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Phys. Rev. B **84**, 115304 — Published 12 September 2011

DOI: 10.1103/PhysRevB.84.115304

Coexistence of super-Poissonian mechanisms in quantum dots with ferromagnetic leads

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Abstract

Spin-resolved noise correlations of electronic currents are investigated through a Coulomb-blockaded two-level quantum dot coupled to two ferromagnetic leads. By changing the bias voltage or/and modulating the spin polarization of leads, different types of dynamical mechanisms are formed to exhibit a sub- to super-Poissonian statistics crossover. An optimized electron correlation is obtained in the specifical regime where several dynamical mechanisms of electron bunching are coexistent. It is found that the shot noise spectrum of spin currents can act as a signal for differentiating which mechanism is responsible for the charge super-Poissonian transport statistics. In addition, the positive cross correlation for two spin species is also predicted.

PACS numbers: 72.70.+m, 72.25.-b, 73.23.Hk, 85.75.-d

I. INTRODUCTION

Quantum fluctuations of electronic current provide a great deal of insights into electronic transport mechanisms and have recently boosted great interest in down-scaling mesoscopic structure [1, 2]. For a system far from equilibrium, the unavoidable shot noise of currents due to granularity of electron charge, obeying no longer the fluctuation-dissipation theorem simply, allows us to access information not available from measurements of the average currents, such as probing the correlations and entanglement of electrons [3–6].

It is well known that noninteracting fermions exhibits suppressed particle number fluctuations relative to the classical Poissonian value (sub-Poissonian statistics) due to the Pauli's exclusion principle [7]. To obtain super-Poissonian statistics in electronic systems, introduction of additional correlated mechanisms is needed, such as the Andreev reflection [8], entanglement [3, 6], and quantum coherence [4, 9, 10]. For a quantum dot (QD), which is a strong confined system, Coulomb interactions between electrons play a unique role in the statistics property. Generally, the repulsive Coulomb interactions result in a further reduction of the shot noise [11]. Under some specifical conditions, however, they can also generate a super-Poissonian characteristics. Three typical bunching mechanisms have been suggested to generate the super-Poissonian statistics. Firstly, the dynamical channel blockade (DCB) was first proposed in a multi-level QD [12], and then extended to tunable situations [13, 14]. Within the time intervals of the electrons in the lower level occasionally tunneling out of the QD, several consecutive tunneling events (bunching of electrons) through the upper level are followed. The second mechanism is the spin accumulation on the QD, which follows from a detailed unbalance of electrons with opposite spin for strongly asymmetric magnetic single-electron-transistors [15]. Thirdly, the dynamical spin blockade (DSB) stems from a spin-dependent bunching of tunneling events, which may appear even in the absence of spin accumulation on the dot [16–18]. In this case, the essential point is that the spin minority with a slow tunneling rate modulates the fast transport of the spin majority.

The three types of mechanisms mentioned above were studied separately in different systems and attributed to different physical origins. The first dynamical mechanism is close associated with the inter-level Coulomb interaction and the other two with the interplay of intra-level Coulomb interaction as well as finite spin polarization. In a realistic system, both of intra- and inter-level interactions along with spin polarization can coexist. Therefore, an interesting issue is to study various super-Poissonian statistics and their mixed effect in the same setup, and further to identify them. On the other hand, these super-Poissonian behaviors are spin-dependent and the study on spin-resolved current correlations [19] is expected to provide additional information about electronic bunching. For instance, spin-current shot noise was suggested to determine the spin unit of quasi-particle [20], or to examine the spin coherence from spin-orbit interactions [21]. Moreover, it was proposed recently to probe attractive or repulsive interactions between electrons even in the lack of charge super-Poissonian behavior [22].

In this paper, we consider a Coulomb-blockaded QD with two levels coupled to two ferromagnetic (FM) leads. This simple model allows us to reach not only various dynamical mechanisms for electron bunching separately, but also their combination by a controllable means. Moreover, the spin-current shot noise is analyzed and the positive cross correlation for two spin channels is predicted in specifical dynamical regimes. It is shown that one can identify various dynamical mechanisms of super-Poissonian by the probe of their spin-current shot noise. The organization of the rest of the paper is as follows. In Sec. II the theoretical model and combination-generation approach for calculation of spin-related shot noise are presented, and in Sec. III the calculation results are discussed in detail for various dynamical mechanisms and their combination. A short summary is given in the last section.

