Coherent control of spin tunneling in a driven spin–orbit coupled bosonic triple well

Yuxin Luo\(^1\), Jia Yi\(^2\), Wenjuan Li\(^1\), Xin Xie\(^1\), Yunrong Luo\(^1\) and Wenhua Hai\(^1\)

\(^1\) Key Laboratory for Matter Microstructure and Function of Hunan Province, and Key Laboratory of Low-dimensional Quantum Structures and Quantum Control of Ministry of Education, School of Physics and Electronics, Hunan Normal University, Changsha 410081, China

\(^2\) School of Management, Hunan University of Information Technology, Changsha 410151, China

E-mail: lyr_1982@hunnu.edu.cn

Received 21 January 2022, revised 2 April 2022
Accepted for publication 7 April 2022
Published 3 May 2022

Abstract

We investigate the coherent control of spin tunneling for a spin–orbit (SO) coupled boson trapped in a driven triple well. In the high-frequency limit, the quasienergies of the system are obtained analytically and the fine energy band structures are shown. By regulating the driving parameters, we reveal that the directed spin-flipping or spin-conserving tunneling of an SO-coupled boson occurs along different pathways and in different directions. The analytical results are demonstrated by numerical simulations and good agreements are found. Further, an interesting scheme of quantum spin tunneling switch with or without spin-flipping is presented. The results may have potential applications in the design of spintronic devices.

Keywords: coherent control, spin tunneling, spin–orbit coupling, triple well

(Some figures may appear in colour only in the online journal)

1. Introduction

Spin–orbit (SO) coupling is the interaction between spin and motion of a particle, which plays a crucial role in many important physical phenomena, e.g. spin-Hall effect [1], topological insulator [2], and the persistent spin helix [3]. For ultracold atoms, SO coupling can be created by utilizing the interaction between laser and atom, which yields Abelian or non-Abelian gauge fields for ultracold atoms in the dressed two hyperfine atomic internal states. As is well-known, the electrons are fermions in the material, but ultracold atoms may be bosons. Therefore, SO-coupled ultracold bosonic gases will lead to novel SO physics that has not been explored in solid materials.

Recently, SO coupling of ultracold atomic gases has been realized in experiments [4–9], which provides a brand new platform to investigate SO coupling physics, due to the unprecedented tunability of experimental parameters. A number of research works have focused on the interesting dynamics of SO-coupled ultracold atoms, for instance, quantum dynamics of SO-coupled Bose–Einstein condensates (BECs) in a double well [10–13], coherent control of an SO-coupled atom in a double-well potential [14], controlling stable spin tunneling in a non-Hermitian double-well system [15], Anderson localization of SO-coupled ultracold atomic gases in an optical lattice [16], Macroscopic Klein tunneling in SO-coupled BECs [17], Landau–Zener transition in an SO-coupled BEC [18], spin dynamics of SO-coupled BECs [19], nonequilibrium dynamics of two-component bosons in an optical lattice [20], controlling localization and directed motion of an SO-coupled single atom in a biparticle lattice [21], Bloch oscillation dynamics of an SO-coupled ultracold atomic gas in an optical lattice [22], quantum tunneling of an SO-coupled ultracold atom in an optical lattice with an impurity [23], controlling second-order tunneling of an SO-coupled atom in optical lattices [24], and so on.

As mentioned above, many works focus on the tunneling dynamical properties of SO-coupled ultracold atomic gases held in a double well or an optical lattice. To the best of our knowledge, the coherent control of spin tunneling for SO-coupled ultracold atoms confined in a triple well is rarely investigated. However, the triple-well model is an important one to study the coherent control of spin tunneling and is a bridge between the double-well and optical lattice models for a better understanding of the spin tunneling dynamics of SO-coupled ultracold atoms in the quantum wells. Thus, it
motivates us to investigate the quantum spin tunneling in an SO-coupled bosonic triple-well system.

In this paper, we theoretically study the coherent control of spin tunneling for an SO-coupled boson confined in a driven triple well. In the high-frequency limit, the quasienergies of the SO-coupled ultracold atomic triple-well system are analytically obtained and the quasienergy spectra are shown. By adjusting the strength of the time-dependent driving field, we can manipulate the directed spin-flipping or spin-conserving tunneling of an SO-coupled boson along different pathways and in different directions. Further, we present an intriguing scheme of a quantum spin switch for transporting an SO-coupled boson accompanied with or without spin-flipping from well 1 to well 3. These results may be useful for the design of spintronic devices [25].

