Complete population transfer between nonresonant tunneling states induced by a train of laser pulses

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Complete population transfer between nonresonant tunneling states in a double-well system is predicted to occur when a train of *bright* or *dark* laser pulses drives transitions to an excited state in the system. The transition induces the phase shift of the wave packet involved in the transition and alters successive evolutions of the wave packets of the nonresonant states as a result of quantum interference, resulting in complete tunneling. [S1050-2947(99)00304-2]

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I. INTRODUCTION

One of the most fundamental problems of quantum mechanics is the quantum-dynamical tunneling occurring in different physical, chemical, and biological systems, and the quantum-dynamical tunneling phenomena in the presence of a laser field have received increasing attention in recent years [1–7]. In contrast to the numerous publications concerning the interaction of a two-state system with a single pulse or pulse trains [8–11], there are only a limited number of studies of interactions between a three-state system and pulse trains, which are particularly devoted to the so-called quantum Zeno effect [12–14]. The measurements by laser pulses in the quantum Zeno system destroy the coherence of the states involved in the transition through relaxations of the excited state.

In the present paper, we present a theoretical study of the coherent population transfer between nonresonant tunneling states in a double-well system under the influence of a train of laser pulses, which coherently changes the tunneling evolutions as a result of quantum interference. It is shown that the complete population transfer between nonresonant tunneling states can be realized by the excitation of the tunneling states with a train of *bright* or *dark* laser pulses and this can be explained by the phase shift of the probability amplitude of the tunneling state involved in the transition, which alters the successive evolution of the wave packets of the nonresonant states, resulting in complete tunneling.

II. DOUBLE-WELL SYSTEM AND PULSE EXCITATIONS

The quantum-dynamical tunneling can be observed in systems whose potential energy has at least two minima, with the form of the potential being locally symmetric. If the system is initially located in one of the wells of the potential, it will undergo tunneling oscillations having period $T_{\mathrm{tun}} = \pi/\kappa$, where κ is the tunneling coefficient [15,16]. Recently, Kilin, *et al.* [17] proposed a new scheme to dynamically suppress the tunneling by a laser field in the system of a molecule that is placed in an appropriate host medium. In effect, the laser field suppresses the tunneling by removing the degeneracy of the ground states. Recently, we adapted their idea to a semiconductor double-well system and predicted dynamical suppression and enhancement of the tunneling, which were due to the dynamic (ac) Stark shift of the states involved in the laser excitations [18].

Here we consider dynamical control of quantum tunneling by a train of laser pulses for an asymmetric double-well potential system as shown in Fig. 1. At low temperature the system can be modeled by a four-state system, each state being the lowest-energy state for corresponding local wells of asymmetric double-well potentials for the ground (V_g) and excited (V_e) states. We assume that the excited state of the right well is far from laser excitations. Therefore, three states $|a\rangle$, $|b\rangle$, and $|c\rangle$ are involved in the analysis of the

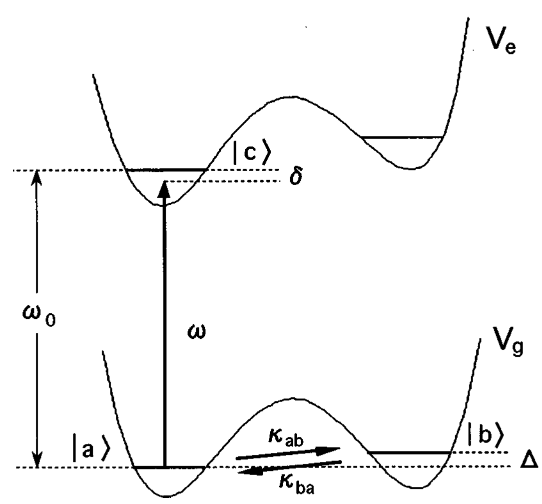

FIG. 1. Asymmetric double potential system investigated in the present paper. The states $|a\rangle$ and $|b\rangle$ in the ground state V_g are coupled by the tunneling through the potential barrier, and the ground state $|a\rangle$ and the excited state $|c\rangle$ are coupled by laser field $E(t) = E_\omega f(t)\sin \omega t$.

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model system. States $|a\rangle$ and $|b\rangle$ in the ground state V_g are coupled by the tunneling through the potential barrier, and ground state $|a\rangle$ and excited state $|c\rangle$ are coupled by laser field $E(t)=E_\omega f(t)\sin\omega t$ with amplitude E_ω, envelope $f(t)$, and frequency ω.

