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Shot noise properties of electron transport through an interacting multi-terminal quantum dots system

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Abstract
We study the correlations of tunneling currents through an interacting quantum dots (QDs) system composed of a top single QD and a bottom qubit with purely capacitive coupling within a quantum master approach. We find that the super-Poissonian current noise of the qubit near resonance, which is a signature of coherent tunneling within the transport qubit for asymmetrical contact couplings, is strongly dependent on non-equilibrium transport through the top QD with different coupling configurations. For pure-dephasing coupling, such a super-Poissonian feature is asymmetrically washed out by increasing coupling strength showing obvious qubit level position dependence with finite bias and temperature, while for orthogonal coupling we can almost symmetrically lower the double peak to a double minimum by increasing coupling strength or adjusting the ratio of the top QD contact couplings in the large bias limit, indicating the transition from coherent tunneling to sequential tunneling.

(Some figures in this article are in colour only in the electronic version)

1. Introduction

Current fluctuations in mesoscopic systems provide much information about electron correlations not available by average current measurements [1, 2]. This has led to intensive experimental [3–8] and theoretical [9–18] studies on the current noise and even higher moment of the current correlator in a variety of open quantum systems, revealing the interplay between Fermi statistics, coherence and Coulomb interaction, which is crucial in solid state quantum information processing.

One such system is quantum dots in a Coulomb blockade regime with multi-terminals, especially involving two or more separate QDs interacting via long range Coulomb forces. Electron transport through one terminal is correlated with transfer events in another terminal nearby. The corresponding negative or positive correlations are accessed by shot noise and cross-correlation measurements, indicating the role of various dissipation mechanisms. However, large adjustable capacitive coupling strength comparative to other system parameters, e.g. intra or interdot charge energy in one conductor, is the key in such experiments. Recently, several groups addressed this problem in different QD configurations [19–21], where a floating metallic top gate provides strong purely capacitive coupling with high controllability. This makes it possible to study two-electron correlation such as non-equilibrium electron transport effects on currents through other dots, two qubit operation in semiconductor QDs and current cross-correlation between two parallel QDs in a controllable way.

In this work, we study the non-equilibrium environment effect on electron statistics of a qubit current. We present a counting statistics formulism based on a Markovian quantum master equation (QME) with lowest order perturbation in contact tunnelings following the line of [25–30]. In this approach, all Coulomb interactions in the subsystem, a transport qubit capacitively coupling to a single dot, are treated nonperturbatively within a many-particle basis. As a signature of coherent tunneling, the super-Poissonian current noise of the qubit in the vicinity of resonance [6, 23, 24] is

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strongly affected by electron transport through the top QD via capacitive coupling. For pure-dephasing coupling, the double peak feature is asymmetrically washed out related to qubit level positions with finite bias and temperature; for orthogonal coupling, because of the same energy absorption and emission rate due to charge fluctuation of the top QD, we can almost symmetrically lower the double peak to a double minimum by increasing the coupling strength or adjusting the relative value of the two top QD tunneling rates.

2. Physical model and counting statistics formulation

We consider the system without a spin degree of freedom for simplicity as sketched in figure 1, which consists of an electrically isolated top single QD and bottom qubit, attached to source and drain leads respectively. Electrons flowing from the bottom source to the bottom drain through the qubit are affected by the top current via long range Coulomb interaction.

This model is described by the Hamiltonian $\hat{\mathcal{H}} = \hat{\mathcal{H}}_b + \hat{\mathcal{H}}_t + \hat{\mathcal{H}}_{leads} + \hat{\mathcal{H}}_{tun} + \hat{\mathcal{H}}_{int}$. Here, the first three terms on the right represent the bottom qubit, top single QD and lead Hamiltonians respectively, given by

$$\hat{\mathcal{H}}_b = \varepsilon_{bl}\hat{a}^\dagger_{bl}\hat{a}_{bl} + \varepsilon_{br}\hat{a}^\dagger_{br}\hat{a}_{br} + (\Omega\hat{a}^\dagger_{br}\hat{a}_{bl} + \text{h.c.}) + U\hat{a}^\dagger_{bl}\hat{a}_{bl}\hat{a}^\dagger_{br}\hat{a}_{br}, \quad (1)$$

