Fano effect on shot noise through a Kondo-correlated quantum dot

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Transport in a semiopen Kondo-correlated quantum dot is mediated through more than one quantum state. Using the Keldysh technique and the equation of motion method, we study the shot noise S for a wide range of source-drain voltages V_{sd} within a model incorporating the additional states as a background continuum, demonstrating the importance of the Fano interference. In the absence of the interference, the noise is revealed to be a probe of the second moment of the local density of states, and our theory reproduces the well-known peak structure around the Kondo temperature in the S-V_{sd} curve. More significantly, it is found that taking account of the background transmission, the voltage dependence of the noise exhibits rich peak-dip line shapes, indicating the presence of the Fano effect. We further demonstrate that due to its two-particle nature, the noise is more sensitive to the quantum interference effect than the simple current.

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Since the work of Esaki and Tsu,[1] mesoscopic physics has been a subject of increasing experimental and theoretical interest for its many potential applications in the not-too-distant future. The study of mesoscopic devices is also of crucial importance in enhancing our understanding of fundamental many-body phenomena. One of the most striking paradigms exhibiting many-body correlations in mesoscopic physics is the Kondo effect in quantum dots (QDs) coupled to external leads. The effect, first discovered in bulk metals slightly doped with magnetic impurities,[2] was predicted to take place in a spin-polarized QD in 1988,[3] and indeed it was observed ten years later,[4] leading to subsequent extensive investigations. While early studies were concentrated on the limiting cases of closed and open dots, recent focus has shifted to the intermediate regime where the dot-lead coupling is weak enough to maintain a charge quantization and strong enough to activate a multilevel transmission. In this regime, transport becomes more complicated and interesting. For example, interference of transmitted electronic waves through different dot levels leads to a Fano-shaped resonance.[5,6] Experiments[7,8] measuring the transmission phase in QDs and subsequent theoretical activities[9,10] have already demonstrated the importance of invoking transmission mediated through more than one quantum state in QDs. The same is true if a QD is strongly Kondo correlated because the dot should be more open to the leads in order to reach the unitary Kondo limit, which has not yet attracted sufficient attention. Multilevel transport readily occurs in a Kondo dot of large size,[6] where irregularities in the shape and defects in the two-dimensional (2D) geometry can give rise to a sequence of additional quasibound levels localized at separate sites within the dot area. A gate voltage can shift them across the electrochemical potential and they become occupied. These quasibound states arise in the dot's energy spectrum each time the extended wave function of the highest occupied level in the dot is unable to adiabatically adjust to the space variations of the confining potential from source to drain, and participate in transport provided that they are in the vicinity of the electrochemical potential.

Out of all the transport probes, shot noise is a powerful experimental tool to investigate correlations in transport processes through nanostructures beyond the capabilities of dc measurements.[11] The study of shot noise through strongly correlated QD devices in the Kondo regime has recently received increasing attention.[12–24] For a QD coupled to two leads, shot noise was addressed for a wide range of voltages from high-voltage Coulomb-blockade regime to the low-voltage Kondo regime.[12] Shot noise in a Kondo QD near the unitary limit was studied with an emphasis on the effective fractional charge.[22,23] Slave-boson mean-field (SBMF) results were also reported for serial and parallel two-QD systems[15] and for a three-terminal QD in the Kondo regime.[16] These investigations have demonstrated that sensitivity to quantum statistics, interference, and interactions between electrons has made noise measurement an effective tool to probe the Kondo physics in QDs, yielding information that is not accessible by current measurements. However, researches on the impact of a multilevel transmission on the shot noise through a Kondo dot are still missing. This point is important because, normally, it is very difficult to ensure a pure Kondo transport through QDs from the experimental point of view.

In this paper, we shall investigate the shot noise through a Kondo dot in which additional transmission channels different from the Kondo-resonant channel are activated. We handle the problem by assuming that the additional channels can be effectively treated as a background of continuum states in the dot. These states hybridize with the electronic states in the leads, but not with the Kondo state in the dot. Such a treatment captures most of the physics contained in the strong coupling Kondo regime, already providing a theoretical interpretation[10] of the unusual phase evolution in the Kondo QDs measured by the Aharonov–Bohm (AB) interferometry.[8] Here, our purpose is to study the shot noise as a function of the source-drain voltage in the whole regime. It is shown that the Fano effect, resulting from interference between the Kondo state and the continuum, manifests itself in the voltage dependence of the noise, which exhibits rich peak-dip shapes, rather than in the corresponding current behavior.

