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Letter

Tunneling-assisted coherent population transfer and creation of coherent superposition states in triple quantum dots

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Abstract
A scheme is proposed for coherent population transfer and creation of coherent superposition states assisted by one time-dependent tunneling pulse and one time-independent tunneling pulse in triple quantum dots. Time-dependent tunneling, which is similar to the Stokes laser pulse used in traditional stimulated Raman adiabatic passage, can lead to complete population transfer from the ground state to the indirect exciton states. Time-independent tunneling can also create double dark states, resulting in the distribution of the population and arbitrary coherent superposition states. Such a scheme can also be extended to multiple quantum dots assisted by one time-dependent tunneling pulse and more time-independent tunneling pulses.

Keywords: triple quantum dots, stimulated Raman adiabatic passage, double dark states

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

1. Introduction

Coherent population transfer and the creation of coherent superposition states play a crucial role in modern quantum physics [1, 2] because of their very interesting applications to electromagnetically induced transparency (EIT) and coherent population trapping (CPT) [3, 4], quantum information and computing [5] and chemical reaction dynamics [6]. Coherent population transfer and creation of coherent superposition states can be realized by the π-pulse [1, 2] and chirped pulse [7, 8] techniques and stimulated Raman adiabatic passage (STIRAP) [1, 2, 9, 10], etc. The π-pulse method is very sensitive to small variations of pulse areas, and the chirped pulse method needs a carefully controlled chirp rate, so neither of these two methods is robust. In contrast, the STIRAP technique does not have these difficulties and has been widely used in the area of coherent population transfer and creation of coherent superposition states. The simplest model of STIRAP is a three-level Λ system [11, 12]. Later, it was found that multilevel systems can also be used to realize coherent population transfer [13–15] and coherent superposition states [16, 17]. However in these STIRAP studies two-photon resonance is required in order to eliminate the excited state from the transfer process and obtain high transfer efficiency. Besides, remarkable enhancement of population transfer can be realized by spontaneously generated coherence (SGC) [18], but such a system requires rigorous conditions of two near degenerate levels and non-orthogonal corresponding dipole matrix elements, which are rarely met in real atomic systems.

On the other hand, double quantum dots (DQDs), which can be fabricated by the self-assembled dot growth method [19], are now receiving extensive attention; they have similar properties to atomic vapors but with the advantage of flexible design and controllable interference strength. In such DQDs, an electron can be excited with a laser field in one dot, then the excited electron can tunnel to the second dot by applying an external voltage [20, 21]. Therefore, with a laser field and the tunneling, DQDs can be treated as a three-level system, in which quantum coherence and interference can be induced. Several experimental works on DQDs have been performed and analyzed. For example, direct observation of molecular resonance and quantum coupling in DQDs manipulated by electric fields has been reported [22, 23]. Such external electric fields can also control interdot electron tunneling [24] and the exciton fine structure [25]. Optical control of entanglement [26], single spin memory [27] and coherent two-electron spin [28] have all been demonstrated in DQDs.

Coherent population transfer has been studied in three-level [29] and multiple-level [30] DQD systems. In these schemes the use of a Stokes pulse or electromagnetic pulse can induce a single dark state. However, the two-photon resonance condition is needed so as to avoid populating the direct exciton state and obtain a high transfer efficiency. When a dark state is coherently coupled to another level by a coupling laser, double dark states are obtained [31, 32]. With the help of the double dark states, the optical responses of atoms can be well controlled and manipulated [33–36]. Recently, double dark states have also been investigated in triple quantum dots (TQDs) [37, 38]. Such TQDs have been fabricated in many processes [39–42] and electron and hole confinement, as well as the intermediate band of such quantum dot molecules, have been studied [43, 44]. Motivated by these works, in this paper we propose to use TQDs to realize coherent population transfer and the creation of coherent superposition states via one time-dependent tunneling pulse and one time-independent tunneling pulse. The advantages of this system are mainly in the following three points. First, population transfer from the ground state to the indirect exciton states results from the time-dependent tunneling, which is similar to the process of STIRAP but without the need for a Stokes laser field. Second, complete population transfer and creation of coherent superposition states can be obtained, even though the condition of two-photon resonance is not satisfied (this is because that time-independent tunneling can create double dark states and compensate the two-photon resonance). Third, compared with the atomic system, the energy scales and physical features of QDs can be flexible designed, not only by their composition but also by the externally applied voltages, which makes QDs an ideal system for experimental and theoretical investigations.

2. Model and equations

We consider TQDs as consisting of three vertically (in the growth direction) stacked self-assembled InAs QDs, as shown in figure 1(a). The QDs are grown with a thin tunnel barrier of GaAs/AlGaAs, thus the electrons can coherently tunnel between the three dots. By controlling the thickness of the QDs, they can have different optical transition energies and can be optically addressable with a resonant laser frequency. The QDs are incorporated into a Schottky diode, so that by adjusting the voltage bias each QD is charged with a single electron. When a gate voltage is not applied, the conduction-band electron energy levels are out of resonance, therefore, electron tunneling between the neighboring QDs is quite weak (figure 1(b)). On the contrary, when a gate voltage is applied, the conduction-band electron energy levels are in resonance and electron tunneling between the neighboring QDs becomes very strong (figure 1(c)). Hole tunneling is neglected due to the off-resonance of the valence-band energy levels in the latter situation.

