non-Hermitian system, even in a Hermitian system. The beam splitting is very useful in optical waveguides and coherent quantum state preparation, especially in quantum sensing and metrology.

As shown in the following, by exploiting the fast shortcuts population transfer in a non-Hermitian Hamiltonian system, one can overcome some decoherence processes and achieve a perfect quantum state preparation. By properly choosing the driving field the maximum coherence and beam splitting can be obtained exactly. The creation of maximum coherence and beam splitting based on counterdiabatic driving or non-Hermitian STA is faster than that based on AP. But the counterdiabatic driving cost more coupling to speed up adiabatic evolution, while non-Hermitian STA modulates non-Hermitian terms without increasing coupling. Therefore, non-Hermitian STA can work in an arbitrarily small coupling and has more advantages.

The above STA in quantum systems can be moved into optical waveguide systems if just substituting time by spatial coordinates [34, 35], which favors a more compact device integration. Furthermore, non-Hermitian terms can be introduced and controlled by ion doping and laser pumping [36], and thus non-Hermitian STA can be implemented in non-Hermitian optical waveguides [36, 37].

In this paper, we firstly present the creation of maximum coherence and beam splitting of AP in a two-level Hermitian system. Secondly, the theoretical description of the counterdiabatic driving creation of maximum coherence and beam splitting in the Hermitian system is given. Then creation of maximum coherence and beam splitting of STA in a non-Hermitian system is described in detail and the physical implementation is considered in optical waveguides systems. The main conclusion is summarized in the end.

Results and discussion

First, for convenience of description, we use a two-level system to describe a directional coupler of two waveguides, and briefly introduce the conventional AP in a simple two-level Hermitian system, whose coherent excitation is described by the Schrödinger equation [38],

$$i\hbar\partial_t c(t) = H_0(t)c(t), \quad (1)$$

where the vector $c(t) = [c_{g1}(t), c_{g2}(t)]^T$ comprises two probability amplitudes of states $|g1\rangle$ and $|g2\rangle$. The Hamiltonian under the rotating wave approximation (RWA) is written as

$$H_0(t) = \frac{\hbar}{2} \begin{pmatrix} 0 & \Omega(t) \\ \Omega(t) & 2\delta(t) \end{pmatrix}, \quad (2)$$

where $\Omega(t)$ is the Rabi frequency, which characterizes the strength of the field-induced coupling between states $|g1\rangle$ and $|g2\rangle$, and $\delta(t)$ is the detuning between field frequency $\omega(t)$ and the transition frequency $\omega_0(t)$, $\delta(t) = \omega_0(t) - \omega(t)$.

The instantaneous eigenstates of this time-dependent Hamiltonian $H_0(t)$ are

$$\begin{aligned} |-\rangle &= \cos\theta|g1\rangle - \sin\theta|g2\rangle, \\ |+\rangle &= \sin\theta|g1\rangle + \cos\theta|g2\rangle, \end{aligned} \quad (3)$$

with eigenvalues

$$\hbar\varepsilon_{\pm}(t) = \frac{\hbar}{2} \left[\delta(t) \pm \sqrt{\Omega^2(t) + \delta^2(t)} \right], \quad (4)$$

and the mixing angle $\theta(t)$ is defined as $\tan 2\theta(t) = \Omega(t)/\delta(t)$. The states $|-\rangle$ and $|+\rangle$ are the adiabatic states. They are connected with the diabatic vector $c(t)$ via the rotation matrix

$$R(\theta) = \begin{pmatrix} \cos\theta & \sin\theta \\ -\sin\theta & \cos\theta \end{pmatrix}, \quad (5)$$

as $c(t) = R(\theta(t))a(t)$, where $a(t) = [a_-(t), a_+(t)]^T$ contains the probability amplitudes of the adiabatic states $|-\rangle$ and $|+\rangle$. Similarly, the Schrödinger equation in the adiabatic basis reads

$$i\hbar\partial_t a(t) = H_a(t)a(t), \quad (6)$$

where

$$H_a(t) = \hbar \begin{pmatrix} \varepsilon_-(t) & -i\dot{\theta}(t) \\ i\dot{\theta}(t) & \varepsilon_+(t) \end{pmatrix}. \quad (7)$$

Here, an overdot means a time derivative.

The adiabatic condition $|\dot{\theta}(t)| \ll \varepsilon_+(t) - \varepsilon_-(t) = \varepsilon(t)$ means the nonadiabatic coupling in the Hamiltonian $H_a(t)$ is negligible compared with the eigenenergy difference. If the system is initially prepared in $|+(0)\rangle = |g1\rangle$, then under the adiabatic condition, this system will always maintain on the instantaneous eigenstate $|+(t)\rangle$. Then by controlling the final mixing angle θ_f according to Eq. (3), one can achieve an arbitrary target state. This means that

the AP makes the initial population in state $|g1\rangle$ be exactly split into two parts respectively populated in states $|g1\rangle$ and $|g2\rangle$, with a population ratio of $\tan^2 \theta_f$. It is needed to be pointed out that the above creation passage of maximum coherence and beam splitting is limited by the adiabatic condition $|\dot{\theta}(t)| \ll \varepsilon(t)$, which requires a slow evolution. As an example, we numerically demonstrate this AP for creation of maximum coherence or beam splitting under a Gaussian shape pulse. The mixing angle is required to satisfy $\theta(0) \rightarrow \pi/2$ and $\theta(t_f) \rightarrow \pi/4$, which means

$$\frac{\Omega(t)}{\delta(t)} \Big|_{t=0} \rightarrow 0^-, \quad \frac{\Omega(t)}{\delta(t)} \Big|_{t=t_f} \rightarrow -\infty. \quad (8)$$