According to our model, a QD coupled to two leads L and R is described by the Hamiltonian, $H=H_r+H_D+H_T+H_0$, where

$$H_r = \sum_{k,\sigma,\alpha} \varepsilon_k C^{\dagger}_{k\sigma\alpha} C_{k\sigma\alpha}, \quad (1)$$

$$H_D = \sum_{\sigma} \varepsilon_d d^{\dagger}_{\sigma} d_{\sigma} + \frac{U}{2} \sum_{\sigma} n_{d\sigma} n_{d\bar{\sigma}}, \quad (2)$$

$$H_T = \sum_{k,\sigma,\alpha} (V_\alpha C^{\dagger}_{k\sigma\alpha} d_\sigma + V^*_\alpha d^{\dagger}_\sigma C_{k\sigma\alpha}), \quad (3)$$

$$H_0 = \sum_{k,k',\sigma} (V_0 C^{\dagger}_{k\sigma L} C_{k'\sigma R} + V^*_0 C^{\dagger}_{k'\sigma R} C_{k\sigma L}). \quad (4)$$

$C_{k\sigma\alpha}(C^{\dagger}_{k\sigma\alpha})$ and $d_\sigma(d^{\dagger}_\sigma)$ are the destruction (creation) operators in the $\alpha(=L,R)$ lead and in the dot, respectively. The part $H_r+H_D+H_T$ is the usual single-level Anderson Hamiltonian, which can faithfully give rise to the Kondo-resonant transmission through the dot. The last term H_0 models electronic transport through the background of continuum states in the dot, activated in the semiopen regime. Without magnetic fields, the system possesses a time-reversal symmetry and, the phase difference ϕ between the resonant and the background components is either 0 or π. Therefore, one can take the tunneling matrix elements as $V_\alpha \in \mathbb{R}$ and $V_0=|V_0|e^{i\phi} =|V_0|\cos\phi$ without loss of generality. In the absence of the background transmission, the intrinsic Anderson hybridization of the dot level due to tunnel coupling to the leads is $\Gamma=\Gamma_L+\Gamma_R$, with $\Gamma_{L/R}=\pi\rho|V_{L/R}|^2$, where ρ is the density of states in the leads. In this work, we choose symmetric coupling $V_L=V_R$ and take $\rho=1/2D$, where D is the half band width. Formally, this model also pertains to an AB interferometer with an embedded QD.[25]

The current and its temporal fluctuation in the α lead are defined as

$$I_\alpha \equiv \langle \hat{I}_\alpha(t)\rangle \equiv -e\langle \partial \hat{N}_\alpha(t)/\partial t\rangle = -\frac{e}{i\hbar}\langle[\hat{N}_\alpha(t),H]\rangle, \quad (5)$$

$$S_{\alpha\alpha}(t) \equiv \langle[\Delta\hat{I}_\alpha(t),\Delta\hat{I}_\alpha(0)]_+\rangle = \langle[\hat{I}_\alpha(t),\hat{I}_\alpha(0)]_+\rangle - 2I^2_\alpha, \quad (6)$$

with notations $\hat{N}_\alpha(t)\equiv\Sigma_{k,\sigma}C^{\dagger}_{k\sigma\alpha}(t)C_{k\sigma\alpha}(t)$ and $\Delta\hat{I}_\alpha(t) \equiv I_\alpha-\hat{I}_\alpha(t)$. Equation (6) is a symmetrized definition of current noise, and nonsymmetrized noise has also been used in the literature.[26,27] The two definitions are equivalent in the zero-frequency limit. While the current only includes the simple off-diagonal one-body nonequilibrium Green functions (GFs) of the Keldysh type,[28] the noise involves the troublesome two-body GFs, which, in the first step, are decoupled and reduced to one-body ones, e.g.,

$$\begin{aligned}&\langle C^{\dagger}_{k\sigma\alpha}(t)d_\sigma(t)C^{\dagger}_{k'\sigma'\alpha}(0)d_{\sigma'}(0)\rangle\\ &=\langle C^{\dagger}_{k\sigma\alpha}(t)d_\sigma(t)\rangle\langle C^{\dagger}_{k'\sigma'\alpha}(0)d_{\sigma'}(0)\rangle\\ &\quad+\delta_{\sigma\sigma'}\langle C^{\dagger}_{k\sigma\alpha}(t)d_{\sigma'}(0)\rangle\langle d_\sigma(t)C^{\dagger}_{k'\sigma'\alpha}(0)\rangle,\end{aligned}$$