With the resonant coupling of a laser field that transmits QD 1, an electron is excited in QD 1. Then, with the tunneling, the electron can be transferred to QD 2 and QD 3. Thus the TQD structure can be treated as a four-level system (figure 1(d)): the ground state $|0\rangle$, where there is no excitation in any QD, the direct exciton state $|1\rangle$, where the electron and hole are both in the first QD, the indirect exciton state $|2\rangle$, where the electron is in the second dot and the hole remains in the first dot, and the indirect exciton state $|3\rangle$, where the electron is in the third dot and the hole remains in the first dot.

In the rotating-wave approximation, the expression of $H(t)$ under the coupling of the pump and tunneling pulses can be written as

$$H = \hbar\begin{pmatrix} 0 & -\Omega_{\rm p}(t) & 0 & 0 \\ -\Omega_{\rm p}(t) & \delta_{\rm p} & -T_1(t) & 0 \\ 0 & -T_1(t) & \delta_{\rm p}-\omega_{12} & -T_2 \\ 0 & 0 & -T_2 & \delta_{\rm p}-\omega_{13} \end{pmatrix}. \tag{1}$$

Here, $\omega_{ij}=\omega_i-\omega_j$ with $\hbar\omega_i$ being the energy of the state $|i\rangle$ $(i=0,1,2,3)$. $\delta_{\rm p}=\omega_{10}-\omega_{\rm p}$ denotes the single-photon detuning ($\omega_{\rm p}$ is the frequency of the pump pulse). In TQDs the energy splittings ω_{12} and ω_{23} depend on the effective confinement potential and are much smaller than ω_{10}, and $\omega_{13}=\omega_{12}+\omega_{23}$. For simplicity, $\delta_{\rm p}$, ω_{12} and ω_{23} are time-independent. $\Omega_{\rm p}(t)$ is the Rabi frequency of the pump pulse, $T_1(t)$ is the time-dependent tunneling pulse and T_2 is the time-independent tunneling. By varying the bias voltage, both tunnelings can be switched off and on smoothly.

At any time t, the state vector can be written as

$$|\Psi(t)\rangle = a_0(t)|0\rangle + a_1(t)|1\rangle + a_2(t)|2\rangle + a_3(t)|3\rangle. \tag{2}$$

The time evolution of the probability amplitude $A(t)=[a_0(t),a_1(t),a_2(t),a_3(t)]^T$ can be described by the Schrödinger equation

$$\frac{\rm d}{{\rm d}t}A(t) = -\frac{\rm i}{\hbar}H(t)A(t) - \Lambda A(t), \tag{3}$$

where Λ is a dissipative process which contains two elements: the spontaneous decay process and the pure dephasing. Substituting equations (1) and (2) into equation (3), we can obtain the following dynamical equations for atomic probability amplitudes in the interaction picture

$${\rm i}\dot{a}_0 = -\Omega_{\rm p}a_1, \tag{4a}$$

$${\rm i}\dot{a}_1 = -\Omega_{\rm p}a_0 - T_1a_2 + (\delta_{\rm p}-{\rm i}\gamma_1)a_1, \tag{4b}$$