In what follows, we will be interested in the time evolution of the populations in states $|a\rangle$, $|b\rangle$, and $|c\rangle$ under the influence of laser pulses. The time evolution of the wave function of the system $\psi(t)$ is treated by the time-dependent Schrödinger equation $i\hbar\partial\psi(t)/\partial t=H\psi(t)$ with Hamiltonian [18],

$$H=\hbar\kappa(|a\rangle\langle b|+|b\rangle\langle a|)+\mu E(t)(|a\rangle\langle c|+|c\rangle\langle a|)+\hbar\Delta|b\rangle\langle b|+\hbar\omega_0|c\rangle\langle c|, \quad (1)$$

where $\hbar\kappa$ is tunneling (coupling) energy determined by matrix elements between states $|b\rangle$ and $|a\rangle$, μ is the dipole transition moment, $\hbar\Delta$ is the energy difference between states $|a\rangle$ and $|b\rangle$, and $\hbar\omega_0$ is the energy separation between states $|a\rangle$ and $|c\rangle$. Expressing the eigenstate of the system as a linear combination of the isolated three states, i.e., $\psi(t)=c_a(t)|a\rangle+c_b(t)|b\rangle+c_c(t)|c\rangle$, we can obtain a set of coupled differential equations for $c_a(t)$, $c_b(t)$, and $c_c(t)$. For convenience, the energy of state $|a\rangle$ is chosen to be zero.

Substituting $\psi(t)$, H, and $E(t)$ into the Schrödinger equation, we can obtain a coupled equation as follows:

$$\frac{dc_a(t)}{dt}=-i\kappa_{ab}c_b(t)-i\Omega_\omega f(t)\sin(\omega t+\theta)c_c(t), \quad (2a)$$

$$\frac{dc_b(t)}{dt}=-i\Delta c_b(t)-i\kappa_{ba}a(t), \quad (2b)$$

$$\frac{dc_c(t)}{dt}=-i\omega_0 c_c(t)-i\Omega_\omega f(t)\sin(\omega t+\theta)c_a(t), \quad (2c)$$

where $\Omega_\omega=\mu E_\omega/\hbar$, κ_{ab} (κ_{ba}) is the tunneling coefficient between state $|a\rangle$ ($|b\rangle$) and state $|b\rangle$ ($|a\rangle$) due to the interwell tunneling. Ω_ω represents the interaction frequency, or Rabi frequency, for the transition between states $|a\rangle$ and $|c\rangle$ induced by the laser field.

We assume throughout this paper that $\kappa_{ab}=\kappa_{ba}=\kappa$, and the energy separation ω_0 is much larger than the tunneling frequency 2κ, i.e., $\omega_0=1000\kappa$. The phase of the laser field is chosen to be $\theta=0$, because the qualitative characteristics of the phenomena presented below are not changed by θ. We also assume that the laser pulses have a hyperbolic secant envelope and the pulse repetition time should be shorter than the tunneling period so as to neglect the evolution of the tunneling oscillations during the pulse excitations. The hyperbolic secant 2π laser pulse is more convenient to control the quantum tunneling rather than the rectangular pulse, because the populations in excited state $|c\rangle$ after each pulse can be neglected even for off-resonant pulses. The hyperbolic secant pulses can change only the phase of the wave packet in state $|a\rangle$ without changing its amplitude. The rectangular pulse, on the other hand, induces large populations in state $|c\rangle$ except near resonance. This prevents us from strictly investigating the interference effects between wave packets in states $|a\rangle$ and $|b\rangle$.

