Optically induced shot noise enhancement and suppression in a Λ-type double-dot system

Tong-Guang Cheng, Hai-Feng Lü, and Yong Guo*
Department of Physics and Key Laboratory of Atomic and Molecular NanoSciences, Ministry of Education, Tsinghua University, Beijing 100084, People's Republic of China
(Received 19 November 2007; revised manuscript received 11 April 2008; published 6 May 2008)

In the sequential tunneling regime, we study the noise properties of a Λ-type double-dot system driven by an optical field. It is found that the zero-frequency shot noise can be tuned between sub- and super-Poissonian characteristics due to the dynamical channel blockade induced by the optical field. When the chemical potential is within the ground and the excited states in the right dot, the Fano factor is enhanced to a super-Poissonian value at a weak optical field but suppressed to a sub-Poissonian value at a strong optical field. The increase in the dot-dot hopping can lead to two peaks of the Fano factor when the absolute value of the detuned frequency approaches the strength of the dot-dot hopping. Moreover, the Fano factor at the two peaks approximate to 2 at a large dot-dot hopping.

DOI: 10.1103/PhysRevB.77.205305 PACS number(s): 73.63.Kv, 72.70.+m, 78.67.Hc

The study of shot noise has become an emerging topic in mesoscopic systems[1] since it characterizes the correlation between charge transport events and can provide more information than the mean current itself.[2–5] Especially, the shot noise's behavior in single-level and two-level quantum dots with electron bunching has both experimentally[6–8] and theoretically[9–13] received noticeable attention. These investigations indicate that a super-Poissonian shot noise results from a dynamical channel blockade, which depends not only on the Coulomb interaction but also on the coupling strengths between the dot system and the leads. Recently, the noise properties in a two-level quantum dot irradiated by an optical field were studied and the emergence of super-Poissonian noise in the dynamical channel blockade regime was reported.[14] It was found that the optical field can suppress the dynamical channel blockade and reduce the Fano factor from a super-Poissonian to sub-Poissonian value since it pumps electrons from a lower to an upper energy level.

Compared to a two-level single dot, the hopping between the ground states of a double dot in the Λ-type system (a three-level structure with two ground states and one excited state) provides another degree of freedom. In quantum optics, such three-level systems are usually used to realize coherent trapping and electromagnetically induced transparency (EIT).[15] These effects have been predicted in Λ-type dot systems[16,17] and can be measured through an optically controlled current.[18–20] An optically controlled current is promising for the design of devices that can apply technologies of optical manipulation in real atom systems to semiconductor materials. Recently, the current of a three-level system under the influence of an optical field is studied when the chemical potential is below two energy levels in the right dot.[21] The current shows two symmetric or asymmetric peaks when the field is resonant or nonresonant. A crossover from a three-level to a two-level system is found when the detuned frequency increases. Moreover, population inversion can be observed by varying the frequency or the strength of the optical field.

In this paper, we investigate the noise properties of a Λ-type double-dot system driven by an optical field. It is found that the effect of the optical field is quite different from that in a two-level single dot.[14] In our double-dot system, the dynamical channel blockade is induced by the optical field because the occupation of the excited state can occur only when the optical field exists. However, at a strong optical field, it will be suppressed by optical pumping. Therefore, the optical field can induce and suppress the dynamical channel blockade at the same time.

The system we consider (see Fig. 1) can be realized by two lateral quantum dots modulated by the anisotropic two-dimensional harmonic potentials. The large charging energy (~meV) forbids more than one electron in the system at one time. The ground states of the two dots are the lowest energy levels of the two harmonic potentials, respectively. The excited state of the left dot can be manipulated so that the energy difference between the excited states of the two dots ($\sim 10^{-1}$ meV) is much larger than the tunneling energy T_c (about 1 μeV).[17] Therefore, electrons are allowed to tunnel only between the two ground states. The energy difference in the right dot corresponds to a frequency of 100 GHz, which belongs to the microwave band. The transition between the ground state of the left dot and the excited state of the right dot by the optical field is neglected.[21] The total Hamiltonian of the system is $H=H_{Lead}+H_{\mathrm{DQD}}+H_T+H_I$, where