$$\hat{\mathcal{H}}_t = \varepsilon_t\hat{a}^\dagger_t\hat{a}_t, \quad (2)$$

$$\hat{\mathcal{H}}_{leads} = \sum_{\alpha,k}\varepsilon_{\alpha k}\hat{c}^\dagger_{\alpha k}\hat{c}_{\alpha k} \quad (3)$$

where $\hat{a}^\dagger_i(\hat{a}_i)$ is the creation (annihilation) operator of an electron with energy ε_i inside the QD $i \in$ bl, br, t and $\hat{c}^\dagger_{\alpha k}(\hat{c}_{\alpha k})$ is the same operator for a noninteracting electron with energy $\varepsilon_{\alpha k}$ in leads $\alpha \in$ bl, br, tl, tr. Ω is the interdot tunneling coupling and U is the interdot charge energy for occupying both dots.

The Hamiltonian of tunneling coupling between leads and dots reads

$$\hat{\mathcal{H}}_{tun} = \sum_{k;\alpha=br,bl}(t_{\alpha k}\hat{a}^\dagger_\alpha\hat{c}_{\alpha k} + \text{h.c.}) + \sum_{k;\alpha=tr,tl}(t_{\alpha k}\hat{a}^\dagger_t\hat{c}_{\alpha k} + \text{h.c.}) \quad (4)$$

with $t_\alpha(\alpha = \text{br, bl, tr, tl})$ the tunneling amplitude, whose influence can be fully characterized by the spectral density $\Gamma_\alpha(\varepsilon) = 2\pi\sum_k |t_{\alpha k}|^2\delta(\varepsilon - \varepsilon_k)$. In the wide-band limit, this spectral density Γ_α is independent of energy, which implies constant tunneling amplitudes and density of states n_α in leads α.

The Coulomb interaction between top QD and bottom qubit has the bilinear form $\hat{\mathcal{H}}_{int} = \frac{g}{2}\hat{n}_t \cdot \hat{\sigma}$, where g is the coupling strength, $\hat{n}_t$ is the charge number operator of the top QD, and $\hat{\sigma}$ is a bottom qubit operator. In the following, we adopt a basis of many-particle states for the electrically isolated QDs and set the average energy of the left and right bottom QDs as the zero energy point of the total system. By diagonalizing the bottom qubit and assuming large interdot Coulomb repulsion so that double occupancy of the bottom QDs is impossible, we get the bonding and antibonding states with energy E_{b1}, E_{b2} (in our setting, $E_{b1} = -E_{b2}$) and its creation (annihilation) operator $\hat{\Psi}^\dagger_i(\hat{\Psi}_i)$. Thus, the operator $\hat{\sigma}$ can be expressed as $\hat{\sigma}_z = \hat{\Psi}^\dagger_{b2}\hat{\Psi}_{b2} - \hat{\Psi}^\dagger_{b1}\hat{\Psi}_{b1}$ for pure-dephasing coupling and $\hat{\sigma}_x = \hat{\Psi}^\dagger_{b2}\hat{\Psi}_{b1} + \hat{\Psi}^\dagger_{b1}\hat{\Psi}_{b2}$ for orthogonal coupling [31, 32]. Charge number fluctuations of the top QD induce only pure dephasing in the qubit since $\hat{\sigma}_z$ commutes with $\hat{\mathcal{H}}_b$, while orthogonal coupling $\hat{\sigma}_x$, which is noncommutative with $\hat{\mathcal{H}}_b$, causes energy relaxation in the qubit.

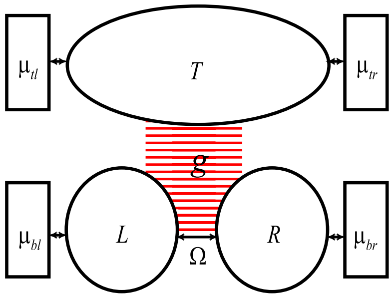

Figure 1. Sketch of three quantum dot systems: the top single QD (labeled T) and bottom double dot (labeled L and R) with mutual coupling Ω are weakly connected to four leads of Fermi energy $\mu_\alpha(\alpha = \text{tl, tr, bl, br})$, coupled via long range Coulomb interaction g.