II. MODEL AND FORMALISM

Consider a two-level single QD with finite Coulomb interactions U within each level and U' between levels, sandwiched by two FM electrodes via tunneling coupling. The schematic diagram is depicted in Fig. 1(a), where only dot levels ε_i $(i=1,2)$ contribute to electronic transport and the other levels $\varepsilon_i+U(U')$ (not shown in the schematic) for strong Coulomb interactions are far out of the bias window $eV=\mu_L-\mu_R$. The Hamiltonian can be written as

$$\begin{aligned} H &= \sum_{\alpha k,\sigma} \varepsilon_{\alpha k\sigma} a^\dagger_{\alpha k\sigma} a_{\alpha k\sigma} + \sum_{\sigma,i=1,2} (\varepsilon_i c^\dagger_{i\sigma} c_{i\sigma} + \frac{U}{2} n_{i\sigma} n_{i\bar{\sigma}}) + \sum_{\sigma,\sigma'} \frac{U'}{2} n_{1\sigma} n_{2\sigma'} \\ &\quad + \sum_{\alpha k, i\sigma} (t_\alpha a^\dagger_{\alpha k\sigma} c_{i\sigma} + H.c.), \end{aligned} \tag{1}$$

where $a^\dagger_{\alpha k\sigma}$ $(c^\dagger_{i\sigma})$ is the creation operator for electrons with spin σ in the $\alpha = L, R$ leads (i-th dot level), spin $\bar{\sigma}$ is opposite to σ, and $n_{i\sigma} = c^\dagger_{i\sigma} c_{i\sigma}$ is the occupation operator. The last term denotes the tunneling Hamiltonian between the dot and electron reservoirs with level-independent coupling strength t_α.

It is assumed that the tunneling coupling is so weak that the electronic transport is dominated by the sequential tunneling. Starting from the Liouville-von Neumann equation under the second-order Born and Markovian approximations [23], we derive the dynamical evolution

$$\frac{\partial \rho_n(t)}{\partial t} = \sum_{\alpha,\sigma,m\neq n,} \rho_m(W^{\alpha\sigma+}_{n\leftarrow m} + W^{\alpha\sigma-}_{n\leftarrow m}) - \rho_n(W^{\alpha\sigma+}_{m\leftarrow n} + W^{\alpha\sigma-}_{m\leftarrow n}), \tag{2}$$

for the diagonal element of QD density matrix $\rho_n(t)$, which stands for the probability of finding the dot state $|n\rangle$ at time t. Here, $W^{\alpha\sigma+}_{m\leftarrow n} = \Gamma_{\alpha\sigma} f_\alpha(\epsilon_m - \epsilon_n) \left|\langle m| c^\dagger_{i\sigma} |n\rangle\right|^2$ and $W^{\alpha\sigma-}_{m\leftarrow n} = \Gamma_{\alpha\sigma}[1 - f_\alpha(\epsilon_n - \epsilon_m)] \left|\langle m| c_{i\sigma} |n\rangle\right|^2$ are the transition rates of electrons tunneling from lead α into the QD (superscript $+$) and vice versa (superscript $-$), accompanied with the dot state changed from $|n\rangle$ to $|m\rangle$ according to selection rule $\langle m| c^\dagger_{i\sigma}(c_{i\sigma}) |n\rangle$. ϵ_n is the eigenenergy of state $|n\rangle$ and $f_\alpha(\varepsilon) = 1/[1 + e^{(\varepsilon-\mu_\alpha)/k_B T_\alpha}]$ is the Fermi-Dirac distribution function with $\mu_{L(R)} = \pm eV/2$ as the chemical potential and T_α as the temperature in lead α. $\Gamma_{\alpha\sigma} = 2\pi\Omega_{\alpha\sigma} |t_\alpha|^2$ characterizes the linewidth function with $\Omega_{\alpha\sigma}$ as the spin-resolved density of states at the Fermi level. The spin polarization of the FM leads is defined as $\xi = (\Omega_\uparrow - \Omega_\downarrow)/(\Omega_\uparrow + \Omega_\downarrow)$. Thus, the linewidth functions can be expressed as $\Gamma_{L\uparrow(\downarrow)} = \Gamma_L(1 \pm \xi)$, and $\Gamma_{R\uparrow(\downarrow)} = \Gamma_R(1 \pm \xi)$ for the parallel (P) alignment of the lead's magnetizations while $\Gamma_{R\uparrow(\downarrow)} = \Gamma_R(1 \mp \xi)$ for the antiparallel (AP) alignment.

