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Coherent Revival of Tunneling

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We introduce a new tunneling effect by a driving field, referred to as coherent revival of tunneling (CRT), corresponding to complete tunneling (transmission coefficient = 1) which revives from the circumstance of total reflection (transmission coefficient ≈ 0) by application of an appropriate perpendicular high-frequency ac field. To illustrate CRT, we simulate electron transport through fishbone-like quantum dot arrays by using single-particle Green's functions along with Floquet theory, and we explore the corresponding current-field amplitude characteristics as well as current-polarization characteristics. In regard to the two characteristics, we show that CRT exhibits entirely different features, when compared with coherent destruction of tunneling and photon assisted tunneling. We also discuss two practical conditions for experimental realization of CRT.

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I. INTRODUCTION

Tunneling in a time-dependent field is a fundamental subject in quantum mechanics and has been investigated for more than fifty years^{1–10}. Tunneling phenomena by a driving field have been considered in many different contexts in chemistry^{3,4,11} and physics including photon-assisted tunneling (PAT)^{1,2,5,10,12–22}, coherent destruction of tunneling (CDT)^{6,23,24}, dynamical localization⁹, and coherent quantum ratchets^{25,26}. For example, PAT has been widely explored in superconductor-insulator-superconductor tunnel junctions^{1,12}, semiconductor nanostructures^{13,14}, quantum dots^{15,16}, superlattices¹⁷, and molecular junctions^{10,18–22}. CDT has been experimentally observed in an Er:Yb-doped glass⁷ and in cold atoms within a double-well potential⁸.

In this Letter, we introduce a new phenomenon, coherent revival of tunneling (CRT), where total transmission can be achieved by an appropriate perpendicular high-frequency ac field. In addition, we show that CRT exhibits completely different current-polarization characteristics from PAT as well as distinct current-field amplitude characteristics from CDT. To demonstrate CRT, we consider the following device setup as shown in Fig. 1 (a): a quantum dot array coupled to two electrodes (the source and the drain) and exposed to a laser field. We will analyze light-driven electron transport through quantum dot arrays based on four different configurations **I–IV** in Fig. 1(b). **I** is used for a demonstration of CDT while **II–IV** are for the demonstrations of CRT.

II. MODEL AND METHOD

—In order to develop the model Hamiltonian, we adopted several approximations including (i) the lead-dot model, (ii) the nearest neighbor tight-binding model and the electric dipole approximation for the quantum dot arrays in a laser field, (iii) a non-interacting electron model for the electrodes, (iv) thermal equilibrium for the electrons in the electrodes, and (v) the wide-band limit

FIG. 1: (a) A quantum dot array coupled to two electrodes (the source and drain electrodes) is exposed to a laser field propagating along the z direction with polarization angle β and frequency ω . Quantum dot arrays are placed in the xy -plane and the blue double headed arrow stands for laser polarization. (b) Quantum dot arrays **I–IV**. The distance between two nearest neighbor quantum dots is d and the coupling is $\Delta_{nn'}$.

approximation for the coupling function. These approximations have been extensively employed in the previous studies^{10,21,23–25,27–30}.

In the lead-dot model, the total system Hamiltonian $H(t)$ is decomposed into the dot Hamiltonian $H_{\text{dot}}(t)$, the lead Hamiltonian H_{lead} , and the dot-lead couplings $H_{\text{dot-lead}}$, *i.e.*, $H(t) = H_{\text{dot}}(t) + H_{\text{lead}} + H_{\text{dot-lead}}$. For the quantum dot arrays, we assume that each dot is placed in the x - y plane, the distance between two nearest neighbor dots is d , and the laser field propagates along the z direction in Fig. 1. Moreover, the dot Hamiltonian is treated within the nearest neighbor tight-binding model and the electric dipole approximation, *i.e.*,

$$H_{\text{dot}}(t) = \sum_n (E_0 - e\mathbf{r}_n \cdot \mathbf{E}(t))|n\rangle\langle n| - \sum_{nn'} \Delta_{nn'}|n\rangle\langle n'|, \quad (1)$$