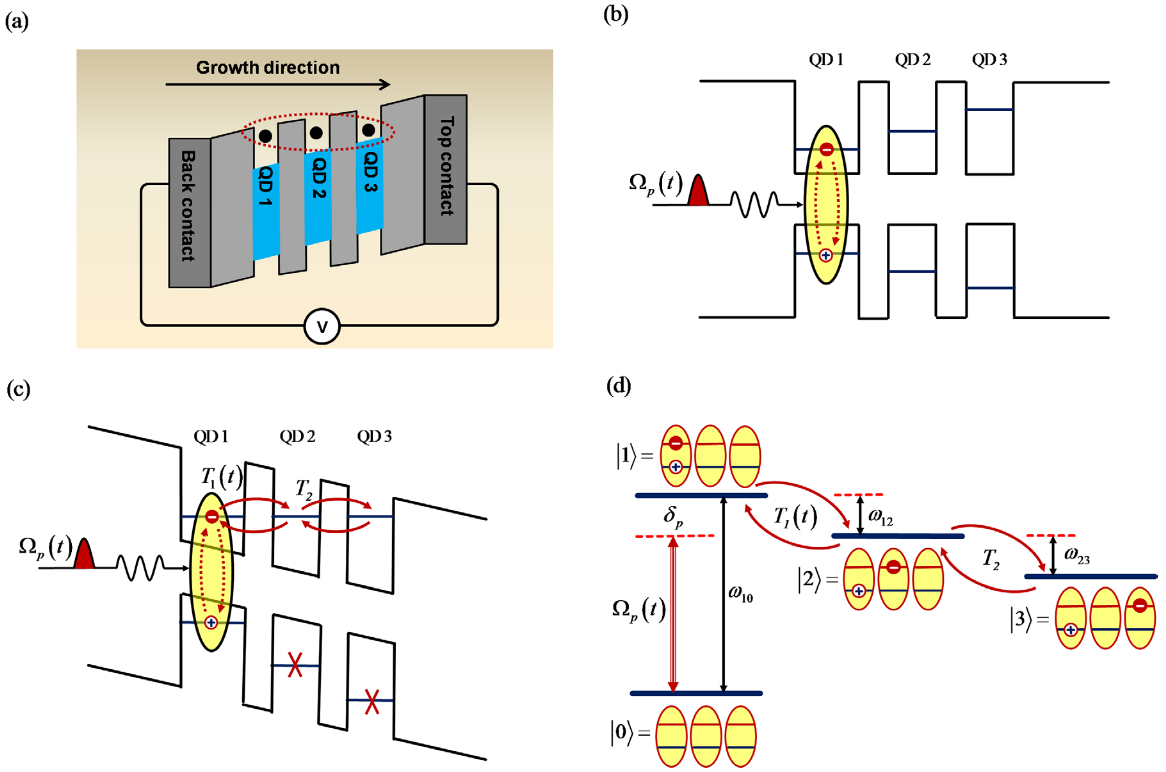

Figure 1. (a) Schematic of the setup of the TQDs. The pump field transmits QD 1. (b) Schematic of the band structure without a gate voltage. (c) Schematic of the band structure with a gate voltage. (d) Schematic of the level configuration of a TQD system.

$$\mathrm{i}\dot{a}_2 = -T_1a_1 - T_2a_3 + (\delta_\mathrm{p} - \delta_1 - \mathrm{i}\gamma_2)a_2, \quad (4c)$$

$$\mathrm{i}\dot{a}_3 = -T_2a_2 + (\delta_\mathrm{p} - \delta_1 - \delta_2 - \mathrm{i}\gamma_3)a_3. \quad (4d)$$

Here $\gamma_i = \frac{1}{2}\Gamma_{i0} + \gamma_{i0}^\mathrm{d}$ $(i = 1-3)$ is the typical effective decay rate, with Γ_{i0} being the radiative decay rate of populations from $|i\rangle \to |0\rangle$ and γ_{i0}^d being the pure dephasing rates.

It can be obtained that one of the eigenvalues of equation (1) will be zero in the condition that

$$\delta_{\mathrm{p}\pm} = \omega_{12} + \frac{\omega_{23} \pm \sqrt{\omega_{23}^2 + 4T_2^2}}{2}. \quad (5)$$

The eigenvector can be written as

$$|\Psi_\mathrm{dark}\rangle_\pm = \cos\theta_\pm|0\rangle - \sin\theta_\pm(\cos\varphi_\pm|2\rangle + \sin\varphi_\pm|3\rangle), \quad (6)$$

and mixing angles θ and φ are

$$\tan\theta_\pm = \frac{\Omega_\mathrm{p}(t)}{T_1(t)}\sqrt{1 + \frac{T_2^2}{(\delta_{\mathrm{p}\pm} - \omega_{13})^2}}, \quad (7a)$$

$$\tan\varphi_\pm = \frac{T_2}{\delta_{\mathrm{p}\pm} - \omega_{13}}. \quad (7b)$$

Here, the mixing angle θ is similar to the conventional one defined in STIRAP of a Λ atomic system. The mixing angle θ is not only related to the value of $\Omega_\mathrm{p}(t)/T_1(t)$, but it is also modified by the time-independent factor $\sqrt{1 + T_2^2/(\delta_{\mathrm{p}\pm} - \omega_{13})^2}$. On the other hand, the additional mixing angle φ is related to detuning and time-independent tunneling. In TQDs, both time-dependent tunneling and time-independent tunneling are modified by the bias voltage.

Equation (6) indicates that the state $|\Psi_\mathrm{dark}\rangle_\pm$ has no component of the state $|1\rangle$. The transition from state $|1\rangle$ to the other states is the main source of possible spontaneous emission, and $|\Psi_\mathrm{dark}\rangle_\pm$ will be immune to the possible spontaneous emission. That is, $|\Psi_\mathrm{dark}\rangle_\pm$ are dark states corresponding to $\delta_{\mathrm{p}\pm}$.

If the time-dependent tunneling pulse $T_1(t)$ precedes the pump pulse $\Omega_\mathrm{p}(t)$, and they overlap in the process, as time progresses from $-\infty$ to ∞ the value of $\Omega_\mathrm{p}(t)/T_1(t)$ rises from 0 to ∞; consequently the mixing angle θ rises from zero to $\pi/2$. As a result, the adiabatic state $|\Psi_\mathrm{dark}\rangle_\pm$ starting in the bare state $|0\rangle$ will end in the coherent superposition state $|\Psi_\mathrm{dark}\rangle_+ = -(\cos\varphi_+|2\rangle + \sin\varphi_+|3\rangle)$ and $|\Psi_\mathrm{dark}\rangle_- = -(\cos\varphi_-|2\rangle + \sin\varphi_-|3\rangle)$. Then with a suitable value for time-independent tunneling T_2 and the detunings, any value of φ can be realized. So either an arbitrary superposition of $|2\rangle$ and $|3\rangle$ or a single state can be established at the end of the process. On the other hand, almost complete population transfer from the ground state $|0\rangle$ to the indirect exciton states $|2\rangle$ and $|3\rangle$ is possible without ever populating the direct exciton state $|1\rangle$.

