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Optical phase control of electron transport in coupled quantum dots

Sha-sha Ke and Gao-xiang Li

Department of Physics, Huazhong Normal University, Wuhan 430079, People's Republic of China

E-mail: gaox@phy.ccnu.edu.cn

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Abstract

The properties of electron transport through an effective closed three-level structure in an asymmetric double quantum dot system are studied. It is demonstrated that the relative phase between two laser fields can strongly modulate the current through the system, which is similar to the magnetic flux controlled coherent transport in an Aharnov–Bohm interferometer. Two different transport regimes are considered, one is the regime of the chemical potential μ_R of the right lead between the ground state ε_2 and the excited state ε_3 of the right dot, the other is the regime of $\mu_R < \varepsilon_2$. In both regimes, the current can be tuned to zero by appropriately choosing the phases of the driving lasers, which can be used as an optically controlled current switch. In the regime of $\mu_R < \varepsilon_2$, the current peak approaches zero because the electron is nearly trapped into the ground state of the left quantum dot. While in the regime of $\varepsilon_2 < \mu_R < \varepsilon_3$, $I = 0$ occurs when the electron is trapped in the dark state, a superposition of the two quantum dot ground states.

1. Introduction

Ultrasmall semiconductor quantum dots (QDs), also called 'artificial atoms', are analogous to real atoms and possess many intrinsic characteristics of atomic physics [1]. Since the features of the dots are controllable, the study of external influences on QDs has become an important topic [2–7]. To date, the effect of time-varying external fields on the transport in QD devices has attracted considerable attention and many interesting phenomena ranging from photon-assisted tunneling [3, 4] to charge or spin pumping [5–7] have been reported. In experiment, the transport spectroscopy has been measured in coupled QDs under microwave field [4]. Recently, it has become desirable to transfer concepts from quantum optics to semiconductor electronics and spintronics [8]. For example, the Dicke effect [9], quantum beats [10], electromagnetically induced transparency (EIT) [11, 12], optical bistability [13], and so on, have been proposed and performed in a transport QD device by measuring its current and noise properties. Sánchez [14] has studied the resonance fluorescence and shot noise in a two-level QD, which is embedded in a phonon bath and irradiated by a time-dependent ac field. Different combinations of sub-and super-Poissonian electron and phonon Fano factors are shown in various regimes. Brandes and Renzoni [11] have proposed a transport mechanism through tunnel-coupled quantum dots based on the coherent population trapping effect (CPT). It is found that the CPT effect in such a device can be used to determine interdot dephasing rates and provide a sensitive control of a current switch. Liao *et al* [15] investigated electron tunneling through a three-level system in an asymmetric double quantum dot irradiated by an external field. Very recently Chu [12] suggested a possible scheme to realize a three-level structure in a QD system.

On the other hand, optical coherent phenomena and phase control in multilevel systems have been well studied in quantum optics in the last two decades [16–20]. Buckle *et al* [17] pointed out that in three or four level systems with closed interaction contour, changing the relative phase of the laser fields has a critical effect upon the temporary or stationary atomic populations. Phase controlled quantum interference in two-color atomic photoionization was reported [21]. Nakajima *et al* [22] presented a theory of modulating ionization by controlling the phases of incident laser fields. Up to now, the controlling quantum dynamics of matter with light by coherent phase is still a fascinating subject of broad scientific and technology [23–25]. For example, very recently, coherent phase control of resonance mediated three-photon

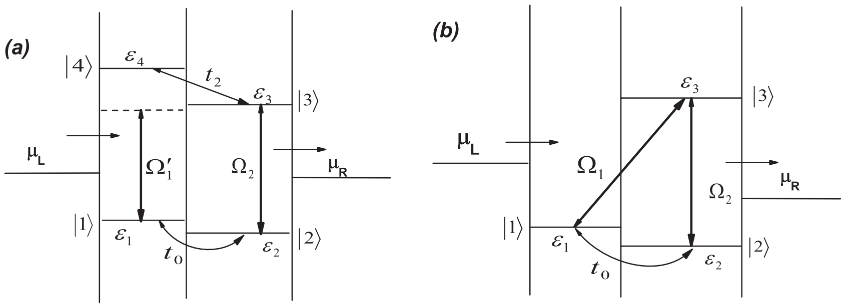

Figure 1. (a) Schematic view of two coupled two-level QD system where the two external fields Ω_1' and Ω_2 irradiate on the device and lead to the transitions $|1\rangle \leftrightarrow |4\rangle$ and $|2\rangle \leftrightarrow |3\rangle$ respectively. And the frequency ω_1 of the first laser Ω_1' is assumed to be far detuned from the transition $|1\rangle \leftrightarrow |4\rangle$. (b) An effective three-level QD system where the laser field with frequency ω_2 and Rabi frequency Ω_2 drives the transition $|2\rangle \leftrightarrow |3\rangle$, and the other laser field with frequency ω_1 and coupling constant $\Omega_1 (= \frac{\Omega_1' t_2}{\Delta_{34}})$ interacts with the transition $|1\rangle \leftrightarrow |3\rangle$.

absorption [23] and phase control of a coherent superposition of degenerate states [24, 25] were studied both theoretically and experimentally [23]. Besides, the Aharnov–Bohm (AB) effect as the most well-known and most extensively studied in solid state physics [26] has attracted considerable attention. In an AB ring interferometer, magnetic-flux-dependent current can be observed due to the two-path interference through the two arms. With the development of nanotechnologies, QDs are usually embedded in one or both arms to study the Fano effect and related physics, e.g. the problem of coherent transport through tunnel-coupled QD systems [27–29]. In the QD device, the coherent phase is usually formed and controlled by an external magnetic flux. However, from the view of experiment, it is technologically difficult to confine a strong magnetic field to such a small region of QDs, this may be an obstacle for future quantum information and computation. Therefore, it is desirable to replace the magnetic field by a more easily controllable source of the coherent phase. On the other hand, abundant research has been performed to investigate the interaction between the light and the nanostructures [30–32]. It is a natural idea to use the laser field phase to control the coherent transport in the QD-based devices.

In the present work, we propose a coherent field phase controlled transport based on an asymmetric double quantum dot system, in which an effective three-level Λ closed structure can be formed. In section 2 we propose a coupled QD device driven by two laser fields and give the theoretical framework. How to form the closed system is discussed and the effective model Hamiltonian is derived by adiabatic elimination. In section 3 we study the phase-dependent transport in such a solid-based device. We separately discuss two different transport regimes, i.e. the chemical potential in the right reservoir lies below both energy levels of the right dot and between them. Finally, section 4 gives the conclusion.