In the strong Coulomb blockade case of U $(U') \gg \delta\epsilon = \epsilon_2 - \epsilon_1$, at most one electron can occupy the QD. There are five probable states whose probabilities can be denoted as a density vector $\rho(t) = (\rho_0, \rho_{1\uparrow}, \rho_{1\downarrow}, \rho_{2\uparrow}, \rho_{2\downarrow})^T$. Under this representation, Eq. (2) can be rewritten as $d\rho(t)/dt = \mathbf{M}\rho(t)$ with

$$M = \sum_\alpha \begin{pmatrix} -\sum\limits_{i\sigma} W^{\alpha+}_{i\sigma\leftarrow 0} & W^{\alpha-}_{0\leftarrow 1\uparrow} & W^{\alpha-}_{0\leftarrow 1\downarrow} & W^{\alpha-}_{0\leftarrow 2\uparrow} & W^{\alpha-}_{0\leftarrow 2\downarrow} \\ W^{\alpha+}_{1\uparrow\leftarrow 0} & -W^{\alpha-}_{0\leftarrow 1\uparrow} & 0 & 0 & 0 \\ W^{\alpha+}_{1\downarrow\leftarrow 0} & 0 & -W^{\alpha-}_{0\leftarrow 1\downarrow} & 0 & 0 \\ W^{\alpha+}_{2\uparrow\leftarrow 0} & 0 & 0 & -W^{\alpha-}_{0\leftarrow 2\uparrow} & 0 \\ W^{\alpha+}_{2\downarrow\leftarrow 0} & 0 & 0 & 0 & -W^{\alpha-}_{0\leftarrow 2\downarrow} \end{pmatrix}. \tag{3}$$

Its formal solution is $\rho(t) = \rho^{(0)} e^{\mathbf{M}t}$, where vector $\rho^{(0)}$ is the steady-state solution satisfying equation $\mathbf{M}\rho^{(0)} = 0$ with normalization condition $\sum\limits_n \rho^{(0)}_n = 1$. The spin-resolved stationary currents, given by $I_\alpha^\sigma = -e\sum_{m,n}\rho_m(W_{n\leftarrow m}^{\alpha\sigma+} - W_{n\leftarrow m}^{\alpha\sigma-})$, can be expressed as $I_\alpha^\sigma = e\sum_\upsilon[\mathbf{I}_\alpha^\sigma\rho^{(0)}]_\upsilon$. Here, the summation goes over all vector elements υ, and $\mathbf{I}_\alpha^\sigma$ is the matrix of current operator, whose nonzero elements are

$$\mathbf{I}_{12}^{\alpha\uparrow} = W_{0\leftarrow 1\uparrow}^{\alpha-}, \mathbf{I}_{14}^{\alpha\uparrow} = W_{0\leftarrow 2\uparrow}^{\alpha-}, \mathbf{I}_{21}^{\alpha\uparrow} = -W_{1\uparrow\leftarrow 0}^{\alpha+}, \mathbf{I}_{41}^{\alpha\uparrow} = -W_{2\uparrow\leftarrow 0}^{\alpha+}$$

for $\mathbf{I}_\alpha^\uparrow$, and

$$\mathbf{I}_{13}^{\alpha\downarrow} = W_{0\leftarrow 1\downarrow}^{\alpha-}, \mathbf{I}_{15}^{\alpha\downarrow} = W_{0\leftarrow 2\downarrow}^{\alpha-}, \mathbf{I}_{31}^{\alpha\downarrow} = -W_{1\downarrow\leftarrow 0}^{\alpha+}, \mathbf{I}_{51}^{\alpha\downarrow} = -W_{2\downarrow\leftarrow 0}^{\alpha+} \tag{4}$$

for $\mathbf{I}_\alpha^\downarrow$.