$$H_{Lead}=\sum_{k,\eta=L,R}\epsilon_{k\eta}c_{k\eta}^{\dagger}c_{k\eta},$$

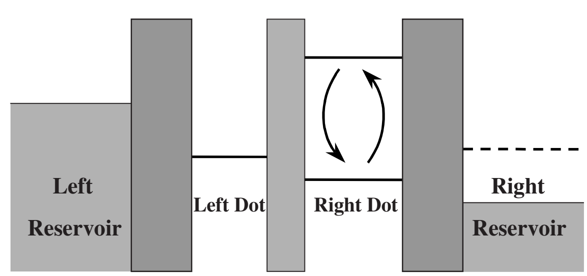

FIG. 1. Schematic view of a Λ-type double-dot system driven by an optical field (Ref. 21). The solid and the dashed horizontal lines indicate two regions where the chemical potential of the right lead is below and within the two levels in the right dot. The curved arrows indicate the driving direction of the optical field.

$$H_{\mathrm{DQD}} = E_L|L\rangle\langle L| + E_R|R\rangle\langle R| + E_e|E\rangle\langle E| + T_c(|L\rangle\langle R| + |R\rangle\langle L|),$$

$$H_I = -\frac{\Omega}{2}(|E\rangle\langle R|e^{-i\omega t}|R\rangle\langle E|e^{i\omega t}),$$

$$H_T = \sum_k [V_{kL}c_{kL}^\dagger|0\rangle\langle L| + V_{kR}c_{kR}^\dagger(|0\rangle\langle R| + |0\rangle\langle E|)] + \mathrm{H.c.} \tag{1}$$

Here, $|0\rangle$ represents the empty state (no electrons in both dots). $|L\rangle$, $|R\rangle$, and $|E\rangle$ correspond to one electron in the ground state of the left or the right dot and the excited state of the right dot, respectively. H_{Lead} describes the electron reservoirs (leads), H_{DQD} is the Hamiltonian of the dot system, and H_T represents the tunneling between the leads and the dot system. The interaction between the electrons and the optical field is expressed in H_I within the rotating-wave approximation,[15] where Ω indicates the coupling strength between the optical field and the quantum dot system. The energy difference between the ground states of the two dots, i.e., the band offset $\Delta_E \equiv E_L - E_R$, is modulated by the gate voltage. The detuned frequency $\Delta_\omega \equiv \omega - (E_e - E_R)$ can be tuned by the frequency of the optical field or the energy level of the excited state by manipulating the anisotropy of the potential of the right dot.

By using the method of quantum rate equation,[22–26] we obtain the stationary current and the zero-frequency shot noise. The states of the three-level system can be represented by a column $\mathbf{P} = (\rho_0, \rho_L, \rho_R, \rho_E, \mathrm{Re}[\rho_{LR}], \mathrm{Im}[\rho_{LR}], \mathrm{Re}[\rho_{LE}], \mathrm{Im}[\rho_{LE}], \mathrm{Re}[\rho_{RE}], \mathrm{Im}[\rho_{RE}],)^T$, of which the components are the density matrix elements of the quantum dot system,

$$\rho_i = \langle i|\rho|i\rangle, \quad i = 0, L, R, E,$$

$$\rho_{LR} = \langle L|\rho|R\rangle, \quad \rho_{LE} = \langle L|\rho|E\rangle, \quad \rho_{RE} = \langle R|\rho|E\rangle. \tag{2}$$

The rate equations for the density matrix elements are

$$\dot{\rho}_0 = -\sum_{i=L,R,E} \Gamma_i^+\rho_0 + \Gamma_i^-\rho_i,$$

$$\dot{\rho}_L = \Gamma_L^+\rho_0 - \Gamma_L^-\rho_L + iT_c(\rho_{LR} - \rho_{RL}),$$

$$\dot{\rho}_R = \Gamma_R^+\rho_0 - \Gamma_R^-\rho_R + iT_c(\rho_{RL} - \rho_{LR}) + 0.5i\Omega(\rho_{RE} - \rho_{ER}),$$

$$\dot{\rho}_E = \Gamma_E^+\rho_0 - \Gamma_E^-\rho_E + 0.5i\Omega(\rho_{ER} - \rho_{RE}),$$

