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Carrier-envelope phase control electron transport in an asymmetric double quantum dot irradiated by a few-cycle pulse

Wen-Xing Yang [a,b,*], Ai-Xi Chen [c], Hao Guo [a], Yanfeng Bai [a], Ray-Kuang Lee [b]

[a] Department of Physics, Southeast University, Nanjing 210096, China
[b] Institute of Photonics Technologies, National Tsing-Hua University, Hsinchu 300, Taiwan
[c] Department of Applied Physics, School of Basic Science, East China Jiaotong University, Nanchang 330013, China

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ABSTRACT

We propose a theoretical scheme to coherently control the transport of a single electron in an asymmetric double-quantum-dot system. The single-electron transport originates from the intrinsic interplay between the externally applied few-cycle pulse and the inter-dot tunneling. Solving the equations of motion for dot-density matrix, we reveal numerically that the current exhibits a significant dependence on the carrier-envelope phase (CEP) of the few-cycle pulse, which is similar to the magnetic flux controlled coherent transport in an Aharnov–Bohm (AB) interferometer. As a result, by varying the CEP of the pulse one can suppress or enhance the current either instantaneously or periodically. Our results illustrate the potential to utilize few-cycle pulses for excitation in quantum dot systems through the CEP control, as well as a guidance in the design for possible experimental implementations.

1. Introduction

In recent years, it has been realized that the phase of electromagnetic (EM) fields can provide a powerful technique to control or manipulate matters. The method of phase control is associated with the coherent interaction between EM fields and matters, which has already been applied to a variety of different systems [1–7]. Phase dependent phenomena are demonstrated by controllable population dynamics [3], suppressed spontaneous emission [4], and coherent transport in coupled tunneling systems [5,6].

Instead of atomic and molecular systems, solid-state media, especially the quantum dot (QD) structures, have attracted significant research attention due to their potential applications in the development of novel optoelectronic devices and solid-state quantum information science [8]. In QD systems, to manipulate electrons an external magnetic flux is applied firstly to implement possible coherent phase controls. However, it is technologically difficult to confine a strong magnetic field within a very small region, which might be a crucial obstacle for its future applications in quantum processing and computation. Therefore, it is a reasonable choice to replace the magnetic field by some source more easy to control. A natural idea is to use the phase of applied EM fields to control the coherent transport between QD systems, which has revealed many interesting phenomena ranging from photon-assisted tunneling to charge/spin pumping [9–17]. In a recent experiment, microwave spectroscopy has been measured in coupled QDs [18], and the photon-assisted resonances, which involve the emission or absorption of a microwave photon, are found when applying a modulated gate voltage. Some other systems based on time-dependent influences also give rise to promising physical phenomena and applications [19–23].

Meanwhile, tremendous progress in the generating of ultra-short pulses in the few-cycle regime allows one to explore a new class of phase-dependent phenomena. The relative phase difference between the carrier wave and the pulse envelope, which is the so-called carrier-envelope phase (CEP), has many distinctive observable features. Both experimental and theoretical studies have revealed that the CEP of few-cycle pulses indeed plays an important role in the light-matter interactions [24–40]. In this work, we propose a scheme for controlling single-electron transport in an asymmetric double QD system. The coherent transport is externally controlled by applying a few-cycle pulse with an adjustable CEP. We demonstrate that the single-electron transport is in fact the result of the intrinsic interplay between the external few-cycle pulse and the inter-dot tunnel coupling. In particular, we further show that the current can be periodically suppressed or enhanced by modulating the CEP. Our study provides an efficient tool to manipulate the quantum dynamics in QD systems with an adjustable CEP.

2. Model and equations of motion

The present QD system is given in Fig. 1. This device is composed of two different QDs (marked as the left (L) and right

* Corresponding author. Tel.: +86 2552090605-8407
E-mail addresses: wenxingyang@seu.edu.cn, wenxingyang2@126.com (W.-X. Yang).