2. Analytical solutions and quasienergy spectra in the high-frequency limit

We consider an SO-coupled ultracold atom confined in a driven triple well and the Hamiltonian of this system reads [21, 24]

\[
\hat{H}(t) = -\nu (\hat{a}_1^\dagger e^{-i\gamma t}\hat{a}_2 + \hat{a}_2^\dagger e^{-i\gamma t}\hat{a}_3 + \text{H.c.}) + \frac{\Omega}{2} \sum_{j=1}^{3} (\hat{n}_j - \hat{\epsilon}_j) + \sum_{j=1}^{3} [\hat{\epsilon}_j(\hat{n}_j - \nu^2) - \hat{\epsilon}_j(\hat{n}_j + \nu^2)],
\]

(1)

where \(\hat{a}_j^\dagger = (\hat{a}_{j,1}^\dagger, \hat{a}_{j,2}^\dagger)\) and \(\hat{a}_j = (\hat{a}_{j,1}, \hat{a}_{j,2})^T\) (the superscript T denotes the transpose), \(\hat{a}_{j,1}, \hat{a}_{j,2}\) is the annihilation (creation) operator of a pseudospin-\(1/2\) atom in the \(j\)th \((j = 1, 2, 3)\) well. \(\nu\) is the tunneling rate in the absence of SO coupling, \(\gamma\) denotes the strength of SO coupling, \(\hat{\epsilon}_j\) is the usual Pauli operator, the parameter \(\Omega\) denotes the effective Zeeman-field strength, and \(\hat{n}_j = \hat{a}_j^\dagger \hat{a}_j\) is the number operator. \(\hat{a}_j(\nu^2)\) and \(\hat{\epsilon}_j(\nu^2)\) are the periodic driving fields applied to well 1 and well 3 respectively [26, 27]. Here \(\hat{\epsilon}_j\) and \(\hat{\epsilon}_j\) are the driving amplitudes, and \(\omega\) is the frequency of the driving field. Throughout this paper, \(\hbar = 1\) and the parameters \(\nu, \Omega, \hat{\epsilon}_1, \hat{\epsilon}_2, \omega\) are normalized in units of the reference frequency \(\omega_0 = 0.1E_s\) with \(E_s = k_B^2/2m\) being 22.5 KHz being the single-photon recoil energy [4], and time \(t\) is measured in units of \(1/\omega_0\). In the experiment, the system parameters can be adjusted in a wide range as follows [13, 15, 21, 28]: \(\nu \sim \omega_0, \hat{\epsilon}_{1,2,3} \sim \omega \in [0, 100]G_{\omega_0}, \Omega \sim \omega\).

We employ the Fock state basis \(|\sigma, 0, 0\rangle\) (or \(|0, 0, \sigma\rangle\) or \(|0, \sigma, 0\rangle\)) to denote the state of a pseudospin-\(1/2\) atom occupying the well 1 (or well 2 or well 3) and no atom in the other wells, the quantum state of the SO-coupled ultracold atomic system can be expanded as

\[
|\psi(t)\rangle = a_1(t)|\uparrow, 0, 0\rangle + a_2(t)|\downarrow, 0, 0\rangle + a_3(t)|0, 1, 0\rangle + a_4(t)|0, 0, 1\rangle + a_5(t)|0, 0, 0\rangle + a_6(t)|0, 0, \downarrow\rangle,
\]

(2)

where \(a_j(t) (k = 1, 2, ..., 6)\) represents the probability amplitude of the boson being in the Fock state \(|\sigma, 0, 0\rangle\) or \(|0, \sigma, 0\rangle\) or \(|0, 0, \sigma\rangle\) (e.g. \(a_1(t)\) represents the probability amplitude of the boson being in Fock state \(|\uparrow, 0, 0\rangle\) or \(|\downarrow, 0, 0\rangle\)). The corresponding probability of the Fock state reads \(P_k(t) = |a_k(t)|^2\), which satisfies the normalization condition \(\sum_{k=1}^{6} P_k(t) = 1\). Inserting equations (1) and (2) into the time-dependent Schrödinger equation \(i\hbar \frac{\partial |\psi(t)\rangle}{\partial t} = \hat{H}(t)|\psi(t)\rangle\), one can obtain the coupled equations