$$\dot{\rho}_{LR} = iT_c(\rho_L - \rho_R) + 0.5i\Omega\rho_{LE} - \left[\frac{1}{2}(\Gamma_L^- + \Gamma_R^-) - i\Delta_E\right]\rho_{LR},$$

$$\dot{\rho}_{LE} = -\left[\frac{1}{2}(\Gamma_L^- + \Gamma_E^-) + i(\Delta_\omega - \Delta_E)\right]\rho_{LE} + 0.5i\Omega\rho_{LR} - iT_c\rho_{RE},$$

$$\dot{\rho}_{RE} = -\left[\frac{1}{2}(\Gamma_R^- + \Gamma_E^-) + i\Delta_\omega\right]\rho_{RE} + 0.5i\Omega(\rho_R - \rho_E) - iT_c\rho_{LE}, \tag{3}$$

where $\Gamma_\eta^\pm = \Gamma_\eta f_\eta^\pm(E_\eta)$, $\Gamma_\eta(\eta = L, R) = 2\pi\Sigma_k|V_{k\eta}|^2\delta(E_\eta - \epsilon_{k\eta})$, and $\Gamma_E = 2\pi\Sigma_k|V_{kR}|^2\delta(E_e - \epsilon_{kR})$ are the dot-lead couplings. $f_{L/R}^\pm(E) = \{1 + \exp[\pm(E - \mu_{L/R})/k_BT]\}^{-1}$ is the Fermi distribution function in the left or the right reservoir, $\pm$ represents the population of the occupied or the empty state of the leads. $\mu_{L/R}$ ($\mu_L = -\mu_R = eV/2$) is the chemical potential of the left or the right lead with V, the voltage of the bias. We set $\Gamma_L = \Gamma_R = \Gamma_E = \Gamma$ as the unit of energy. Equation (3) can be rewritten in the following matrix form:[26]

$$\frac{d\mathbf{P}}{dt} = \mathbf{MP}(t). \tag{4}$$

Therefore, the statistical averaging of any time-dependent operator $\hat{A}(t)$ becomes

$$\langle \hat{A}(t)\rangle = \mathrm{Tr}[\hat{\Lambda}\rho] = \sum_k [\mathbf{AP}(t)]_k = \sum_k [\mathbf{A}(t)\mathbf{P}^{(0)}]_k = \sum_k [\mathbf{AT}(t)\mathbf{P}^{(0)}]_k, \tag{5}$$

where $\mathbf{A}$ is the matrix form of the operator $\hat{A}$, $\mathbf{P}^{(0)}$ is the steady state solution of Eq. (3), and $\mathbf{T}(t) = \exp(\mathbf{M}t)$ governs the time evolution of $\mathbf{P}$. It is well known that the noise power spectra can be expressed as the Fourier transformation of the current-current correlation function

$$S_{\eta\eta'}(\omega) = 2\langle \hat{I}_\eta(t)\hat{I}_{\eta'}(0)\rangle_\omega - 2\langle \hat{I}_\eta\rangle_\omega\langle \hat{I}_{\eta'}\rangle_\omega. \tag{6}$$

Moreover, the current-current correlation function can be expressed in the $\mathbf{P}$ representation as follows:

$$\langle \hat{I}_\eta(t)\hat{I}_{\eta'}(0)\rangle = \theta(t)\sum_k [\mathbf{\Gamma}_\eta\mathbf{T}(t)\mathbf{\Gamma}_{\eta'}\mathbf{P}^{(0)}]_k + \theta(-t)\sum_k [\mathbf{\Gamma}_{\eta'}\mathbf{T}(-t)\mathbf{\Gamma}_\eta\mathbf{P}^{(0)}]_k, \tag{7}$$

where $\mathbf{\Gamma}_\eta$ is the matrix form of the current operator.[27] Finally, the current-current correlation in the ω space is

$$\langle \hat{I}_\eta(t)\hat{I}_{\eta'}(0)\rangle_\omega = \sum_{\lambda,k}\left[\frac{\mathbf{\Gamma}_\eta\mathbf{P}_\lambda\mathbf{\Gamma}_{\eta'}}{-i\omega - \lambda}\mathbf{P}^{(0)}\right]_k + \left[\frac{\mathbf{\Gamma}_{\eta'}\mathbf{P}_\lambda\mathbf{\Gamma}_\eta}{i\omega - \lambda}\mathbf{P}^{(0)}\right]_k. \tag{8}$$

Here, $\mathbf{P}_\lambda$ is a projector operator associated with the eigenvalue λ and $\mathbf{M} = \Sigma_\lambda \lambda\mathbf{P}_\lambda$.