where the first term on the right-hand side corresponds to disconnected diagrams that cancel out the term $2I^2_\alpha$ of Eq. (6). This decoupling scheme (or variations of it) has already been employed in literature[14–20,24,29] and can be justified as follows. (i) The d electrons and C electrons are two different types of electrons, having different spatial configurations and less overlap of their wave functions, so that their dynamical correlations are expected to be small. (ii) Since U is large, the most possible population of the d level is that with only one electron because of the strong effective Coulomb repulsion, which, in turn, reduces the two-body correlations between d electrons and C electrons. Therefore, like the current, the noise is also pinned down to off-diagonal one-body GFs corresponding to the states in the dot and in the α lead. In the second step, with the help of the equations of motion (EOMs) of the operators and applying the analytical continuation rules in a complex time contour,[28] all off-diagonal one-body GFs are cast in terms of diagonal ones, i.e., the dot GF G_d and the bare GF $g_{k\sigma\alpha}$ in the α lead. The lesser, greater, retarded, and advanced components of $g_{k\sigma\alpha}$ are $g^<_{k\sigma\alpha}(t) = i/\hbar f_\alpha(\varepsilon_k)e^{-(i/\hbar)\varepsilon_k t}$, $g^>_{k\sigma\alpha}(t)=i/\hbar[f_\alpha(\varepsilon_k)-1]e^{-(i/\hbar)\varepsilon_k t}$, and $g^{r,a}_{k\sigma\alpha}(t)=\mp i/\hbar\,\theta(\pm t)e^{-(i/\hbar)\varepsilon_k t}$, where $f_\alpha(\varepsilon_k)$ is the Fermi distribution function in the α lead. Although the noise $S_{\alpha\alpha}(t)$ typically includes current fluctuations of any frequencies, in this work, we concentrate on the zero-frequency contribution $S_{\alpha\alpha}=\lim_{\omega\to 0}\int dt e^{i\omega t}S_{\alpha\alpha}(t)$. Therefore, the time-translational invariance is restored for both $S_{\alpha\alpha}$ and I_α, and it is advantageous for all GFs to go over to energy variables. Finally, assuming quasielastic transport, we make use of the current conservation, $I_L=-I_R$, which implies $S_{LL}=S_{RR}=-S_{LR}=-S_{RL}$, and then express the lesser and greater dot GFs $(G^<_d, G^>_d)$ by the retarded one (G^r_d).[19,20] At zero temperature, where the thermal noise vanishes and the remaining noise is called shot noise, we obtain very compact expressions for the current and for the shot noise as follows:

$$\begin{aligned}I=\frac{2e}{h}\int_0^{V_{sd}} d\varepsilon\{a-\mathrm{Im}[\tilde{\Gamma}G^r_d(\varepsilon,V_{sd})]\\ +\mathrm{Im}[(Z+1)\tilde{\Gamma}G^r_d(\varepsilon,V_{sd})]\},\end{aligned} \quad (7)$$

$$\begin{aligned}S=\frac{4e^2}{h}\int_0^{V_{sd}} d\varepsilon\{a-a^2-\mathrm{Im}^2[Z\tilde{\Gamma}G^r_d(\varepsilon,V_{sd})]\\ -\mathrm{Im}[(2a-1)Z\tilde{\Gamma}G^r_d(\varepsilon,V_{sd})]\},\end{aligned} \quad (8)$$

where $a=4\lambda/(1+\lambda)^2$ is the background transmission probability (BTP) though the additional continuum states in the dot, $\tilde{\Gamma}=\Gamma/(1+\lambda)$ is the Anderson hybridization renormalized by these states, $\lambda=\pi^2\rho^2|V_0|^2$ is a dimensionless measure of the background tunneling, and $Z=2a-1+2i\sqrt{a(1-a)}\cos\phi$. Derivations were performed for the source-drain voltage bias V_{sd} taking place in one lead only with the other lead being grounded. The current expression [Eq. (7)] is exact and composed of three distinct parts. The first term in Eq. (7) is the current mediated through the continuum states. The second term corresponds to the usual current expression for the single-level Anderson model, but with renormalized parameters. The third term is the interference between the two current components, which includes multiple-scattering events. Noteworthy, the shot-noise expression [Eq. (8)] can be regrouped in the renormalized noninteracting form $\mathcal{T}(1-\mathcal{T})$,[30] where $\mathcal{T}$ is the total transmission probability through the device, i.e., the integrand of Eq. (7). This regrouping is a direct consequence of the decoupling scheme and the fact that the electronic interaction U is restricted to the dot, allowing for simple interpretations in terms of transmission probabilities. Thus, the noise probes the second moment of the local density of states on the dot, while the current only measures the first one.