And as shown in Fig. 1a, the Rabi frequency and detuning for the field are

$$\begin{aligned} \Omega(t) &= \Omega_0 \exp \left[-\frac{(t - t_f/2 - \tau)^2}{\sigma^2} \right], \\ \delta(t) &= \delta_0 \exp \left[-\frac{(t - t_f/2 + \tau)^2}{\sigma^2} \right]. \end{aligned} \quad (9)$$

In Eq. (9), Ω_0 and δ_0 are the peak values of Rabi frequency Ω and detuning δ , and σ is full width at half maximum. The τ is separation time between detuning and Rabi frequency, which guarantees to satisfy the boundary conditions in Eq. (8), and t_f is the operation time. In the creation process of maximum coherence and beam splitting, the adiabatic condition is always satisfied, if one can check the time evolution of $\varepsilon(t)/|\dot{\theta}(t)|$ (Fig. 1b) where this ratio is much larger than 1. Thus the state $|g1\rangle$ is equally split into state $|g1\rangle$ and state $|g2\rangle$ (Fig. 1c).

Although the AP can be used to obtain a precise quantum state preparation, it requires an extremely long operation time to avoid the influence of nonadiabatic effect in the evolution process. To speed up the slow evolution process, researchers have proposed the counterdiabatic driving [14, 16, 17], which is one of STA methods. The addition of an additional counterdiabatic term [14] can effectively suppress the nonadiabatic effect, which is written as

$$H_{cd}(t) = \hbar \begin{pmatrix} 0 & i\dot{\theta}(t) \\ -i\dot{\theta}(t) & 0 \end{pmatrix}. \quad (10)$$

And then the total Hamiltonian is constructed as

$$H = H_0 + H_{cd} = \frac{\hbar}{2} \begin{pmatrix} 0 & \sqrt{\Omega^2 + 4\dot{\theta}^2} e^{-i\alpha} \\ \sqrt{\Omega^2 + 4\dot{\theta}^2} e^{i\alpha} & 2\delta \end{pmatrix}, \quad (11)$$

where $\alpha(t) = -\arctan 2\dot{\theta}(t)/\Omega(t)$. By using the unitary transformation,

$$U(t) = \begin{pmatrix} e^{-i\frac{\alpha}{2}} & 0 \\ 0 & e^{i\frac{\alpha}{2}} \end{pmatrix}, \quad (12)$$

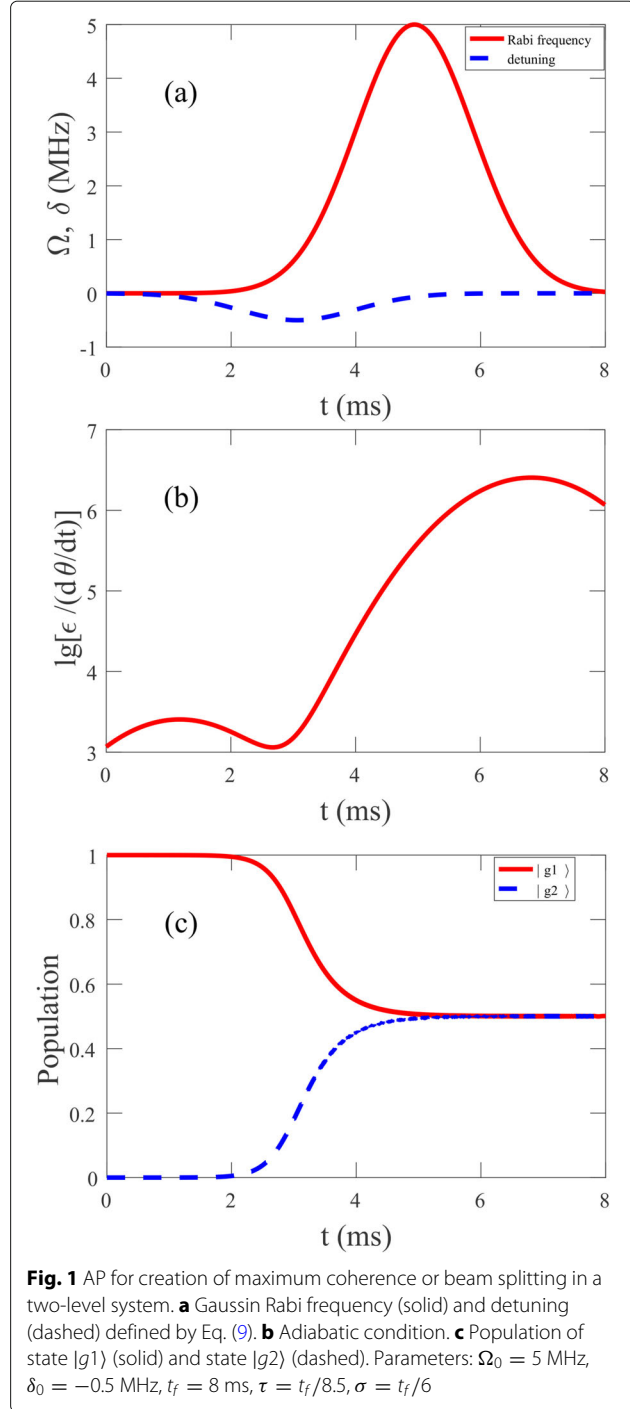


Fig. 1 AP for creation of maximum coherence or beam splitting in a two-level system. **a** Gaussian Rabi frequency (solid) and detuning (dashed) defined by Eq. (9). **b** Adiabatic condition. **c** Population of state $|g1\rangle$ (solid) and state $|g2\rangle$ (dashed). Parameters: $\Omega_0 = 5$ MHz, $\delta_0 = -0.5$ MHz, $t_f = 8$ ms, $\tau = t_f/8.5$, $\sigma = t_f/6$

modified Hamiltonian, $H_m(t) = U^{-1} H U - i\hbar U^{-1} \dot{U}$, can be calculated as

$$H_m = \frac{\hbar}{2} \begin{pmatrix} -\dot{\alpha} & \sqrt{\Omega^2 + 4\dot{\theta}^2} \\ \sqrt{\Omega^2 + 4\dot{\theta}^2} & 2\delta + \dot{\alpha} \end{pmatrix}. \quad (13)$$