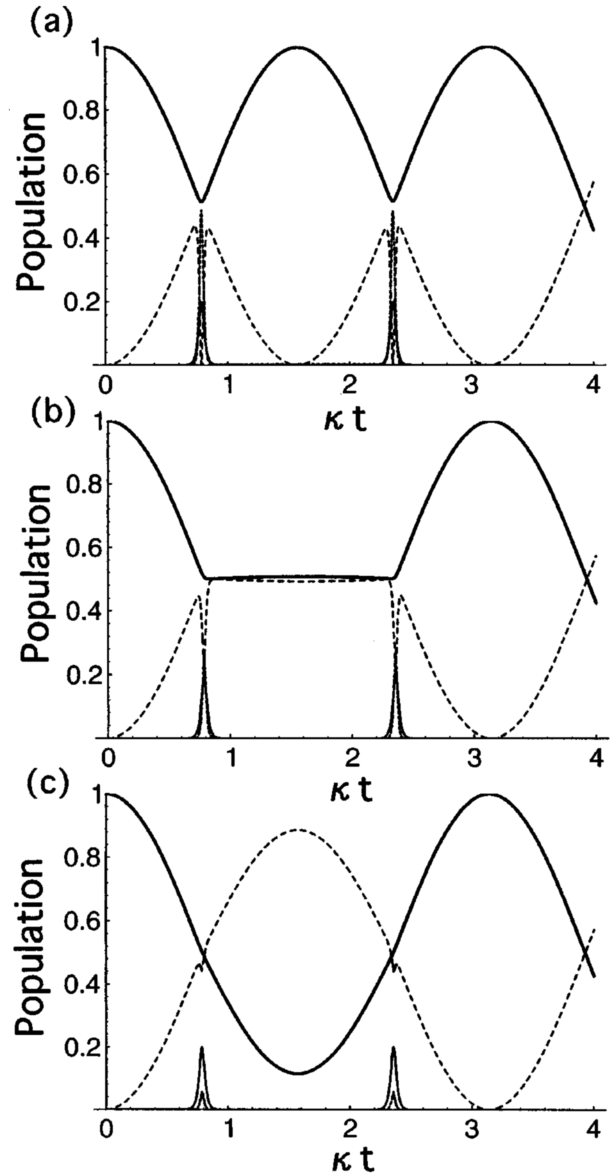

FIG. 2. Modifications of the tunneling oscillations by a sequence of two laser pulses. The populations $|c_a|^2$, $|c_b|^2$, and $|c_c|^2$ are shown by solid, dashed, and dotted lines, respectively, and the pulse envelope is shown by the thin solid line. The parameters used in the calculations are $\omega=\omega_0=1000\kappa$, $\Omega_\omega=100\kappa$, and $\tau_p=0.02$. The first and the second pulses are applied at $\tau=\tau_{\rm tun}/4$ and $\tau=3\tau_{\rm tun}/4$, respectively. (a) $\delta=0$, (b) $\delta=50\kappa$, and (c) $\delta=150\kappa$.

III. PHASE SHIFT BY PULSE EXCITATIONS

The coupled equations can be solved numerically by use of the Runge-Kutta algorithm. We introduce a dimensionless time $\tau=\kappa t$, and hence the dimensionless tunneling time is given by $\tau_{\rm tun}\equiv\kappa T_{\rm tun}=\pi$. Direct numerical integration of Eqs. (2) is performed with an initial condition $c_b(\tau=0)=1$, $c_a(\tau=0)=c_c(\tau=0)=0$.

Let us begin with a case where the tunneling states $|a\rangle$ and $|b\rangle$ have the same energy, i.e., $\Delta=0$. In Figs. 2(a)–2(c), we show typical numerical results that show the modifications of the tunneling oscillations by a sequence of two laser pulses, which have the envelope $f(\tau)=\mathrm{sech}(\tau/\tau_p)$. The populations of states $|a\rangle$, $|b\rangle$, and $|c\rangle$ are shown as a function of the dimensionless time τ together with the envelope of laser pulses. The populations $|c_a|^2$, $|c_b|^2$, and $|c_c|^2$ are shown by solid, dashed, and dotted lines, respectively, and the pulse envelope by a thin solid line. The parameters used in the calculations are $\omega=\omega_0=1000\kappa$, $\Omega_\omega=100\kappa$, and $\tau_p=0.02$. The pulse area is chosen to be 2π, i.e., $\int_{-\infty}^{+\infty}\Omega_\omega f(\tau)d\tau=2\pi$. The first and the second pulses are applied at $\tau=\tau_{\rm tun}/4$ and $\tau=3\tau_{\rm tun}/4$, respectively.

In the case of resonant excitation ($\delta\equiv\omega_0-\omega=0$), the phase of the tunneling oscillations is reversed (π out of phase) at each pulse, as shown in Fig. 2(a). The population in state $|a\rangle$ is completely transferred to state $|c\rangle$ at the peak of the laser pulse and it returns to state $|a\rangle$ just after the pulse because of 2π pulse. A drastic change in the tunneling oscillations is seen for the detuned excitation ($\delta=50\kappa$), as shown in Fig. 2(b). After the first laser pulse, the tunneling oscillations are completely suppressed and then the second laser pulse restores the free tunneling oscillations. When the detuning is further increased ($\delta=150\kappa$), the oscillations approach the free tunneling oscillations, as shown in Fig. 2(c). As will be described below, the characteristics of the tunneling oscillations for the negative detunings ($\delta<0$) are slightly different from those for the positive detunings ($\delta>0$), because the calculations are performed without the rotating-wave approximation (RWA).