The shot noise spectrum can be calculated by the Fourier transform of the current-current correlation function

$$S_{\alpha,\alpha'}^{\sigma,\sigma'}(\omega) = 2\int_{-\infty}^{\infty} dt e^{i\omega t}[\langle \hat{I}_\alpha^\sigma(t)\hat{I}_{\alpha'}^{\sigma'}(0)\rangle - \langle \hat{I}_\alpha^\sigma\rangle\langle \hat{I}_{\alpha'}^{\sigma'}\rangle]. \tag{5}$$

Following the widely applied combination-generation approach [16, 24, 25], we have

$$\langle \hat{I}_\alpha^\sigma(t)\hat{I}_{\alpha'}^{\sigma'}(0)\rangle = \theta(t)\sum_\upsilon[\mathbf{I}_\alpha^\sigma e^{\mathbf{M}t}\mathbf{I}_{\alpha'}^{\sigma'}\rho^{(0)}]_\upsilon + \theta(-t)\sum_\upsilon[\mathbf{I}_{\alpha'}^{\sigma'} e^{-\mathbf{M}t}\mathbf{I}_\alpha^\sigma\rho^{(0)}]_\upsilon, \tag{6}$$

where $\theta(t)$ is the Heavisider function. Substituting Eq.(6) into Eq.(5) and performing Fourier transform, the expression for the spin-resolved shot noise spectrum is obtained as

$$S_{\alpha,\alpha'}^{\sigma,\sigma'}(\omega) = \delta_{\alpha,\alpha'}\delta_{\sigma,\sigma'}S_{\alpha\sigma}^{Sch} - 2e^2\sum_\upsilon[\mathbf{I}_\alpha^\sigma\mathbf{S}\mathbf{E}_+\mathbf{S}^{-1}\mathbf{I}_{\alpha'}^{\sigma'}\rho^{(0)} + \mathbf{I}_{\alpha'}^{\sigma'}\mathbf{S}\mathbf{E}_-\mathbf{S}^{-1}\mathbf{I}_\alpha^\sigma\rho^{(0)}]_\upsilon, \tag{7}$$

which is the same as Eq. (16) of Ref. [25]. Here $S_{\alpha\sigma}^{Sch} = 2eI_\alpha^\sigma$ is the self-correlation Schottky noise, $\mathbf{S}$ is a matrix whose columns are eigenvectors of matrix $\mathbf{M}$, and $\mathbf{E}_\pm$ is a diagonal matrix whose diagonal elements are $\mathbf{E}_\pm^{(nn)} = \frac{1}{\lambda_n \pm i\omega}$ for $\lambda_n \neq 0$ and $\mathbf{E}_\pm^{(nn)} = 0$ for $\lambda_n = 0$ with λ_n as the n-th eigenvalue of matrix $\mathbf{M}$.

The shot noise definition for $I_\alpha^{c(s)} = I_\alpha^\uparrow \pm I_\alpha^\downarrow$ gives the charge-current shot noise as $S^c = \sum_{\sigma=\uparrow,\downarrow}(S^{\sigma\sigma} + S^{\sigma\bar{\sigma}})$ and the spin-current shot noise as $S^s = \sum_{\sigma=\uparrow,\downarrow}(S^{\sigma\sigma} - S^{\sigma\bar{\sigma}})$. In this paper, we are interested in the zero frequency shot noise $S_{\alpha,\alpha'}^{\sigma,\sigma'}(0) = S_{\alpha,\alpha'}^{\sigma,\sigma'}(\omega = 0)$. It can be easily shown that auto-correlations and cross-correlations satisfy $S^{\sigma\sigma'} = S_{LL}^{\sigma\sigma'}(0) = S_{RR}^{\sigma\sigma'}(0) = -S_{LR}^{\sigma\sigma'}(0) = -S_{RL}^{\sigma\sigma'}(0)$.

III. RESULTS AND DISCUSSION

We focus on the shot noise in the region of $\mu_L > \varepsilon_2 > \varepsilon_1 \sim \mu_R$, as shown in Fig. 1(a), for which the Fermi distribution functions at temperature $k_BT_\alpha \ll \delta\epsilon$ reduce to be $f_L(\varepsilon_1)$

$= f_L(\varepsilon_2) = 1$ and $f_R(\varepsilon_2) = 0$. For ε_1 close to $\mu_R = -eV/2$, define x to be $f_R(\varepsilon_1) = 1/[1+e^{(\varepsilon_1+eV/2)/k_BT_R}]$, which can be tuned by the bias voltage. It spans from $x=0$ to $x=1$ with changing μ_R well below ε_1 to well above ε_1.