Noting that the summation over k of the zero eigenvalue is exactly canceled by the term $\langle \hat{I}_\eta\rangle\langle \hat{I}_{\eta'}\rangle$, the eventual expression for the noise power spectrum is

$$S_{\eta\eta'}(\omega) = \delta_{\eta\eta'}S_\eta^{\mathrm{Sch}} + 2\sum_{k,\lambda\neq 0}\left\{\frac{[\mathbf{\Gamma}_\eta\mathbf{P}_\lambda\mathbf{\Gamma}_{\eta'}\mathbf{P}^{(0)}]_k}{-i\omega - \lambda} + \frac{[\mathbf{\Gamma}_{\eta'}\mathbf{P}_\lambda\mathbf{\Gamma}_\eta\mathbf{P}^{(0)}]_k}{i\omega - \lambda}\right\}, \tag{9}$$

where $S_\eta^{\mathrm{Sch}} = \Sigma_k|[\mathbf{\Gamma}_\eta\mathbf{P}^{(0)}]_k|$ is the frequency-independent Schottky noise originating from the self-correlation of a given tunneling event with itself, which the double-time correlation function Eq. (7) cannot contain. Here, we focus on the zero-frequency shot noise. Due to current conservation, in the two-terminal system, $S_{LL}(0) = S_{RR}(0) = S_{LR}(0) = S_{RL}(0)$.[28] We use the expression $S = (1/4)[S_{LL}(0) + S_{RR}(0)$

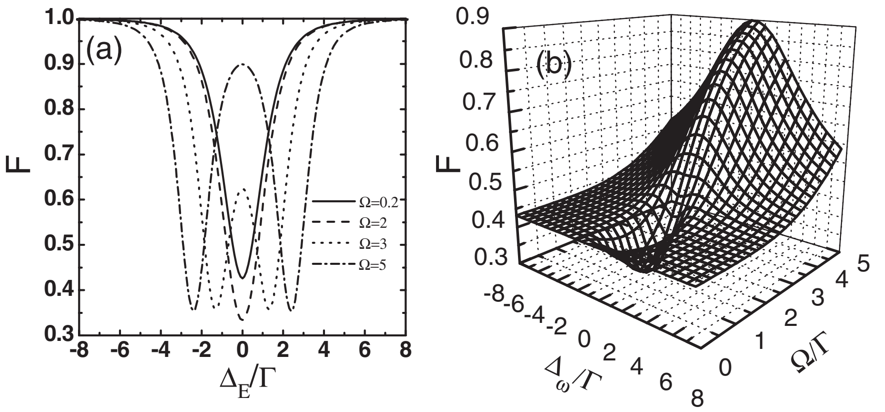

FIG. 2. (a) The Fano factor F as a function of the band offset Δ_E at different Rabi frequencies when the detuned frequency Δ_ω equals zero. (b) The Fano factor F as a function of the detuned frequency Δ_ω at different Rabi frequencies when the band offset Δ_E equals zero. $T_c=\Gamma$.

$+S_{LR}(0)+S_{RL}(0)]$, and the Fano factor is defined as $F=S/2eI$.

In the following discussion, we focus on two regions where the chemical potential of the right lead μ_R is below and within the two energy levels in the right dot. The chemical potential of the left lead is set to be far beyond the energy level of the left dot, i.e., $\mu_L-E_L\gg k_BT$ and $f_L^+(E_L)\approx 1$. Since the energy spacing in the right dot reaches the scale of 0.1 meV, the thermal width ($\sim k_BT$) of the Fermi distribution function can be neglected. Hence, $f^+(E_e)$ approaches zero whether μ_R is below or within the two energy levels in the right dot. Moreover, the value of $f_R^+(E_R)$ mostly depends on whether the chemical potential of the right lead is below or beyond the ground state energy E_R. Therefore, when $\mu_R<E_R<E_e$, $f_R^+(E_R)\approx f_R^+(E_e)\approx 0$, and when $E_R<\mu_R<E_e$, $f_R^+(E_R)\approx 1$ while $f_R^+(E_e)\approx 0$.