The modified Rabi frequency $\Omega_m(t)$ and detuning $\delta_m(t)$ are obtained as

$$\Omega_m(t) = \sqrt{\Omega^2(t) + 4\dot{\theta}^2(t)}, \quad (14)$$

$$\delta_m(t) = \delta(t) + \dot{\alpha}(t). \quad (15)$$

To demonstrate the advantages of counterdiabatic driving for accelerating quantum evolution, we set the same Rabi frequency and detuning defined by Eq. (9) with the same amplitude of Rabi frequency and detuning ($\Omega_0 = 5$ MHz and $\delta_0 = -0.5$ MHz) as shown in Fig. 1a, but for a very short operation time, $t_f = 1 \mu\text{s}$. According Eq. (15), one can see that the amplitudes of the modified Rabi frequency and detuning obviously increase (Fig. 2b). This means that the counterdiabatic driving costs more coupling to achieve an accurate adiabatic evolution (Fig. 2c). Here, an operation time of $1 \mu\text{s}$ is an example, in fact, the theoretically speeding operation time can be arbitrarily short. In addition, it is necessary to point out that there is a phase difference $\alpha(0)$ between states $|g1\rangle$ and $|g2\rangle$ at the initial moment.

As for a non-Hermitian system, there is gain or loss of state population. One can nullify the nonadiabatic coupling by controlling gain or loss instead of the counterdiabatic term in a Hermitian system. For this goal, two nonzero terms are added in the Hamiltonian $H_0(t)$, and we obtain

$$H^\Gamma(t) = \frac{\hbar}{2} \begin{pmatrix} -i\Gamma_1(t) & \Omega(t) \\ \Omega(t) & 2\delta(t) + i\Gamma_2(t) \end{pmatrix}, \quad (16)$$

where $\Gamma_1(t)$ and $\Gamma_2(t)$ are both real numbers and represent the loss rate of the state $|g1\rangle$ and the gain rate of the state $|g2\rangle$ respectively.

As illustrated in AP, the Hamiltonian $H_a^\Gamma(t)$ in the basis $|\pm\rangle$ has the form

$$H_a^\Gamma(t) = \hbar \begin{pmatrix} A_{11} & A_{12} \\ A_{21} & A_{22} \end{pmatrix}, \quad (17)$$

where

$$\begin{aligned} A_{11} &= \varepsilon_- - \frac{i}{2} (\Gamma_1 \cos^2 \theta - \Gamma_2 \sin^2 \theta), \\ A_{12} &= -\frac{i}{4} (\Gamma_1 + \Gamma_2) \sin 2\theta - i\dot{\theta}, \\ A_{21} &= -\frac{i}{4} (\Gamma_1 + \Gamma_2) \sin 2\theta + i\dot{\theta}, \\ A_{22} &= \varepsilon_+ - \frac{i}{2} (\Gamma_1 \sin^2 \theta - \Gamma_2 \cos^2 \theta). \end{aligned} \quad (18)$$

The $\varepsilon_\pm(t)$ are again defined by Eq. (4). We assume that initially the system only populates in state $|g1\rangle$. According to the initial condition $\theta(0) = \pi/2$ and Eq. (3), the initial state coincides completely with the state $|+\rangle$. To nullify the nonadiabatic coupling, it requires that the nondiagonal

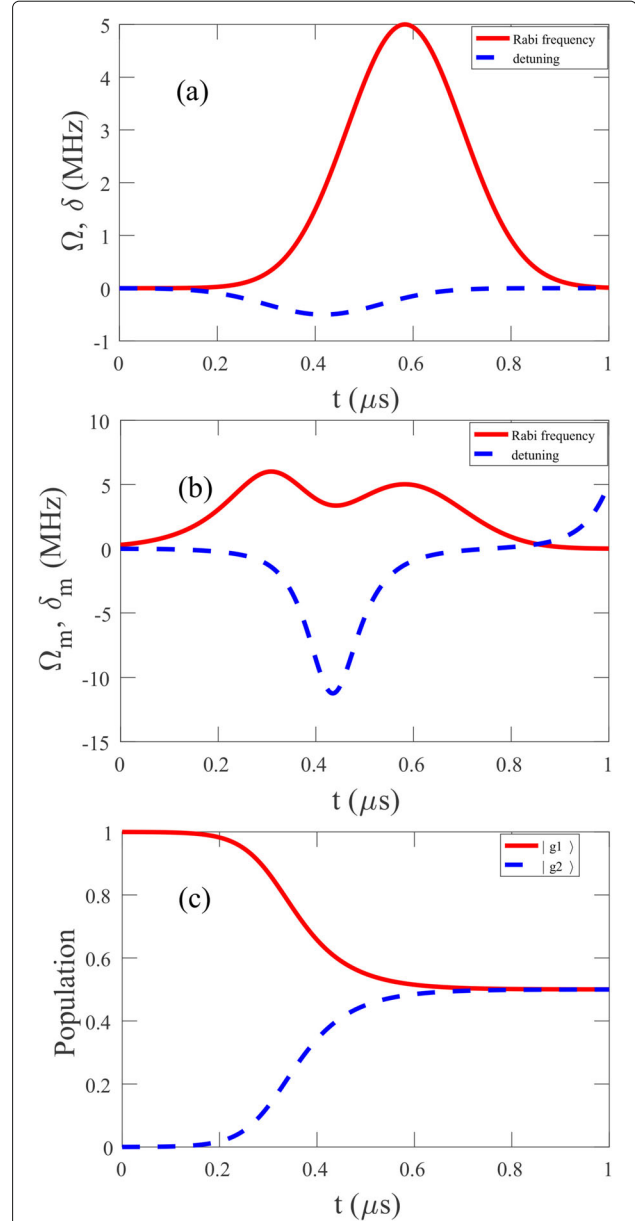


Fig. 2 STA for creation of maximum coherence or beam splitting. **a** Original Rabi frequencies (solid) and detuning (dashed) defined by Eq. (9). **b** Modified Rabi frequencies (solid) and detuning (dashed) defined by Eq. (15). **c** Population of state $|g1\rangle$ (solid) and state $|g2\rangle$ (dashed). Parameters: $\Omega_0 = 5$ MHz, $\delta_0 = -0.5$ MHz, $t_f = 1 \mu\text{s}$, $\tau = t_f/12$, $\sigma = t_f/6$