In Figs. 2(a) and 2(b), we plot charge-current Fano factor $F^c = S^c/(2eI^c)$ (solid lines) and spin-current Fano factor $F^s = S^s/(2eI^c)$ (dashed lines) as a function of x for different spin polarizations of the leads. For the nonmagnetic leads ($\xi = 0$), a prominent characteristic is that the charge-current Fano factor F^c starts with a sub-Poissonian statistics $F^c < 1$ and increases monotonically with x to maximal $F^c = 2$ at $x = 1$, whereas the spin-current Fano factor F^s keeps to the Poissonian value $F^s = 1$, independently of x. Interestingly, F^c and F^s merge at $x = x_0$, which is determined by $S^{\sigma\bar{\sigma}} = 0$, yielding $(2\Gamma_L+\Gamma_R){x_0}^2+3\Gamma_R x_0-2\Gamma_R = 0$. For $x > x_0$, we have $F^c > F^s$, indicating that the current correlation $S^{\sigma\bar{\sigma}}$ between opposite spins becomes attractive. For the FM leads ($\xi \neq 0$), the Fano factor shows quite different variation with x. In the P alignment [Fig. 2(a)], the increase of ξ lifts both F^c and F^s curves and makes them faster beyond the Poissonian value, but the cross-point of F^c and F^s is fixed at x_0. More complicated behaviors appear in the AP alignment, e.g., the cross-point x_0 is ξ-dependent and vanishes for larger ξ, as shown in Fig. 2(b). To clarify the underlying physics, we discuss in what follows two limit cases of $x = 0$ and $x = 1$.

A. Double level transport regime: $x = 0$

At the regime of $x = 0$ (i.e., $\varepsilon_1 > \mu_R$), the electronic transport through either level ε_1 or ε_2 is allowed. The corresponding Fano factors in the P and AP alignments are plotted in Fig. 3(a) as a function of polarization ξ with symmetric coupling $\Gamma_L = \Gamma_R$. In the P case, the spin-resolved shot noises are obtained as

$$S^{\uparrow\uparrow(\downarrow\downarrow)} = eI\frac{1\pm\xi}{(4\Gamma_L+\Gamma_R)^2}\left[\frac{16\Gamma_L^2}{1\mp\xi}+4\Gamma_L\Gamma_R+{\Gamma_R}^2\right], S^{\sigma\bar{\sigma}} = -eI\frac{4\Gamma_L\Gamma_R}{(4\Gamma_L+\Gamma_R)^2}. \tag{8}$$

For $\xi = 0$, it follows that the Fano factor of charge current $F^c = 1 - \frac{8\Gamma_L\Gamma_R}{(4\Gamma_L+\Gamma_R)^2} < 1$ is of sub-Poissonian type, the spin-current Fano factor $F^s = 1$ is of Poissonian type, and the opposite-spin correlation $S^{\sigma\bar{\sigma}}$ is repulsive. If $4\Gamma_L$ is replaced with Γ_L in Eq. (8), it will reproduce the result without interactions. The quadrupling of Γ_L arises from the presence of strong Coulomb interactions within and between levels. Although there are two levels ε_1 and ε_2 within the bias window, only one electron can occupy on the QD. As a result, there are four possible spins ($|1\uparrow\rangle$, $|1\downarrow\rangle$, $|2\uparrow\rangle$, and $|2\downarrow\rangle$) that can enter an empty dot, each of them with Γ_L, whereas for a singly-occupied dot there is only one possible spin to tunnel out with Γ_R. For the FM leads ($\xi \neq 0$), F_P^c and F_P^s in the P alignment increase monotonically with ξ, and F_P^c develops from a sub-Poissonian type to a super-Poissonian type, exhibiting ξ-induced bunching of electrons. The transition of sub- to super-Poissonian Fano factor happens at a threshold of spin polarization, $\xi_1 = \sqrt{\Gamma_R/(4\Gamma_L + \Gamma_R)}$. Such a change stems mainly from the contribution of $S^{\uparrow\uparrow}$, for which the DSB for spin majority is formed.