In this study, the top QD and bottom qubit constitute the subsystem of interest, whose Hamiltonian is $\hat{\mathcal{H}}_s = \hat{\mathcal{H}}_b + \hat{\mathcal{H}}_t + \hat{\mathcal{H}}_{int}$, and all the Coulomb interaction terms are fully taken into account, while tunneling between dots and leads is treated perturbatively to the lowest order. This weak tunneling approximation referring to sequential tunneling is reasonable when Γ_α is much smaller than voltage bias eV or temperature $k_B T$ [33]. By tracing out the leads degree of freedom, the quantum master equation for the subsystem density operator $\hat{\rho}_s$ to the second order perturbation theory reads [33–35]

$$\frac{d}{dt}\hat{\rho}_s(t) = -i\check{\mathcal{L}}_s\hat{\rho}_s(t) - \int_{t0}^{t} dt'\,\check{\mathcal{L}}_{tun}G_s(t,t')G_{leads}(t,t')\check{\mathcal{L}}_{tun}\hat{\rho}_s(t') \quad (5)$$

where the Liouville operators are $\check{\mathcal{L}}_s = [\hat{\mathcal{H}}_s, \cdot]$ and $\check{\mathcal{L}}_{tun} = [\hat{\mathcal{H}}_{tun}, \cdot]$, and propagators are defined as $G_s(t,t') = \exp[-i\check{\mathcal{L}}_s(t-t')]$ and $G_{leads}(t,t') = \exp[-i\check{\mathcal{L}}_{leads}(t-t')]$. We set $e = \hbar = 1$ and $n_\alpha = \pi^{-1}$.

Since our interest is stationary state properties and only sequential tunneling processes are considered, we can apply the Markov approximation which means replacing $\hat{\rho}_s(t')$ with $\hat{\rho}_s(t)$ and setting $t' \to -\infty$ in equation (5). In order to get a compact form, we introduce the following left and right action superoperators as done in [29]: $\check{\Psi}^{(L)}\cdot = \hat{\Psi}\cdot$ and $\check{\Psi}^{(R)}\cdot = \cdot\hat{\Psi}$. Then we get

$$\dot{\hat{\rho}}_s = \check{M}\hat{\rho}_s \quad (6)$$

where the generator of QME $\check{M}$ is

$$\check{M} = -i\check{\mathcal{L}}_s + \sum_{\alpha=bl,br,tl,tr}(-\check{\Pi}^\alpha + \check{\Sigma}^\alpha_+ + \check{\Sigma}^\alpha_-). \quad (7)$$

$\check{\mathcal{L}}_s$ describes the isolated coherent subsystem dynamics and the dissipation term is split into a non-diagonal part leaving the number of electrons inside the subsystem unchanged,

$$\check{\Pi}^{\alpha} = \sum_i (\check{\Psi}_i^{\dagger(\mathrm{L})}\check{\Psi}_{i,\alpha}^{(+,\mathrm{L})} + \check{\Psi}_i^{\dagger(\mathrm{R})}\check{\Psi}_{i,\alpha}^{(-,\mathrm{R})} + \mathrm{h.c.}) \tag{8}$$

and two diagonal parts responsible for increasing or decreasing the number of electrons in the subsystem

$$\begin{aligned} \check{\Sigma}_{+}^{\alpha} &= \sum_i (\check{\Psi}_i^{\dagger(\mathrm{L})}\check{\Psi}_{i,\alpha}^{(-,\mathrm{R})} + \check{\Psi}_i^{(\mathrm{R})}\check{\Psi}_{i,\alpha}^{\dagger(-,\mathrm{L})}), \\ \check{\Sigma}_{-}^{\alpha} &= \sum_i (\check{\Psi}_i^{\dagger(\mathrm{R})}\check{\Psi}_{i,\alpha}^{(+,\mathrm{L})} + \check{\Psi}_i^{(\mathrm{L})}\check{\Psi}_{i,\alpha}^{\dagger(+,\mathrm{R})}). \end{aligned} \tag{9}$$