terms are zero, which means

$$\Gamma_1(t) + \Gamma_2(t) = -\frac{4\dot{\theta}(t)}{\sin 2\theta}. \quad (19)$$

During the evolution, the system follows the state $|+\rangle$, and the state $|-\rangle$ has never received any populations.

We obtain the evolution of the amplitudes of the adiabatic state by using Eq. (6)

$$a_-(t_f) = 0, \quad (20a)$$

$$a_+(t_f) = \exp \left\{ -i \int_0^{t_f} \varepsilon_+(t') - \frac{i}{2} [\Gamma_1(t') \sin^2 \theta(t') - \Gamma_2(t') \cos^2 \theta(t')] dt' \right\}, \quad (20b)$$

where t_f is the final time of the evolution process. In order to guarantee the final state $|a_+(t_f)|^2 = 1$, we require that the real part in Eq. (20b) is zero. It means

$$\int_0^{t_f} [\Gamma_1(t') \sin^2 \theta(t') - \Gamma_2(t') \cos^2 \theta(t')] dt' = 0. \quad (21)$$

Eq. (21) ensures that the total population at the end of the process is conserved. In particular, the total population in the system remains constant throughout the process, if the Γ_1 and Γ_2 satisfy

$$\Gamma_1(t) \sin^2 \theta(t) - \Gamma_2(t) \cos^2 \theta(t) = 0. \quad (22)$$

So we obtain

$$\begin{aligned} \Gamma_1(t) &= -\frac{2\dot{\theta}(t)}{\tan \theta(t)}, \\ \Gamma_2(t) &= -2\dot{\theta}(t) \tan \theta(t). \end{aligned} \quad (23)$$

So far, we have obtained the exact analytic solutions to equations about loss rate $\Gamma_1(t)$ of the state $|g1\rangle$ and gain rate $\Gamma_2(t)$ of the state $|g2\rangle$. From Eq. (23), it can be seen that there are $\Gamma_1(t)$ and $\Gamma_2(t)$, which satisfy Eq. (23), to ensure that the total population is conservative, regardless of any set of field parameters ($\Omega(t)$ and $\delta(t)$) selected.

After obtaining the relation between $\Gamma_{1,2}$ and mixing angle, the corresponding field parameters can be selected according to our purposes. Now, efficiency of creation of maximum coherence and beam splitting based on non-Hermitian STA in a Gaussian model described by Eq. (9) is investigated as well as that of a Hermitian STA. We set the same parameters ($\Omega_0 = 5$ MHz, $\delta_0 = -0.5$ MHz, $t_f = 1$ μ s, $\tau = t_f/10$, and $\sigma = t_f/6$) as Fig. 2. In our method, we introduce a non-Hermitian term (Fig. 3b) to cancel the nonadiabatic coupling, and Rabi frequency and detuning (Fig. 3a) are neither modified nor increased. And the accurate and rapid creation of maximum coherence or beam splitting is obtained in Fig. 3c. Same as Hermitian STA, non-Hermitian STA is three-order or even arbitrarily faster than AP. Non-Hermitian STA which is better than Hermitian one can be applied under arbitrarily small coupling.

In fact, the above creation of maximum coherence or beam splitting is not limited to a Gaussian model, and the non-Hermitian STA beam splitting can be realized in other models, such as the widely used Allen-Eberly (AE) model [39]. The Rabi frequency and detuning in the AE