A qualitative explanation of the tunneling characteristics described above can be given by a simple analytical consideration. In the case of the resonant excitation ($\delta=0$) shown in Fig. 2(a), the system has a superposition of states $|a\rangle$ and $|b\rangle$: $|\psi\rangle=(1/\sqrt{2})(i|a\rangle+|b\rangle)$ just before the first pulse excitation. Since the quantum tunneling can be neglected during the laser pulse, the time evolution of states $|a\rangle$ and $|c\rangle$ by laser excitation is simply given in the following matrix form in the RWA [19]:

$$\begin{pmatrix} c_a(\tau_+) \\ c_c(\tau_+) \end{pmatrix} = \begin{pmatrix} \cos(\Omega_\omega/2\kappa)\tau & i\sin(\Omega_\omega/2\kappa)\tau \\ i\sin(\Omega_\omega/2\kappa)\tau & \cos(\Omega_\omega/2\kappa)\tau \end{pmatrix} \begin{pmatrix} c_a(\tau_-) \\ c_c(\tau_-) \end{pmatrix}, \tag{3}$$

where we define the time-dependent Rabi frequency $\Omega_\omega(\tau)=\Omega_\omega f(\tau)$. For the initial conditions immediately before the first pulse, i.e., $c_a(\tau_-)=i(1/\sqrt{2})$ and $c_c(\tau_-)=0$ at $\tau=\tau_{\rm tun}/4$, we obtain $c_a(\tau_+)=-c_a(\tau_-)=i(1/\sqrt{2})e^{-i\pi}$ and $c_c(\tau_+)=0$ immediately after the laser excitation. This means that the phase of the wave packet in state $|a\rangle$ is altered by π by the pulse excitation. The successive time evolution of states $|a\rangle$ and $|b\rangle$ after the first pulse is then given by $|c_b(\tau)|^2=(1/2)(1+\sin\tau)$, which is π out of phase with respect to the free tunneling oscillations. The numerical result shown in Fig. 2(a) is completely consistent with this interpretation.

The population evolutions for the detuned excitations ($\delta\neq0$) are also simply interpreted by considering that each laser pulse induces the phase alteration ϕ of the wave packet in state $|a\rangle$. Keeping in mind that $c_a(\tau_+)=i(1/\sqrt{2})e^{-i\phi}$ and $c_b(\tau_+)=1/\sqrt{2}$, the time evolution of the population in state $|b\rangle$ between the two pulses is given by $|c_b(\tau)|^2=(\frac{1}{2})(1-\cos\phi\sin\tau)$ immediately after the first pulse. The complete suppression of the tunneling oscillations, i.e., $|c_a(\tau)|^2=|c_b(\tau)|^2=\frac{1}{2}$ between the pulses in Fig. 2(b), is achieved for $\phi=\pm\pi/2$. We can infer that the amount of phase change ϕ approaches zero (or 2π) as the detuning δ increases. This behavior is confirmed by use of the trajectories of the Bloch vector $u=c_a^*c_b+c_ac_b^*$, $v=-i(c_a^*c_b-c_ac_b^*)$, and $w=c_ac_a^*-c_bc_b^*$ on the unit sphere (Bloch sphere). The phase shift obtained as a function of the detuning δ is shown in Fig. 3. It is worth noting that the phase of the wave packet in state $|a\rangle$ is altered from $0\sim2\pi$ (or $-\pi\sim\pi$) by changing the detuning from infinite negative to infinite positive. As expected from above discussions, the resonant excitation exactly induces the phase shift of π and for the detuning of $\delta=\Omega_\omega/2$, the phase shift becomes $\pi/2$.