The auxiliary annihilation operator is given by

$$\check{\Psi}_{i,\alpha}^{(\pm)} = \sum_j T_{\alpha,i}^{*} T_{\alpha,j} f_{\alpha}^{(\pm)}(\check{\mathcal{L}}_{\mathrm{s}})\Psi_j \tag{10}$$

where $f_{\alpha}^{(+)}(\check{\mathcal{L}}_{\mathrm{s}}) = 1 - f_{\alpha}^{(-)}(\check{\mathcal{L}}_{\mathrm{s}}) = (\exp(\check{\mathcal{L}}_{\mathrm{s}} - \mu_\alpha)/k_{\mathrm{B}}T + 1)^{-1}$ is the Fermi distribution function with μ_α the Fermi energy of lead α, and the tunneling matrix element $T_{\alpha,i}$ can be got in the many-particle basis of separate top or bottom dots. Here, we neglect the level renormalization contributions which only change the bare Bohr frequencies of the subsystem.

In order to get full information about transport through the subsystem, we study the probability distribution $P(\vec{n}, t)$ that $\vec{n}$ electron transfers are measured during a time interval t, which is related to the generating function (GF) $G(\vec{\chi}, t)$ and the cumulant generating functions (CGF) $S(\vec{\chi}, t)$ as

$$G(\vec{\chi}, t) = \mathrm{e}^{-S(\vec{\chi},t)} = \sum_{\vec{n}} P(\vec{n}, t)\mathrm{e}^{\mathrm{i}\vec{n}\cdot\vec{\chi}} \tag{11}$$

where $\vec{n}$ is a vector with component n_α^{+} or n_α^{-} (the number of electrons transferring into or out of the subsystem through lead α) and $\vec{\chi}$ is the corresponding counting field vector. The GF is obtained by tracing the generate operator (GO) $\hat{g}(\vec{\chi}, t)$ over the subsystem degree of freedom $G(\vec{\chi}, t) = \mathrm{Tr}_{\mathrm{s}}\{\hat{g}(\vec{\chi}, t)\}$, while the GO satisfies the evolution equation derived from equation (6) in the trajectory picture [25, 27]:

$$\begin{aligned} \frac{\mathrm{d}}{\mathrm{d}t}\hat{g}(\vec{\chi}, t) &= \check{\mathcal{W}}(\vec{\chi})\hat{g}(\vec{\chi}, t), \\ \check{\mathcal{W}}(\vec{\chi}) &= -\mathrm{i}\check{\mathcal{L}}_{\mathrm{s}} + \sum_{\alpha=\mathrm{bl,br,tl,tr}} (-\check{\Pi}^{\alpha} + \mathrm{e}^{\mathrm{i}\chi_\alpha^{+}}\check{\Sigma}_{+}^{\alpha} + \mathrm{e}^{\mathrm{i}\chi_\alpha^{-}}\check{\Sigma}_{-}^{\alpha}). \end{aligned} \tag{12}$$

Counting begins after subsystems have reached the stationary state, i.e. $\hat{g}(\vec{\chi}, t = 0) = g^{\mathrm{st}}$. From the CGF we can obtain all the (zero-frequency) cumulants $C_n = -\frac{1}{t}(-\mathrm{i}\partial_{\vec{\chi}})^n S(\vec{\chi}, t)|_{\vec{\chi}\to 0, t\to\infty}, n = 1, 2\ldots$. If only one certain lead is concerned, e.g. bottom right lead with counting fields $\vec{\chi} := (\chi_{\mathrm{br}}^{+} = -\chi, \chi_{\mathrm{br}}^{-} = \chi)$, the first three cumulants are related to the average current, the current noise and the skewness, which characterize the distribution, e.g. a Fano factor $F = C_2/C_1$ bigger (smaller) than 1 means super-(sub-)Poissonian noise. Cross-correlations can also be obtained with multi-type counting fields.