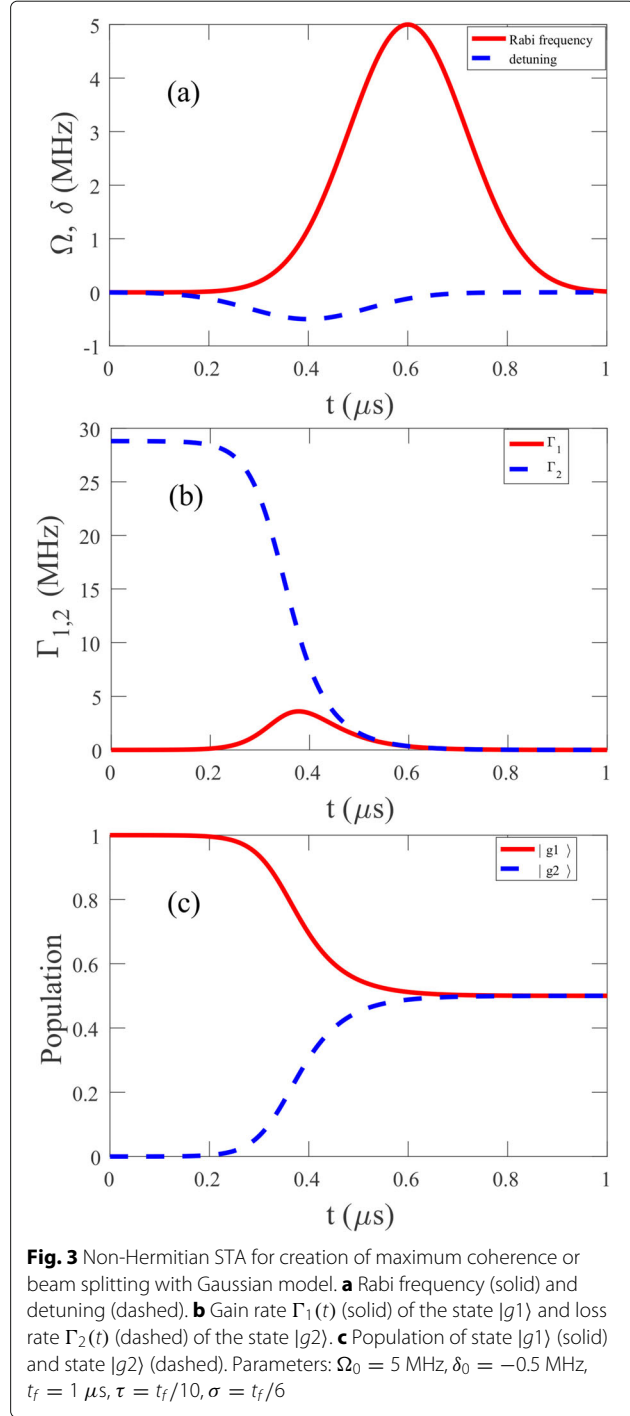


Fig. 3 Non-Hermitian STA for creation of maximum coherence or beam splitting with Gaussian model. **a** Rabi frequency (solid) and detuning (dashed). **b** Gain rate $\Gamma_1(t)$ (solid) and loss rate $\Gamma_2(t)$ (dashed) of the state $|g2\rangle$. **c** Population of state $|g1\rangle$ (solid) and state $|g2\rangle$ (dashed). Parameters: $\Omega_0 = 5$ MHz, $\delta_0 = -0.5$ MHz, $t_f = 1$ μ s, $\tau = t_f/10$, $\sigma = t_f/6$

model are

$$\begin{aligned} \Omega(t) &= \Omega_0 \operatorname{sech} \left(\frac{t - t_f/2}{\sigma} \right), \\ \delta(t) &= \delta_0 \left[1 - \tanh \left(\frac{t - t_f/2}{\sigma} \right) \right]. \end{aligned} \quad (24)$$

Here the detuning is a transformed hyperbolic tangent function, which is a little different from that used in the

traditional AE model, in order to ensure that the boundary conditions (8) are satisfied. Under the parameters $\Omega_0 = 5$ MHz, $\delta_0 = -0.5$ MHz, $t_f = 1 \mu s$, and $\sigma = t_f/12$, a perfect creation of maximum coherence and beam splitting are achieved in short time, as shown in Fig. 4c.

With a further discussion, our suggested STA technique in non-Hermitian systems could be physically realized in waveguide optics [36, 37]. A PT-symmetric directional

coupler consisting of two waveguides, waveguide $|L\rangle$ and waveguide $|R\rangle$, placed in proximity with propagation constants $\beta_L(z)$ and $\beta_R(z)$ is considered. The refractive index or geometry of the two waveguides are allowed to vary along the propagation direction z . Light is coupled into the device at input plane $z = z_i (= 0)$ and out at output plane $z = z_f$. The waveguides $|L\rangle$ and $|R\rangle$ here substitute for the above quantum states $|g1\rangle$ and $|g2\rangle$, and time t in Eq. (16) is replaced by direction z . Light propagation in the coupler is described by coupled-mode equation [40, 41]

$$i \frac{d}{dz} \begin{pmatrix} A_L \\ A_R \end{pmatrix} = \begin{pmatrix} \beta_L(z) - i\gamma_L(z) & \kappa(z) \\ \kappa(z) & \beta_R(z) + i\gamma_R(z) \end{pmatrix} \begin{pmatrix} A_L \\ A_R \end{pmatrix}, \quad (25)$$

where A_L and A_R are mode amplitudes of two waveguides, κ is coupling coefficient, γ_L and γ_R are respectively the loss of waveguide $|L\rangle$ and the gain of waveguides $|R\rangle$. The relation between the mismatch $\Delta = (\beta_R - \beta_L)/2$ and width difference δW of two waveguides can be approximated by a linear relation [35, 42], and the coupling coefficient scales exponentially with the waveguide separation $d(z)$ according to the law $\kappa(z) = \kappa_0 \exp[-k(d(z) - d_0)]$ [35, 40] (where κ_0 and k are two constant parameters to be determined from the fabrication process). A practical consideration of experimental realization of this system can be Ti in-diffused Fe:LiNbO₃ optical waveguides which have recently been applied to demonstrate PT-symmetric directional couplers in the visible spectrum [36]. The loss rate γ_L is dependent on the number of Fe²⁺ dopants and can be controlled easily. The laser pumping through the two-wave mixing process can provide the gain γ_R [36]. Thus, non-Hermitian STA can be experimentally verified in optical waveguide systems. Better yet, although we focus on equal mode splitting of optical waveguides, our technique can be applied to arbitrary ratio mode splitting and can be applied in more fields.

Conclusion

In this paper, we have presented the fast creation of maximum coherence and beam splitting via shortcuts to adiabaticity technique in Hermitian and non-Hermitian systems. In a Hermitian system, one achieved a fast adiabatic creation of maximum coherence and beam splitting by introducing an additional counterdiabatic term to eliminate the nonadiabatic coupling in evolution. Theoretical analysis and numerical results verified that shortcuts to adiabaticity is three-order or even arbitrarily faster than conventional adiabatic passage, but this realization is at the cost of increasing coupling. Instead, a novel non-Hermitian STA scheme proposed here can be used to achieve perfect creation of maximum coherence and beam splitting without increasing coupling. By