It is worth pointing out that without the time-independent tunneling T_2 the system will reduce to a three-level DQD, and complete population transfer from the ground state $|0\rangle$ to the indirect exciton state $|2\rangle$ can be obtained in the condition that the two-photon resonance ($\delta_\mathrm{p} = \omega_{12}$) is satisfied. In our TQD system with time-independent tunneling complete population transfer

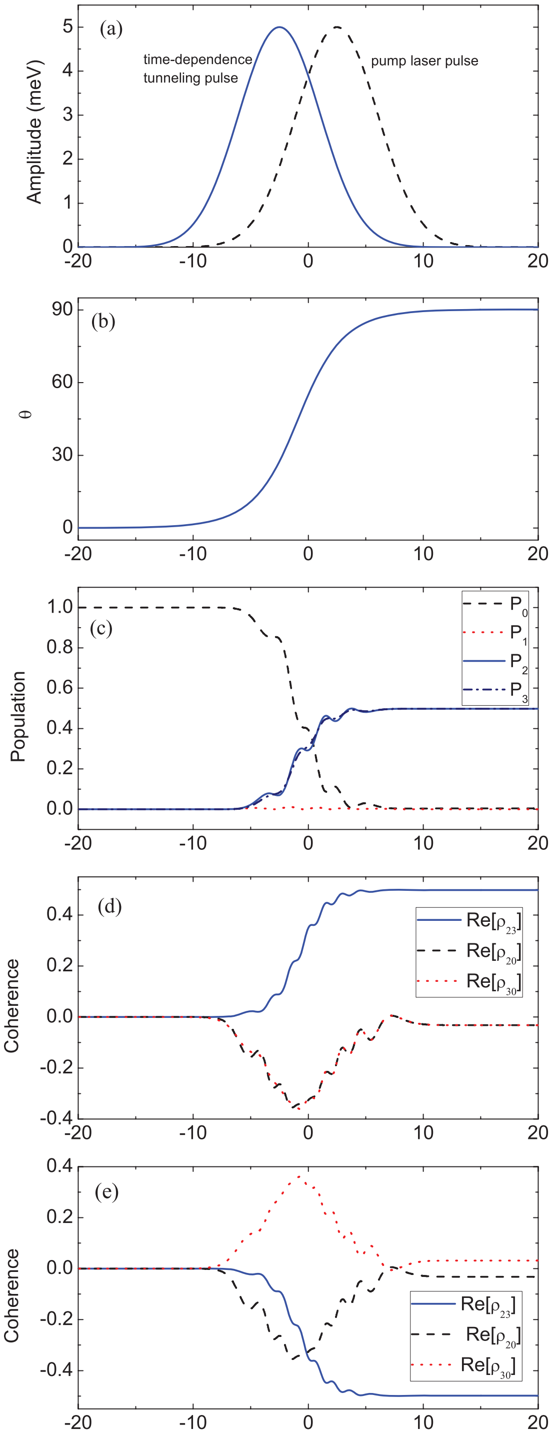

Figure 2. (a) The pump pulse and the time-dependent tunneling pulse. (b) The mixing angle θ for $\hbar\delta_p = \pm 2$ meV. (c) The time evolutions of population $P_i = |a_i|^2$ $(i = 0 - 3)$ for $\hbar\delta_p = \pm 2$ meV. (d) The real part of the coherence dynamics $\rho_{02,03,23}$ for $\hbar\delta_p = 2$ meV. (e) The real part of the coherence dynamics $\rho_{02,03,23}$ for $\hbar\delta_p = -2$ meV. Other parameters are $\hbar\Omega_p = \hbar T_l = 5$ meV, $\hbar T_2 = 2$ meV, $\hbar\omega_{12} = \hbar\omega_{23} = 0$, $T = 5$ ps, $\tau = 2.5$ ps, $\hbar\gamma_1 = 0.01$ meV and $\gamma_2 = \gamma_3 = 10^{-3}\gamma_1$.

can be obtained even when the condition of two-photon resonance is not satisfied. Furthermore, via time-independent tunneling, creation of coherent superposition states can be obtained.

3. Results and discussions

In what follows, we show a numerical simulation to illustrate the above analytic solutions. The pulses $\Omega_p(t)$ and $T_l(t)$ are shown in figure 2(a) and they are expressed as

$$\Omega_p(t) = \Omega_p \exp\left[-(t-\tau)^2/T^2\right], \tag{8a}$$

$$T_l(t) = T_l \exp\left[-(t+\tau)^2/T^2\right]. \tag{8b}$$

Here Ω_p and T_l are their peak values, respectively, T is the pulse duration and $T_l(t)$ precedes $\Omega_p(t)$ with a 2τ delay time between them. In the following numerical analysis, the initial population is assumed to be in state $|0\rangle$, that is $a_0(-\infty) = 1$, $a_{1,2,3}(-\infty) = 0$. Other realistic values of parameters are according to [37] and references therein.

In the case of $\omega_{12} = \omega_{23} = 0$, a suitable pump pulse detuning is $\delta_{p\pm} = \pm T_2$ according to equation (5). The time evolutions of mixing angles θ and φ can be obtained from equation (7), and φ has a fixed value $\varphi_\pm = \pi/4$, while the mixing angle θ is relative to the evolution of the pump pulse $\Omega_p(t)$ and the time-dependent tunneling pulse $T_l(t)$. We show the time evolutions of mixing angle θ in figure 2(b). As can be seen, the mixing angle θ arises smoothly from zero to $\pi/2$ for $\delta_{p\pm} = \pm T_2$. Therefore, the ultimate superposition state $|\Psi_{dark}\rangle_\pm$ has the form of $|\Psi_{dark}\rangle_+ = -(|2\rangle + |3\rangle)/\sqrt{2}$ for $\delta_{p+} = T_2$ and $|\Psi_{dark}\rangle_- = -(|2\rangle - |3\rangle)/\sqrt{2}$ for $\delta_{p-} = -T_2$. This means that the target states $|\Psi_{dark}\rangle_\pm$ are maximally coherent superpositions of $|2\rangle$ and $|3\rangle$, and $|\Psi_{dark}\rangle_+$ is orthogonal to $|\Psi_{dark}\rangle_-$.