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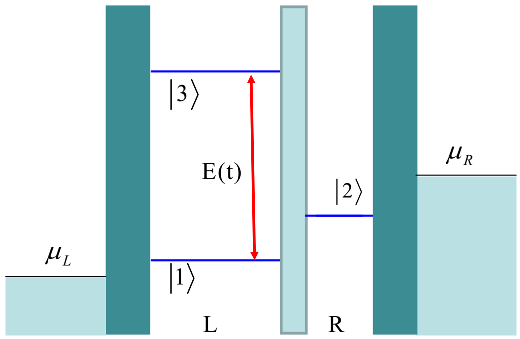

Fig. 1. Schematic diagram of a three-level system which consists of the ground state $|1\rangle$ and the first excited state $|3\rangle$ in the left dot, and the ground state $|2\rangle$ in the right dot in an asymmetric double QD structure. The ground state $|1\rangle$ in the left dot is coupled resonantly to the excited state $|3\rangle$ by a few-cycle pulse $E(t)$.

(R) dots) and two normal metal leads. The size of the right dot is assumed to be much smaller than that of the left one, therefore the energy spacing of it is much larger than that of the left one. The ground state $|1\rangle$ and the first excited state $|3\rangle$ of the left dot and the other ground state $|2\rangle$ of the right dot form a three-level system. Such kind of QD configuration can be realized by a quantum heterostructure consisting of different semiconductor materials (GaAs/AlGaAs). Metal gates could be deposited on the top of a GaAs/AlGaAs heterostructure with a two-dimensional electron gas about 100 nm below the surface [8]. Generically, Coulomb charging energies are of a few meV and they are the largest energy scales. Hence the states associated with more than one additional transport electron can be safely neglected [41–43]. Typical energy difference between the ground state $|1\rangle$ and the first excited state $|3\rangle$ is of the order of several meV [8]. The states $|1\rangle$ and $|3\rangle$ are both coupled to the state $|2\rangle$ via the tunneling through the barrier. The transition between the states $|1\rangle$ and $|3\rangle$ themselves is resonantly driven by a few-cycle pulse characterized by the electric field [25] $E(t)$, which is further defined via the vector potential, $E(t) = -\partial A(t)/\partial t$. This definition guarantees that the vector potential $A(t)$ vanishes at $t = \pm\infty$, or equivalently, the pulse area under $E(t)$ over the entire pulse duration becomes exactly zero [44], which excludes unphysical results. Specifically we assume that the vector potential has a Gaussian temporal envelope, i.e.,

$$A(t) = A_0 e^{-(2\ln 2)(t-2\tau)^2/\tau^2}\sin(\omega t + \varphi), \tag{1}$$

where A_0 is the peak of the pulse envelope, ω is the carrier frequency ($2\pi/\omega$ corresponds to an optical oscillation cycle time), and φ is the CEP. The CEP describes the offset of the peak laser pulse relative to the peak position of the envelope. In Eq. (1), we have assumed that the temporal intensity profile is Gaussian temporal envelope with pulse duration τ (full width at half maximum (FWHM)).

In the present system, the effective Hilbert space can be thought of as being spanned by four basis states: $|0\rangle$ (no electron in both dots), $|1\rangle$ (one electron in the ground 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 left dot). The total Hamiltonian of the system can be expressed in the form [8]

$$H = H_{D0} + H_{DL} + H_{DT} + H_B + H_{ep}. \tag{2}$$

The first term is the free Hamiltonian of the two QDs,

$$H_{D0} = \sum_{j=1,2,3} \epsilon_j |j\rangle\langle j|, \tag{3}$$

where ϵ_j represents the energy of the state $|j\rangle$ $(j = 1, 2, 3)$. The second term describes the interaction between the left dot and a few-cycle pulse,

$$H_{DL} = -\Omega\xi(t)|1\rangle\langle 3| + H.c., \tag{4}$$

where $\xi(t) = \omega^{-1}\partial[e^{-(2\ln 2)(t-2\tau)^2/\tau^2}\sin(\omega t + \varphi)]/\partial t$. The Rabi frequency for the transition $|1\rangle \leftrightarrow |3\rangle$ is denoted by $2\Omega = \mu_{13}\omega A_0/\hbar$, with μ_{13} being the corresponding transition dipole moment. The third term describes the coupling Hamiltonian of the dots through the tunneling effect,

$$H_{DT} = \kappa_{12}|1\rangle\langle 2| + \kappa_{32}|3\rangle\langle 2| + H.c., \tag{5}$$

where κ_{ij} is the tunneling coefficient between the states $|i\rangle$ and $|j\rangle$. The fourth term represents the interaction between the leads and the dots and the interaction between the photon modes and the left dot,