where e is an elementary charge, $|n\rangle$ stands for the orbital on the n -th quantum dot at position $\mathbf{r}_n$, E_0 is the on-site energy, $\Delta_{nn'}$ is the coupling between two nearest

neighbor dots, and $\mathbf{E}(t) = \mathbf{E}(t + T)$ is a monochromatic linearly-polarized electric field with a period T and angular frequency $\omega = 2\pi/T$. In a practical experimental setup, it is difficult to obtain the same couplings between dots. To model the fluctuation of dot couplings, we assume that $\Delta_{nn'} = \Delta_{n'n} = \Delta_0 + \Delta_1$, where Δ_0 is a constant and Δ_1 is a random variable which describes the variation of coupling strength. The linearly-polarized electric field $\mathbf{E}(t)$ has amplitude E and a polarization angle β in x-y plane, *i.e.*,

$$\mathbf{E}(t) = E(\cos \beta \hat{e}_x + \sin \beta \hat{e}_y) \cos \omega t, \quad (2)$$

where $\hat{e}_x$ and $\hat{e}_y$ stand for unit vectors in x and y directions, respectively.

The two leads are described by a non-interacting electron model, *i.e.*, $H_{\text{lead}} = \sum_{lq} \epsilon_{lq} |lq\rangle \langle lq|$, where $|lq\rangle$ stands for the orbital with energy ϵ_{lq} in the lead l with the mode q , and $l = S$ and D denote the source and the drain, respectively. In addition, assuming that the electrons in the leads are in thermal equilibrium, their average occupation number can be expressed as the Fermi function $f_l(\epsilon) = (1 + e^{(\epsilon - \mu_l)/k_B\theta})^{-1}$ with chemical potential μ_l in the lead l at temperature θ . The dot-lead couplings are modeled as

$$H_{\text{dot-lead}} = \sum_q V_{Sq,u} |Sq\rangle \langle u| + V_{Dq,v} |Dq\rangle \langle v| + \text{H.c.}, \quad (3)$$

where $|u\rangle$ ($|v\rangle$) represent the dot u (v) coupled to the source (drain), and the element of the coupling function can be expressed as $\Gamma_{l,m}(\epsilon) = 2\pi \sum_q |V_{lq,m}|^2 \delta(\epsilon - \epsilon_{lq})$, $(l, m) = (S, u)$ and (D, v) . Consider the symmetric coupling and the wide-band limit approximation, which corresponds to $\Gamma_{S,u}(\epsilon) = \Gamma_{D,v}(\epsilon) = \Gamma = \text{a constant}$. That is, we only keep the imaginary part of the self-energy and set $\Sigma = \frac{-i}{2}(|u\rangle \Gamma \langle u| + |v\rangle \Gamma \langle v|)$ since the real part of the self-energy does not play a crucial role in the analysis.

In a time-periodic field, the time-average current can be computed using^{10,21,24,25,27–29,31}

$$\bar{I} = \frac{2e}{h} \sum_{k=-\infty}^{+\infty} \int d\epsilon \{T_{DS}^{(k)}(\epsilon) f_S(\epsilon) - T_{SD}^{(k)}(\epsilon) f_D(\epsilon)\}, \quad (4)$$

where $T_{DS}^{(k)}(\epsilon) = \Gamma_{D,v}(\epsilon + k\hbar\omega) \Gamma_{S,u}(\epsilon) |G_{vu}^{(k)}(\epsilon)|^2$ reads the transmission of the tunneling electron from the source to the drain with energy ϵ accompanied by k -photon absorption ($k > 0$) or emission ($k < 0$). The k -th component of the retarded Green function $G_{vu}^{(k)}(\epsilon)$ can be derived from the time-independent infinite-dimensional Floquet eigenvalue matrix equation^{10,21,22}. We let $T_{SD}(\epsilon) = \sum_{k=-\infty}^{\infty} T_{SD}^{(k)}(\epsilon)$ and $T_{DS}(\epsilon) = \sum_{k=-\infty}^{\infty} T_{DS}^{(k)}(\epsilon)$, Eq. (4) becomes

$$\bar{I} = \frac{2e}{h} \int_{-\infty}^{\infty} d\epsilon [T_{DS}(\epsilon) f_S(\epsilon) - T_{SD}(\epsilon) f_D(\epsilon)]. \quad (5)$$

In **I–IV**, the generalized parity symmetry $S_{GP} : (\mathbf{r}, t) \rightarrow (-\mathbf{r}, t + \pi/\omega)$ is satisfied³² so we have $T_{SD}(\epsilon) = T_{DS}(\epsilon)$ and the transmission coefficient can be defined as $T(\epsilon) = T_{SD}(\epsilon)$ ²¹. Eq (5) reduces to a Landauer-type formula.