\[
\begin{align*}
\dot{a}_1(t) &= \left[\frac{\Omega}{2} + \epsilon_1 \cos(\omega t)\right]a_1(t) - \nu \cos(\gamma t) a_2(t) + \nu \sin(\gamma t) a_4(t), \\
\dot{a}_2(t) &= \left[-\frac{\Omega}{2} + \epsilon_1 \cos(\omega t)\right]a_2(t) - \nu \sin(\gamma t) a_3(t) - \nu \cos(\gamma t) a_4(t), \\
\dot{a}_3(t) &= -\nu \cos(\gamma t) a_1(t) + a_5(t) - \nu \sin(\gamma t) a_2(t) - \nu \cos(\gamma t) a_4(t), \\
\dot{a}_4(t) &= \nu \sin(\gamma t) a_1(t) - a_3(t) - \nu \cos(\gamma t) a_2(t) - \nu \sin(\gamma t) a_5(t) - \nu \cos(\gamma t) a_6(t), \\
\dot{a}_5(t) &= \nu \sin(\gamma t) a_1(t) + a_3(t) - \nu \cos(\gamma t) a_2(t) - \nu \sin(\gamma t) a_4(t) - \nu \cos(\gamma t) a_6(t), \\
\dot{a}_6(t) &= \nu \sin(\gamma t) a_1(t) - a_4(t) - \nu \cos(\gamma t) a_2(t) + \nu \sin(\gamma t) a_5(t) + \nu \cos(\gamma t) a_4(t).
\end{align*}
\]

(3)

Clearly, equation (3) is very hard to be solved exactly, because of the time-dependent coefficients. However, the spin tunneling dynamics of the SO-coupled ultracold atomic system can be studied analytically in the high-frequency limit [29, 30]. Therefore, we introduce the slowly varying function of time \(b_k(t)(k = 1, 2, ..., 6)\) through the transformation

\[
\begin{align*}
a_{1,2}(t) &= b_{1,2}(t)e^{-i\frac{\Omega}{2}t + \epsilon_1 \cos(\omega t)dt}, \\
a_{3,4}(t) &= b_{3,4}(t)e^{i\frac{\Omega}{2}t - \epsilon_1 \cos(\omega t)dt}, \\
a_{5,6}(t) &= b_{5,6}(t)e^{i\frac{\Omega}{2}t + \epsilon_1 \cos(\omega t)dt}.
\end{align*}
\]

(4)

In the high-frequency limit and by using of the Fourier expansion \(e^{i\omega_0 t} = \sum_{\mu = -\infty}^{\infty} J_\mu(\omega_0) e^{i\mu \omega t}\) and \(e^{i\omega t}e^{i\omega_0 t} = \sum_{\mu = -\infty}^{\infty} \nu \mu J_\mu(\omega_0) e^{i(\mu \omega + \omega_0) t}\), the rapidly oscillating terms of the Fourier expansion with \(n \neq 0\) and \(n' \pm \frac{\Omega}{2} \neq 0\) can be neglected [31] and the equation (3) reduces to the form

\[
\begin{align*}
\dot{b}_1(t) &= -J_1 b_3(t) + J_2 b_4(t), \\
\dot{b}_2(t) &= -J_2 b_1(t) + J_3 b_4(t), \\
\dot{b}_3(t) &= -J_3 b_1(t) - J_4 b_2(t) - J_5 b_3(t) + J_6 b_5(t), \\
\dot{b}_4(t) &= J_4 b_1(t) - J_5 b_2(t) - J_6 b_3(t) - J_7 b_6(t), \\
\dot{b}_5(t) &= -J_6 b_1(t) - J_7 b_3(t) + J_8 b_6(t), \\
\dot{b}_6(t) &= J_7 b_2(t) + J_8 b_4(t) - J_9 b_6(t).
\end{align*}
\]

(5)