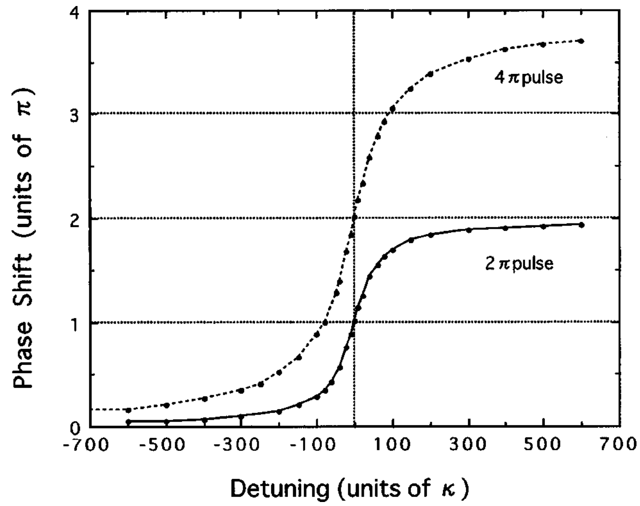

FIG. 3. Phase shift by a 2π pulse excitation as a function of the detuning δ. The phase of the wave packet in state $|a\rangle$ is altered from $0\sim2\pi$ by changing the detuning from infinite negative to infinite positive. The phase shift induced by a 4π pulse is also shown.

The phase shift of the wave packet induced by the pulse excitation plays a crucial role for the complete population transfer between nonresonant tunneling states. The phase shift induced by a 4π pulse is also shown in Fig. 3, which is nearly twice as large as that by 2π pulse but not exactly twice, except for the resonant condition ($\delta=0$). By use of a 4π pulse instead of a 2π pulse, we can obtain larger phase shifts with smaller detunings; the phase shift of $\pi\sim3\pi$ is obtained for $-80\kappa\leqslant\delta\leqslant80\kappa$.

IV. COMPLETE POPULATION TRANSFER

A. Numerical results

We can see from the numerical result shown in Fig. 2(a) that the phase of the tunneling oscillations is reversed by the resonant excitation. Therefore, we infer that when the pulse interval coincides with the half-cycle of the generalized Rabi oscillation $[\Omega=\sqrt{(2\kappa)^2+\Delta^2}]$, the generalized Rabi oscillations with small amplitudes are reversed at each pulse and the oscillations constructively accumulated as a result of quantum interference, resulting in the complete population transfer between nonresonant tunneling states $|a\rangle$ and $|b\rangle$.

This behavior is confirmed by the numerical results shown in Fig. 4(a). In the calculations, it is assumed that the energy difference between states $|a\rangle$ and $|b\rangle$ is chosen to be $\Delta=10\kappa$ and the pulse train has 20 resonant 2π pulses with the pulse envelope $f(\tau)=\Sigma_{n=1}^{20}\,\mathrm{sech}[\tau-\tau_0-(n-1)\tau_i]/\tau_P$. Other parameters are $\Omega_\omega=100\kappa$ and $\tau_i=\kappa(\pi/\Omega)$, τ_0

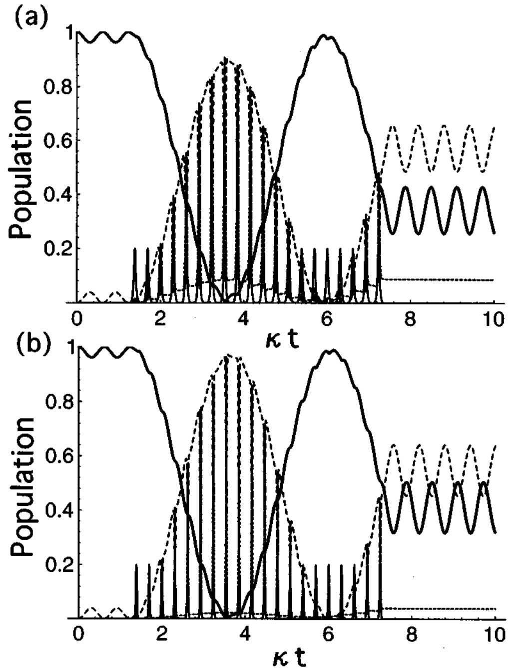

FIG. 4. Complete population transfer between nonresonant tunneling states $|a\rangle$ and $|b\rangle$ by a train of 2π resonant laser pulses with $\tau_i=\kappa\pi/\Omega$ and $\tau_0=4.5\tau_i$. The energy difference between states $|a\rangle$ and $|b\rangle$ is chosen to be $\Delta=10\kappa$. (a) $\Omega_\omega=100\kappa$, $\tau_p=0.02$, (b) $\Omega_\omega=200\kappa$, $\tau_p=0.01$.