Despite their simplicity, Eqs. (7) and (8) are far from being trivial since the calculation scheme for the dot GF $G_d^r(\varepsilon, V_{sd})$ has to include all the necessary many-body correlations in order to describe the Kondo physics. Here, we employ the EOM method[20,31–33] to evaluate $G_d^r(\varepsilon, V_{sd})$ in the Kondo regime. The EOM for $G_d^r(\varepsilon, V_{sd})$ introduces higher-order GFs, whose EOMs produce, in turn, more GFs. One can terminate this hierarchy by decoupling out averages of operators. The latter are then found using the fluctuation-dissipation relationship from the relevant GFs. In this manner, the treatment becomes self-consistent. In Ref. 20, in the limit $U\to\infty$, we have solved $G_d^r(\varepsilon)$ for an equilibrium QD embedded in an AB interferometer by using a formally equivalent Hamiltonian. For a semiopen dot we consider here, the dot GF $G_d^r(\varepsilon, V_{sd})$ becomes strongly voltage dependent since the dot-lead coupling is strong enough for the source-drain bias to drive the dot out of equilibrium. Within the EOM scheme, the voltage V_{sd} enters into $G_d^r(\varepsilon, V_{sd})$ via the Fermi distribution functions in the leads, $f_L(\varepsilon) = 1/[1+\exp\beta\varepsilon]$ and $f_R(\varepsilon) = 1/[1+\exp\beta(\varepsilon - V_{sd})]$, and the signature of Kondo physics is logarithmic divergences arising from integrals of the types:

$$\int_{-D}^{D} d\varepsilon' \frac{f_\alpha(\varepsilon')}{\varepsilon - \varepsilon' + i0^+}$$

and

$$\int_{-D}^{D} d\varepsilon' \frac{f_\alpha(\varepsilon')[G_d^r(\varepsilon')]^*}{\varepsilon - \varepsilon' + i0^+}. \tag{9}$$

Therefore, the Kondo resonance in the local density of states is dramatically affected by V_{sd}. The Kondo temperature is within the EOM $T_K = D\exp[\pi|\varepsilon_d + \mathrm{Re}\,\Sigma_0|/\mathrm{Im}\,\Sigma_0]$, where $\Sigma_0 = -{}_{1/2}\sqrt{a}\Gamma\cos\phi - i\tilde{\Gamma}$ is the noninteracting self-energy of $G_d^r(\varepsilon, V_{sd})$. Our EOM calculation for $G_d^r(\varepsilon, V_{sd})$ is consistent with previous theories[20,32,33] and correctly reproduces qualitative features required by the Kondo effect and by the Fermi-liquid relations. For a quantitative analysis, one should, of course, resort to more advanced techniques, e.g., the numerical renormalization group method. In all our numerical calculations, the reference energy is always set at $E_F = 0$ and the energy cutoff and the intrinsic Anderson hybridization are taken as $D = 1$ and $\Gamma = 0.01$, respectively.

Before we present numerical results on the shot noise as a function of V_{sd}, it is helpful to make some general remarks. At zero temperature, the formula $\mathcal{T}(1-\mathcal{T})$ for interacting QD devices in the Kondo regime was first obtained by the SBMF theory,[13–17,24] where the Anderson Hamiltonian in the limit $U\to\infty$ is rewritten in terms of the slave-boson language and becomes quadratic. Then Wick's theorem can be exactly applied, leading to this formula. Whereas the SBMF theory

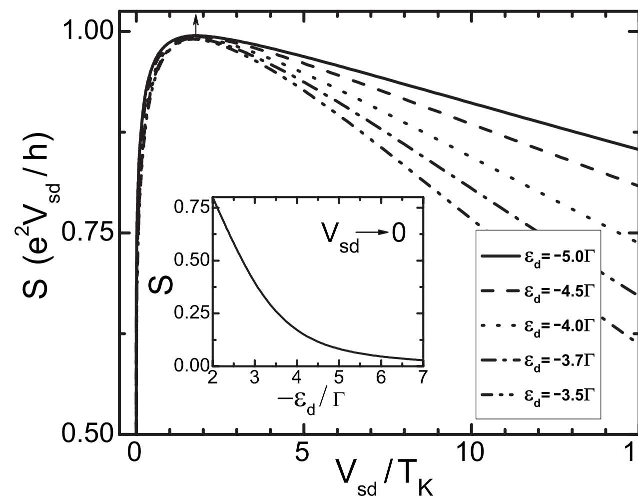

FIG. 1. Shot noise (in units of V_{sd}) through the QD in the absence of background transmissions. Inset: dependence of S/V_{sd} on the dot level ε_d in the limit $V_{sd}\to 0$. All indicated values of ε_d fall in the Kondo regime.