FIG. 2: Current-field amplitude characteristics for $V_{SD} = 0.05\Delta_0$, $k_B\theta = 0$, $\Delta_1 = 0$, $\hbar\omega = 10\Delta_0$, and $\Gamma = 0.1\Delta_0$, where $A = edE$ is proportional to the field amplitude E . Red solid (dashed) lines show the current through **I** in a laser field with a polarization angle $\beta = 0^\circ$ (90°). Green, blue, black solid lines correspond to currents through **II–IV**, respectively, in a laser field with a polarization angle $\beta = 90^\circ$.

III. RESULTS AND DISCUSSIONS

—Consider the following conditions in the experimental setup. We assume a small applied voltage ($eV_{SD} = 0.05\Delta_0$) and symmetric chemical potentials ($\mu_S = E_0 + eV_{SD}/2$, $\mu_D = E_0 - eV_{SD}/2$, and $E_0 = 0$). Fig. 2 shows current-field amplitude characteristics of the triple quantum dots **I** and the fishbone-like quantum dot arrays **II–IV**. Note that in Fig. 2, 3, and 5, we do not consider variation of coupling strength and set $\Delta_1 = 0$. In the array **I**, the current is strongly suppressed at $A \approx 2.4, 5.5$, and $8.6\hbar\omega$ (corresponding to the roots of Bessel functions of the first kind, *i.e.*, $J_0(\frac{A}{\hbar\omega}) = 0$) in a laser field with $\beta = 0^\circ$, where A is proportional to the field amplitude E ($A = edE$). This phenomenon is the so-called coherent destruction of tunneling, which has been extensively studied^{27,30}. On the other hand, the current does not change with A in a laser field with $\beta = 90^\circ$ because $\mathbf{e}_r \cdot \mathbf{E}(t) = 0$. In addition, $T(\epsilon) \approx 1$ between μ_S and μ_D due to one of the conduction channels at $\epsilon = 0$. Consequently, the magnitude of the current can be estimated as $\bar{I} = \frac{2e}{h} \int d\epsilon T(\epsilon) [f_S(\epsilon) - f_D(\epsilon)] \approx \frac{2e}{h} V_{SD} = 0.1e\Delta_0/h$, which is very close to the numerical result (the red dashed line in Fig. 2).

On the other hand, the fishbone-like quantum dot arrays **II–IV** in a perpendicular laser field ($\beta = 90^\circ$) exhibit entirely different current-field amplitude characteristics, *e.g.*, the current reaches the maximum at the roots of the Bessel function, which is complete contrast to the CDT observed in the array **I** in a parallel laser field ($\beta = 0^\circ$). To understand this feature, we consider the transmission spectra of **II** and the corresponding energy levels of the quasi-states. The upper and lower bright lines in Fig. 3(a) correspond to the energies of quasi-states in Fig. 3(b), indicating that the quasi-states provide conduction channels for tunneling electrons. However, the intensity of the transmission spectra along $\epsilon = 0$

FIG. 3: (a) Transmission spectra of **II** under laser fields for $E_0 = 0$, $\hbar\omega = 10\Delta_0$, $\beta = 90^\circ$, $\Delta_1 = 0$, and $\Gamma = 0.1\Delta_0$, where the color denotes the magnitude of transmission on a linear scale. (b) The energies of the quasi-states of **II** in the first Brillouin zone ($-\hbar\omega/2 < \epsilon < \hbar\omega/2$).