3. Results and analyses of fluctuating environment effects

For numerical calculation, we express all the operators in the many-particle basis of the subsystem. A general subsystem state is written as $|ij\rangle = |i\rangle_{\mathrm{b}} \otimes |j\rangle_{\mathrm{t}}$, with $|i\rangle_{\mathrm{b}} \in |0\rangle_{\mathrm{b}}, |+\rangle_{\mathrm{b}}, |-\rangle_{\mathrm{b}}$ denoting the empty, bonding and antibonding states of the bottom qubit and $|j\rangle_{\mathrm{t}} \in |0\rangle_{\mathrm{t}}, |1\rangle_{\mathrm{t}}$ denoting the empty and one singly occupied states of the top QD. The subsystem has six many body states and the density matrix in this Liouville space is a vector with ten elements (six populations and four coherences between $|+0\rangle$ ($|+1\rangle$) and $|-0\rangle$ ($|-1\rangle$)). Therefore we only need to evaluate all the equations in matrix form and the calculation of cumulants becomes an algebraic operation [30].

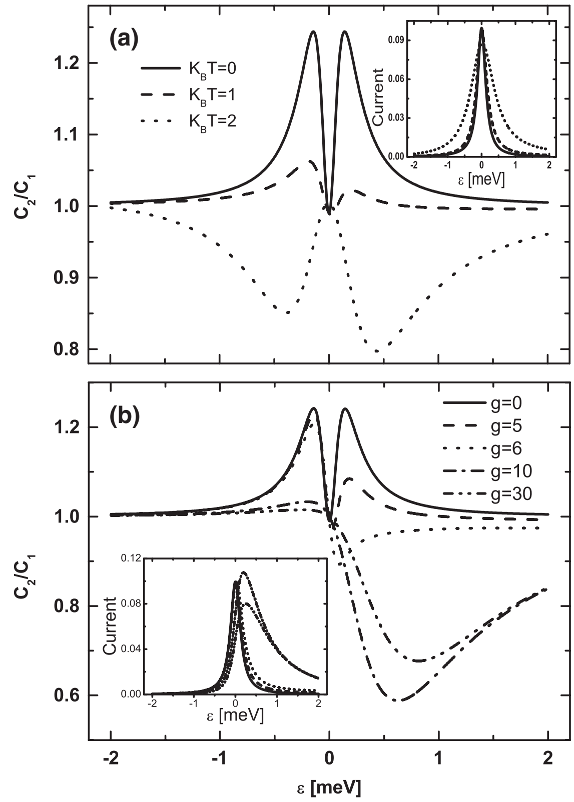

Figure 2. (a) Fano factor C_2/C_1 and current (inset) versus level detuning $\epsilon = \epsilon_{\mathrm{bl}} - \epsilon_{\mathrm{br}}$ without capacitive coupling for different temperatures and $\Gamma_{\mathrm{bl}} = 0.0628$, $\Gamma_{\mathrm{br}} = 0.006\,28$, $\Omega = 0.1$, $\mu_{\mathrm{bl}} = 5$, $\mu_{\mathrm{br}} = -5$ (in meV); (b) effect of pure-dephasing coupling on noise and current (inset) of the bottom qubit. $k_{\mathrm{B}}T = 0.5$ meV, $\Gamma_{\mathrm{tl}} = \Gamma_{\mathrm{tr}} = 1.57$ meV and coupling strength $g = 0, 5, 6, 10, 30$ meV. Other parameters are the same as in (a). The bias applied on the top QD is large enough so that transport is unidirectional.

As pointed out by Kießlich *et al*, super-Poissonian shot noise can be used as an indicator of quantum coherent coupling between two QDs [23]. They found that the Fano factor versus energy level detuning ϵ may exhibit two symmetric peaks close to resonance larger than unity with asymmetric contact couplings, and this structure is asymmetrically washed out in the large bias limit by increasing temperature because of electron–phonon scattering. Without dissipation, however, a similar result can also be found with finite bias by increasing temperature as long as $\mu_{\mathrm{r}} + k_{\mathrm{B}}T < E_{\mathrm{b1}}, E_{\mathrm{b2}} < \mu_{\mathrm{l}} - k_{\mathrm{B}}T$, see figure 2(a). This shows that such a super-Poissonian feature is sensitive to temperature and the molecule level position relative to the Fermi level of the leads. The latter effect is enlarged in our model. Instead of an equilibrium