$$H_B = \sum_{k,\eta=L,R} \varepsilon_{k,\eta} c^\dagger_{k,\eta} c_{k,\eta} + \sum_{q\nu} \omega_q a^\dagger_{q\nu} a_{q\nu} + \sum_k [V_{kL} c^\dagger_{kL}(|0\rangle\langle 1| + |0\rangle\langle 3|) + V_{kR} c^\dagger_{kR}|0\rangle\langle 2| + H.c.] + \sum_{q\nu}(V_{q\nu}|3\rangle\langle 1|a^\dagger_{q\nu} + H.c.), \tag{6}$$

where $c_{k\eta}$ is the annihilation operator of electrons in the lead η $(\eta = L, R)$, and $a_{q\nu}$ is the annihilation operator of photons with momentum q and polarization ν. $V_{k\eta}$ denotes the coupling strength of the interaction between electrons in the QDs and the leads η. For simplicity, we have assumed that both levels in the left dot interact with the lead L via the same strength. The final term

$$H_{ep} = \sum_Q \frac{1}{2} g_Q \sigma_{z1}(a_{-Q} + a^\dagger_Q) + \sum_Q \omega_Q a^\dagger_Q a_Q + \sum_p \frac{1}{2} g_p \sigma_{z2}(a_{-p} + a^\dagger_p) + \sum_p \omega_p a^\dagger_p a_p \tag{7}$$

describes the coupling of phonons to the charge density which has been found to be the dominant interacting mechanism in a single two-level dot and in the double QDs. Here $\sigma_{z1} = (|3\rangle\langle 3| - |1\rangle\langle 1|)$, $\sigma_{z2} = (|2\rangle\langle 2| - |1\rangle\langle 1|)$, $a^\dagger_Q(a^\dagger_p)$ is the creation operator of phonons with frequency $\omega_Q(\omega_p)$, and $g_Q(g_p)$ is the coupling strength of the interaction between electrons and phonons [45].

The transition rates between the states of the left lead and the two states of the left dot are expressed as $\Gamma^\pm_{Lj} = \Gamma_L f^\pm_j(\epsilon_j)$ $(j = 1, 3)$, where $\Gamma_L = 2\pi\sum_k |V_{kL}|^2\delta(\epsilon_j - \varepsilon_{kL})$ and $f^\pm_j(\epsilon) = \{1 + \exp[\pm(\epsilon - \mu_L)/k_BT\}^{-1}$ is the Fermi distribution function of the left reservoir. Here μ_L is the chemical potential, and $\pm$ corresponds to the occupied/empty state of the left lead. More specifically, Γ^+_{Lj} represents the tunneling rate of the transition of electrons from the left lead into the state $|j\rangle$, while Γ^-_{Lj} represents the tunneling rate of the inverse process. Furthermore, here we only consider a low bias configuration where μ_L is well below the two energy levels of the left dot [11,17], hence the Fermi distribution functions can be approximated as $f^+_L(\epsilon_j) = 0$ and $f^-_L(\epsilon_j) = 1$. In this case, the transition rates are further simplified as $\Gamma^-_{Lj} = \Gamma_L$ and $\Gamma^+_{Lj} = 0$.

Now we consider the transition of electrons in the “right” part of the system. We assume that the chemical potential μ_R of the right lead is well above the energy level of the right dot. By similar deduction, the tunneling rate of transition of electrons from the state $|2\rangle$ to the right lead is given by $\Gamma^-_R = 0$, and the tunneling rate of the inverse process is approximated as $\Gamma^+_R = \Gamma_R$ [11,17]. For simplicity, we assume that the dot-lead tunneling rates satisfy $\Gamma_L = \Gamma_R = \Gamma$. Under the assumption of weak coupling between the QDs and the leads, the behavior of the double QDs in the sequential regime can be described in terms of the density operator of the system. After adiabatically eliminating the reservoir operators c_{kR}, c_{kL}, $a_{q\nu}$, a_p, a_Q and employing the Born–Markov approximation, the master equations for the density matrix of the system can be written as [49]

$$\frac{\partial\rho_{33}}{\partial t} = -\gamma_3\rho_{33} - \Gamma\rho_{00} + i\Omega\xi(t)(\rho_{13} - \rho_{31}) + i\kappa_{32}(\rho_{23} - \rho_{32}), \tag{8}$$