FIG. 4: Mechanism of coherent revival of tunneling. Electron transport through **II** in a laser field with $\beta = 90^\circ$ (the green double headed arrow denotes laser polarization). In a high-frequency field, the coupling $\Delta_{nn'}$ in the three dots indicated by the red box can transform to the effective coupling $\Delta_{\text{eff}} = \Delta_{nn'} J_0(\frac{A}{\hbar\omega})$, where $J_0(\frac{A}{\hbar\omega})$ is a Bessel function of the first kind. As $J_0(\frac{A}{\hbar\omega}) = 0$, **II** is nearly equivalent to **I**, and destructive quantum interference caused by the fishbone-like structure vanishes.

changes with the field amplitude A and shows strong transmission suppression at small A . It is intriguing that the quasi-states exist but they do not assist tunneling. The transmission suppression is a manifestation of tunneling with destructive quantum interference (DQI) due to the synergistic effect of resonant states (conduction channels). The synergistic effect of resonant states caused by specific structures in the absence of a driving field has been extensively discussed in the domain of molecular electronics^{33–39}. However, as A reaches the roots of the Bessel function, the strong transmission suppression vanishes and the total transmission revives (The synergistic effect of the conduction channels disappears, and the conduction channel revivals.), indicating that the current reaches the local maximum. As a result, we call this phenomenon *coherent revival of tunneling* (CRT).

It is interesting that the patterns of quasi-state energies in CRT and CDT are completely different (see Fig. A.1 in Appendix A). For **II**, the energies of the five quasi-

states in the first Brillouin zone can be solved in the high-frequency approximation ($\hbar\omega \gg \Delta_0$) and they are $-\Delta_0 \sqrt{2 + 2|J_0(\frac{A}{\hbar\omega})|^2}$, 0 (three degenerate states), and $\Delta_0 \sqrt{2 + 2|J_0(\frac{A}{\hbar\omega})|^2}$ (The details can be found in Appendix B). Numerical results also show that the energies of the three degenerate states are insensitive to the field amplitude and the upper and lower lines never intersect at $\epsilon = 0$. In Fig. 3, the suppression and revival of tunneling are independent from the quasi-state energies of the three degenerate states. On the contrary, CDT results from the fact that the quasi-state energies of **I** intersect with each other at the roots of the Bessel function. In addition to numerical results, for the CRT effect, the closed form expression of Green's functions also indicates that $T(0) \approx 1$ at $J_0(\frac{A}{\hbar\omega}) = 0$ while $T(0) \approx 0$ at $J_0(\frac{A}{\hbar\omega}) = 1$ (see Appendix B). CRT and CDT show extremely different characteristics in the quasi-state energy plot.

The mechanism of CRT can be also understood by the time-average effective Hamiltonian in the high-frequency approximation. In a high-frequency field ($\hbar\omega \gg \Delta_0$), we can make a transformation of the Hamiltonian of **II** to a rotating frame⁴⁰, and the effective coupling in the three dots indicated by the red box in Fig. 4 becomes $\Delta_{\text{eff}} = \Delta_{nn'} J_0(\frac{A}{\hbar\omega})$. As $J_0(\frac{A}{\hbar\omega}) = 0$, **II** is nearly equivalent to **I** so the tunneling with DQI vanishes and the current reaches the maximum. In addition, Fig. 2 shows that the maximum current through **II** (the green line) almost coincides with the current through **I**, which also indicates that **II** reduces to **I** when $J_0(\frac{A}{\hbar\omega}) = 0$. Based on this assessment, in a high-frequency laser field with $\beta = 90^\circ$ and $J_0(\frac{A}{\hbar\omega}) = 0$, **III** and **IV** are equivalent to linear five and seven linear quantum dots, respectively.

The tunneling effect in a time-dependent field can be also investigated by current-polarization characteristics. For example, one-photon assisted tunneling exhibits that

FIG. 5: Current-polarization characteristics of **II-IV** under laser fields for $\hbar\omega = 10\Delta_0$, $A = 2.405\hbar\omega$, $\Delta_1 = 0$, and $\Gamma = 0.1\Delta_0$.