Here, \(J_\lambda = \nu \cos(\gamma t) J_\lambda(\omega_0), J_{2\lambda} = \nu \sin(\gamma t) J_{2\lambda}(\omega_0), J_4 = \nu \cos(\gamma t) J_4(\omega_0)\), and \(J_{5,6} = \nu \sin(\gamma t) J_{5,6}(\omega_0)\) are the renormalized coupling constants and \(J_\lambda(\nu x)\) is the \(n\)-order Bessel function of \(x\). It is worth noting that the spin tunneling dynamics of the original system (1) can be described effectively by equation (5), which is the basis of the following analysis.
Before moving on, we first briefly introduce the Floquet theorem [32, 33]. The basis of this theorem lies in the observation that for a time-dependent Hamiltonian with period \( \tau \), \( \hat{H}(t) = \hat{H}(t + \tau) \), there exists a set of complete bases \( \{ |\psi_p(t)\rangle \} \) which constitute the solutions of Schrödinger equation \( i\frac{\partial |\psi(t)\rangle}{\partial t} = \hat{H}(t)|\psi(t)\rangle \) in the form

\[
|\psi(t)\rangle = |\varphi_p(t)\rangle e^{-i\Delta E_p t}, |\varphi_p(t)\rangle = |\varphi_p(t + \tau)\rangle,
\]

where the time-periodic functions \( |\varphi_p(t)\rangle \) are called Floquet states and \( E_p \) are quasienergies. For the system we are studying whose Hamiltonian is periodic in time with period \( \tau = 2\pi/\omega \), according to the Floquet theorem, the Floquet solution of this system can be constructed as \( |\psi(t)\rangle = |\varphi(t)\rangle e^{-i\Delta E t} \), where \( E \) is the Floquet quasienergy and \(|\varphi(t)\rangle = A_0 e^{-i\frac{\Delta E}{2} t}[|0\rangle, 0] + A_1 e^{-i\frac{\Delta E}{2} t}[|1\rangle, 0] + A_2 e^{-i\frac{\Delta E}{2} t}[|0\rangle, 1] + A_3 e^{-i\frac{\Delta E}{2} t}[|1\rangle, 1] \) is the Floquet state. Based on the superposition principle of quantum mechanics, the non-Floquet state can be obtained by the linear superposition of the Floquet states [15, 34–36], which implies the enhancement or suppression of quantum tunneling. We have introduced the stationary solution of equation (5) as \( b_0(t) = A_0 e^{-i\Delta E t} \) with \( A_k \) being constant. Inserting such a form of \( b_0(t) \) into equation (5), the constant \( A_k \) and Floquet quasienergy \( E \) can be obtained. Here, we only give the Floquet quasienergies as

\[
E_{1,2} = 0, E_{3,4} = \pm \sqrt{\frac{\alpha - \sqrt{\beta}}{2}}, E_{5,6} = \pm \sqrt{\frac{\alpha + \sqrt{\beta}}{2}}.
\]

with the constants

\[
\alpha = 2J_1^2 + J_2^2 + J_3^2 + 2J_4^2 + J_5^2 + J_6^2, \quad \beta = J_1^2 + 4J_2^2(J_2 - J_3)^2 + J_4^2 + 2J_2^2 J_5^2 + 2J_3^2 J_5^2 + 2J_3^2 J_6^2 + J_5^4 - 8J_2^2 J_5 J_6 + 6J_4^2 - 2J_2^2 - 2J_3^2 + J_5^3 - 2J_6^2(J_3 + J_6^2 - J_5^2).
\]

Clearly, the quasienergies depend on the system parameters, except for the degenerate zero quasienergies \( E_1 \) and \( E_2 \).

From equation (7), it is surprisingly found that when the SO coupling strength \( \gamma \) is an integer or half integer, the constant \( \beta \) will equal to zero. This will result in the quasienergies \( E_3 = E_5 = -\sqrt{\alpha/2} \) and \( E_4 = E_6 = \sqrt{\alpha/2} \), which means that the new quasienergy degeneracy occurs. To our knowledge, the degeneracy of energy levels generally implies quantum decoherence, so it will lead to the intriguing phenomenon of quantum tunneling, for instance, the selective coherent destruction of tunneling (SCDT) popularly occurs at the degeneracy (crossing) point of the partial energy levels [34–36] and coherent destruction of tunneling (CDT) commonly happens at the collapse (crossing) point of all energy levels [37, 38]. As an example, we fix the parameters \( \gamma = 0.5, \omega = 50, \Omega = 100, \nu = 1 \), and (a) \( \varepsilon_2 = 5.1356/\omega \), (b) \( \varepsilon_2 = 2\omega \) to plot the quasienergy spectra with quasienergy as a function of the driving parameter \( \varepsilon_1/\omega \), as shown in figures 1(a)–(b), respectively. Here, the circle points denote the analytical results based on equation (7) and the solid curves label the numerical results from the original model (1). The two sets of parameters lead to \( E_3 = E_5 = -|J_3(\frac{\alpha}{2})| \) and \( E_4 = E_6 = |J_3(\frac{\alpha}{2})| \). Here, circles denote the numerical results and solid curves label the numerical correspondences. Hereafter, any parameter adopted in the figures is dimensionless.