$\mu_R<E_R<E_e$: In this regime, $f_R^+(E_R)\approx f_R^+(E_e)\approx 0$. Figure 2(a) shows the Fano factor as a function of the band offset at different Rabi frequencies when the optical field is at the resonance case. With an increase in Ω, the Fano factor undergoes a transition from a global minimum to a local maximum at $\Delta_E=0$. When the coupling strength with the optical field Ω is smaller than both the dot-dot hopping T_c and the dot-lead coupling Γ, the global minimum of the Fano factor locates at $\Delta_E=0$. When Ω increases, two symmetric valleys of the Fano factor occur where the band offset roughly equals one-half the value of Ω. For the case of a strong optical field, the value of the peak between the valleys enhances and approaches a Poissonian value. The influence of the detuned frequency is depicted in Fig. 2(b). Different from the band offset, the Fano factor monotonously increases or decreases with an increase in $|\Delta_\omega|$ at a weak or a strong optical field, respectively.

The competition among the coupling strength Ω, the dot-dot hopping T_c, and the tunneling rate Γ results from the correlation of two transport channels for electrons to flow through the system. One is along the left lead $\rightarrow E_L\rightarrow E_R\rightarrow$ the right lead, and the other is along the left lead $\rightarrow E_L\rightarrow E_R\rightarrow E_0\rightarrow$ the right lead. In the limit of the nondriving case ($\Omega=0$), the excited state in the right dot does not contribute to the transport, and the three-level system reduces to two coupled single-level quantum dots with a dot-dot hopping T_c. Hence, the Fano factor becomes[26]

$$F=1-\frac{8T_c^2\Gamma_L\Gamma_R(3\Gamma_L\Gamma_R+\Gamma_R^2+T_c^2)}{[4T_c^2(2\Gamma_L+\Gamma_R)+\Gamma_L\Gamma_R^2]^2}. \tag{10}$$

At a low coupling strength Ω, the transport through the excited state is much weaker than the one through the ground state in the right dot. In this situation, the population of the excited state decays faster than that of the ground state; therefore, the Fano factor is most strongly suppressed when there is no band offset. With an increase in Ω, the channel through the excited state in the right dot gradually contributes to the transport. Since Coulomb interaction prevents a double occupancy, the two paths have a negative correlation with each other.[23] Therefore, when the two transport channels have the same contribution, the Fano factor is suppressed to its minimal value. With respect to that, the current is proportional to the population in the right dot, the transition from a global minimum to a local maximum of the Fano factor roughly occurs when the population inversion arises (see Fig. 3).

When the strength of the optical field increases, the population of the excited state will exceed the population of the ground state until the band offset approximately reaches one-half of Ω. Therefore, the Fano factor is suppressed to its minimum due to the same contribution of the two transport channels when the band offset roughly equals one-half of Ω.

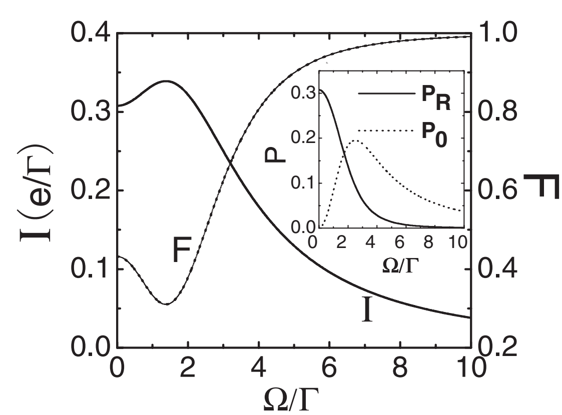

FIG. 3. The current I (solid line) and the Fano factor F (dotted line) as a function of Ω when $\Delta_E=\Delta_\omega=0$ and $T_c=\Gamma$. Inset: Population in the right dot as a function of Ω.