Then we display the time evolution of population $P_i = |a_i|^2$ $(i = 0 - 3)$ in figure 2(c). As figure 2(c) reveals, for both pump pulse detunings ($\delta_{p\pm} = \pm T_2$), all the population is transferred from state $|0\rangle$ to states $|2\rangle$ and $|3\rangle$ with equal value at the end of the process, and state $|1\rangle$ is empty in the whole process. Next we plot the real part of the coherence dynamics ρ_{ij} for $\delta_{p+} = T_2$ and $\delta_{p-} = -T_2$ in figures 2(d) and (e), respectively. The imaginary part is not given, because under the adiabatic condition the imaginary part should be zero in the whole process. It can be seen that Re $[\rho_{23}]$ rises from zero to a maximum value 0.5 during the process, while Re $[\rho_{02}]$ and Re $[\rho_{03}]$ are nearly zero at the end of the process. This reveals that two states $|\Psi_{dark}\rangle_+$ and $|\Psi_{dark}\rangle_-$, which are maximally coherent superpositions of $|2\rangle$ and $|3\rangle$, are created.

To see this more clearly, we give the time evolution of the dressed-state eigenvalues for $\delta_{p-} = T_2$ and $\delta_{p-} = -T_2$, as shown in figures 3(a) and (b), respectively. From the figures, the eigenvalue of $|\Psi_{dark}\rangle_\pm$ stays at zero (solid line) during the whole process, therefore, $|\Psi_{dark}\rangle_\pm$ are dark states or adiabatic states. On the other hand, the non-adiabatic coupling $\dot{\theta}$, which can be obtained by equation (5*a*), is much smaller than the eigenvalue splitting of the adiabatic states and the other three non-adiabatic states $|\varepsilon_{dark} - \varepsilon_{a,b,c}|$. So non-adiabatic coupling between the eigenstates is negligible and the adiabatic condition can be satisfied in both cases [1].

From equation (6), the ratio of the final population probability of state $|2\rangle$ to state $|3\rangle$ is $R = (\cos\varphi/\sin\varphi)^2$. Therefore,

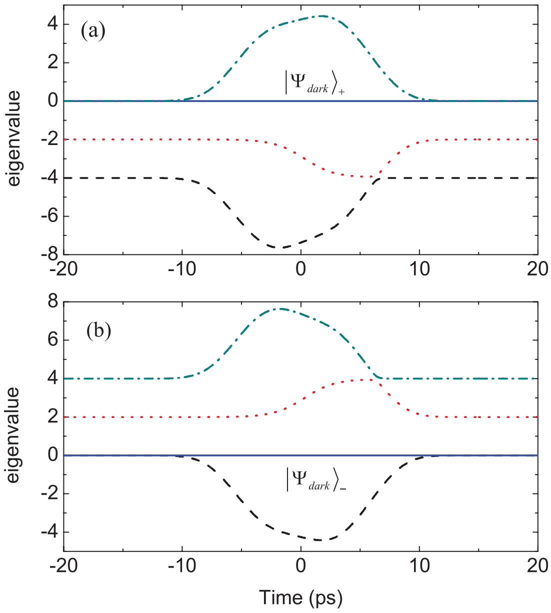

Figure 3. The dark state eigenvalues $\varepsilon_{\mathrm{dark}}$ (solid line) and the other three non-zero eigenvalues $\varepsilon_{a,b,c}$ of the dressed state picture: (a) for $\hbar\delta_{\mathrm{p}} = 2\,\mathrm{meV}$, (b) for $\hbar\delta_{\mathrm{p}} = -2\,\mathrm{meV}$. Other parameters are the same as those in figure 2.

through proper choice of parameters, the value of φ can be precisely controlled; as a consequence, arbitrary population of states $|2\rangle$ and $|3\rangle$, and an arbitrary single or superposition quantum state, can be realized. In figure 4 we give some examples. As can be seen, the population can be transferred from state $|0\rangle$ to states $|2\rangle$ and $|3\rangle$ with a designed value, and in all these cases, state $|1\rangle$ has no population.

As is well known, the STIRAP technique applied in normal Λ-type systems can transfer a population from the initial state to the final one completely, but it is essential that the condition of two-photon resonance is met—otherwise, the excited state will be populated during the transfer process and thus its decay will sharply deteriorate the transfer efficiency. However, in the present case, efficient population transfer without any population in the excited state can be realized even when the two-photon resonance is not satisfied.