![Figure 1. Quasienergy as a function of the driving parameter \( \varepsilon_1/\omega \) for \( \gamma = 0.5, \nu = 1, \omega = 50, \Omega = 100 \), and (a) \( \varepsilon_2 = 5.1356/\omega \), (b) \( \varepsilon_2 = 2\omega \). Here, circles denote the analytical results and solid curves label the numerical correspondences. Hereafter, any parameter adopted in the figures is dimensionless.](image-url)
energy band structure shown in figure 1(a), the quasienergy curves are divided to three energy bands in figure 1(b). The energy gap between two adjacent energy bands is equal to the minimum value $J_2(2) \approx 0.3528$ of quasienergy $E_4$ or $E_6$. Furthermore, from figure 1(b), it can be seen that the two energy-level curves of each band from the numerical results are approximately degenerate (e.g. quasienergy curves of $E_1$ and $E_2$, $E_3$ and $E_5$, $E_4$ and $E_6$), which have some small deviations from analytical results (circles).

3. Directed tunneling and spin switch with or without spin-flipping

Coherent manipulation of quantum tunneling is an interesting research subject in quantum information technologies [39, 40]. Generally, quantum tunneling of ultracold atomic systems relies intensively on the periodic driving external field. Thus, the quantum spin tunneling and transport of an SO-coupled ultracold boson held in a triple well can be manipulated by adjusting the driving parameters. In the section, we will focus on studying the directed tunneling and spin tunneling switch with or without spin-flipping.

3.1. Directed tunneling with or without spin-flipping

For an SO-coupled ultracold atomic triple-well system, an attractive physical problem is how to control the directed spin-flipping or spin-conserving tunneling of an SO-coupled atom from the initial middle well to the left well or to the right well. Based on equation (5), it can be seen that when the SO-coupling strength $\gamma$ is half integer or integer, quantum tunneling with or without spin-flipping can occur. Here, we will try to manipulate the directed spin-flipping tunneling of a spin-up boson initially located in the well 2 as an example. In order to realize the directed spin-flipping tunneling of the spin-up particle from well 2 to well 1, it is easily found from equation (5) that when the effective coupling constants satisfy $J_1 = J_2 = J_3 = 0$ and $J_3 \neq 0$, $b_1$ is only coupled with $b_2$. It means that the SCDT occurs. We take the parameters $\gamma = 0.5$, $\nu = 1$, $\omega = 50$, $\Omega = 100$, $\varepsilon_1 = 2\omega$, $\varepsilon_2 = 5.1356\omega$ satisfying the conditions $J_1 = J_2 = J_3 = 0$ and $J_3 \neq 0$ to plot the time evolutions of the probabilities in figure 2(a). Clearly, the spin-flipping tunneling channel between well 2 and well 3 is closed and the particle performs a Rabi oscillation of spin-flipping along this pathway between well 2 and well 1, namely, the directed spin-flipping tunneling takes place between initial state $|0, \uparrow, 0\rangle$ and state $|\downarrow, 0, 0\rangle$ with tunneling time $\Delta t \approx 4.4$. It means the occurrence of SCDT and the set of parameters corresponds to the degeneracy position of the partial quasienergies in figure 1(a). If we set the same parameters as in figure 2(a) except for $\varepsilon_1 = 5.1356\omega$, which corresponds to the collapse (crossing) point of all the quasienergies in figure 1(a), the quantum tunneling of the particle will be frozen. It means the CDT occurs as shown in figure 2(b).