2. Theoretical framework

2.1. Adiabatic elimination and the effective Hamiltonian

Figure 1(a) shows the system and the expected level configuration. The device is composed of two coupled QDs and two normal metal leads, where the size of the left dot is much smaller than the right one. In this case the energy spacing in the left dot is larger than that of the right one ($\backsim 1/a^2$ and a the length of dot). The coupled quantum dot device considered here for fabrication can be formed by a heterostructure consisting of different semiconductor materials (GaAs/AlGaAs). Metal gates are deposited on top of a GaAs/AlGaAs heterostructure with a 2DEG about 100 nm below the surface [1]. Considering the Coulomb blockade regime, only a single electron is allowed in the two-dot system. Thus there exist five states in the effective Hilbert space of the electronic system: $|0\rangle$ (no electron in both dots), $|1\rangle$ (one electron in the ground state of the left dot), $|4\rangle$ (one electron in the excited state of the left dot), $|2\rangle$ (one electron in the ground state of the right dot), and $|3\rangle$ (one electron in the excited state of the right dot), respectively corresponding to $|0\rangle = |N_1, N_2, N_3, N_4\rangle$, $|1\rangle = |N_1 + 1, N_2, N_3, N_4\rangle$, $|4\rangle = |N_1, N_2, N_3, N_4 + 1\rangle$, $|2\rangle = |N_1, N_2 + 1, N_3, N_4\rangle$, and $|3\rangle = |N_1, N_2, N_3 + 1, N_4\rangle$. Here we only consider the transport in the sequential tunneling regime and there is no spin-rotation symmetry breaking, thus the spin degeneracy is not involved in this case. The Hamiltonian which describes the above model can be written as

$$H = H_{\mathrm{D}} + H_{\mathrm{T}} + H_{\mathrm{L}_1} + H_{\mathrm{L}_2}. \tag{1}$$

Here the Hamiltonian H_{D} representing the energy of QDs is

$$H_{\mathrm{D}} = \sum_{i=1}^{4} \varepsilon_i |i\rangle\langle i|, \tag{2}$$

where $|i\rangle\langle i|$ ($i = 1, 2, 3, 4$) are the energy operators of the QD four states. The Hamiltonian H_{T} describing the system of two dots connected to two leads can be written as

$$H_{\mathrm{T}} = \sum_{k;\alpha\in L,R} \varepsilon_k^{\alpha} c_{\alpha k}^{\dagger} c_{\alpha k} + \sum_{k;\beta=1,4} (V_{Lk} c_{Lk}^{\dagger} s_{\beta} + \mathrm{h.c.}) + \sum_{k;\beta=2,3} (V_{Rk} c_{Rk}^{\dagger} s_{\beta} + \mathrm{h.c.}), \tag{3}$$

where $c_{\alpha k}$ ($c_{\alpha k}^{\dagger}$) is the annihilation (creation) operator in lead $\alpha (\in L, R)$ with wavevector k. The first term describes the energy of the electron reservoirs. The last two terms stand for the tunneling coupling between the QDs and the reservoirs with the operators $s_{\beta} = |0\rangle\langle\beta|$ ($\beta = 1, 2, 3, 4$), and the coupling strength $V_{\alpha k}$ for α (=L and R). In the dipole and rotating-wave approximations, the Hamiltonian H_{L_1} which describes the interaction between electron and external field Ω_1' in the left dot and the tunneling between the two excited states in two QDs is expressed as

$$H_{L_1} = \frac{\Omega_1'}{2}(|4\rangle\langle 1|e^{-i(\omega_1 t+\theta_1)} + |1\rangle\langle 4|e^{i(\omega_1 t+\theta_1)}) + t_2(|4\rangle\langle 3| + |3\rangle\langle 4|). \quad (4)$$

The Hamiltonian H_{L_2} which describes the interaction between electron and external field Ω_2 in the right dot and the tunneling between the two ground states in two QDs is

$$H_{L_2} = \frac{\Omega_2}{2}(|3\rangle\langle 2|e^{-i(\omega_2 t+\theta_2)} + |2\rangle\langle 3|e^{i(\omega_2 t+\theta_2)}) + t_0(|1\rangle\langle 2| + |2\rangle\langle 1|). \quad (5)$$

The field Ω_1' with frequency ω_1 and phase θ_1 drives the transition $|1\rangle \leftrightarrow |4\rangle$ in the left QD while the field Ω_2 with frequency ω_2 is applied to drive the transition $|2\rangle \leftrightarrow |3\rangle$ in the right QD. The tunnel matrix elements t_0 and t_2 determine the strength of the tunneling process and are mainly determined by the interdot distance, whose typical value is 200–500 nm [1]. For the present scheme, the driven lasers are required to be well-separated, so that the respective laser fields couple only with their intended QD. The laser wavelength is of the order of 500–2000 nm and it is well known that the minimum spot size in the far-field regime is equal to the wavelength divided by the numerical aperture of the optical system. Advanced optical and semiconductor techniques are required to realize the scheme. Fortunately, it has been pointed out that lenses made out of flat slabs of left-handed materials (materials with negative permittivity and permeability) [33] could overcome the diffraction limit [34]. Many experiments [35] have shown that lens made of metamaterials can focus light onto an area smaller than a square wavelength. Besides, we may use relatively large quantum dots in our scheme. In the former experiments, large quantum dots with diameters of 400–700 nm were fabricated to study the Kondo effect at low temperatures [36]. Therefore, the present scheme may be realized in the near future by applying a left-handed metamaterial lens to focus light onto an area smaller than a square wavelength or fabricating larger quantum dots.

Considering that the detuning $|\omega_1-(\varepsilon_4-\varepsilon_1)|$ is sufficiently large, the electron in $|1\rangle$ absorbs a photon irradiated by the first laser and it immediately tunnels out from $|4\rangle$ to $|3\rangle$, that is, the electron is hardly able to populate the state $|4\rangle$, which can be adiabatically eliminated [37]. Therefore the two-step electron transition $|1\rangle \leftrightarrow |4\rangle \leftrightarrow |3\rangle$ can be effectively treated as a transition between $|1\rangle$ and $|3\rangle$, and described by an effective Hamiltonian

$$H_{L_1}' = \frac{\Omega_1}{2}|3\rangle\langle 1|e^{-i(\omega_1 t+\theta_1)} + \text{h.c.}, \quad (6)$$

where $\Omega_1 = \Omega_1' t_2/\Delta_{34}$ is the effective coupling constant. Here we define $\Delta_{ij} = \varepsilon_i - \varepsilon_j$ ($i \neq j = 1, 2, 3, 4$). The detailed derivation of the effective Hamiltonian H_{L_1}' is given in the appendix. Therefore the effective Hamiltonian of the whole system can be rewritten as

$$H_{\text{eff}} = H_0 + H_{I_1} + H_{I_2}. \quad (7)$$

Here, the Hamiltonian H_0 describes the electron reservoir contributions and the dot–lead coupling,

$$H_0 = \sum_{k;\alpha\in L,R} \varepsilon_k^\alpha c_{\alpha k}^\dagger c_{\alpha k} + \sum_k (V_{Lk} c_{Lk}^\dagger s_1 + V_{Rk} c_{Rk}^\dagger s_2 + V_{Rk} c_{Rk}^\dagger s_3 + \text{h.c.}). \quad (8)$$