always produces a Lorentzian Kondo peak and is, hence, completely blind to charge fluctuations, our theory has the flexibility to follow the effect of charge fluctuations on the shot noise in the Kondo regime ($\varepsilon_d < -\Gamma$); this effect gradually vanishes as the dot approaches the Kondo limit ($\varepsilon_d \ll -\Gamma$). We emphasize that our EOM treatment of shot noise is a self-consistent procedure since the decoupling scheme used in deriving Eq. (8) and the EOM determination of the dot GF $G_d^r(\varepsilon, V_{sd})$ contain the same high-order many-body correlations.[20,31]

The validity of the EOM method to study the noise is well illustrated in Fig. 1, where we have plotted the voltage-scaled shot noise as a function of the voltage V_{sd} in the case of no background transmissions activated, i.e., $a=0$. From the low-voltage ($V_{sd} < T_K$) Kondo regime to the high-voltage ($V_{sd} > T_K$) Coulomb-blockaded regime, the shot noise exhibits a nonmonotonic dependence on the voltage, with a peak around $(1-2)T_K$ (marked by the arrow). Results are presented for different values of the dot level ε_d, leading to values of T_K within the EOM differing by a factor of 100. Nevertheless, the peak position remains almost unchanged, demonstrating that shot-noise measurements give a straightforward reliable estimate of the Kondo temperature. These features qualitatively agree with the five-method quantitative calculations[12] and the SBMF results.[15] This qualitative validity of the renormalized noninteracting noise formula reinforces the fact that the strongly correlated Kondo resonance can be described by noninteracting quasiparticles with renormalized parameters.[34] At the equilibrium limit $V_{sd}\to 0$, on the condition that a number of symmetries of the Hamiltonian are respected, the system reaches the unitary limit, resulting in $S/V_{sd}=0$.[12,15,22] However, these symmetries are not easily accessible in realistic QDs. In this work, while the $SU(2)$ spin symmetry and the left-right symmetry (requiring equal dot-lead couplings for the left and right leads) are preserved, the particle-hole symmetry ($2\varepsilon_d + U = 0$) is broken because the EOM calculation assumes the infinite Coulomb

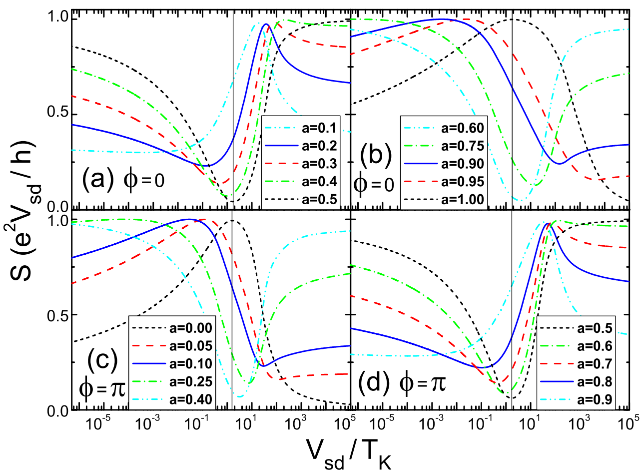

FIG. 2. (Color online) Dependence of the shot noise on the voltage for the level position $\varepsilon_d=-4\Gamma$ when background transmissions through the dot are activated. The four vertical lines mark the peak or dip positions in the curves corresponding to (a) $a=0.5$, (b) $a=1.0$, (c) $a=0.0$, and (d) $a=0.5$, respectively. For either phase difference (0 or π), the voltage dependence of S/V_{sd} dramatically evolves peak-dip structures as the BTP changes from 0 to 1.

interaction, which introduces potential scattering. Therefore, due to a deviation from the unitary limit, we obtain a nonzero S/V_{sd} as $V_{sd}\to 0$, as indicated in the inset of Fig. 1. Here, charge fluctuations on the dot level, which deforms the line shape of the Kondo resonance, also contribute to this deviation; the contribution is small but finite for not big $-\varepsilon_d/\Gamma$ in the Kondo regime. As $-\varepsilon_d/\Gamma$ increases by further lowering the dot level, both the degree of particle-hole asymmetry and charge fluctuations decrease, leading to a complete suppression of S/V_{sd} in the Kondo limit. It is thus clear that the remarkable fractional backscattered charge predicted by Sela *et al.*[22] and nicely confirmed by Zarchin *et al.*[23] is absent in the present case due to the broken particle-hole symmetry and our noninteracting quasiparticle description of the noise. Nonetheless, as we show below, such a description faithfully captures the quantum interference, leading to the distinct Fano–Kondo effect.