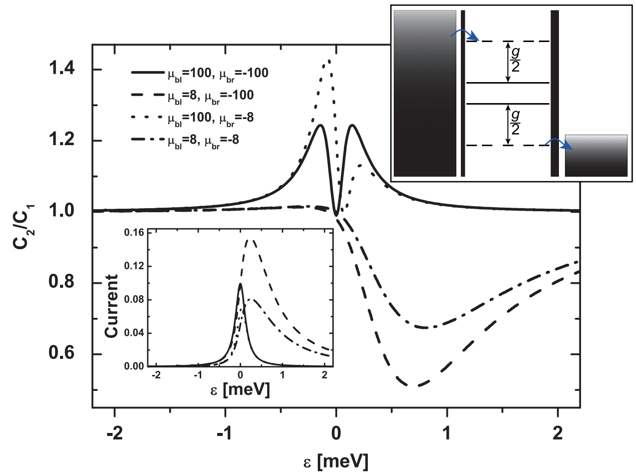

Figure 3. Fano factor C_2/C_1 and current (bottom left inset) versus level detuning $\epsilon = \epsilon_{bl} - \epsilon_{br}$ for different bias configurations. $g = 20$ meV and other parameters are the same as in figure 2(a). Top right inset: sketch of tunneling processes of the bottom qubit where solid lines symbolize two molecule levels which fluctuate by $\frac{g}{2}$ (dashed lines). The arrows represent g-sensitive tunneling processes.

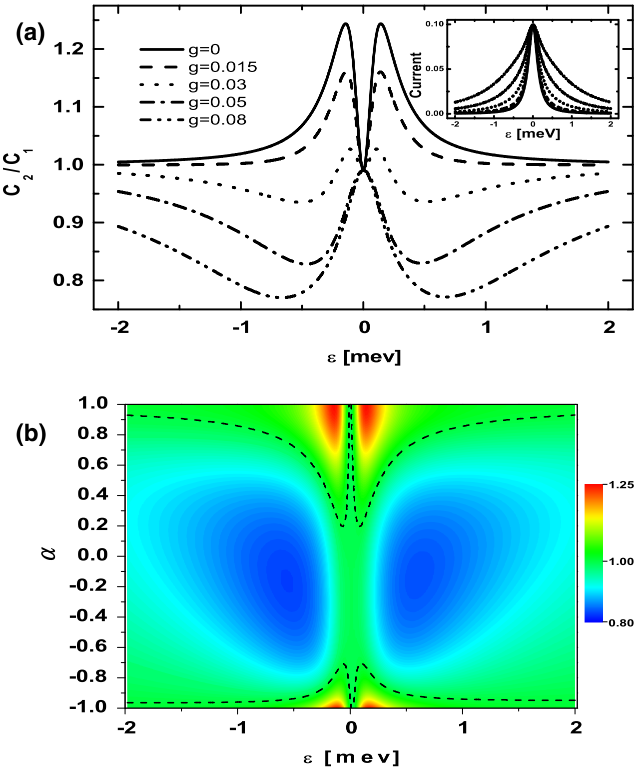

Figure 4. (a) Fano factor C_2/C_1 and current (inset) versus level detuning $\epsilon = \epsilon_{bl} - \epsilon_{br}$ for different orthogonal coupling strengths g. Both the top QD and the bottom qubit are in the large bias limit. Parameters $\Gamma_{tl} = 0.0628$ meV and $\Gamma_{tr} = 1.57$ meV, other parameters are the same as in figure 2(a). (b) Fano factor versus level detuning and the asymmetry of the top QD's contact couplings $\alpha = (\Gamma_{tr} - \Gamma_{tl})/(\Gamma_{tr} + \Gamma_{tl})$. Parameter $\Gamma_{tl} = 1.57$ meV, $g = 0.02$ meV and other parameters are the same as in (a). The black dashed curve corresponds to Fano factor $C_2/C_1 = 1$.

phonon environment, we consider dissipation due to non-equilibrium electron fluctuations in the top QD. Figure 2(b) shows the corresponding Fano factor for pure-dephasing coupling. Increasing the coupling strength g, the double peak is asymmetrically lowered and finally shows obvious sub-Poissonian features for $\epsilon > 0$, but for $\epsilon < 0$ it is still super-Poissonian although its value is close to unity.