The Hamiltonian H_{I_1} represents the energy of QDs and the tunneling between the ground states in two dots,

$$H_{I_1} = \sum_{i=1}^{3} \varepsilon_i |i\rangle\langle i| + t_0(|1\rangle\langle 2| + |2\rangle\langle 1|). \quad (9)$$

The term H_{I_2} which describes the interaction between QDs and driving lasers can be written as

$$H_{I_2} = \frac{\Omega_1}{2}\left[|3\rangle\langle 1|e^{-i(\omega_1 t+\theta_1)} + |1\rangle\langle 3|e^{i(\omega_1 t+\theta_1)}\right] + \frac{\Omega_2}{2}\left[|3\rangle\langle 2|e^{-i(\omega_2 t+\theta_2)} + |2\rangle\langle 3|e^{i(\omega_2 t+\theta_2)}\right]. \quad (10)$$

The field Ω_1 drives the transition between the ground state of the left dot ($|1\rangle$) and the excited state of the right dot ($|3\rangle$) while the field Ω_2 is applied to drive the transition between the ground state of the left dot ($|2\rangle$) and the excited state of the left dot ($|3\rangle$). The three-level loop is completed by tunneling between the two ground states $|1\rangle$ and $|2\rangle$ as sketched in figure 1(b).

2.2. Rate equation

To indicate the laser phase controlled coherent transport, we present a set of rate equations [38–40] to study the quantum transport in such a device in the sequential regime. In order to describe the transport through the QD, the master equations or 'quantum' version of the rate equations were first proposed by Nazarov [39], and later derived microscopically from the Schrödinger equation directly [40] and from the von Neumann equation and superoperators [41], respectively. It gives a good description about the sequential tunneling for weak coupling between the dot and the leads [38]. Actually at low temperature, higher-order transport, such as the cotunneling process and Kondo effect, plays an important role in the transport properties. Especially in the Kondo regime, the electron transport is through the spin–spin exchange interaction between the electrons in the central region and in the leads. To deal with these high-order processes, more sophisticated treatment should be considered, e.g. the equation of motion method [42], non-crossing approximation [43], or numerical renormalization group calculation [44]. Furthermore, the temperature taken in the calculation is of the order of 0.1Γ (Γ is the dot–lead coupling strength and is of the order of 1–10 μeV), which is much higher than the Kondo temperature. Therefore, the effect of the Kondo correlation on transport is not considered here.

Under the assumption of weak coupling between the QD and the leads, and applying the wide-band limit in the two leads, electronic transport through this system in the sequential regime can be described by the quantum rate equations for the dynamical evolution of the density matrix elements, $\rho(t)$ [38]. In the case of infinite Coulomb repulsion, the statistical expectations of the diagonal elements of the density matrix, ρ_i ($i = 0, 1, 2, 3$), give the occupation probabilities of the states of the dot. ρ_0 is the probability of finding the dot unoccupied, $\rho_{1,2,3}$ are the probabilities of finding the lower level occupied in the left dot, the lower and upper level occupied in the right dot, respectively. The off-diagonal density matrix elements describe the coherent superposition of the three levels.

Starting from the effective Hamiltonian $H_{\rm eff}$, we can obtain a set of time-dependent motion equations as follows

$$\begin{aligned}
\dot{\rho}_1 &= -(\Gamma_{\rm L}^+ + \Gamma_{\rm L}^-)\rho_1 - \Gamma_{\rm L}^+\rho_2 - \Gamma_{\rm L}^+\rho_3 + {\rm i}t_0\rho_{12} - {\rm i}t_0\rho_{12}^* \\
&\quad + {\rm i}\frac{\Omega_1}{2}{\rm e}^{-{\rm i}\theta_1}\tilde{\rho}_{13} - {\rm i}\frac{\Omega_1}{2}{\rm e}^{{\rm i}\theta_1}\tilde{\rho}_{13}^* + \Gamma_{\rm L}^+, \\
\dot{\rho}_2 &= -\Gamma_{\rm R_2}^+\rho_1 - (\Gamma_{\rm R_2}^+ + \Gamma_{\rm R_2}^-)\rho_2 - \Gamma_{\rm R_2}^+\rho_3 - {\rm i}t_0\rho_{12} \\
&\quad + {\rm i}t_0\rho_{12}^* + {\rm i}\frac{\Omega_2}{2}{\rm e}^{-{\rm i}\theta_2}\tilde{\rho}_{23} - {\rm i}\frac{\Omega_2}{2}{\rm e}^{{\rm i}\theta_2}\tilde{\rho}_{23}^* + \Gamma_{\rm R_2}^+, \\
\dot{\rho}_3 &= -\Gamma_{\rm R_3}^+\rho_1 - \Gamma_{\rm R_3}^+\rho_2 - (\Gamma_{\rm R_3}^+ + \Gamma_{\rm R_3}^-)\rho_3 - {\rm i}\frac{\Omega_1}{2}{\rm e}^{-{\rm i}\theta_1}\tilde{\rho}_{13} \\
&\quad + {\rm i}\frac{\Omega_1}{2}{\rm e}^{{\rm i}\theta_1}\tilde{\rho}_{13}^* - {\rm i}\frac{\Omega_2}{2}{\rm e}^{-{\rm i}\theta_2}\tilde{\rho}_{23} + {\rm i}\frac{\Omega_2}{2}{\rm e}^{{\rm i}\theta_2}\tilde{\rho}_{23}^* + \Gamma_{\rm R_3}^+, \\
\dot{\tilde{\rho}}_{12} &= it_0\rho_1 - {\rm i}t_0\rho_2 - \left[\tfrac{1}{2}(\Gamma_{\rm L}^- + \Gamma_{\rm R_2}^-) + {\rm i}\Delta\right]\rho_{12} \\
&\quad + {\rm i}\frac{\Omega_2}{2}{\rm e}^{-{\rm i}\theta_2}\tilde{\rho}_{13} - {\rm i}\frac{\Omega_1}{2}{\rm e}^{{\rm i}\theta_1}\tilde{\rho}_{23}^*, \\
\dot{\tilde{\rho}}_{13} &= {\rm i}\frac{\Omega_1}{2}{\rm e}^{{\rm i}\theta_1}\rho_1 - {\rm i}\frac{\Omega_1}{2}{\rm e}^{{\rm i}\theta_1}\rho_3 + {\rm i}\frac{\Omega_2}{2}{\rm e}^{{\rm i}\theta_2}\rho_{12} \\
&\quad - \left[\tfrac{1}{2}(\Gamma_{\rm L}^- + \Gamma_{R3}^-) + {\rm i}\Delta\right]\tilde{\rho}_{13} - {\rm i}t_0\tilde{\rho}_{23}, \\
\dot{\tilde{\rho}}_{23} &= {\rm i}\frac{\Omega_2}{2}{\rm e}^{{\rm i}\theta_2}\rho_2 - {\rm i}\frac{\Omega_2}{2}{\rm e}^{{\rm i}\theta_2}\rho_3 + {\rm i}\frac{\Omega_1}{2}{\rm e}^{{\rm i}\theta_1}\rho_{12}^* \\
&\quad - {\rm i}t_0\tilde{\rho}_{13} - \left[\tfrac{1}{2}(\Gamma_{\rm R_2}^- + \Gamma_{\rm R_3}^-) - {\rm i}\Delta\right]\tilde{\rho}_{23},
\end{aligned} \tag{11}$$