When the coupling of the Kondo dot to the leads increases to a strong regime and background transmissions are activated, interference processes can play a crucial role, changing the simple single-peak structure in the voltage dependence of S/V_{sd} in a striking way, as shown in Fig. 2. We find asymmetric line shapes of the peaks and dips, which indicate the presence of the Fano effect, resulting from the interference between the discrete Kondo channel and the continuum of background transmission. If one ignores the voltage dependence of the dot GF, an intuitive understanding of these complicated line shapes can be reached by analyzing some special cases. (i) At $a=0$ or 1, the integrand of Eq. (8), $\mathcal{T}(1-\mathcal{T})$, is reduced to $0.25-[\mathrm{Im}(\tilde{\Gamma}G_d^r)+0.5]^2$, where $-\mathrm{Im}(\tilde{\Gamma}G_d^r)$ develops the unitary Kondo peak near $\varepsilon=0$ with its half-width at half-height of T_K, resulting in the fact that $\mathcal{T}(1-\mathcal{T})$ exhibits a maximum value of 0.25 at $\varepsilon=T_K$. Thus, one may expect a corresponding peak around T_K in the voltage dependence of the noise, which is indeed observed

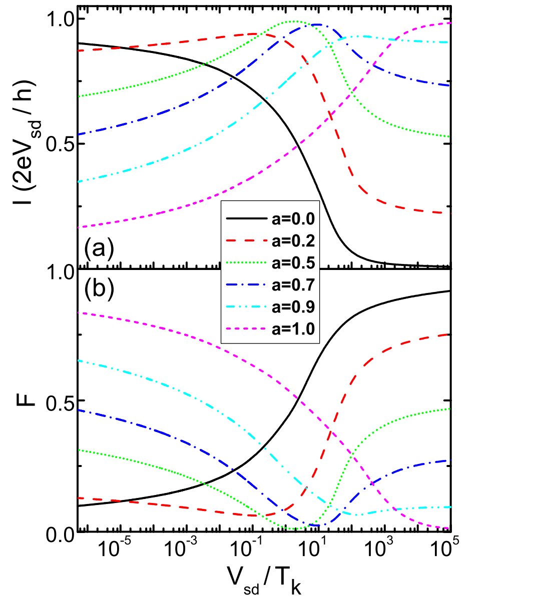

FIG. 3. (Color online) Current and Fano factors as functions of the applied voltage for $\varepsilon_d=-4\Gamma$, $\phi=0$. They exhibit asymmetric peak or dip structure, but by no means the peak-dip shape of a typical Fano effect.

[see the curves corresponding to $a=1$ in Fig. 2(b) and $a=0$ in Fig. 2(c)]. (ii) At $a=0.5$, one has $\mathcal{T}(1-\mathcal{T})=0.25-\mathrm{Re}^2(\tilde{\Gamma}G_d^r)$, where $\mathrm{Re}(\tilde{\Gamma}G_d^r)$ is responsible for the Fano effect and possesses a peak at $\varepsilon=T_K$ with its height of about 0.5. Therefore, $\mathcal{T}(1-\mathcal{T})$ is completely suppressed at $\varepsilon=T_K$. As one might expect, a full Fano dip shows up around T_K in the voltage dependence of S/V_{sd} [see the curves corresponding to $a=0.5$ in Figs. 2(a) and 2(d)]. In other cases, the dependence of the noise on the voltage consists of a complicated peak-dip structure due to the interplay of the imaginary and real parts of the dot GF. Here, we emphasize the similar behaviors of Figs. 2(a) and 2(d), and of Figs. 2(b) and 2(c), which can be understood from the invariance of the noise expression [Eq. (8)] under the transformation: $\{a,\phi=0\}\Rightarrow\{1-a,\phi=\pi\}$. In the equilibrium limit $V_{sd}\to 0$, at $a=0$, the Kondo effect is identified by the suppression of shot noise, although a complete suppression $S/V_{sd}=0$ is not accessible in this work due to the residual potential scattering and charge fluctuations [see Figs. 2(c)]. At finite a, the suppression is violated and, in some cases, even changes to an enhancement. This effect does not at all indicate the absence of the Kondo resonance, but rather is due to the strong Fano interference between the Kondo resonance and the background transmission. In the other limit, of high voltage, $V_{sd}\gg T_K$, the Kondo effect is irrelevant, and the Kondo-resonant channel is effectively Coulomb blockaded, resulting in a very weak average transmission over a scale $0\le\varepsilon\le V_{sd}$ as well as weak interference contributions. Thus, the shot noise in this regime is completely determined by the background transmission, $S/V_{sd}\approx 4a(1-a)e^2/h$, which is independent of the source-drain voltage bias.