In the above calculations, because the time-dependent tunneling pulse precedes the pump pulse, the mixing angle θ is $\pi/2$ at the end of the process. As a result, there is no population remaining in state $|0\rangle$, and the coherent superposition state $|\Psi_{\mathrm{dark}}\rangle_{\pm}$ does not contain the component of state $|0\rangle$. But these can be changed by using fractional STIRAP (F-STIRAP) [12]. As shown in figure 5(a), the time-dependent tunneling pulse and the pump pulse are switched off simultaneously:

$$\Omega_{\mathrm{p}}(t) = \Omega_{\mathrm{p}} \exp(-t^2/T^2), \tag{9a}$$

$$T_1(t) = \begin{cases} T_1 \exp[-(t+10)^2/T^2] & t < -10\,\mathrm{ps} \\ T_1 & -10\,\mathrm{ps} \leqslant t \leqslant 0\,\mathrm{ps} \\ T_1 \exp(-t^2/T^2) & t > 0\,\mathrm{ps}. \end{cases} \tag{9b}$$

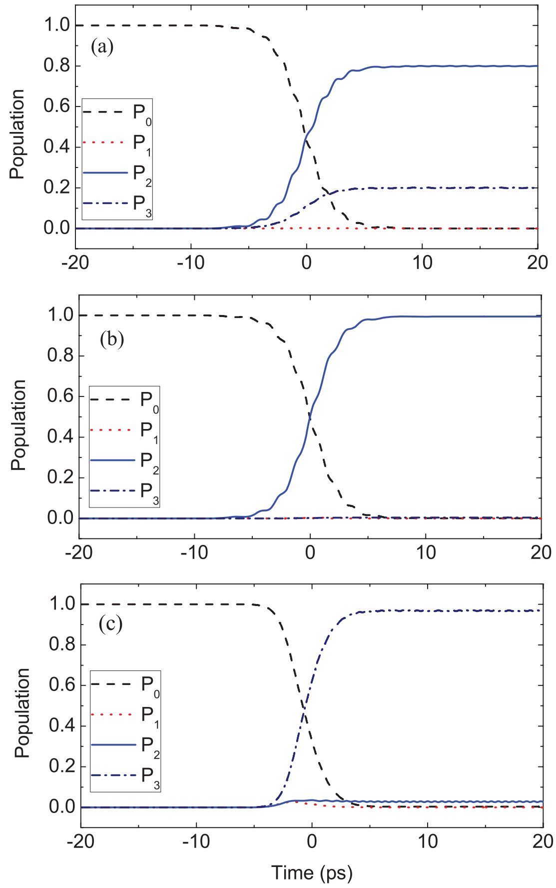

Figure 4. The time evolution of population $P_i = |a_i|^2$ $(i = 0-3)$ for different parameters: (a) $\hbar\Omega_{\mathrm{p}} = \hbar T_1 = 5\,\mathrm{meV}$, $\hbar T_2 = 2\,\mathrm{meV}$, $\hbar\delta_{\mathrm{p}} = -1\,\mathrm{meV}$, $\hbar\omega_{12} = 0$, $\hbar\omega_{23} = 3\,\mathrm{meV}$; (b) $\hbar\Omega_{\mathrm{p}} = \hbar T_1 = 5\,\mathrm{meV}$, $\hbar T_2 = 0.5\,\mathrm{meV}$, $\hbar\delta_{\mathrm{p}} = 0.04\,\mathrm{meV}$, $\hbar\omega_{12} = 0$, $\hbar\omega_{23} = 6\,\mathrm{meV}$; (c) $\hbar\Omega_{\mathrm{p}} = 2\,\mathrm{meV}$, $\hbar T_1 = 9\,\mathrm{meV}$, $\hbar T_2 = 1.7\,\mathrm{meV}$, $\hbar\delta_{\mathrm{p}} = 11\,\mathrm{meV}$, $\hbar\omega_{12} = 1\,\mathrm{meV}$, $\hbar\omega_{23} = 9.71\,\mathrm{meV}$. Other parameters are $T = 5\,\mathrm{ps}$, $\tau = 2.5\,\mathrm{ps}$, $\hbar\gamma_1 = 0.01\,\mathrm{meV}$ and $\gamma_2 = \gamma_3 = 10^{-3}\gamma_1$.

For $\omega_{12} = \omega_{23} = 0$, the detuning of the pump pulse is also $\delta_{\mathrm{p}\pm} = \pm T_2$. Then the final value of the mixing angle goes to $\theta_{\pm} = \arctan\sqrt{2}$ (figure 5(b)). Together with the fixed value of $\varphi_{\pm} = \arctan\left(\pm 1/\sqrt{2}\right)$, the adiabatic state $|\Psi_{\mathrm{dark}}\rangle_{\pm}$ starting in the bare state $|0\rangle$ will end in the coherent superposition state $|\Psi_{\mathrm{dark}}\rangle_{+} = \left(|0\rangle - |2\rangle - |3\rangle\right)/\sqrt{3}$ and $|\Psi_{\mathrm{dark}}\rangle_{-} = \left(|0\rangle - |2\rangle + |3\rangle\right)/\sqrt{3}$. This indicates that the target state is a maximally coherent superposition of $|0\rangle$, $|2\rangle$ and $|3\rangle$.

In order to clearly understand this process, we calculate the time evolution of population $P_i = |a_i|^2$ $(i = 0-3)$ and the real part of the coherence dynamics $\rho_{23,20,30}$ in figures 5(c)–(e). As figure 5(c) reveals, for both pump pulse detunings $(\delta_{\mathrm{p}\pm} = \pm T_2)$, the population is distributed equally in three states $|0\rangle$, $|2\rangle$ and $|3\rangle$ at the end of the process and state $|1\rangle$ is empty during the whole process. From figures 5(d) and (e), it can be seen that $\mathrm{Re}\,[\rho_{23}]$, $\mathrm{Re}\,[\rho_{20}]$ and $\mathrm{Re}\,[\rho_{30}]$ all increase from zero to the maximum value of 1/3 during the process for both $\delta_{\mathrm{p}+} = T_2$ and $\delta_{\mathrm{p}-} = -T_2$. This reveals that the states $|\Psi_{\mathrm{dark}}\rangle_{+}$ and $|\Psi_{\mathrm{dark}}\rangle_{-}$, which are coherent superpositions of $|0\rangle$,