Further, from equation (5) it can also be found when the effective coupling constants satisfy $J_1 = J_3 = J_4 = 0$ and $J_5 \neq 0$, the probability function $b_3$ is only related to function $b_6$, which means the SCDT happens. We fix the same initial conditions and parameters as that in figure 2(a) except for $\varepsilon_1 = 5.1356\omega$ and $\varepsilon_2 = 2\omega$ to plot the time evolutions of the probabilities in figure 2(c). Clearly, the directed spin-flipping tunneling along another pathway between well 2 and well 3 occurs, in which the spin particle performs a spin-flipping Rabi oscillation between state $|0, \uparrow, 0\rangle$ and state $|0, 0, \downarrow\rangle$ with tunneling time $\Delta t \approx 4.4$. It means the SCDT occurs and the set of parameters corresponds to the degeneracy location of the partial quasienergies in figure 1(b). The analytical results (dashed lines) based on equation (5) are in perfect agreement with the numerical results from the accurate model (3) in figure 2. Because the above-mentioned results are related to controlling the directed spin-flipping tunneling of a spin boson, it is similar to the case of manipulating the directed spin-conserving tunneling of the particle.

3.2. Quantum spin tunneling switch with or without spin-flipping

From the above subsection, we find the time-dependent driving field affects dramatically spin tunneling dynamics of this system, thus we can propose a scheme of quantum spin tunneling switch with or without spin-flipping by means of sudden regulating of driving parameters to control the occurrence and suppression of quantum tunneling. The means have been performed in many research works [41–44]. Here, we present a scheme of quantum spin tunneling switch with spin-flipping by adjusting the driving strength as an example, as shown in figure 3. In figure 3(a), we take a spin-up boson initially occupied in well 1 and set the parameters $\gamma = 0.5$, $\nu = 1$, $\omega = 50$, $\Omega = 100$, $\varepsilon_1 = \varepsilon_2 = 5.1356\omega$ corresponding to...
attributed to the effect of CDT. At any given time the collapse (crossing) point of all the quasienergies in figure 1(a). One can see that the spin-up boson is frozen in well 1, due to the CDT effect. At any given time $t = t_1 = 2$, the driving strength $\varepsilon_1$ is changed to $\varepsilon_1 = 2\omega$ and holds this value until the time is $t = t_2 = 6.4$. At the moment, the spin-up particle tunnels completely from state $|\uparrow, 0, 0\rangle$ to state $|0, \downarrow, 0\rangle$. Then we immediately adjust the driving strength $\varepsilon_1$ to $\varepsilon_1 = 5.1356\omega$, such that the state $|0, \downarrow, 0\rangle$ is kept which is attributed to the effect of CDT. At any given time $t = t_3 = 8.4$, we regulate the driving strength $\varepsilon_2$ to $\varepsilon_2 = 2\omega$ and preserve this value until $t = t_4 = 12.8$. At this time, the spin-down particle tunnels completely from state $|0, \downarrow, 0\rangle$ to state $|0, 0, \uparrow\rangle$. Then, we return the driving strength $\varepsilon_2$ to $\varepsilon_2 = 5.1356\omega$, so that the final state $|0, 0, \uparrow\rangle$ is kept, due to the CDT effect. Therefore, the spin-up boson is successfully transported through two spin-flipping tunnels from well 1 to well 3 by adjusting the driving strength $\varepsilon_i$ ($i = 1, 2$) of the driving external fields, namely, the quantum spin tunneling switch with spin-flipping is theoretically realized. From figure 3(a), it is also found that we can manipulate the spin-flipping tunneling of the spin particle from an arbitrarily initial occupied well to any other well. The corresponding spatial distributions of the spin particle at the tuning moments are shown in figure 3(b), where $\Delta t_i$ denotes transferring time between the different populations and the transferring time $\Delta t_1 = \Delta t_2 = t_2 - t_1 = t_4 - t_3 = 4.4$. Similarly, the quantum spin tunneling switch without spin-flipping can also be performed by adjusting the driving parameters.

4. Conclusion and discussion

In conclusion, we have investigated the coherent control of spin tunneling for an SO-coupled boson trapped in a driven triple well. In the high-frequency limit, we analytically obtain the Floquet quasienergies of the system and the fine quasienergy spectra are shown. By modulating the strength of the driving external field, we can control the directed spin-flipping or spin-conserving tunneling of an SO-coupled boson along different pathways and in different directions from the middle well to the left well or to the right well. The analytical results are demonstrated by numerical calculations and perfect agreements between both are shown. Further, based on the combined effect of CDT and SCDT, we propose a diverting scheme of quantum spin tunneling switch in which a spin particle is transported with or without spin-flipping from the initial well 1 to the final well 3. It is worth noting that we also can manipulate the spin-flipping or spin-conversing tunneling of a spin boson from an arbitrarily initial occupied well to any other well. These results may be useful in the design of spintronic devices and quantum information technology.