Here, the noise we study probes the second moment of the local density of states. In order to further demonstrate its advantages in comparison with the usual current measurement, which only reveals information on the first moment of the density of states, we plot in Fig. 3(a) the evolution of the current I/V_{sd} with a variation of the BTP. At $a=0$, the current is maximal as $V_{sd}\rightarrow 0$ and is suppressed for $V_{sd}\gg T_K$, giving rise to the so-called zero bias anomaly.[35] At finite a, the current evolves an asymmetric peak for $\phi=0$ [see Fig. 3(a)] or dip for $\phi=\pi$ (not shown here), but by no means the peak-dip shape of a typical Fano effect. Therefore, it is clear that the noise is more sensitive to the effect of quantum interference than the current.[20] The corresponding behavior of the Fano factor, defined as $F=S/2eI$, as presented in Fig. 3(b), almost complements the current curves. It develops dip structures at finite BTP where the current exhibits peaks. In the absence of the background transmission, the Fano factor is suppressed in the linear response regime $V_{sd}\rightarrow 0$ and increases to the Poissonian value $F=1$ in the high-voltage Coulomb-blockaded regime. At the moment, we have to point out some SBMF studies[13,24] where the Fano factor saturates at a universal value of 0.5 in the large voltage limit, which is nonphysical. Indeed, in the high-voltage Coulomb-blockaded regime, the shot noise is Poissonian $(F=1)$ (Ref. 11) because the transmission probability is very small $(\mathcal{T}\ll 1)$ and correlations become not so important for transmitting electrons.

In summary, we have presented a qualitative study of shot noise through a Kondo-correlated QD, when the dot-lead coupling increases to a strong regime and additional transmissions are activated. In particular, various peak-dip structures of the Fano effect are predicted in the voltage dependence of the noise, which are absent in the corresponding current behavior. The present work can enrich the existing studies (spreading over one decade) of the noise in the strongly correlated Kondo regime. We hope that it will be informative for future experimental effort.