To clearly illustrate the above bunching process with respect to the spin majority, we plot corresponding schematic diagrams in Figs. 1(b)-(d). For simplify, only one level of the dot is plotted there. If a spin-down electron occupies the level [see Fig. 1(b)], the spin-up electron cannot flow through the dot under strong Coulomb blockade. Until the spin-down electron leaves the dot for the right lead, the spin-up electron has a chance to tunnel [see Fig. 1(c)]. Here, the spin-up electron is assumed to be the spin-majority. It means that large numbers of spin-up electrons will flow through the dot within a time interval $t_\uparrow$, which is shorter than that of a spin-down electron due to $t_\sigma \sim 1/\Gamma_\sigma$ and $\Gamma_\uparrow > \Gamma_\uparrow$. The larger the polarization ξ, the more the spin-up electrons of consecutive tunneling are, and at the same time, the shorter $t_\uparrow$ is, leading to bunching effect of the spin majority at certain threshold ξ_1 [see Fig. 1(d)]. This can explain the Fano factor changing from sub- to super-Poissonian type in Fig. 3(a). It is emphasized that this spin-dependent bunching induces not only a super-Poissonian F_P^c but also $F_P^s > 1$ owing to enhancing $S^{\uparrow\uparrow}$. However, the DSB does not change negative $S^{\sigma\bar{\sigma}}$, which is independent of ξ as given by Eq. (8).

In the AP alignment, there exhibits a nonmonotonic behaviors of Fano factor, as shown in Fig. 3(a). The charge-current Fano factor can be obtained analytically as $F_{AP}^c = 1 + 8(4\xi^2 - 3\xi^4 - 1)/(5 + 3\xi^2)^2$. It then follows that $F_{AP}^c > 1$ for $\xi > \xi_2$ with ξ_2 determined by $4\xi_2^2 - 3\xi_2^4 - 1 = 0$. This super-Poissonian behavior is close associated with the joint effect of finite spin polarization and strong Coulomb interactions. Noted that F_{AP}^c is always less than one in the single-level QD without or with interactions [17]. In the AP configuration, there is no DSB-induced bunching, since both spin species now experience equal coupling strengths with only the roles of source and drain exchanged. Further analysis indicates that the weak bunching of tunneling events results from the spin accumulation due to the strong asymmetric tunneling, which is further enhanced by combinative inter- and intra-interactions. Distinguishing significantly from the DSB, this super-Poissonian statistics F_{AP}^c

in the range of $\xi_2 < \xi < 1$ is accompanied with a spin-current shot noise $F^s_{AP} = 1 - 48\xi^2(1-\xi^2)/(5+3\xi^2)^2$ below the Poissonian value and a positive cross noise $S^{\sigma\bar{\sigma}} = (20\xi^2 - 18\xi^4 - 2)/(5+3\xi^2)^2 > 0$. The ξ dependence of $S^{\sigma\bar{\sigma}}$ is the origin of cross-point x_0 of F^c_{AP} and F^s_{AP} that shifts with ξ, as shown in Fig. 2 for the AP case. It is remarked that positive $S^{\sigma\bar{\sigma}}$ has also been reported previously due to "population inversion" for strong asymmetric coupling [22].

B. Single level transport regime: $x = 1$

At $x = 1$ (or $\varepsilon_1 < \mu_R$), only the upper level ϵ_2 is within the bias window and forms the only channel for an electron to pass through the dot. In this regime, the Fano factors F^c and F^s in the P and AP alignments are plotted in Fig. 3(b) as a function of spin polarization ξ in leads. Unlike in the $x = 0$ case, even for the unpolarized leads ($\xi = 0$), F^c exhibits super-Poissonian behavior, $F^c = 1 + \frac{2\Gamma_L}{\Gamma_L+\Gamma_R}$, and the corresponding spin-resolved shot noises is given by

$$S^{\sigma\sigma} = eI[1 + \frac{\Gamma_L}{\Gamma_L + \Gamma_R}], S^{\sigma\bar{\sigma}} = eI\frac{\Gamma_L}{\Gamma_L + \Gamma_R}. \tag{9}$$

Here not only the same-spin correlations but also the opposite-spin correlations are positive, all the spin components deviating remarkably from the Fermi statistics. This bunching effect arises from the spin-dependent DCB, different from the usual DCB in a two-level QD [12, 13] where the opposite-spin currents are uncorrelated, $S^{\sigma\bar{\sigma}} = 0$. When an electron in the lower level (slow channel) is thermally activated out of the dot, the probability of bunching effect through the higher level can happen between electrons with either same spin $S^{\sigma\sigma}$ or opposite spins $S^{\sigma\bar{\sigma}}$ due to random spin of the slow channel. It then can be understood why the same correlated factor $\Gamma_L/(\Gamma_L+\Gamma_R)$ appears in $S^{\sigma\sigma}$ and $S^{\sigma\bar{\sigma}}$ of Eq. (9). This spin-dependent DCB differs from the DSB for which the slow channel is only for spin minority and so the bunching effect arises always for spin majority $S^{\uparrow\uparrow}$. Interestingly, it is found that despite $F^c > 1$, the spin-current noise $F^s = 1$ is of Poissonian-type statistics in the DCB mechanism, which is able to distinguish evidently from the mechanisms associated with the DSB ($F^s > 1$) and the spin accumulation ($F^s < 1$), as shown in Fig. 3(a).