where $\tilde{\rho}_{3j} = \rho_{3j}{\rm e}^{{\rm i}\omega t}$ ($j = 1, 2$) are slowly varying off-diagonal matrix elements of the reduced density operator of the double dots. Here, for simplicity, we assume $\omega_1 = \omega_2 = \omega$ and the laser Ω_2 is resonant with the transition $|2\rangle \leftrightarrow |3\rangle$. But the laser Ω_1 is not resonant with the transition $|1\rangle \leftrightarrow |3\rangle$. Therefore the states $|1\rangle$ and $|2\rangle$ are non-degenerate and the energy differences between them are defined by Δ ($= \varepsilon_1 - \varepsilon_2$). The temperature-dependent tunneling rates are defined as $\Gamma_j^\pm = \sum_\alpha \Gamma_{\alpha j} f_\alpha^\pm(\epsilon_j)$, where $f_\alpha^+(\omega) = [1 + {\rm e}^{(\omega-\mu_\alpha)/k_B T}]^{-1}$ is the Fermi distribution function, $f_\alpha^-(\omega) = 1 - f_\alpha^+(\omega)$, and μ_α is the chemical potential in lead α. In the wide-band limit, the tunneling rates between the reservoirs and the dots are assumed to be energy independent: $\Gamma_\alpha = 2\pi\sum_k |V_k^\alpha|\delta(\varepsilon_\alpha - \varepsilon_k)$. Here, Γ_j^+ (Γ_j^-) describes the tunneling rate of electrons into (out from) the QD.

We can solve equation (11) at the stationary case and subsequently obtain the current (in units of e) flowing from the lead L to the QD in the sequential regime [38],

$$I = \Gamma_{\rm L}^+\rho_0 - \Gamma_{\rm L}^-\rho_1. \tag{12}$$

Due to the current conservation, the stationary current (in units of e) flowing from the QD to the lead R can be written as [38]

$$I = \Gamma_{\rm R_2}^-\rho_2 + \Gamma_{\rm R_3}^-\rho_3 - \Gamma_{\rm R_2}^+\rho_0 - \Gamma_{\rm R_3}^+\rho_0. \tag{13}$$

In the sequential tunneling regime, the transport is mainly determined by the relative height between the chemical potential and the dot levels. From equations (11)–(13), the current can be expressed by the tunneling magnitude $\Gamma_i^\pm$, which is only determined by the sign of $(\mu_\alpha - \varepsilon_i)$ at low temperature but is less related to the bias voltage $(\mu_{\rm L} - \mu_{\rm R})$. Here we consider that the chemical potential in the left lead is much higher than the energy level of state $|1\rangle$ and thus $\Gamma_{\rm L}^+ \simeq \Gamma$ (Γ is taken as the energy unit), which means the electrons in the left lead can tunnel into the left dot but the reverse process is almost prohibited (i.e. $\Gamma_{\rm L}^- \simeq 0$). Therefore, both the laser and the bias voltage can have a strong effect on the non-equilibrium transport in the present device.

3. Results and discussions

Except for the temperature, the electron transport in the present system can be modulated by the following five parameters: the energy difference $\Delta(= \varepsilon_1 - \varepsilon_2)$ between two ground states, laser intensities Ω_1, Ω_2, direct tunneling strength t_0 and the relative phases θ_i ($i = 1, 2$). We mainly focus on the modulation of the phase θ_i for different sets of $(\Delta, t_0, \Omega_1, \Omega_2)$ and discuss two different regions with respect to the location of the chemical potential $\mu_{\rm R}$, i.e. $\mu_{\rm R} < \varepsilon_2$ and $\varepsilon_2 < \mu_{\rm R} < \varepsilon_3$, respectively. For simplicity we assume that the frequencies of two coherent fields are equal and the dot-lead tunneling meets $\Gamma_{\rm L} = \Gamma_{\rm R_2} = \Gamma_{\rm R_3} = \Gamma$ and Γ is taken as the energy unit in the following discussion.

3.1. $\mu_{\rm R} < \varepsilon_2$

We first consider the case of $\mu_{\rm R} < \varepsilon_2$, where the tunneling rates $\Gamma_{\rm L}^+ = \Gamma_{\rm R_2}^- = \Gamma_{\rm R_3}^- \approx \Gamma$ and $\Gamma_{\rm L}^- = \Gamma_{\rm R_2}^+ = \Gamma_{\rm R_3}^+ \approx 0$ at low temperature. In this case, electrons can tunnel out from both the excited and ground state of the right QD. According to equations (12) and (13), the expression of the stationary current can be simplified as

$$I = \Gamma\rho_0 = \Gamma(\rho_3 + \rho_2). \tag{14}$$

Figure 2 shows the results for the current I as a function of the energy difference Δ for different Rabi frequencies Ω_1; here we have chosen $\Omega_2 = 20$, $t_0 = 0.5$, and $\theta_1 = \theta_2 = 0$. It can be seen that the current exhibits a two-peak structure and the two peaks are located at $\Delta = \pm\Omega_2/2$. With the increase of Ω_1 from 0.0 to 1.6, the position of the two peaks holds the line, while the structure of the two-peak changes from symmetric to asymmetric type. In particular, the left peak disappears when Ω_1 is equal to $2t_0$. In order to make these phenomena clear, we discuss the system in a convenient basis for the Hilbert space constructed by the unperturbed state $|1\rangle$ and dressed states $|\pm\rangle$. The dressed states $|\pm\rangle$, originating from the interaction of the right QD with the strong coherent laser field Ω_2, can be easily derived from the eigenvalue equation, $\frac{\Omega_2}{2}({\rm e}^{-{\rm i}\theta_2}|3\rangle\langle 2| + {\rm e}^{{\rm i}\theta_2}|2\rangle\langle 3|)|\pm\rangle = \lambda_\pm|\pm\rangle$, with eigenvalues $\lambda_\pm = \pm\Omega_2/2$ and eigenstates $|\pm\rangle = \frac{1}{\sqrt{2}}(|3\rangle \pm {\rm e}^{{\rm i}\theta_2}|2\rangle)$. In the basis of $|1\rangle$ and $|\pm\rangle$, the Hamiltonian $H_{I_1} + H_{I_2}$ in the rotating frame of the laser frequency ω_2 becomes

$$\begin{aligned}
H_{I_1} + H_{I_2} &= \Delta|1\rangle\langle 1| + \frac{\Omega_2}{2}(|+\rangle\langle +| - |-\rangle\langle -|) \\
&\quad + (g_+|1\rangle\langle +| + g_-|1\rangle\langle -| + {\rm h.c.}),
\end{aligned} \tag{15}$$