$=4.5\tau_i$. In Fig. 4(a), we can see that nearly complete ($\sim$92%) population transfer from state $|b\rangle$ to $|a\rangle$ is achieved by eight pulses. The incomplete population transfer is mainly due to population transfer ($\sim$8%) to state $|c\rangle$. This can be reduced ($\sim$2%) by use of steep 2π pulses as shown in Fig. 4(b), where we used the pulse width (amplitude), which is half (twice) of that in Fig. 4(a). It should be noted that the complete transfer behavior hardly depends on τ_0 and this gives us flexibilities for the practical experiments.

B. Theoretical considerations

In our previous paper [20], the complete population inversion of a two-state system by a train of nonresonant optical pulses is proposed with two methods: the phase reversal method and the interval pulse method. The matrix multiplication procedures for the interval pulse method used in [20] can be utilized for the present system with some modifications.

During the intervals of the pulse excitations, the system undergoes free tunneling oscillations. For arbitrary initial amplitudes $c_a(0)$ and $c_b(0)$, the solution of the equation of motion can be written in the matrix form

$$\begin{pmatrix} c_a(\tau) \\ c_b(\tau) \end{pmatrix} = \begin{pmatrix} A & B \\ B & A^* \end{pmatrix} \begin{pmatrix} c_a(0) \\ c_b(0) \end{pmatrix} = M_I \begin{pmatrix} c_a(0) \\ c_b(0) \end{pmatrix}, \quad (4)$$

where the asterisk indicates a complex conjugate and the matrix elements are

$$A=\cos(\Omega/2\kappa)\tau - i\frac{\Delta}{\Omega}\sin(\Omega/2\kappa)\tau, \quad B=i\frac{2\kappa}{\Omega}\sin(\Omega/2\kappa)\tau, \quad (5,6)$$

where $\Omega=\sqrt{(2\kappa)^2+\Delta^2}$. If the phase shift by the pulse excitation is given by ϕ, the transfer matrix is given by

$$M_\phi = \begin{pmatrix} e^{i\phi/2} & 0 \\ 0 & e^{-i\phi/2} \end{pmatrix}. \quad (7)$$

Here, we consider a simple case that the pulse repetition time (or time interval) is equal to half of the period of the generalized Rabi frequency, i.e., $\tau_i=\kappa\pi/\Omega$, and the phase shift by the pulse excitations is $\phi=\pm\pi$. Substituting the values of τ_i and ϕ into Eqs. (4) and (7), we obtain

$$M_I = -i\frac{1}{\Omega}\begin{pmatrix} \Delta & 2\kappa \\ 2\kappa & -\Delta \end{pmatrix}, \quad M_\phi = \pm i \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}. \quad (8,9)$$

The transfer matrix M_u for one pulse and one interval can be obtained by matrix multiplication of M_I and M_ϕ, i.e., $M_u=M_\phi \cdot M_I$, and, therefore, the transfer matrix for a train of N pulses can be written in the form [20]

$$M_u^N = (\pm 1)^N \begin{pmatrix} \cos(N\varphi) & \sin(N\varphi) \\ -\sin(N\varphi) & \cos(N\varphi) \end{pmatrix}, \quad (10)$$

where

$$\varphi = \tan^{-1}\left(\frac{2\kappa}{\Delta}\right). \quad (11)$$

To obtain the complete population transfer, we let $\cos(N\varphi)=0$, and combining this with Eq. (11), we obtain

$$\frac{2\kappa}{\Delta} = \tan\frac{\pi}{2N}(2\lambda+1), \quad (12)$$

where λ is integers. At large Δ, the number of pulses N, in order to achieve the complete transfer, increases approximately linearly with Δ.

As seen in Fig. 4(a), the complete population transfer is achieved at eight pulses, which agrees quite well with $N=\pi/2\tan^{-1}(2\kappa/\Delta)\approx 7.96$. It is worth noting that the complete population transfer between nonresonant tunneling states by means of sequential laser pulses can be considered as a temporal version of the phase matching between phase-mismatched waveguides induced by spatial periodic (or grating) structures.

V. DETUNED EXCITATIONS

In this section we consider the complete population transfer between nonresonant tunneling states by a train of *detuned* ($\delta\neq 0$) laser pulses. In Fig. 5, we show the phase shift by nonresonant excitations as a function of the amplitude of the pulse Ω_ω with a constant pulse width $\tau_p=0.02$ for $\delta=\pm 50\kappa$, $\pm 100\kappa$, and $\pm 200\kappa$. The positive (negative) phase shifts are obtained for negative (positive) detunings. For the small detuning $\delta=\pm 50\kappa$, the numerical results are shown only for $2n\pi$ ($n=1\sim 4$) pulses, because the populations transferred to state $|c\rangle$ cannot be neglected for other pulse areas. The amount of the phase shift monotonically increases as the laser amplitude Ω_ω is increased. The population transfer to state $|c\rangle$ is less than 2% for all numerical results obtained for $\delta=\pm 100\kappa$ and $\pm 200\kappa$.