If Ω is far beyond both T_c and Γ, photon-assisted Rabi oscillations dominate the tunneling processes and lead to a Poissonian noise. The reason is that fast oscillations in the right dot will result in the randomness of tunneling events. Moreover, when the intensity of the optical field is fixed, as the band offset increases and exceeds far beyond 0.5 Ω, the system will gradually approach the coupled single-level dots due to the faster decay rate of the population of the excited state than that of the ground state. Hence, the Fano factor will approach a Poissonian value at a large band offset. The influence of the detuned frequency is different from the band offset. The reason is that it can transfer the population from the excited state to the ground state, whereas the band offset suppresses the population of both. The effect of the optical field decays with an increase in the detuned frequency, and the Fano factor approaches the value of two coupled single-level dots in series (see Fig. 3).

$E_R < \mu_R < E_e$: In this region, $f_R^+(E_R) \approx 1$ while $f_R^+(E_e) \approx 0$, and a dynamical channel blockade occurs for a two-level system.[14] Different from the case of a single two-level dot, in our considered system, the dynamical channel blockade not only results from the Coulomb interaction but is also induced by the optical field. In the absence of an optical field, the transport can only occur due to the thermal excitation through the ground state of the right dot, leading to a Poissonian shot noise. When Ω is smaller than both T_c and Γ, the occupation of the ground state blocks the transport through the excited state, therefore, the dynamical channel blockade occurs,[14,22,25] which leads to a super-Poissonian noise. With an increase in Ω, the pumping that drives electrons from the ground state to the excited state suppresses the dynamical channel blockade, thus resulting in a reduction in the Fano factor. When Ω becomes sufficiently large, the Fano factor can be suppressed to a sub-Poissonian value.[14]

Figures 4(a) and 4(b) indicate the influence of the band offset and the detuned frequency on the Fano factor at different coupling strengths under the influence of an optical field. It is found that the Fano factor shows a minimum at the point with $\Delta_\omega=0$ or $\Delta_E=0$ and have two symmetric peaks when Δ_ω or Δ_E is nonzero. Since the band offset and the detuned frequency can increase the population of the ground state relative to that of the excited state, the super-Poissonian peaks of the Fano factor occur due to strengthening of the dynamical channel blockade. If either the band offset or the detuned frequency becomes even larger, the transport through the excited state is severely suppressed, leading to a Poissonian noise because the current is restricted to occur only through thermal excitation. With an increase in Ω, the positions of the two peaks gradually move from the origin point and their values rapidly decay. When Ω is large enough, no apparent peaks appear and the Fano factor is suppressed to be sub-Poissonian.

Now, we investigate the Fano factor as a function of the detuned frequency at different dot-dot hopping T_c. The result is shown in Fig. 5. When T_c is smaller than Γ, the two peaks approximately locate at $\Delta_\omega = \pm\Gamma$. When T_c exceeds the value of Γ, the two peaks almost exit at $\Delta_\omega = \pm T_c$. The value of the peaks increase with T_c and have an asymptotic value of 2. The behavior of the Fano factor can be explained as follows: At a small T_c, the dot-lead coupling Γ dominates the evolution of the density matrix of the double-dot system $\rho \sim \exp(Mt)$.[26] The dot-lead couplings contribute a decay rate of Γ^{-1}, which results in the width of energy levels ($\sim\Gamma$) due to the uncertainty relation. If the detuned frequency is within this width, the Fano factor will increase with the detuned frequency since the possibility of pumping electrons from the ground state to the excited state is suppressed, and at the same time the dynamical blockade is strengthened. However, if the detuned frequency exceeds this width, the Fano factor will trend to a Poissonian value. The reason is that the cur-