Acknowledgments

This work was supported by the Scientific Research Foundation of Hunan Provincial Education Department under Grants No. 21B0063 and No. 18C0027, the Hunan Provincial Natural Science Foundation of China under Grants No. 2021JJ30435 and No. 2017JJ3208, and the National Natural Science Foundation of China under Grant No. 11747034.

References

[1] Kato Y K, Myers R C, Gossard A C and Awschalom D D 2004 Observation of the spin Hall effect in semiconductors Science 306 1910
[2] Bernevig B A, Hughes T L and Zhang S C 2006 Quantum spin Hall effect and topological phase transition in HgTe quantum wells Science 314 1757
[3] Koralek J D, Weber C P, Orenstein J, Bernevig B A, Zhang S C, Mack S and Awschalom D D 2009 Emergence of the persistent spin helix in semiconductor quantum wells Nature 458 610
[4] Lin Y J, Jiménez-García K and Spielman I B 2011 Spin–orbit-coupled Bose–Einstein condensates Nature 471 83
[5] Wang P J, Yu Z Q, Fu Z K, Miao J, Huang L H, Chai S J, Zhai H and Zhang J 2012 Spin–orbit coupled degenerate Fermi gases Phys. Rev. Lett. 109 095301
[6] Cheuk L W, Sommer A T, Hadzibabic Z, Yefsah T, Bakr W S and Zwierlein M W 2012 Spin-injection spectroscopy of a spin–orbit coupled Fermi gas Phys. Rev. Lett. 109 095302
[7] Zhang J et al 2012 Collective dipole oscillations of a spin–orbit coupled Bose–Einstein condensate Phys. Rev. Lett. 109 115301
[8] Huang L, Meng Z, Wang P, Peng P, Zhang S L, Chen L, Li D, Zhou Q and Zhang J 2016 Experimental realization of two-dimensional synthetic spin–orbit coupling in ultracold Fermi gases Nat. Phys. 12 540
[9] Wu Z, Zhang L, Sun W, Xu X T, Wang B Z, Ji S C, Deng Y, Chen S, Liu X J and Pan J W 2016 Realization of two-dimensional spin–orbit coupling for Bose–Einstein condensates Science 354 83
[10] Zhang D, Fu L, Wang Z and Zhu S 2012 Josephson dynamics of a spin–orbit-coupled Bose–Einstein condensate in a double-well potential Phys. Rev. A 85 043609
[11] Garcia-March M A, Mazzarella G, Dell’Anna L, Juliá-Díaz B, Salasnich L and Polli A 2014 Josephson physics of spin–orbit-coupled elongated Bose–Einstein condensates Phys. Rev. A 89 063607
[12] Citro R and Naddeo A 2015 Spin–orbit coupled Bose–Einstein condensates in a double well Eur. Phys. J. Spec. Top. 224 503
[13] Yu Z and Xue J 2014 Selective coherent spin transportation in a spin–orbit-coupled bosonic junction Phys. Rev. A 90 033618
[14] Wang W, Dou F and Duan W 2017 Coherent control of spin–orbit-coupled atom in a double-well potential Eur. Phys. J. D 71 294
[15] Luo Y, Wang X, Luo Y, Zhou Z, Zeng Z and Luo X 2020 Controlling stable tunneling in a non-Hermitian spin–orbit coupled bosonic junction New. J. Phys. 22 093041
[16] Zhou L, Pu H and Zhang W 2013 Anderson localization of cold atomic gases with effective spin–orbit interaction in a quasiperiodic optical lattice Phys. Rev. A 87 023625
[17] Zhang D, Xue Z, Yan H, Wang Z and Zhu S 2012 Macroscopic Klein tunneling in spin–orbit-coupled Bose–Einstein condensates Phys. Rev. A 85 013616
[18] Olson A J, Wang S, Niffenegger R J, Li C, Greene C H and Chen Y 2014 Tunable Landau–Zener transitions in a spin–orbit-coupled Bose–Einstein condensate Phys. Rev. A 90 013616
[19] Mardonov S, Modugno M and Sherman E Y 2015 Dynamics of spin–orbit coupled Bose–Einstein condensates in a random potential Phys. Rev. Lett. 115 180402
[20] Ng H T 2015 Nonequilibrium dynamics of spin–orbit-coupled lattice bosons Phys. Rev. A 92 043634