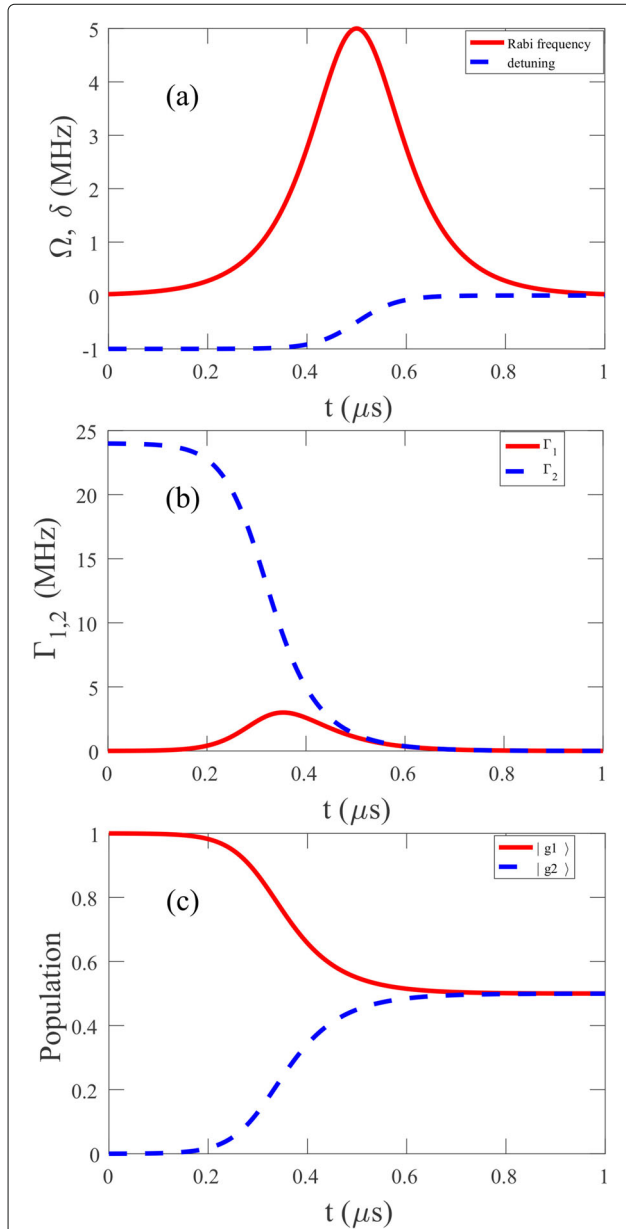


Fig. 4 Non-Hermitian STA for creation of maximum coherence or beam splitting with AE model. **a** Rabi frequency (solid) and detuning (dashed). **b** Gain rate $\Gamma_1(t)$ (solid) of the state $|g1\rangle$ and loss rate $\Gamma_2(t)$ (dashed) of the state $|g2\rangle$. **c** Population of state $|g1\rangle$ (solid) and state $|g2\rangle$ (dashed). Parameters: $\Omega_0 = 5$ MHz, $\delta_0 = -0.5$ MHz, $t_f = 1 \mu s$, $\sigma = t_f/12$

tuning non-Hermitian terms and field parameters, it is convenient to control the dynamic evolution of the system. Moreover, even for an arbitrarily small coupling between two states, non-Hermitian STA succeeds in perfect creation of maximum coherence and beam splitting, but counterdiabatic driving fails. The feasibility of non-Hermitian STA method has demonstrated in Gaussian and Allen-Eberly models and its experimental realization in optical waveguide systems is described. The fast quantum state manipulation based on shortcuts to adiabaticity demonstrated here possesses unique advantages in quantum sensing and metrology in a noisy environment.

Abbreviations

STA: Shortcuts to adiabaticity; AP: adiabatic passage; STIRAP: stimulated Raman adiabatic passage; PT: parity-time; LZ: Landau-Zener; AE: Allen-Eberly

Acknowledgements

X. C. thanks the Ramón y Cajal grant (RYC-2017-22482) and the Program for Eastern Scholar.

Author's contributions

All authors contributed equally in compiling this paper. All authors read and approved the final manuscript.

Funding

NSFC (No. 11674342, No. 11474193, and No. 91536220), SKLDDQP (No. KF201707), NKRD (No. 2016YFA0301504), SMSTC (No. 18010500400 and No. 18ZR1415500).

Availability of data and materials

Upon request to the authors.

Competing interests

The authors declare that they have no competing interests.

Author details

¹International Center of Quantum Artificial Intelligence for Science and Technology (QuArtist) and Department of Physics, Shanghai University, 200444 Shanghai, China. ²State Key Laboratory of High Field Laser Physics and CAS Center for Excellence in Ultra-intense Laser Science, Shanghai Institute of Optics and Fine Mechanics, Chinese Academy of Sciences, 201800, Shanghai, China. ³The Key Laboratory of Quantum Optics, Shanghai Institute of Optics and Fine Mechanics, Chinese Academy of Sciences, 201800, Shanghai, China. ⁴State Key Laboratory of Low Dimensional Quantum Physics, Tsinghua University, 100084, Beijing, China. ⁵Department of Physical Chemistry, University of the Basque Country, 48080, Bilbao, Spain.