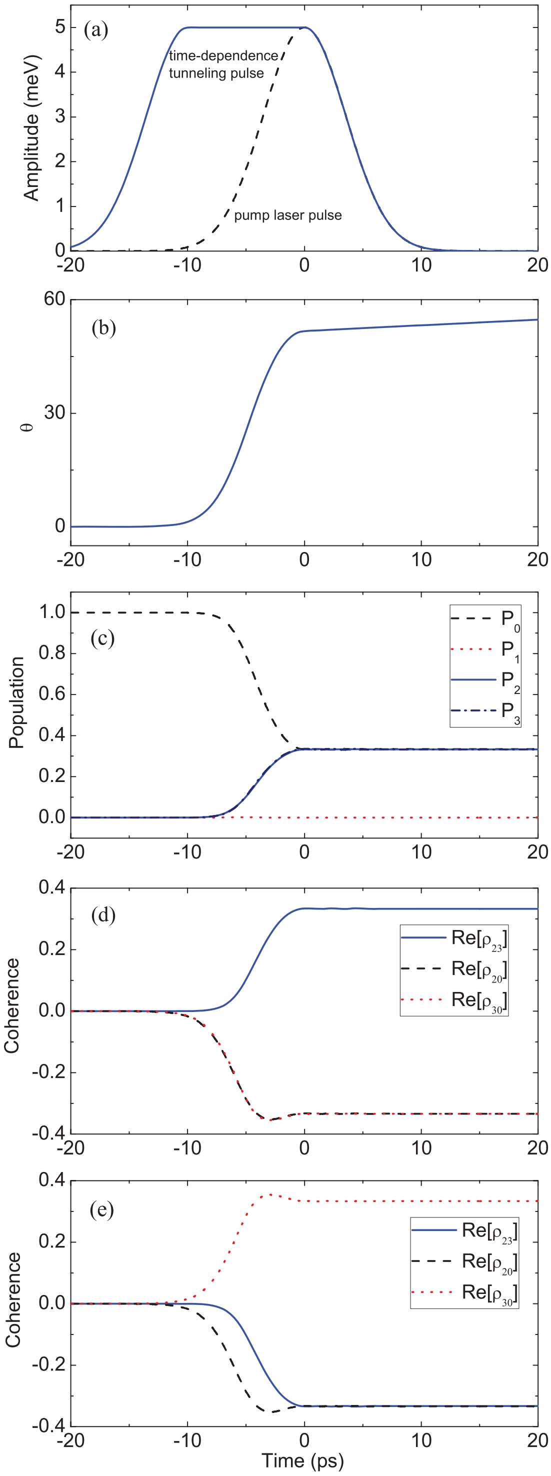

Figure 5. (a) The pump pulse and the time-dependent tunneling pulse. (b) The mixing angle θ for $\hbar\delta_p = \pm 2$ meV. (c) The time evolutions of population $P_i = |a_i|^2$ $(i = 0 - 3)$ for $\hbar\delta_p = \pm 2$ meV. (d) The real part of the coherence dynamics $\rho_{02,03,23}$ for $\hbar\delta_p = 2$ meV. (e) The real part of the coherence dynamics $\rho_{02,03,23}$ for $\hbar\delta_p = -2$ meV. Other parameters are the same as those in figure 2.