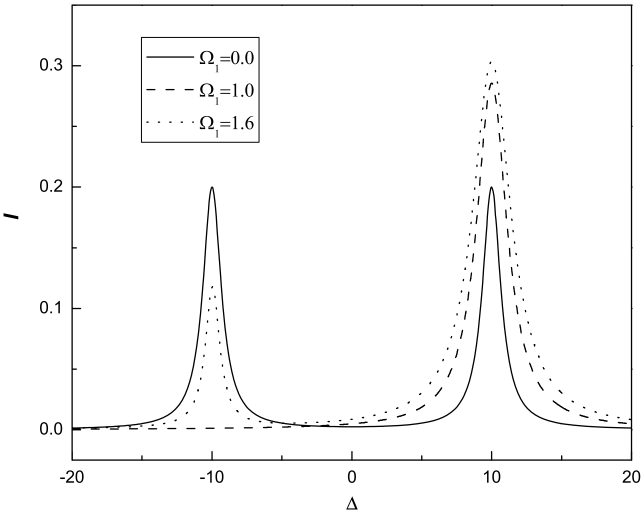

Figure 2. Current I as a function of energy difference Δ at $\theta_1 = \theta_2 = 0$, $\Omega_2 = 20$, and $t_0 = 0.5$ for different Rabi frequencies $\Omega_1 = 0.0$ (solid line), $\Omega_1 = 1.0$ (dash–dotted line), and $\Omega_1 = 1.6$ (dotted line). All coefficients are scaled with Γ.

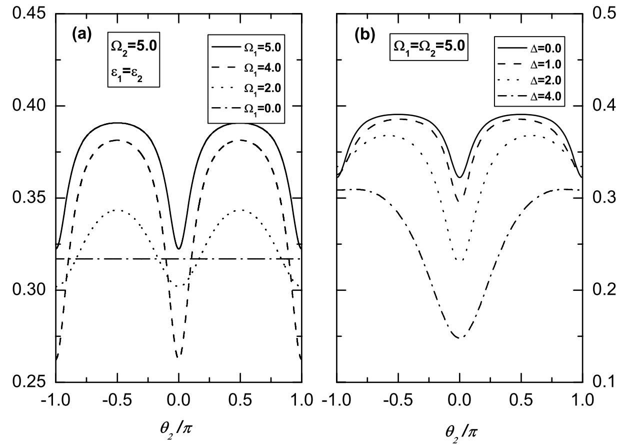

Figure 3. Current I as a function of phase θ_2 at $\theta_1 = 0.0$, $t_0 = 2.0$, and (a) $\Omega_2 = 5.0$, $\Delta = 0$ for different Rabi frequencies Ω_1; (b) $\Omega_1 = \Omega_2 = 5.0$ for different energy differences Δ. All coefficients are scaled with Γ.

here $g_+ = \frac{1}{2\sqrt{2}}(\Omega_1 e^{i\theta_1} + 2t_0 e^{i\theta_2})$, and $g_- = \frac{1}{2\sqrt{2}}(\Omega_1 e^{i\theta_1} - 2t_0 e^{i\theta_2})$. In the interaction picture, the above Hamiltonian reads as

$$\tilde{H}_I = g_+|1\rangle\langle+|e^{i(\Delta-\frac{\Omega_2}{2})t} + g_-|1\rangle\langle-|e^{i(\Delta+\frac{\Omega_2}{2})t} + \text{h.c.} \quad (16)$$

If we assume the laser field Ω_2 is a strong coherent laser field, i.e. $\Omega_2 \gg |g_\pm|$, then when we tune the energy difference Δ to be equal to $\Omega_2/2$, the transition $|1\rangle \leftrightarrow |+\rangle$ is resonant but $|1\rangle \leftrightarrow |-\rangle$ is far-off-resonant. Therefore we can ignore the highly oscillating terms $g_-|1\rangle\langle-|e^{i(\Delta+\frac{\Omega_2}{2})t} + h.c.$ in equation (16) within the effective rotating-wave approximation. That is to say, in the case of $\Delta = \Omega_2/2$, the electron can only transmit from state $|1\rangle$ to state $|+\rangle$ and consequently tunnel out from state $|+\rangle$ and then a peak appears at $\Delta = \Omega_2/2$. However, there are two transition channels from states $|1\rangle$ to $|+\rangle$, one is the electron transition driven by the laser field Ω_1 with the strength of $\Omega_1 e^{i\theta_1}/2\sqrt{2}$, the other is the single electron tunneling determined by the term of $t_0 e^{i\theta_2}/\sqrt{2}$. These two channels interfere with each other and the interference can be constructive or destructive depending on the relative phases θ_1 and θ_2. On the contrary, if the energy difference Δ is tuned to $-\Omega_2/2$, the main channels of the electron transport are from states $|1\rangle$ to $|-\rangle$ and subsequently tunnel out from state $|-\rangle$ because $|1\rangle \leftrightarrow |-\rangle$ is resonant and $|1\rangle \to |+\rangle$ is far-off-resonant. Therefore a peak appears at $\Delta = -\Omega_2/2$. There also exist two transition channels from states $|1\rangle$ and $|-\rangle$ and these two channels interfere with each other. One is the electron transition driven by the laser field Ω_1 with the strength of $\Omega_1 e^{i\theta_1}/2\sqrt{2}$, the other is the single electron tunneling determined by the term of $t_0 e^{i(\theta_2+\pi)}/\sqrt{2}$. As shown in figure 2, for $\Omega_1 = 0$ two symmetric peaks appear in the current spectrum with the driving coherent field Ω_2 as depicted in [15]. There only exists one transition channel, i.e, the single electron tunneling between the states $|\pm\rangle$ and $|1\rangle$ without the coherent field Ω_1 and the corresponding strength are equal, i.e. $|g_+| = |g_-| = t_0/\sqrt{2}$, so the two symmetric peaks appear around $\Delta = \pm\Omega_2/2$. However, if $\Omega_1 \neq 0$, we can get $g_\pm = \frac{1}{2\sqrt{2}}(\Omega_1 \pm 2t_0)$ for $\theta_1 = \theta_2 = 0$, which indicates that the constructive interference between the two transition channels between the states $|1\rangle$ and $|+\rangle$, and destructive interference among the other two transition channels between the states $|1\rangle$ and $|-\rangle$ happen. Therefore for $\theta_1 = \theta_2 = 0$, $t_0 = 0.5$, $\Omega_1 = 1.6$, $|g_+| = 0.919$, and $|g_-| = 0.212$, the currents exhibit asymmetric peaks and the height of the right peak is higher than the left one as shown by a dotted line in figure 2. Furthermore, the left peak disappears corresponding to the point where the current is almost equal to zero when $\Omega_1 = 2t_0 = 1$ as shown by a dash–dotted line in figure 2. The reason is that there exists complete destructive interference between the two channels between states $|-\rangle$ and $|1\rangle$ so that $g_- = 0$. In the case of $I = 0$, the electron is trapped in the state $|1\rangle$. The trapped electron cannot tunnel out of the ground state of the left QD. Moreover, this state is protected from incoming electrons because of the Coulomb blockade. In the double dot case discussed here, the peaks of current can be nearly tuned to zero because of a complete destructive interference determined by relative phases θ_1 and θ_2. This effect may serve as an optical phase controlled current switch [11].