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[1] L. Esaki and R. Tsu, IBM J. Res. Dev. **14**, 61 (1970).
[2] A. C. Hewson, *The Kondo Problem to Heavy Fermions* (Cambridge University Press, Cambridge, 1993).
[3] L. I. Glazman and M. E. Raikh, JETP Lett. **47**, 452 (1988); T. K. Ng and P. A. Lee, Phys. Rev. Lett. **61**, 1768 (1988).
[4] D. Goldhaber-Gordon, H. Shtrikman, D. Mahalu, D. Abusch-Magder, U. Meirav, and M. A. Kastner, Nature (London) **391**, 156 (1998); S. M. Cronenwett, T. H. Oosterkamp, and L. P. Kouwenhoven, Science **281**, 540 (1998).
[5] J. Göres, D. Goldhaber-Gordon, S. Heemeyer, M. A. Kastner, H. Shtrikman, D. Mahalu, and U. Meirav, Phys. Rev. B **62**, 2188 (2000); I. G. Zacharia, D. Goldhaber-Gordon, G. Granger, M. A. Kastner, Y. B. Khavin, H. Shtrikman, D. Mahalu, and U. Meirav, *ibid.* **64**, 155311 (2001); A. C. Johnson, C. M. Marcus, M. P. Hanson, and A. C. Gossard, Phys. Rev. Lett. **93**, 106803 (2004).
[6] P. Stefański, A. Tagliacozzo, and B. R. Bułka, Phys. Rev. Lett. **93**, 186805 (2004).
[7] A. Yacoby, M. Heiblum, D. Mahalu, and H. Shtrikman, Phys. Rev. Lett. **74**, 4047 (1995); E. Schuster, E. Buks, M. Heiblum, D. Mahalu, V. Umansky, and H. Shtrikman, Nature (London) **385**, 417 (1997); M. Avinun-Kalish, M. Heiblum, O. Zarchin, D. Mahalu, and V. Umansky, *ibid.* **436**, 529 (2005).
[8] Y. Ji, M. Heiblum, D. Sprinzak, D. Mahalu, and H. Shtrikman, Science **290**, 779 (2000); Y. Ji, M. Heiblum, and H. Shtrikman, Phys. Rev. Lett. **88**, 076601 (2002).
[9] A. Aharony, O. Entin-Wohlman, and Y. Imry, Phys. Rev. Lett. **90**, 156802 (2003); C. A. Büsser, G. B. Martins, K. A. Al-Hassanieh, A. Moreo, and E. Dagotto, Phys. Rev. B **70**, 245303 (2004); D. I. Golosov and Y. Gefen, *ibid.* **74**, 205316 (2006); V. Kashcheyevs, A. Schiller, A. Aharony, and O. Entin-Wohlman, *ibid.* **75**, 115313 (2007); C. Karrasch, T. Hecht, A. Weichselbaum, Y. Oreg, J. von Delft, and V. Meden, Phys. Rev. Lett. **98**, 186802 (2007); C. Karrasch, T. Hecht, A. Weichselbaum, J. von Delft, Y. Oreg, and V. Meden, New J. Phys. **9**, 123 (2007); P. G. Silvestrov and Y. Imry, *ibid.* **9**, 125 (2007).
[10] K. Kang, M.-S. Choi, and S. Lee, Phys. Rev. B **71**, 045330 (2005).
[11] Ya. M. Blanter and M. Büttiker, Phys. Rep. **336**, 1 (2000); Yu. V. Nazarov, *Quantum Noise in Mesoscopic Physics* (Kluwer, Dordrecht, 2003).
[12] Y. Meir and A. Golub, Phys. Rev. Lett. **88**, 116802 (2002).
[13] B. Dong and X. L. Lei, J. Phys.: Condens. Matter **14**, 4963 (2002).
[14] R. López and D. Sánchez, Phys. Rev. Lett. **90**, 116602 (2003).
[15] R. López, R. Aguado, and G. Platero, Phys. Rev. B **69**, 235305 (2004).
[16] D. Sánchez and R. López, Phys. Rev. B **71**, 035315 (2005).
[17] B. H. Wu, J. C. Cao, and K. H. Ahn, Phys. Rev. B **72**, 165313 (2005).
[18] G. B. Zhang, S. J. Wang, and L. Li, Phys. Rev. B **74**, 085106 (2006).
[19] T. F. Fang and S. J. Wang, J. Phys.: Condens. Matter **19**, 026204 (2007).
[20] T. F. Fang, S. J. Wang, and W. Zuo, Phys. Rev. B **76**, 205312 (2007).
[21] A. O. Gogolin and A. Komnik, Phys. Rev. Lett. **97**, 016602 (2006).
[22] A. Golub, Phys. Rev. B **73**, 233310 (2006); **75**, 155313 (2007); E. Sela, Y. Oreg, F. von Oppen, and J. Koch, Phys. Rev. Lett. **97**, 086601 (2006).
[23] O. Zarchin, M. Zaffalon, M. Heiblum, D. Mahalu, and V. Umansky, arXiv:0711.4552 (unpublished).
[24] H. F. Lu and Y. Guo, Phys. Rev. B **76**, 045120 (2007).
[25] See Refs. 10 and 20, and references therein.
[26] E. Onac, F. Balestro, B. Trauzettel, C. F. J. Lodewijk, and L. P. Kouwenhoven, Phys. Rev. Lett. **96**, 026803 (2006).
[27] H.-A. Engel and D. Loss, Phys. Rev. Lett. **93**, 136602 (2004).

[28] H. Haug and A.-P. Jauho, *Quantum Kinetics in Transport and Optics of Semiconductors* (Springer, Berlin, 1998).
[29] B. H. Wu, J. C. Cao, and K.-H. Ahn, Phys. Rev. B **74**, 195316 (2006).
[30] J. X. Zhu and A. V. Balatsky, Phys. Rev. B **67**, 165326 (2003).
[31] H. G. Luo, Z. J. Ying, and S. J. Wang, Phys. Rev. B **59**, 9710 (1999).
[32] O. Entin-Wohlman, A. Aharony, and Y. Meir, Phys. Rev. B **71**, 035333 (2005).
[33] V. Kashcheyevs, A. Aharony, and O. Entin-Wohlman, Phys. Rev. B **73**, 125338 (2006).
[34] L. Vaugier, A. A. Aligia, and A. M. Lobos, Phys. Rev. Lett. **99**, 209701 (2007).
[35] The zero bias anomaly usually pertains to the maximal differential conductance dI/dV_{sd} at zero bias due to the Kondo effect. In fact, I/V_{sd} is quite similar in the voltage dependence to dI/dV_{sd}.