This super-Poisson behavior can be understood by dynamical channel blockade [36]. Because of a strong Coulomb blockade and asymmetric couplings ($\Gamma_{bl} > \Gamma_{rl}$), one of the transport channels is much more efficient than the other. As a result, once an electron occupies a state belonging to the slow transport channel, it takes more time to tunnel-out of the dots so that another one can tunnel into them. Thus in the time series electrons seem to transfer in bunches. Charge fluctuation of the top QD makes two states belonging to fast and slow transport channels respectively in the bottom qubit sensitive to g due to pure-dephasing coupling, which means the energy level of each state fluctuates between two values separated by $\frac{g}{2}$ (figure 3, top right inset). For $\epsilon > 0$ with increasing g, the g-sensitive processes strongly change the tunnel-in rate of the slow type transport channel and the tunnel-out rate of the fast one asymmetrically so that one cannot distinguish the two types of transport channels, which contributes to the transition from super-Poissonian to sub-Poissonian; however, for $\epsilon < 0$, the partition of two types of transport channels still exists, which leads to the asymmetric weakening of the double peaks. This effect is more obvious when there is only one type of (slow or fast) g-sensitive process lying in the transport window. In figure 3, we even find an enhanced super-Poissonian noise when $E_{b1} - \frac{g}{2} < \mu_{br} < E_{b1}, E_{b2}, E_{b2} + \frac{g}{2} < \mu_{bl}$ for $\epsilon < 0$ where only tunnel-out rate of the slow type is reduced.

In addition to fluctuating the levels of the qubit as pure-dephasing coupling, orthogonal coupling induces relaxation similar to electron–phonon scattering [22–24]. Interestingly, we can almost symmetrically lower the double peaks to a double minimum by increasing the coupling strength g in the large bias limit, indicating the crossover between coherent and sequential tunneling regimes, see figure 4(a). This implies the same rate of energy absorption and emission due to charge fluctuation of the top QD taking place during transport through the qubit, which resolves the dynamical channel blockade symmetrically for $\epsilon > 0$ and $\epsilon < 0$. In contrast to pure-dephasing coupling, the crossover between these two regimes is also sensitive to the contact couplings of the top QD in orthogonal coupling as shown in figure 4(b) for fixed coupling Γ_{tl}, g and temperature T. Increasing Γ_{tr}, the transition of the tunneling regime is recognized from coherent for small Γ_{tr} to sequential for Γ_{tr} comparative to Γ_{tl} and finally coherent again for large Γ_{tr}, which implies that the current noise of the qubit is strongly dependent on the non-equilibrium transport conditions of the top QD.

4. Conclusions

We have calculated the counting statistics of an interacting quantum dots system consisting of a top QD and a bottom qubit with capacitive coupling based on a lowest order perturbative Markovian quantum master equation. We investigated the effect of a non-equilibrium environment (the top QD with two leads) on electron transport through the bottom qubit near resonance. By increasing the pure-dephasing coupling strength, we find asymmetrical weakening of the double peaks, strongly related to the qubit energy level position referring to the Fermi level of the leads, which suggests a transition from coherent tunneling to sequential tunneling. With proper bias, we even find enhanced super-Poissonian noise. In contrast, for orthogonal coupling super-Poissonian noise can be almost symmetrically lowered to the sub-Poissonian regime with increasing coupling strength because of the same rate of energy absorption and emission due to charge fluctuation of the top QD. This noise is also sensitive to the ratio of the top QD tunneling rates for certain coupling strengths in the large bias limit. For further theoretical study, it would be interesting to investigate the effect of cotunneling and the possible non-Markovian signature in our model [37] and to investigate the effect of non-equilibrium entanglement between two electrically separated qubits with different coupling configurations [38]. We hope our findings may encourage more experimental work in terms of shot noise measurements in similar setups such as those presented in [8, 20, 21].

Acknowledgments

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