In the discussion above, we have studied different super-Poissonian behaviors dominated separately by only one of the dynamical mechanisms. In order to optimize bunching of electrons, one can combine these dynamical mechanisms as shown in Fig. 3(b) for finite spin polarization. Here, all the Fano factors start from the DCB behavior of $F^c = 2$ and $F^s = 1$ at $\xi = 0$, and increase monotonically with ξ for both P and AP configurations. Especially, after $\xi \sim 0.5$, $F_P^{c(s)}$ in the P case shows a dramatic enhancement and becomes diverge at full polarization $\xi = 1$, while $F_{AP}^{c(s)}$ in the AP case increases relatively slowly and maximizes to $F_{AP}^{c(s)} = 3$ at $\xi = 1$. The former is governed by a joint contribution of the DCB and DSB, while the latter is governed by that of the DCB and weak spin accumulation. In these regimes, there are the optimized correlations between electrons, giving rise to the super-Poissonian statistics of charge current and spin current as well as positive cross correlations between the currents of two spin species at the same time. They arise from the interplay of intra- and inter-level Coulomb interactions together with the spin block effect. The joint super-Poissonian statistics is characterized by $F^s_{P(AP)} > 1$ and $F^s_{P(AP)} < F^c_{P(AP)}$ (or $S^{\sigma\bar{\sigma}} > 0$), which is significantly distinct from $F_P^s > F_P^c$ in the DSB mechanism and from $F_{AP}^s < 1$ in the spin accumulation mechanism. Thus, probing of spin-current shot noise F^s can give important information to understand different dynamical mechanisms of super-Poissonian statistics.

IV. SUMMARY

We have investigated the spin-resolved noise correlations of electronic currents through a double-level QD coupled to two FM leads by taking into account the interplay of strong intra- and inter-level Coulomb interactions on the QD and spin polarization of the leads. By changing x (equivalently the bias voltage), spin polarization, and the relative orientation of lead's magnetizations, we can realize different types of charge super-Poissonian statistics separately, and their combination for which electron correlations are optimized due to the simultaneous formation of several super-Poissonian mechanisms. Interestingly, it is shown that one can identify different dynamical mechanisms of super-Poissonian statistics by probing spin-current shot noise F^s (either its Poissonian type or/and its relative magnitude compared to the charge-current shot noise F^c). In addition, the positive cross correlation for two spin species is predicted for either the spin-dependent DCB or the spin accumulation mechanism.

V. ACKNOWLEDGMENTS

This work was supported by Program for New Century Excellent Talents in University, and NSF-China under Grant No.10974058 and No. 60825402. D.Y.X. acknowledges support from the State Key Program for Basic Researches of China under Grants No. 2011CB922103 and No. 2010CB923400.

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FIG. 1: (Color online) (a) Schematic view of a two-level QD coupled to two FM leads. Fermi distribution function $x = 1/[1+e^{(\varepsilon_1-\mu_R)/k_BT_R}]$ weighs the tunneling rate Γ_R between incoming and outgoing electrons when ε_1 is close to $\mu_R = -eV/2$. (b)-(d) Formation process of spin majority bunching, where faster transport of spin-up electrons is modulated by slower tunneling of the spin-down electron.

FIG. 2: (Color online) Fano factors F^c of charge current and F^s of spin current as a function of x for different spin polarizations of the leads in the (a) P and (b) AP magnetization configuration.

FIG. 3: (Color online) Fano factors $F_P^{c(s)}$ in the P configuration and $F_{AP}^{c(s)}$ in the AP one as a function of spin polarization ξ at (a) $x = 0$ and (b) $x = 1$.