We also notice that the current curves depend on the phase θ_2 for different laser field strength Ω_1 and energy difference Δ in the case of $\theta_1 = 0$ respectively as shown in figures 3(a) and (b). The current oscillates with the phase θ_2, which demonstrates that the laser field phase can act as a magnetic-flux-type AB phase. It should be noticed that the laser phase is used here to control the coherent transport, which seems to play the same role as the magnetic flux in an AB ring. However, it is quite different from the AB interferometer. Here the interaction between the electrons and the laser is time-dependent, while in an AB ring the non-equilibrium is only due to the bias voltage. But they have the same physical reason why both of them exhibit phase mediated quantum interference. With the decrease of Ω_1 from $\Omega_1 = \Omega_2$, the current modulation $\Delta I (= I_{max} - I_{min})$ increases. When Ω_1

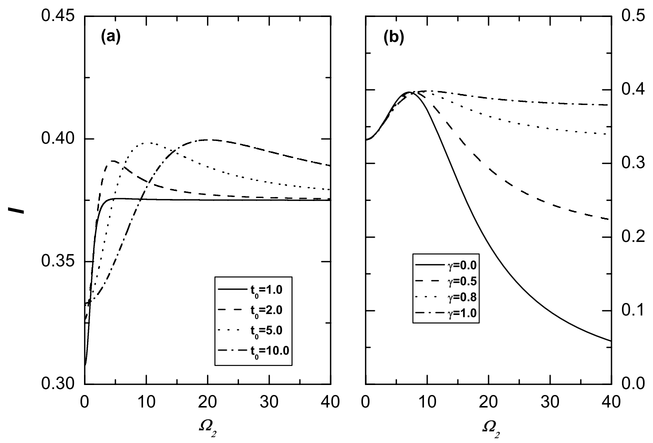

Figure 4. Dependence of the current I on Rabi frequency Ω_2 at $\theta_1 = 0.0$, $\theta_2 = \frac{\pi}{2}$, $\Delta = 0.0$, (a) $\Omega_1 = \Omega_2$ for different tunneling strength t_0; (b) $t_0 = 0.5$, $\Omega_1 = r\Omega_2$ for different r. All coefficients are scaled with Γ.

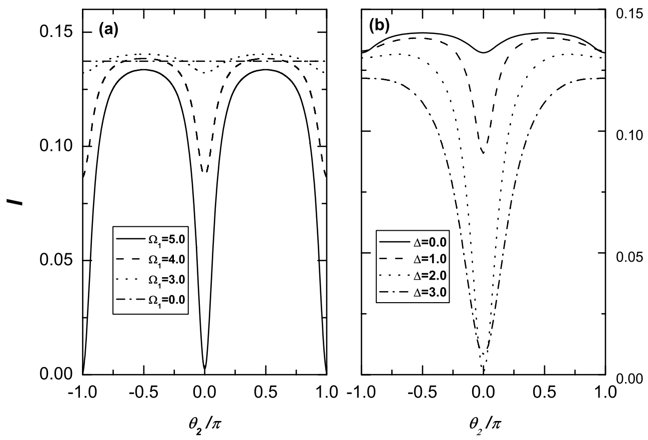

Figure 5. Dependence of current I on change of phase θ_2 (a) for different Rabi frequencies Ω_1 with $\theta_1 = 0$, $\Delta = 0$, $t_0 = 2.0$, and $\Omega_2 = 5.0$; (b) for different energy differences Δ with $\theta_1 = 0$, $t_0 = 2.0$, $\Omega_1 = 3.0$, and $\Omega_2 = 5.0$. All coefficients are scaled with Γ.

decreases more, ΔI begins to decrease and reach zero at $\Omega_1 = 0$ as shown in figure 3(a). That is to say, the phase θ_2 cannot control the current without the field Ω_1. This is because the closed loop is broken and we find the current expression as $4t_0^2(2t_0^2 + \Gamma^2 + \Omega_2^2)\Gamma[24t_0^4 + (\Gamma^2 + \Omega_2^2)^2 + 2t_0^2(7\Gamma^2 + 3\Omega_2^2)]^{-1}$ for $\Omega_1 = 0$, which is independent of the relative phase. In the case of $\Omega_1 = \Omega_2$, the oscillation period changes from π to 2π when the energy difference Δ changes from $\Delta = 0$ to $\Delta \neq 0$. Furthermore the minimum value of I decreases with the increase of Δ as shown in figure 3(b).

Let us now turn our attention to the effect of the tunneling strength t_0 on the current. Figure 4 shows the dependence of the current on the Rabi frequency Ω_2 for various coupling strengths t_0 and laser asymmetry $r(= \Omega_1/\Omega_2)$. In the strong field limit, I approaches a certain value $3/8\Gamma$ for $\Omega_1 = \Omega_2$ and the resonance peak still appears near $\Omega_2 = 2t_0$. Furthermore, as shown in figure 4(b), r can modulate the current behavior in quite a different way in the strong field limit. In this case, the current can be simplified as

$$I \simeq \left[\frac{4r^2 + 4r^4 + r^6}{4 + 8r^2 + 9r^4 + 3r^6} + \frac{1}{4 + 8r^2 + 9r^4 + 3r^6} \times \left(\frac{t_0}{\Omega_2}\right)^2\right]\Gamma, \quad (17)$$

from which one can prove that I decreases with the increase of Ω_2 and approaches 0.375Γ for $r = 1$ in the limit of $\Omega_2 \gg \Gamma$, as shown by a dot-dashed line in figure 4(b).