FIG. 6: Current-field amplitude characteristics of **IV** for $V_{SD} = 0.05\Delta_0$, $k_B\theta = 0$, $\hbar\omega = 4\Delta_0$, and $\Gamma = 0.1\Delta_0$. The red line corresponds to current through **IV** with uniform couplings ($\Delta_1 = 0$) while the black lines correspond to current through **IV** with six sets of random couplings ($|\Delta_1| < 0.3\Delta_0$).

$\bar{I} \propto \cos^2 \beta^{22}$, *i.e.*, the current reaches the maximum at $\beta = 0^\circ$ and the minimum at $\beta = 90^\circ$. It can be understood that electron tunneling driven by electric field oscillation is most efficient when the directions of electron tunneling and electric field oscillation are the same. However, the fishbone-like quantum dot arrays show different features in current-polarization amplitude characteristics. Fig. 5 shows that the maximum current through **II-IV** occurs at $\beta = 90^\circ$ (the directions of current and polarization are perpendicular) while the nearly zero current occurs at $\beta = 0^\circ$. The nearly zero current at $\beta = 0^\circ$ originates from CDT while the maximum current at $\beta = 90^\circ$ results from CRT. Moreover, the width of the current peak in Fig. 5 becomes narrower as the length of the fishbone structure increases. This feature results from tunneling with DQI enhanced by the repeated cross structures.

In practice, various experimental conditions need to be considered. Here we discuss two of these experimental conditions: (i) variation of coupling strength between the dots and (ii) leakage currents caused by a breakdown electric field. First, it is extremely difficult to obtain the same couplings between dots. In order to model the variation of coupling strength between dots, we assume that

Δ_1 is a random variable and $|\Delta_1| < 0.3\Delta_0$. Fig. 6 shows current through **IV** with uniform couplings (the red line) and with six sets of random couplings (the black lines). We find that even in the intermediate frequency regime ($\hbar\omega = 4\Delta_0$) the peaks of the red and black lines all still occur close to the roots of the Bessel function, indicating that the feature of CRT is not sensitive to variation of coupling strength, and CRT could be observed in the fishbone-like quantum dot arrays with imperfect fabrication. Note that variation of coupling strength may enhance the maximum current due to the breaking of the generalized parity symmetry. As the generalized parity symmetry no longer holds, zero bias current can be driven by a time-periodic field, leading to current enhancement or suppression³². Second, when the strong electric field oscillation is parallel to the direction of tunneling current, source-to-drain leakage current may happen due to electric field breakdown of the quantum dots (The breakdown electric field of excellent insulators such as Al_2O_3 is about 1 V/nm⁴¹). However, for CRT, source-to-drain leakage current can be avoided because the electric field oscillation is perpendicular to the direction of source-to-drain current.

IV. CONCLUSION

To summarize, we have presented the concepts of CRT and discussed the corresponding current-field amplitude and current-polarization characteristics. The results imply that the fishbone-like quantum dot arrays are good candidates for the observation of CRT. We would like emphasize the distinction between CRT and CDT: (i) CRT and CDT occurs in a laser field with the polarization angle 90° and 0° , respectively. (ii) As $J_0(\frac{A}{\hbar\omega}) = 0$, CRT exhibits current enhancement (complete tunneling) while CDT shows current suppression. (iii) No energy level crossing of quasi-states occurs in CRT. For a practical experimental setup, we find that CRT is not sensitive to reasonable variation of coupling strength between dots. Moreover, the condition $\beta = 90^\circ$ for CRT can avoid unexpected tunneling mechanisms such as source-to-drain leakage currents. Note that the Coulomb interaction and spin-orbit coupling may influence the CRT effect. In this study, we explored the CRT effect in the single-electron tunneling regime (the low-bias condition). In this condition, Coulomb interaction should not play a crucial role^{42,43}. In addition, according to a previous study⁴⁴, the spin-orbit interaction is much smaller than Δ_0 in double quantum dots. As a result, it is reasonable to neglect the spin-orbit interaction in order to facilitate the theoretical analysis and provide a clear explanation of CRT mechanism. We hope that this study motivates additional investigations into CRT as well as other novel tunneling effects driven by time-dependent fields in various contexts.