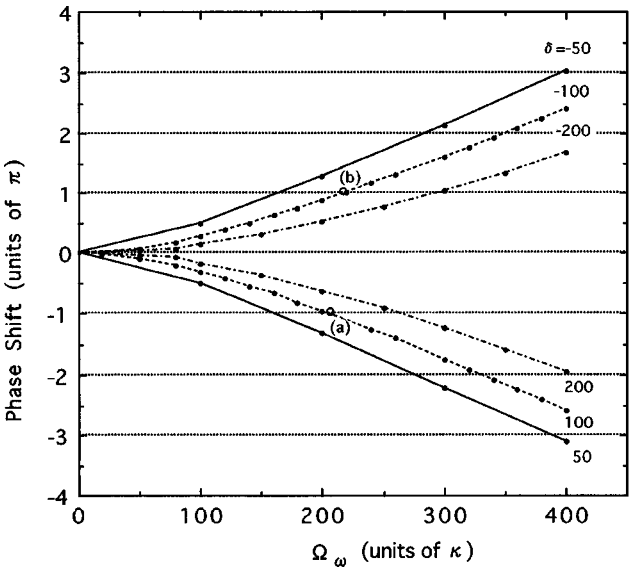

FIG. 5. Phase shift by a nonresonant pulse excitation as a function of the amplitude of the pulse Ω_ω with a constant pulse width $\tau_p=0.02$ for $\delta=\pm 50\kappa$, $\pm 100\kappa$, and $\pm 200\kappa$.

Keeping in mind that the resonant 2π pulse excitation induces the phase shift of π, resulting in the complete population transfer, we infer that the detuned pulses can also achieve the complete population transfer if they bring about π phase shift. For an example, the π phase shift for $\delta=\pm 100\kappa$ is obtained at $\Omega_\omega=205\kappa$ [indicated by (a)] and 216κ [indicated by (b), respectively]. We can show such results in Figs. 6(a) and 6(b), where the complete population transfers are obtained at the expected values of $\Omega_\omega=205\kappa$ and 216κ for $\delta=100\kappa$, -100κ with a fixed $\Delta=15\kappa$. The number of pulses to obtain the complete population transfer depends on the sign of the detuning (ten pulses for positive detuning of $\delta=100\kappa$ and 17 pulses for negative detuning of $\delta=-100\kappa$) and cannot be determined by Eq. (10). If Δ is positive (negative), the complete transfer for negative (positive) detunings requires larger numbers of pulses than for positive (negative) detunings. This may be due to the complex coherent accumulation effects of the tunneling evolutions, which cannot be interpreted by the simple analysis presented in Sec. IV B.

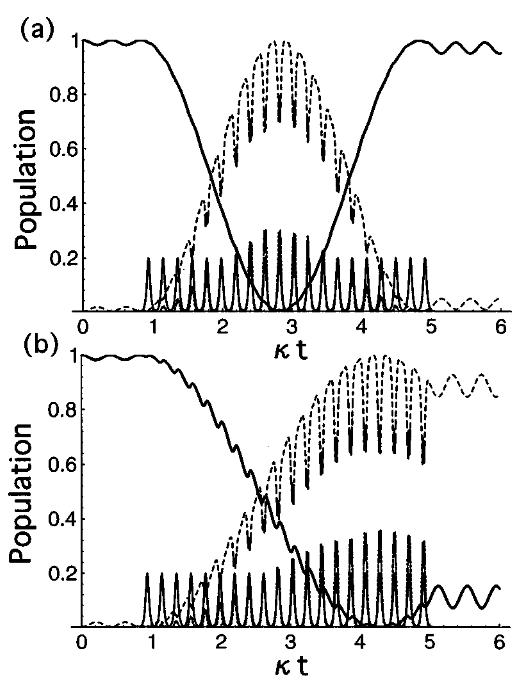

FIG. 6. Complete population transfer between nonresonant tunneling states $|a\rangle$ and $|b\rangle$ by a train of detuned laser pulses. Parameters used are $\Delta=10\kappa$, $\tau_i=\kappa\pi/\Omega$, and $\tau_0=4.5\tau_i$. (a) $\delta=100\kappa$ and $\Omega_\omega=205\kappa$, (b) $\delta=-100\kappa$ and $\Omega_\omega=216\kappa$.