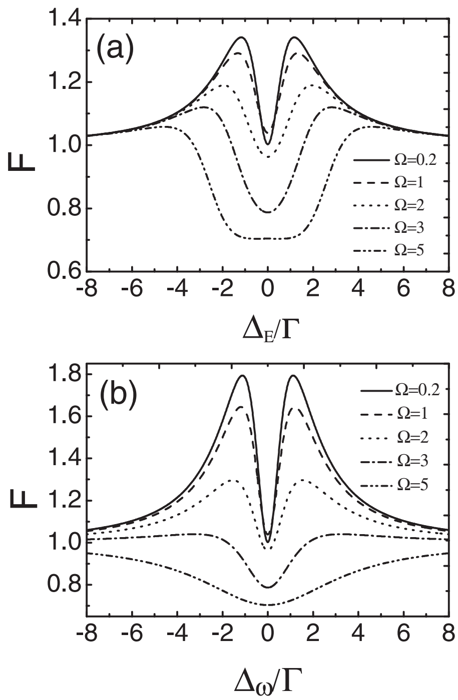

FIG. 4. (a) The Fano factor F as a function of the band offset Δ_E at different Rabi frequencies when the detuned frequency Δ_ω equals zero. (b) The Fano factor F as a function of the detuned frequency Δ_ω at different Rabi frequencies when the band offset Δ_E equals zero. $T_c=\Gamma$.

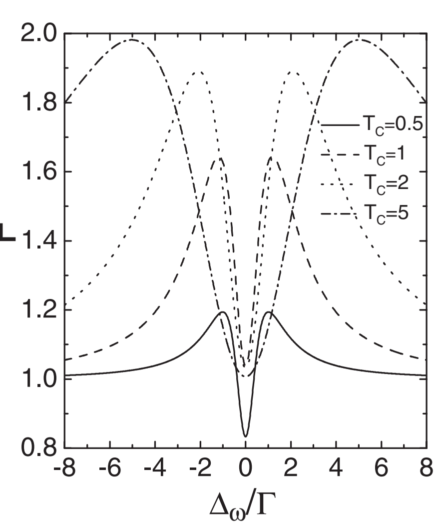

FIG. 5. The Fano factor F as a function of the detuned frequency Δ_ω at different dot-dot hopping T_c. $\Delta_E=0$ and $\Omega=\Gamma$.

rent will be restricted to occur only through thermal excitation. Consequently, the peaks of the Fano factor approximately locate at $\Delta_\omega = \pm \Gamma$. With an increase in T_c, electrons gradually localize in the two ground states, and the coherent tunneling between two dots dominates the evolution of the density matrix. Due to the coherent tunneling, there exist two paths of pumping electrons from the ground states to the excited state, which can be seen from the nondiagonal density matrix elements ρ_{EL} and ρ_{ER}. Although the ground state of the left dot does not have a direct transition to the excited state, the electrons in it can be indirectly pumped to the excited state by first hopping to the right dot. If $\Delta_\omega=0$, the two paths are distinguishable since they are apart from each other in the time domain. When Δ_ω approaches $\pm T_c$, the two paths incline to simultaneously pump the electrons to the excited state. Although the current is strongly suppressed due to the localization in the ground states, electrons tend to reach and leave the excited state in pairs within the time interval of empty ground states. Therefore, the Fano factor asymptotically approaches 2. If Δ_ω becomes even larger, the two paths resume to be apart in the time domain and are more greatly suppressed, which leads to a Poissonian Fano factor.

In conclusion, we study the shot noise behavior in a Λ-type double-dot system driven by an optical field. The super-Poissonian Fano factor is found, which can be explained by the dynamical channel blockade. Different from the case in a single two-level quantum dot, it is the optical field that induces the dynamical channel blockade. The Fano factor is enhanced when the coupling strength with the optical field is weak but suppressed for a strong optical field. The results also indicate that the Fano factor has a strong dependence on dot-dot hopping T_c. When T_c is large enough, the Fano factor shows two peaks approximately located at $\Delta_\omega = \pm T_c$, which can be used as a probe for the dot-dot hopping of strong coupling dots. It is further revealed that with an increase in T_c, the values of the peaks approaches 2. Our study on the optically induced dynamical channel blockade indicates a regime wherein the relationship between Coulomb blockade and the coherent pumping quite differs from that in a two-level system.

This work was supported by the National Natural Science Foundation of China (Grant No. 10774083), Tsinghua Basic Research Foundation (JCpy2005058), the Specialized Research Fund for the Doctoral Program of Higher Education (Grant No. 2006003047), and the 973 Program (Grant No. 2006CB605105).