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Appendix A: Transmission Spectrum and Quasi-State Energy Plot of I

FIG. A.1: (a) Transmission spectra of **I** under laser fields for $E_0 = 0$, $\hbar\omega = 10\Delta_0$, $\beta = 0^\circ$, $\Delta_1 = 0$, and $\Gamma = 0.1\Delta_0$, where the colors denote the magnitude of transmission on a linear scale. (b) The energies of the quasi-states of **I** in the first Brillouin zone ($-\hbar\omega/2 < \epsilon < \hbar\omega/2$). As $J_0(\frac{A}{\hbar\omega}) = 0$, the three quasi-state energy curves cross each other and the transmission becomes zero at $A \approx 2.40, 5.52$, and $8.65\hbar\omega$, leading to coherent destruction of tunneling.

Appendix B: Closed Form Expression for Quasi-State Energy and Green's Function of II

FIG. B.1: Fishbone-like quantum dot array **II**. The numbers denote the n -th quantum dot, *e.g.*, 1 stands for the 1st quantum dot.

In order to understand the origin of coherent revival of tunneling (CRT) and the curves of the quasi-state energies in Fig. 3, we solve the time-dependent Schrödinger equation of **II** in the high-frequency limit. According to Eq. (1)

and (2), the Hamiltonian of the system **II** in a laser field with the polarization angle $\beta = 90^\circ$ is

$$H_{\text{II}}(t) = \begin{matrix} & |1\rangle & |2\rangle & |3\rangle & & |4\rangle & & |5\rangle \\ \begin{matrix} \langle 1| \\ \langle 2| \\ \langle 3| \\ \langle 4| \\ \langle 5| \end{matrix} & \begin{pmatrix} E_0 & \Delta_0 & 0 & & 0 \\ \Delta_0 & E_0 & \Delta_0 & & \Delta_0 \\ 0 & \Delta_0 & E_0 & & 0 \\ 0 & \Delta_0 & 0 & E_0 + A \cos(\omega t) & 0 \\ 0 & \Delta_0 & 0 & 0 & E_0 - A \cos(\omega t) \end{pmatrix} \end{matrix}, \quad (\text{B1})$$

where $|n\rangle$ stands for the orbital on the n -th quantum dot (see Fig. B.1) and $A = edE$. Note that for simplicity we adopt $\Delta_{n,n'} = \Delta_0$.

The system **II** is coupled to the two electrodes, and the effect of the electrodes can be modeled as the self-energy term $\Sigma = \frac{-i}{2}(|1\rangle\Gamma\langle 1| + |3\rangle\Gamma\langle 3|)$, where Γ is a coupling function. To get the closed form expression of the quasi-state energies, we neglect the contribution of the self-energy since $\Gamma \ll \Delta_0$, *i.e.*, set $\Gamma = 0$, and transform the Hamiltonian in Eq. (B1) into a rotating frame

$$H_{\text{II}}^{\text{rot}}(t) = U_{\text{rot}}^\dagger(t) H_{\text{II}}(t) U_{\text{rot}}(t) - i\hbar U_{\text{rot}}^\dagger(t) \frac{dU_{\text{rot}}(t)}{dt} \\ = \begin{pmatrix} E_0 & \Delta_0 & 0 & 0 & 0 \\ \Delta_0 & E_0 & \Delta_0 & \Delta_0 \exp(-i\frac{A}{\hbar\omega} \sin(\omega t)) & \Delta_0 \exp(i\frac{A}{\hbar\omega} \sin(\omega t)) \\ 0 & \Delta_0 & E_0 & 0 & 0 \\ 0 & \Delta_0 \exp(i\frac{A}{\hbar\omega} \sin(\omega t)) & 0 & E_0 & 0 \\ 0 & \Delta_0 \exp(-i\frac{A}{\hbar\omega} \sin(\omega t)) & 0 & 0 & E_0 \end{pmatrix}, \quad (\text{B2})$$

where

$$U_{\text{rot}}(t) = \begin{pmatrix} 1 & 0 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 & 0 \\ 0 & 0 & 1 & 0 & 0 \\ 0 & 0 & 0 & \exp(-i\frac{A}{\hbar\omega} \sin(\omega t)) & 0 \\ 0 & 0 & 0 & 0 & \exp(i\frac{A}{\hbar\omega} \sin(\omega t)) \end{pmatrix}. \quad (\text{B3})$$