3.2. $\varepsilon_2 < \mu_R < \varepsilon_3$

Next, we will discuss current transport in the regime of $\varepsilon_2 < \mu_R < \varepsilon_3$ as shown in figure 1(b) where the tunneling rates $\Gamma_L^+ = \Gamma_{R_2}^+ = \Gamma_{R_3}^- \approx \Gamma$ and $\Gamma_L^- = \Gamma_{R_2}^- = \Gamma_{R_3}^+ \approx 0$ at low temperature. According to equations (12) and (13), the expression current can be simplified as

$$I = \Gamma\rho_0 = \Gamma(\rho_3 - \rho_0). \quad (18)$$

In this case, the electron can tunnel out to the right lead only from the excited state $|3\rangle$. From equations (11) and (18), the concise expressions of steady state current are obtained as

$$I = \frac{\Gamma^4\{(\Omega_2^2+\Omega_1^2)^2(\frac{\Delta^2}{4}+t_0^2) - [2t_0\Omega_1\Omega_2\cos\theta_2 + \frac{\Delta}{2}(\Omega_2^2-\Omega_1^2)]^2\}}{f(\Omega_1, \Omega_2, \Gamma, t_0, \theta_2, \Delta)}, \quad (19)$$

where $f(\Omega_1, \Omega_2, \Gamma, t_0, \theta_2, \Delta)$ is the function of the parameters Ω_1, Ω_2, Γ, t_0, θ_2, and Δ and we assume $\theta_1 = 0$ for simplicity. Figure 5(a) shows the current I versus phase θ_2 for different Rabi frequencies Ω_1, where $\theta_1 = 0$, $\Delta = 0$, $t_0 = 2.0$, and $\Omega_2 = 5.0$. The steady state current exhibits oscillations with the phase θ_2 and reaches its minimum value at the point $\theta_2 = k\pi$ $(k = 0, \pm 1, \pm 2 \ldots)$. With the increase of Ω_1 from zero to Ω_2, the minimum value of current decreases gradually until it goes down to zero when $\Omega_1 = \Omega_2 = \Omega$. $I = 0$ occurs when the electron is trapped into a coherent superposition of the two ground states $|1\rangle$ and $|2\rangle$ that decoupled from the light. It is convenient to work in an 'uncoupled state' representation: $|B\rangle = \frac{1}{\sqrt{2}}(e^{i\theta_1}|1\rangle + e^{i\theta_2}|2\rangle)$, $|D\rangle = \frac{1}{\sqrt{2}}(e^{-i\theta_2}|1\rangle - e^{-i\theta_1}|2\rangle)$, and $|3\rangle$. In this basis, the Hamiltonian $H_{I_1} + H_{I_2}$ in the rotating frame of the laser frequency ω_2 becomes

$$\begin{aligned} H_{I_1} + H_{I_2} = &\frac{\Omega}{\sqrt{2}}(|B\rangle\langle 3| + |3\rangle\langle B|) \\ &+ t_0\cos(\theta_1 - \theta_2)(|B\rangle\langle B| - |D\rangle\langle D|) \\ &+ \frac{t_0}{2}(e^{-2i\theta_1} - e^{-2i\theta_2})|B\rangle\langle D| + \frac{t_0}{2}(e^{2i\theta_1} - e^{2i\theta_2})|D\rangle\langle B|. \end{aligned} \quad (20)$$

It is now apparent that if $\theta_2 = \theta_1 + k\pi (k = 0, \pm 1, \pm 2\cdots)$, the last two terms of the above Hamiltonian are equal to zero. Therefore in the condition of $\theta_2 = \theta_1 + k\pi (k = 0, \pm 1, \pm 2\cdots)$, only $|B\rangle$ couples to the light and excitation of the electron from $|B\rangle$ to the excited state $|3\rangle$ with a subsequent decay into $|B\rangle$ and $|D\rangle$ gradually pumps all the population into $|D\rangle$. The reason is that the electron in the state $|D\rangle$ is decoupled from the light and cannot be excited again. That is to say, in the long time limit, we will find that the electron can be trapped in two different superposition states, the so-called dark states, depending on the relative phases θ_1 and θ_2. One is the state of $\frac{1}{\sqrt{2}}(|1\rangle - |2\rangle)$ for $\theta_1 = \theta_2 = 0$, the other is the state of $\frac{1}{\sqrt{2}}(|1\rangle + |2\rangle)$ for $\theta_1 = 0$ and $\theta_2 = \pm\pi$ as shown by a solid line in figure 5(a). It has been found that atomic dark states are extraordinary stable against a number of perturbations. In the case of quantum dots, a trapped electron cannot tunnel out of the coherent superposition of two ground states. Moreover, this superposition is protected from incoming electrons because of a Coulomb blockade. In this case, the effect as a sudden appearance of $I = 0$ can be used as an optically controlled current switch. Brandes *et al* [11] have pointed out that the dark resonance effect(i.e. the effect of $I = 0$) appears for $\delta_R = 0$ and the Raman detuning $\delta_R = 0$ is defined as $\delta_R = \varepsilon_2 - \varepsilon_1$. Different from [11], in our system the current can be tuned to zero by choosing appropriate relative phases θ_1 and θ_2 as discussed above. Furthermore, we must point out that the physical reason for the sudden appearance of $I = 0$ here is completely different from the reason for the current peak disappearing and the corresponding current is nearly equal to zero, as discussed in section 3.1. For $\mu_R < \epsilon_2 < \epsilon_3$, due to the complete destructive interference between the two transition channels, the electron is trapped into the ground state of the left QD and the current peak disappears. In contrast, for $\epsilon_2 < \mu_R < \epsilon_3$, the appearance of $I = 0$ is because the electron is trapped into a dark state that decoupled from the light.

We also discuss I as a function of phase θ_2 for various energy differences Δ as shown in figure 5(b). Here we have chosen $\theta_1 = 0$, $t_0 = 2.0$, $\Omega_1 = 3.0$, and $\Omega_2 = 5.0$. For $\Omega_1 \neq \Omega_2$, the energy difference Δ, which can be controlled by the external gate voltage, can be used to change the modulation of the phase on the current. The current can reach its minimum value at the point of $\theta_2 = 0$ and the minimum value of I can be expressed as

$$I_{\min} = 4t_0^2(\Omega_1^2 - \Omega_2^2)^2\Gamma\{f(t_0, \Omega_1, \Omega_2) + 256\Delta t_0^3\Omega_1\Omega_2 + 32\Delta t_0(\Omega_1 + \Omega_2)(2\Delta t_0 - \Omega_1\Omega_2)\}^{-1}. \quad (21)$$

From the expression of $I_{\min}$, the value of $I_{\min}$ decreases with the increase of energy difference Δ as shown in figure 5(b).