Received: 17 November 2019 Accepted: 11 June 2020

Published online: 11 July 2020

References

- Shore, B. W., Bergmann, K., Kuhn, A., Schieman, S., Oreg, J., Eberly, J. H.: Laser-induced population transfer in multistate systems: A comparative study. *Phys. Rev. A*. **45**, 5297–5300 (1992). <https://doi.org/10.1103/PhysRevA.45.5297>
- Melinger, J., Gandhi, S. R., Hariharan, A., Goswami, D., Warren, W.: Adiabatic population transfer with frequency-swept laser pulses. *J. Chem. Phys.* **101**(8), 6439–6454 (1994). <https://doi.org/10.1063/1.468368>
- Vitanov, N., Halfmann, T., Shore, B., Bergmann, K.: Laser-induced population transfer by adiabatic passage techniques. *Annu. Rev. Phys. Chem.* **52**, 763–809 (2001). <https://doi.org/10.1146/annurev.physchem.52.1.76>
- Malinovsky, V., Krause, J.: General theory of population transfer by adiabatic rapid passage with intense, chirped laser pulses. *Eur. Phys. J. D*. **14**(2), 147–155 (2001). <https://doi.org/10.1007/s10053017021>
- Kotru, K., Butts, D. L., Kinast, J. M., Stoner, R. E.: Large-area atom interferometry with frequency-swept Raman adiabatic passage. *Phys. Rev. Lett.* **115**, 103001 (2015). <https://doi.org/10.1103/PhysRevLett.115.103001>
- Kotru, K., Brown, J. M., Butts, D. L., Kinast, J. M., Stoner, R. E.: Robust Ramsey sequences with Raman adiabatic rapid passage. *Phys. Rev. A*. **90**, 053611 (2014). <https://doi.org/10.1103/PhysRevA.90.053611>
- Beterov, I. I., Saffman, M., Yakshina, E. A., Zhukov, V. P., Tretyakov, D. B., Entin, V. M., Ryabtsev, I. I., Mansell, C. W., MacCormick, C., Bergamini, S., Fedoruk, M. P.: Quantum gates in mesoscopic atomic ensembles based on adiabatic passage and Rydberg blockade. *Phys. Rev. A*. **88**, 010303 (2013). <https://doi.org/10.1103/PhysRevA.88.010303>
- Hardy, C. J., Edelstein, W. A., Vatis, D.: Efficient adiabatic fast passage for NMR population inversion in the presence of radiofrequency field inhomogeneity and frequency offsets. *J. Magn. Reson.* **66**(3), 470–482 (1986). [https://doi.org/10.1016/0022-2364\(86\)90190-3](https://doi.org/10.1016/0022-2364(86)90190-3)
- Laane, J.: *Frontiers and Advances in Molecular Spectroscopy*. Elsevier Science, Amsterdam (2017). <https://books.google.co.jp/books?id=DeDWDgAAQBAJ>. Accessed 19 June 2020
- Schmidgall, E. R., Eastham, P. R., Phillips, R. T.: Population inversion in quantum dot ensembles via adiabatic rapid passage. *Phys. Rev. B*. **81**, 195306 (2010). <https://doi.org/10.1103/PhysRevB.81.195306>
- Feynman, R. P., Vernon Jr, F. L., Hellwarth, R. W.: Geometrical representation of the Schrödinger equation for solving maser problems. *J. Appl. Phys.* **28**(1), 49–52 (1957). <https://doi.org/10.1063/1.1722572>
- Torrentegui, E., Ibanez, S., Martinez-Garaot, S., Modugno, M., del Campo, A., Guery-Odelin, D., Ruschhaupt, A., Chen, X., Gonzalo Muga, J.: Shortcuts to adiabaticity. *Advances in Atomic Molecular and Optical Physics*. **62**, 117–169 (2013). <https://doi.org/10.1016/B978-0-12-408090-4.00002-5>
- Chen, X., Ruschhaupt, A., Schmidt, S., del Campo, A., Guery-Odelin, D., Muga, J. G.: Fast optimal frictionless atom cooling in harmonic traps: Shortcut to adiabaticity. *Phys. Rev. Lett.* **104**, 063002 (2010). <https://doi.org/10.1103/PhysRevLett.104.063002>
- Chen, X., Lizuain, I., Ruschhaupt, A., Guery-Odelin, D., Muga, J. G.: Shortcut to adiabatic passage in two- and three-level atoms. *Phys. Rev. Lett.* **105**, 123003 (2010). <https://doi.org/10.1103/PhysRevLett.105.123003>
- Chen, X., Muga, J. G.: Engineering of fast population transfer in three-level systems. *Phys. Rev. A*. **86**, 033405 (2012). <https://doi.org/10.1103/PhysRevA.86.033405>
- Berry, M. V.: Transitionless quantum driving. *J. Phys. A Math. Theor.* **42**(36), 365303 (2009). <https://doi.org/10.1088/1751-8113/42/36/365303>
- del Campo, A.: Shortcuts to adiabaticity by counterdiabatic driving. *Phys. Rev. Lett.* **111**, 100502 (2013). <https://doi.org/10.1103/PhysRevLett.111.100502>
- Masuda, S., Rice, S. A.: Fast-forward assisted STIRAP. *J. Phys. Chem. A*. **119**(14), 3479–3487 (2015). <https://doi.org/10.1021/acs.jpca.5b00525>. PMID: 25775133
- Bason, M. G., Viteau, M., Malossi, N., Huillery, P., Arimondo, E., Ciampini, D., Fazio, R., Giovannetti, V., Mannella, R., Morsch, O.: High-fidelity quantum driving. *Nat. Phys.* **8**(2), 147–152 (2012). <https://doi.org/10.1038/NPHYS217>
- Zhang, J., Shim, J. H., Niemeyer, I., Taniguchi, T., Teraji, T., Abe, H., Onoda, S., Yamamoto, T., Ohshima, T., Isoya, J., Suter, D.: Experimental implementation of assisted quantum adiabatic passage in a single spin. *Phys. Rev. Lett.* **110**, 240501 (2013). <https://doi.org/10.1103/PhysRevLett.110.240501>
- An, S., Lv, D., del Campo, A., Kim, K.: Shortcuts to adiabaticity by counterdiabatic driving for trapped-ion displacement in phase space. *Nat. Commun.* **7** (2016). <https://doi.org/10.1038/ncomms12999>
- Du, Y.-X., Liang, Z.-T., Li, Y.-C., Yue, X.-X., Lv, Q.-X., Huang, W., Chen, X., Yan, H., Zhu, S.-L.: Experimental realization of stimulated Raman shortcut-to-adiabatic passage with cold atoms. *Nat. Commun.* **7** (2016). <https://doi.org/10.1038/ncomms1247>
- Guery-Odelin, D., Ruschhaupt, A., Kiely, A., Torrentegui, E., Martinez-Garaot, S., Muga, J. G.: Shortcuts to adiabaticity: Concepts, methods, and applications. *Rev. Mod. Phys.* **91**(4) (2019). <https://doi.org/10.1103/RevModPhys.91.04500>
- Ibanez, S., Martinez-Garaot, S., Chen, X., Torrentegui, E., Muga, J. G.: Shortcuts to adiabaticity for non-Hermitian systems. *Phys. Rev. A*. **84**, 023415 (2011). <https://doi.org/10.1103/PhysRevA.84.023415>
- Toroso, B. T., Della Valle, G., Longhi, S.: Non-Hermitian shortcut to adiabaticity. *Phys. Rev. A*. **87**, 052502 (2013). <https://doi.org/10.1103/PhysRevA.87.052502>