By using $e^{ia \sin(b)} = \sum_{m=-\infty}^{\infty} J_m(a) e^{imb}$ and the high-frequency approximation, we can derive time-average effective Hamiltonian

$$\bar{H}_{\text{II}}^{\text{rot}} = \begin{pmatrix} E_0 & \Delta_0 & 0 & 0 & 0 \\ \Delta_0 & E_0 & \Delta_0 & \Delta_0 J_0(\frac{A}{\hbar\omega}) & \Delta_0 J_0(\frac{A}{\hbar\omega}) \\ 0 & \Delta_0 & E_0 & 0 & 0 \\ 0 & \Delta_0 J_0(\frac{A}{\hbar\omega}) & 0 & E_0 & 0 \\ 0 & \Delta_0 J_0(\frac{A}{\hbar\omega}) & 0 & 0 & E_0 \end{pmatrix} \quad (\text{B4})$$

The quasi-state energies and their corresponding eigenvectors can be solved from the effective matrix equation

$$\sum_n \langle n' | \bar{H}_{\text{II}}^{\text{rot}} | n \rangle \phi_{n\nu} = q_\nu \phi_{n'\nu}. \quad (\text{B5})$$

From Eq. (B4) and (B5), as $E_0 = 0$, we can derive the quasi-state energies $q_1 = -\Delta_0 \sqrt{2 + 2|J_0(\frac{A}{\hbar\omega})|^2}$, $q_2 = q_3 = q_4 = 0$ (three degenerate states), and $q_5 = \Delta_0 \sqrt{2 + 2|J_0(\frac{A}{\hbar\omega})|^2}$. q_1 and q_5 correspond to the lower and upper curves in Fig 3(b), respectively. q_2 , q_3 , and q_4 correspond to the line along $\epsilon = 0$ in Fig. 3(b). Moreover, as $A = 0$, we can derive that $J_0(\frac{A}{\hbar\omega}) = 1$

and $q_1 = -q_5 = -2\Delta_0$. On the other hand, when $A \rightarrow \infty$ or $\frac{A}{\hbar\omega}$ satisfies the roots of the Bessel function ($J_0(\frac{A}{\hbar\omega}) = 0$), we can obtain $q_1 = -q_5 = -\sqrt{2}\Delta_0$. The two results are in agreement with the behavior of the curves in Fig 3(b).

In the zero-bias limit, the time-average current can be approximated as $\bar{I} \approx \frac{2e}{h} T(\mu) V_{\text{SD}} = \frac{2e}{h} T(0) V_{\text{SD}}$, where $\mu = 0$ is the chemical potential of the electrodes and V_{SD} is the source-drain voltage. The transmission coefficient $T(\epsilon)$ and the retarded Green's function $G_{nn'}^{\text{R}}(\epsilon)$ for a tunneling electron with energy ϵ can be computed via

$$T(\epsilon) = \Gamma^2 |G_{nn'}^{\text{R}}(\epsilon)|^2, \quad (\text{B6})$$

$$G_{nn'}^{\text{R}}(\epsilon) = \sum_\nu \frac{\phi_{n\nu} \phi_{n'\nu}}{\epsilon - q_\nu}, \quad (\text{B7})$$

where n and n' correspond to the dots connected to the electrodes. For the system **II**, n and $n' = 1$ or 3 .