Comparing Fig. 6 with Fig. 5, we can see that the population transfer to state $|c\rangle$ can be safety neglected and the population transfer between the nonresonant states is more completely achieved by detuned laser pulses.

VI. EXCITATIONS BY DARK PULSES

As shown in Fig. 7, a train of *dark* resonant pulses with an appropriate pulse interval and pulse area also induces the complete population transfer between nonresonant states. Complete population transfer between detuned state $|a\rangle$ and $|b\rangle$ by means of a train of 20 *dark* pulses. The dark pulse envelope is given by $1-\Sigma_{n=1}^{20}\,\mathrm{sech}[\tau-\tau_0-(n-1)\tau_i]/\tau_p$, $\Delta=40\kappa$, and $\Omega_\omega=100\kappa$. When the laser field is given by Ω_ω, state $|a\rangle$ splits into two dressed states, which are shifted by $\Omega_\omega/2$ from the original position [21,17,18]. Therefore, the effective detuning between states $|a\rangle$ and $|b\rangle$ is given by $\Delta'=\Omega_\omega/2-\Delta$ and the complete population transfer is expected by a train of 2π dark pulses with the pulse interval $\tau_i=\kappa\pi/\Omega'$, where $\Omega'=\sqrt{(2\kappa)^2+\Delta'^2}$. An example of such a case is shown in Fig. 7 where the complete population transfer from state $|c\rangle$ to state $|a\rangle$ is achieved at nine-dark pulses. By the dressing field, the populations of states $|a\rangle$ and $|c\rangle$ are strongly mixed and they show rapid oscillations with an equal time-averaged population. Here we assume that relaxation effects from state $|c\rangle$ are not present or can be neglected. In other words, this restricts the applicability of the results to the case of short pulses and high repetition rates compared with the relaxation times of the system.

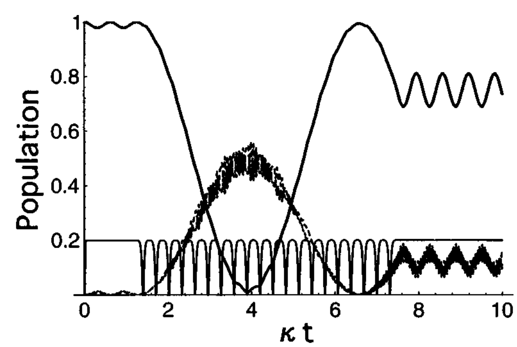

FIG. 7. Complete population transfer between nonresonant tunneling states $|a\rangle$ and $|b\rangle$ by a train of dark resonant pulses. The dark pulse envelope is given by $1-\Sigma_{n=1}^{20}\,\mathrm{sech}[\tau-\tau_0-(n-1)\tau_i]/\tau_p$. The energy difference between states $|a\rangle$ and $|b\rangle$ is chosen to be $\Delta=40\kappa$ and other parameters are $\Omega_\omega=100\kappa$, $\tau_i=\kappa\pi/\Omega$, and $\tau_0=4.5\tau_i$.

VII. CONCLUSIONS

Complete population transfer between nonresonant tunneling states in a double-well system is predicted to occur when a train of *bright* or *dark* laser pulses drives transitions to an excited state in the system. The pulse excitation induces the phase shift of the wave packet involved in the transition, which alter the successive evolutions of the wave packets of the nonresonant states as a result of quantum interference, resulting in complete tunneling.

The complete population transfer can be obtained for both resonant and detuned pulse trains with the repetition rate being half the period of the generalized Rabi oscillations. Furthermore, it is shown that a train of the dark pulses also induces the complete population transfer between nonresonant tunneling states. To control the quantum tunneling by short laser pulses, it might be useful not only to investigate the interaction between laser field and quantum structures, but also, in view of practical point, to modulate the THz oscillations in semiconductor coupled quantum wells.

To conclude, it should be pointed out that since the tunneling interaction between two well states is formally identical to a coherent coupling, the system studied in the present paper is equivalent to an atomic three-level system coupled by two coherent fields and most of the phenomena predicted in the present paper can be observed experimentally, particularly in view of the recent advances in laser technology, including effective pulse shaping [22] and producing trains of equally spaced identical pulses with repetition times of the order of 10–100 ps [23,24].

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