26. Torosov, B. T., Della Valle, G., Longhi, S.: Non-Hermitian shortcut to stimulated Raman adiabatic passage. *Phys. Rev. A*. **89**, 063412 (2014). <https://doi.org/10.1103/PhysRevA.89.063412>
27. Bender, C. M., Boettcher, S.: Real Spectra in Non-Hermitian Hamiltonians Having \mathcal{PT} Symmetry. *Phys. Rev. Lett.* **80**, 5243–5246 (1998). <https://doi.org/10.1103/PhysRevLett.80.5243>
28. Bender, C. M.: Making sense of non-Hermitian Hamiltonians. *Rep. Prog. Phys.* **70**(6), 947–1018 (2007). <https://doi.org/10.1088/0034-4885/70/6/r03>
29. Uzdin, R., Guenther, U., Rahav, S., Moiseyev, N.: Time-dependent Hamiltonians with 100% evolution speed efficiency. *J. Phys. A Math. Theor.* **45**(41), 415304 (2012). <https://doi.org/10.1088/1751-8113/45/41/415304>
30. Reyes, S. A., Olivares, F. A., Morales-Molina, L.: Landau-Zener-Stuckelberg interferometry in \mathcal{PT} -symmetric optical waveguides. *J. Phys. A Math. Theor.* **45**(44), 444027 (2012). <https://doi.org/10.1088/1751-8113/45/44/444027>
31. Uzdin, R., Moiseyev, N.: Rapid azimuthal rotation in the Hermitian and non-Hermitian Landau–Zener problem. *J. Phys. A Math. Theor.* **45**(44), 444033 (2012). <https://doi.org/10.1088/1751-8113/45/44/444033>
32. Ibanez, S., Muga, J. G.: Adiabaticity condition for non-Hermitian Hamiltonians. *Phys. Rev. A*. **89**, 033403 (2014). <https://doi.org/10.1103/PhysRevA.89.033403>
33. Marr, C., Beige, A., Rempe, G.: Entangled-state preparation via dissipation-assisted adiabatic passages. *Phys. Rev. A*. **68**(3) (2003). <https://doi.org/10.1103/PhysRevA.68.03381>
34. Martinez-Garaot, S., Tseng, S.-Y., Muga, J. G.: Compact and high conversion efficiency mode-sorting asymmetric Y junction using shortcuts to adiabaticity. *Opt. Lett.* **39**(8), 2306–2309 (2014). <https://doi.org/10.1364/OL.39.00230>
35. Tseng, S.-Y., Wen, R.-D., Chiu, Y.-F., Chen, X.: Short and robust directional couplers designed by shortcuts to adiabaticity. *Opt. Express*. **22**(16), 18849–18859 (2014). <https://doi.org/10.1364/OE.22.018849>
36. Rueter, C. E., Makris, K. G., El-Ganainy, R., Christodoulides, D. N., Segev, M., Kip, D.: Observation of parity-time symmetry in optics. *Nat. Phys.* **6**(3), 192–195 (2010). <https://doi.org/10.1038/NPHYS151>
37. Guo, A., Salamo, G. J., Duchesne, D., Morandotti, R., Volatier-Ravat, M., Aimez, V., Siviloglou, G. A., Christodoulides, D. N.: Observation of \mathcal{PT} -symmetry breaking in complex optical potentials. *Phys. Rev. Lett.* **103**, 093902 (2009). <https://doi.org/10.1103/PhysRevLett.103.093902>
38. Vasilev, G. S., Vitanov, N. V.: Coherent excitation of a two-state system by a Gaussian field. *Phys. Rev. A*. **70**, 053407 (2004). <https://doi.org/10.1103/PhysRevA.70.053407>
39. Allen, L., Eberly, J. H.: Optical Resonance and Two-level Atoms, Dover books on physics and chemistry. Dover, New York (1987). <https://books.google.co.jp/books?id=1q0ae-XNmWwC>. Accessed 19 June 2020
40. Okamoto, K.: Fundamentals of Optical Waveguides, Electronics & Electrical. Elsevier Science, Amsterdam (2006). <https://books.google.co.jp/books?id=0SNXgnx6un4C>. Accessed 19 June 2020
41. Yariv, A.: Coupled-mode theory for guided-wave optics. *IEEE J. Quantum Electron.* **9**(9), 919–933 (1973). <https://doi.org/10.1109/JQE.1973.1077767>
42. Syahriar, A., Schneider, V. M., Al-Bader, S.: The design of mode evolution couplers. *J. Light. Technol.* **16**(10), 1907–1914 (1998). <https://doi.org/10.1109/50.721079>

Publisher's Note

Springer Nature remains neutral with regard to jurisdictional claims in published maps and institutional affiliations.

Submit your manuscript to a SpringerOpen[®] journal and benefit from:

- Convenient online submission
- Rigorous peer review
- Open access: articles freely available online
- High visibility within the field
- Retaining the copyright to your article

Submit your next manuscript at ► [springeropen.com](https://www.springeropen.com)