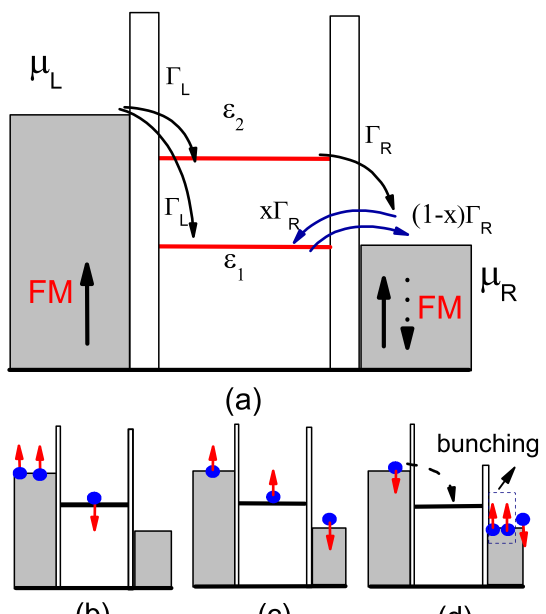

Figure 1 LF13250B 15Aug2011

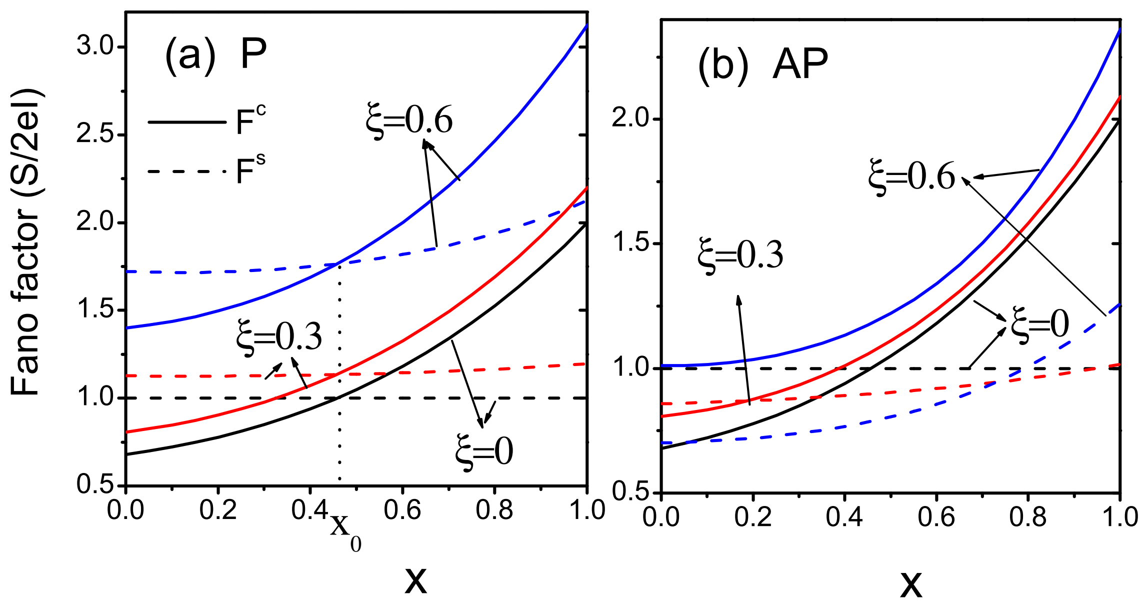

Figure 2 LF13250B 15Aug2011

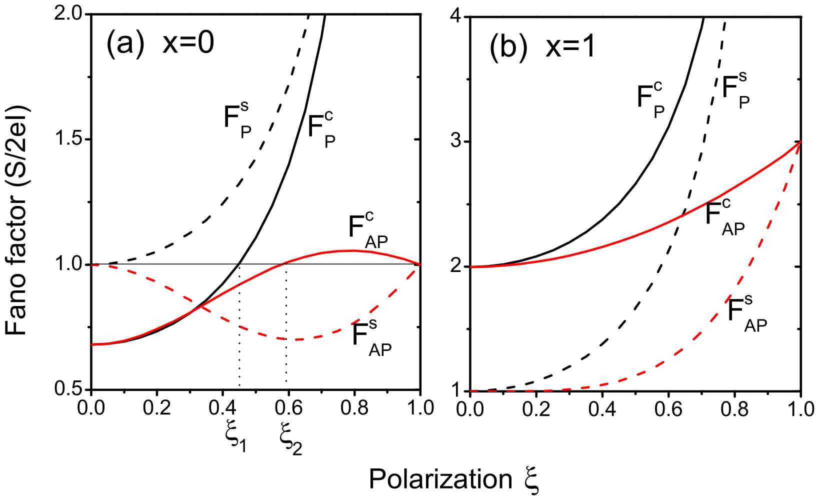

Figure 3 LF13250B 15Aug2011