*Corresponding author; guoy66@tsinghua.edu.cn

[1] Ya. M. Blanter and M. Büttiker, Phys. Rep. **336**, 1 (2000).
[2] R. Zhu and Y. Guo, Appl. Phys. Lett. **90**, 232104 (2007); J. Appl. Phys. **102**, 083706 (2007).
[3] L. Qin, Y. Guo, and J. L. Zhu, Phys. Lett. A **366**, 124 (2007).
[4] G. Iannaccone, G. Lombardi, M. Macucci, and B. Pellegrini, Phys. Rev. Lett. **80**, 1054 (1998).
[5] E. Onac, F. Balestro, B. Trauzettel, C. F. J. Lodewijk, and L. P. Kouwenhoven, Phys. Rev. Lett. **96**, 026803 (2006).
[6] P. Barthold, F. Hohls, N. Maire, K. Pierz, and R. J. Haug, Phys. Rev. Lett. **96**, 246804 (2006).
[7] S. S. Safonov, A. K. Savchenko, D. A. Bagrets, O. N. Jouravlev, Y. V. Nazarov, E. H. Linfield, and D. A. Ritchie, Phys. Rev. Lett. **91**, 136801 (2003).
[8] S. Gustavsson, R. Leturcq, B. Simovič, R. Schleser, P. Studerus, T. Ihn, K. Ensslin, D. C. Driscoll, and A. C. Gossard, Phys. Rev. B **74**, 195305 (2006).
[9] Ya. M. Blanter and M. Büttiker, Phys. Rev. B **59**, 10217 (1999).
[10] A. Thielmann, M. H. Hettler, J. König, and G. Schön, Phys. Rev. B **71**, 045341 (2005).
[11] A. Cottet, W. Belzig, and C. Bruder, Phys. Rev. B **70**, 115315 (2004).
[12] W. Belzig, Phys. Rev. B **71**, 161301(R) (2005).
[13] J. Aghassi, Axel Thielmann, M. H. Hettler, and G. Schön, Phys. Rev. B **73**, 195323 (2006); F. Bodoky, W. Belzig, and C. Bruder, *ibid.* **77**, 035302 (2008).
[14] R. Sánchez, G. Platero, and T. Brandes, Phys. Rev. Lett. **98**, 146805 (2007).
[15] Marlan O. Scully and M. Suhail Zubairy,*Quantum Optics* (Cambridge University Press, Cambridge, 1997).
[16] T. Brandes and F. Renzoni, Phys. Rev. Lett. **85**, 4148 (2000).
[17] W. D. Chu, S. Q. Duan, and J. L. Zhu, Appl. Phys. Lett. **90**, 222102 (2007).
[18] T. H. Oosterkamp, T. Fujisawa, W. G. van der Wiel, K. Ishibashi, R. V. Hijman, S. Tarucha, and L. P. Kouwenhoven, Nature (London) **395**, 873 (1998).
[19] H. Qin, A. W. Holleitner, K. Eberl, and R. H. Blick, Phys. Rev. B **64**, 241302(R) (2001).
[20] J. R. Petta, A. C. Johnson, C. M. Marcus, M. P. Hanson, and A. C. Gossard, Phys. Rev. Lett. **93**, 186802 (2004).
[21] Y. Y. Liao, D. S. Chuu, and Y. N. Chen, Phys. Rev. B **75**, 125325 (2007).
[22] G. Kießlich, A. Wacker, and E. Schöll, Phys. Rev. B **68**, 125320 (2003).
[23] I. Djuric, B. Dong, and H. L. Cui, IEEE Trans. Nanotechnol. **4**, 71 (2004).
[24] B. Dong, H. L. Cui, and X. L. Lei, Phys. Rev. B **69**, 035324 (2004).
[25] I. Djuric, B. Dong, and H. L. Cui, Appl. Phys. Lett. **87**, 032105 (2005).
[26] I. Djuric, B. Dong, and H. L. Cui, J. Appl. Phys. **99**, 063710 (2006).
[27] S. Hershfield, J. H. Davies, P. Hyldgaard, C. J. Stanton, and J. W. Wilkins, Phys. Rev. B **47**, 1967 (1993).
[28] Here, we use the same convention of the sign in the expression of the current as that of Ref. 26, in which the sign is absorbed in the current operator $\mathbf{\Gamma}$ to keep the stationary current for both leads (I_L and I_R) always positive.