To compute $T(\epsilon)$ and $G_{nn'}^{\text{R}}(\epsilon)$, the self-energy contributed from the electrodes has to be considered. q_ν and $\phi_{n\nu}$ in Eq. (B7) can be solved by

$$\sum_n \langle n' | \bar{H}_{\text{II}}^{\text{rot}} + \Sigma | n \rangle \phi_{n\nu} = q_\nu \phi_{n'\nu}. \quad (\text{B8})$$

As $\Gamma \neq 0$, the closed form of q_ν and $\phi_{n\nu}$ are very complicated. However, when $\Gamma \ll \Delta_0$, the quasi-state energies can be approximately expressed as $q_1 \approx -\Delta_0 \sqrt{2 + 2|J_0(\frac{A}{\hbar\omega})|^2}$, $q_2 \approx q_3 \approx q_4 \approx 0$, and $q_5 \approx \Delta_0 \sqrt{2 + 2|J_0(\frac{A}{\hbar\omega})|^2}$. The asymptotic values of q_ν

TABLE I: The Asymptotic Values of q_ν and $\phi_{n\nu}$ for $\Gamma \neq 0$ and $\Gamma \ll \Delta_0$

ν	$J_0(\frac{A}{\hbar\omega}) = 1$			$J_0(\frac{A}{\hbar\omega}) = 0$		
	q_ν	$\phi_{1\nu}$	$\phi_{3\nu}$	q_ν	$\phi_{1\nu}$	$\phi_{3\nu}$
1	$-2\Delta_0 - i\Gamma/8$	$\frac{1}{2\sqrt{2}}$	$\frac{1}{2\sqrt{2}}$	$-\sqrt{2}\Delta_0 - i\Gamma/4$	$1/2$	$1/2$
2	$-i\Gamma/2$	$\frac{1}{\sqrt{2}}$	$-\frac{1}{\sqrt{2}}$	$-i\Gamma/2$	$\frac{1}{\sqrt{2}}$	$-\frac{1}{\sqrt{2}}$
3	$-i\Gamma/4$	$1/2$	$1/2$	0	0	0
4	0	0	0	0	0	0
5	$2\Delta_0 - i\Gamma/8$	$\frac{1}{2\sqrt{2}}$	$\frac{1}{2\sqrt{2}}$	$\sqrt{2}\Delta_0 - i\Gamma/4$	$1/2$	$1/2$

and $\phi_{n\nu}$ for $J_0(\frac{A}{\hbar\omega}) = 1$ and 0 are listed in Table I. The sum of eigenvalues is equal to the trace of $\overline{H}_{\text{II}}^{\text{rot}} + \Sigma$, i.e., $\sum_n \langle n | \overline{H}_{\text{II}}^{\text{rot}} + \Sigma | n \rangle = -i\Gamma$. Note that in Table I we only show the real part of $\phi_{n\nu}$ since $\Re(\phi_{n\nu}) \gg \Im(\phi_{n\nu})$.

According to Table I, when $J_0(\frac{A}{\hbar\omega}) = 1$, we can derive

$$G_{13}^{\text{R}}(0) = \sum_{\nu} \frac{\phi_{1\nu}\phi_{3\nu}}{-q_\nu} = \frac{-i\Gamma}{128\Delta_0^2 + \frac{\Gamma^2}{2}} \approx 0, \quad (\text{B9})$$

which indicates that $T(0) \approx 0$ and the time-average current is nearly zero in the absence of a laser field. This result is intriguing because (i) the state $\nu = 4$ has no contribution to transmission, and (ii) the synergistic effect of the states $\nu = 2$ and 3 leads to zero transmission. In other words, the existence of the states in a conductor

does not guarantee that they can play as conduction channels.

For $J_0(\frac{A}{\hbar\omega}) = 0$, we can derive

$$G_{13}^{\text{R}}(0) = \sum_{\nu} \frac{\phi_{1\nu}\phi_{3\nu}}{-q_\nu} = \frac{-i\Gamma}{16\Delta_0^2 + \frac{\Gamma^2}{2}} + \frac{i}{\Gamma} \approx \frac{i}{\Gamma}. \quad (\text{B10})$$

Substitute Eq. (B10) into Eqs. (B6) and (B7), one can derive $T(0) \approx 1$ and current reaches the maximum. In addition, as $J_0(\frac{A}{\hbar\omega}) = 0$, the dots $|4\rangle$ and $|5\rangle$ are decoupled from the other dots. Consequently, the synergistic effect of the states $\nu = 2$ and 3 vanishes, and the state $\nu = 2$ becomes an efficient conduction channel. This is the origin of CRT.

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