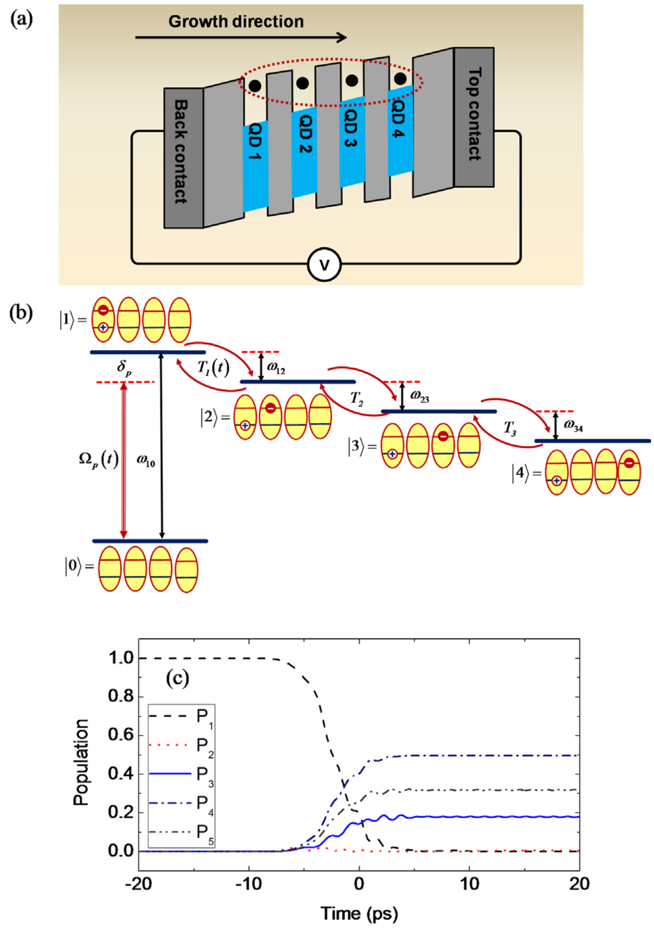

Figure 6. (a) Schematic of the setup of the four-QD system. The pump field transmits the first QD. (b) Schematic of the level configuration of the four-QD system. (c) The time evolution of population $P_i = |a_i|^2$ $(i = 0 - 4)$ of the four-QD system. The parameters are $\hbar\Omega_p = \hbar T_1 = 4$ meV, $\hbar T_2 = 3$ meV, $\hbar T_3 = 4$ meV, $\hbar\delta_p = 0$, $\hbar\omega_{12} = -5$ meV, $\hbar\omega_{23} = \hbar\omega_{34} = 0$, $T = 5$ ps, $\tau = 2.5$ ps, $\hbar\gamma_1 = 0.01$ meV and $\gamma_2 = \gamma_3 = \gamma_4 = 10^{-3}\gamma_1$.

$|2\rangle$ and $|3\rangle$, are created. Furthermore, with suitable parameters, arbitrary superposition of $|0\rangle$, $|2\rangle$ and $|3\rangle$ can be established. Also, the population can be distributed in states $|0\rangle$, $|2\rangle$ and $|3\rangle$ with arbitrary ratio, without ever populating state $|1\rangle$.

Additionally, the TQDs can be extended to multiple QDs (MQDs). Here we put forward a MQD system with four QDs, as shown in figure 6(a). In such system, the pump laser pulse transmits the first QD, and with the help of three tunneling pulses the electrons can transfer from the first QD to the others, creating a five-level system (figure 6(b)). The pump laser pulse $\Omega_p(t)$ and the tunneling pulse $T_1(t)$ are time-dependent, while the other tunneling pulses T_2 and T_3 are time-independent. δ_p is the detuning of the pump laser pulse and ω_{12}, ω_{23} and ω_{34} are the energy splittings. Then the Hamiltonian can be given

$$H' = \hbar \begin{pmatrix} 0 & \Omega_p(t) & 0 & 0 & 0 \\ \Omega_p(t) & -\delta_p & T_1(t) & 0 & 0 \\ 0 & T_1(t) & -(\delta_p - \omega_{12}) & T_2 & 0 \\ 0 & 0 & T_2 & -(\delta_p - \omega_{12} - \omega_{23}) & T_3 \\ 0 & 0 & 0 & T_3 & -(\delta_p - \omega_{12} - \omega_{23} - \omega_{34}) \end{pmatrix}. \quad (10)$$

For simplicity, ω_{23} and ω_{34} are set to zero. In the condition that $\delta_{p\pm} = \omega_{12} \pm \sqrt{T_2^2 + T_3^2}$, two dark states can be given

$$|\Psi_{\rm dark}\rangle_{\pm} = \sin\phi_{\pm}|0\rangle - \cos\phi_{\pm}\frac{1}{\sqrt{(\delta_{p\pm} - \omega_{12})^2 + T_2^2 + T_3^2}} \times \left(\Omega_p|2\rangle + (\delta_{p\pm} - \omega_{12})|3\rangle + T_3|4\rangle\right), \quad (11)$$

with

$$\tan\phi_{\pm} = \frac{T_1 T_2}{\sqrt{2}(\delta_{p\pm} - \omega_{12})\Omega_p}. \quad (12)$$

Here Ω_p and T_1 are the peak values of pump laser pulse and the tunneling pulse $T_1(t)$, respectively. By adjusting the parameters, on one hand, the population can be transferred from $|0\rangle$ to $|2\rangle$, $|3\rangle$ and $|4\rangle$ without populating $|1\rangle$. On the other hand, the desired single quantum state or superposition states can be constructed. We give one example of a numerical simulation in figure 6(c). In MQDs, population can be transferred to more states, and multi-component coherent superposition states can be created in the same manner.

We should note that the dynamics of single-electron transport in a linear array of tunnel-coupled quantum dots has been studied [45], and complete coherent transfer of electron wavepacket between the two ends of the array was achieved. But in the present paper we show that time-dependent tunneling can lead to complete population transfer from the ground state to the indirect exciton states. Furthermore, we also show that time-independent tunneling can create double dark states; therefore, it is possible to realize the distribution of the population and the arbitrary coherent superposition states.

In our theoretical simulations above, we do not include radiative decay or pure dephasing. This is because that the realistic values for radiative decay and pure dephasing are much smaller than the Rabi frequency of the pump field and the tunneling. Such dissipative processes do not break the adiabatic condition and have nearly no effect on the time evolution of population and coherence dynamics.

4. Conclusions

In conclusion, we have shown that it is possible to realize coherent population transfer and create coherent superposition states in TQDs via one time-dependent tunneling pulse and one time-independent tunneling pulse. Time-dependent tunneling can lead to complete population transfer from the ground state to the indirect exciton states, and time-independent tunneling can result in the distribution of the population and the arbitrary coherent superposition states. The function of time-dependent tunneling is similar to the Stokes laser pulse used in the traditional STIRAP, while the function of time-independent tunneling is to create double dark states and compensate the two-photon resonance. We have also extended the system to MQDs assisted by one time-dependent tunneling pulse and more than one time-independent tunneling pulse. The scheme of such semiconductor nanostructures has essential applications in novel optoelectronic devices, such as quantum phase gates and quantum computers.

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