[21] Luo Y, Lu G, Kong C and Hai W 2016 Controlling spin-dependent localisation and directed transport in a bipartite lattice Phys. Rev. A 93 043609
[22] Ji W, Zhang K, Zhang W and Zhou L 2019 Bloch oscillations of spin–orbit-coupled cold atoms in an optical lattice and spin-current generation Phys. Rev. A 99 023604
[23] Luo X, Yang B, Cui J, Guo Y, Li L and Hu Q 2019 Dynamics of spin–orbit-coupled cold atomic gases in a Floquet lattice with an impurity J. Phys. B: At. Mol. Opt. Phys. 52 085301
[24] Luo X, Zeng Z, Guo Y, Yang B, Xiao J, Li L, Kong C and Chen A 2021 Controlling directed atomic motion and second-order tunneling of a spin–orbit-coupled atom in optical lattices Phys. Rev. A 103 043615
[25] Z’tur I, Fabian J and Sarma S D 2004 Spintronics: fundamentals and applications Rev. Mod. Phys. 76 323
[26] Li L, Luo X, Lu X, Yang X and Wu Y 2015 Coherent destruction of tunneling in a lattice array with a controllable boundary Phys. Rev. A 91 063804
[27] Zeng Z, Li L, Yang B, Xiao J and Luo X 2020 Coherent control of dissipative dynamics in a periodically driven lattice array Phys. Rev. A 102 012221
[28] Chen Y A, Nascimbène S, Aidelsburger M, Atala M, Trotzky S and Bloch I 2011 Controlling correlated tunneling and superexchange interactions with ac-driven optical lattices Phys. Rev. Lett. 107 210405
[29] Blanes S, Casas F, Oteo J A and Ros J 2009 The Magnus expansion and some of its applications Phys. Rep. 470 151
[30] Thimmel B, Nalbach P and Terzidis O 1999 Rotating wave approximation: systematic expansion and application to coupled spin pairs Eur. Phys. J. B 9 207
[31] Zou M, Lu G, Hai W and Zou R 2013 Quantum manipulation of tunneling for two bosons held in a driven triple-well J. Phys. B: At. Mol. Opt. Phys. 46 045004
[32] Shirley J H 1965 Solution of the Schrödinger equation with a Hamiltonian periodic in time Phys. Rev. 138 A979
[33] Sambe H 1973 Steady states and quasienergies of a quantum-mechanical system in an oscillating field Phys. Rev. A 7 2203
[34] Lu G, Fu L, Hai W, Zou M and Guo Y 2015 Directed selective-tunneling of bosons with periodically modulated interaction Phys. Lett. A 379 947
[35] Luo Y, Lu G, Tao J and Hai W 2015 Transparent control of three-body selective destruction of tunneling via unusual states J. Phys. B: At. Mol. Opt. Phys. 48 015002
[36] Luo Y, Hai K, Zou M and Hai W 2017 Coherent control of quasi-degenerate stationary-like states via multiple resonances Sci. Rep. 7 211
[37] Grossmann F, Dittrich T, Jung P and Hänggi P 1991 Coherent destruction of tunneling Phys. Rev. Lett. 67 516
[38] Liu J, Hai W and Zhou Z 2013 Coherent control via interplay between driving field and two-body interaction in a double well Phys. Lett. A 377 3078
[39] Monroe C, Meekhof D, King B and Wineland D 1996 A Schrödinger cat superposition state of an atom Science 272 1131
[40] Král P, Thanopoulos I and Shapiro M 2007 Colloquium: coherently controlled adiabatic passage Rev. Mod. Phys. 79 53
[41] Weiss C and Jinasundera T 2005 Coherent control of mesoscopic tunneling in a Bose–Einstein condensate Phys. Rev. A 72 053626
[42] Cerefield C E 2007 Quantum control and entanglement using periodic driving fields Phys. Rev. Lett. 99 110501
[43] Hai K, Hai W and Chen Q 2010 Controlling transport and entanglement of two particles in a bipartite lattice Phys. Rev. A 82 053412
[44] Luo X, Huang J and Lee C 2011 Coherent destruction of tunneling in a lattice array under selective in-phase modulations Phys. Rev. A 84 053847