In figure 6 we illustrate the curves of the Rabi-frequency-dependent current. For $\theta_1 = 0$, $\theta_2 = \pi$, $\Delta = 0$, and $\Omega_1 = \Omega_2 = \Omega$, the current can be written as

$$I = \frac{2t_0^2\Omega^2\Gamma}{(\Omega^2 + 3t_0^2)^2 + 7t_0^4 + 4\Gamma^2 t_0^2}. \quad (22)$$

It is found that the current increases with the increase of Ω in the regime of $\Omega < \sqrt{2t_0\sqrt{4t_0^2 + \Gamma^2}}$ and reaches its maximum value $\frac{4t_0\sqrt{4t_0^2+\Gamma^2}\Gamma}{(2\sqrt{4t_0^2+\Gamma^2}+3t_0)^2+7t_0^2+4\Gamma^2}$ at the point of $\Omega = \sqrt{2t_0\sqrt{4t_0^2 + \Gamma^2}}$. After reaching its maximum, I decreases sharply as shown in figure 6(a). Figure 6(b) illustrates the dependence of the current on the Rabi frequency Ω for the different asymmetry $r(= \frac{\Omega_1}{\Omega_2})$, where $t_0 = 5.0$, $\Omega_2 = \Omega$, $\Omega_1 = r\Omega$, and other parameters are the same as those in figure 6(a). In this condition, the current can be expressed as

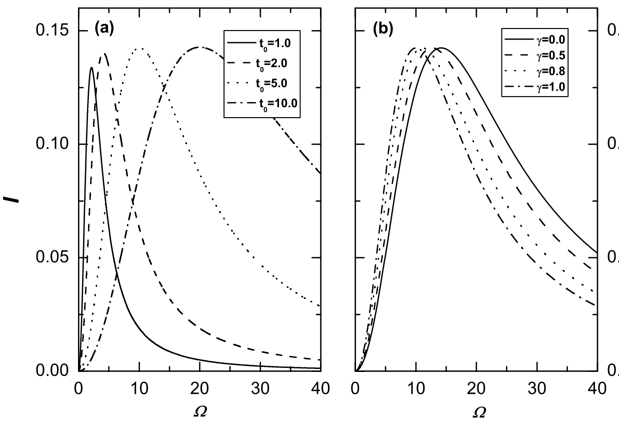

Figure 6. (a) Dependence of the current I on Rabi frequency Ω for different tunneling strengths t_0 with $\Omega_1 = \Omega_2 = \Omega$. (b) Dependence of the current I on Rabi frequency Ω for different Rabi frequencies Ω_1 with $\Omega_1 = r\Omega_2 = r\Omega$ and $t_0 = 5.0$. Other parameters $\theta_1 = 0$, $\theta_2 = \frac{\pi}{2}$, $\Delta = 0.0$, are the same in the two figures. All coefficients are scaled with Γ.

$$I = \frac{4(1 + r^2)t_0^2\Omega^2\Gamma}{[\Omega^2(1 + r^2) + 6t_0^2]^2 + 16t_0^2\Gamma^2 + 28t_0^4}, \quad (23)$$

the maximum value of which is at the point of $\Omega = 2\sqrt{\frac{t_0\sqrt{\Gamma^2+4t_0^2}}{1+r^2}}$. For $t_0 = 5.0$, and $r = 0.0, 0.5, 0.8, 1.0$, the locations of the maximum value of current I are respectively at $\Omega \simeq 14.0$, 12.7, 11.0, and 10.0. And the current maximum values are the same and equal to 0.1424 for different asymmetry r as shown in figure 6(b).

4. Conclusion

In conclusion, we have studied the coherent field phase control of transport through an effective closed three-level structure in an asymmetric double quantum dots system. It is shown that the current spectrum can be sensitively modulated by the relative phases of the laser fields, which indicates a quantum interference similar to the magnetic flux controlled coherent transport in an AB ring interferometer. The results are separately discussed in two different transport regimes, with respect to the location of the chemical potential μ_R. In the regime of $\mu_R < \varepsilon_2$, we discuss the phase modulation on the current for energy difference Δ. It is shown that the two-peak structure appears and the location and height of the current peak can be modulated by choosing different relative phases θ_1 and θ_2. Importantly, the peak of the current can be tuned to nearly zero because the electron is nearly trapped into the ground state of the left QD, due to a complete destructive interference determined by relative phases θ_1 and θ_2. This effect can serve as an optically phase controlled current switch. In the other regime of $\varepsilon_2 < \mu_R < \varepsilon_3$, we find that the two coupled quantum dots can be trapped into a dark state $|D\rangle = \frac{1}{\sqrt{2}}(e^{-i\theta_2}|1\rangle - e^{-i\theta_1}|2\rangle)$, which is decoupled from the light, and here we choose $\Omega_1 = \Omega_2$. That the electron is trapped in a decoupled coherent superposition corresponds to the effect of the current disappearance. Due to its easier manoeuvrability, the laser phase controlled coherent transport provides another feasible mode and a potential application in nanoelectronics.

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Appendix

Here we give the concrete methods to adiabatically eliminate the level $|4\rangle$. The Hamiltonian H_{L_1}, in the interaction picture through the unitary transformation $\exp[iH_D t]$, is given by

$$\tilde{H}_{L_1} = t_2|4\rangle\langle 3|e^{i\Delta_{43}t} + \frac{\Omega_1'}{2}|4\rangle\langle 1|e^{-i(\omega_1 t+\Delta_{14}t+\theta_1)} + \text{h.c.}, \quad (A.1)$$

where $\Delta_{ij} = \varepsilon_i - \varepsilon_j$ $(i \neq j = 1, 2, 3, 4)$. In the limit when $|\omega_1 - \Delta_{41}| \gg \Omega_1$ and the coupling strength $V_{4k} \ll V_{ik}$ $(i = 1, 2, 3)$, the Hamiltonian $\tilde{H}_{L_1}$ consists of highly oscillating terms and to a good approximation we finally obtain the effective Hamiltonian $\tilde{H}_{L_1}' = -i\tilde{H}_{L_1}(t)\int \tilde{H}_{L_1}(t')dt'$ [37], given by

$$\tilde{H}_{L_1}' = \frac{\Omega_1}{2}|3\rangle\langle 1|e^{-i(\omega_1 t+\theta_1)}e^{i\Delta_{31}t} + \text{h.c.}, \quad (A.2)$$

where $\Omega_1 = \Omega_1' t_2/\Delta_{34}$ is the effective coupling constant. Then the Hamiltonian $\tilde{H}_{L_1}'$, in the Schrödinger picture through the unitary transformation $\exp[-iH_D t]$, is given by

$$H_{L_1}' = \frac{\Omega_1}{2}|3\rangle\langle 1|e^{-i(\omega_1 t+\theta_1)} + \text{h.c.} \quad (A.3)$$

That is to say, in the limit of $|\omega_1 - \Delta_{41}| \gg \Omega_1$, the two-step electron transition $|1\rangle \leftrightarrow |4\rangle \leftrightarrow |3\rangle$ can be effectively treated as a transition between $|1\rangle$ and $|3\rangle$.

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