Tunable Kondo-Luttinger systems far from equilibrium

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We theoretically investigate the nonequilibrium current through a quantum dot coupled to one-dimensional electron leads, utilizing a controlled frequency-dependent renormalization group approach. We compute the nonequilibrium conductance for large bias voltages and address the interplay between decoherence, Kondo entanglement, and Luttinger physics. The combined effect of large bias voltage and strong interactions in the leads, known to stabilize two-channel Kondo physics, results in nontrivial modifications in the conductance. Interestingly, these unusual changes in the conductance persist in the presence of a finite channel asymmetry.

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I. INTRODUCTION

Understanding strongly correlated quantum systems far from equilibrium is an outstanding challenge in condensed-matter physics. Many of the theoretical approaches that have been proven successful in treating strong correlations are inadequate once the system is driven out of equilibrium. Quantum dot devices provide an ideal setting to study transport under nonequilibrium conditions, as they constitute comparatively simple model systems with high tunability.[1–7] Kondo physics plays a crucial role in understanding their transport properties.[8,9] It has been shown that several effects in these devices will suppress or modify the Kondo screening, such as dissipation and the electron-electron interaction in Luttinger-liquid quantum wires that couple to the dot.[10–16] In this paper, we study nonequilibrium currents across quantum dots in nanosettings involving Kondo entanglement and Luttinger physics.[17]

We consider a quantum dot in the Kondo regime coupled to one-dimensional (1D) reservoir leads described by the Luttinger theory. This model exhibits either a one-channel Kondo (1CK) or a two-channel Kondo (2CK) ground state,[14,16] as the Luttinger parameter K is decreased; the control parameter corresponds to the interaction strength in the 1D leads.[17] The nonequilibrium properties of this system were addressed only in an exactly solvable limit[18] or in the linear (low-bias) region.[19] The full crossover in the nonequilibrium conductance between the 1CK and 2CK fixed points,[14] with much relevance to experiments, has not yet been addressed. In particular, interactions in 1D wires are expected to result in a peculiar nonequilibrium transport.[20,21]

Here, we apply a nonequilibrium renormalization group (RG) method[2,22] to tackle these issues. We calculate the conductance for bias voltages large compared to the relevant Kondo scales. We identify signatures of intermediate 2CK behavior in the RG flow for all $K<1$ which strongly modify the conductance profile. The low-temperature conductance is nonuniversal in the sense that it does not depend on V/T_{2CK} only, where T_{2CK} is the relevant 2CK scale.

II. MODEL HAMILTONIAN AND EQUILIBRIUM PROPERTIES

In our setup, the Kondo Hamiltonian is given by

$$H = H_{lead} + H_K,$$

$$H_K = \sum_{\alpha,\alpha',k,k',\sigma,\sigma'} J_{\alpha,\alpha'} \mathbf{S} \cdot c^{\dagger}_{\alpha',k',\sigma'} \tau_{\sigma'\sigma} c_{\alpha,k,\sigma}, \tag{1}$$

where $\alpha,\alpha' = \mathrm{L},\mathrm{R}$ denote the left/right lead, H_{leads} describes two—left (L) and right (R)—Luttinger liquid leads with Luttinger parameter $0<K<1$ (Refs. 14 and 16) and Fermi energies (chemical potentials) $\mu_{\mathrm{L/R}} = \pm V/2$, $c^{\dagger}_{\alpha,k,\sigma}$ is the electron creation operator for the lead α, $\mathbf{S}$ is the spin-1/2 operator on the dot, and τ are Pauli matrices. In terms of the pseudofermion operators f_{γ}, $\mathbf{S} = \frac{1}{2}\Sigma_{\gamma\gamma'} f^{\dagger}_{\gamma} \tau_{\gamma\gamma'} f_{\gamma'}$. In the Kondo regime, the number operator of the pseudofermions satisfy the constraint: $Q = \Sigma_{\gamma} f^{\dagger}_{\gamma} f_{\gamma} = 1$. Here, we denote by $g_{\mathrm{LR}} = N_0 J_{\mathrm{LR}}$ and $g_{\mathrm{LL}} = g_{\mathrm{RR}} = N_0 J_{\mathrm{LL}} = N_0 J_{\mathrm{RR}}$ the dimensionless interlead and intralead Kondo couplings, respectively,[8,9] where $N_0 = \frac{1}{2D_0}$ and D_0 is an ultraviolet cutoff (approximately few kelvins).

In equilibrium (or zero-bias voltage, $V=0$), the RG analysis results in two infrared fixed points:[14,16] the 1CK and 2CK fixed points. In the former case, all Kondo couplings, g_{LR} and g_{LL}, are relevant under RG transformation and flow toward strong coupling, such that the two leads can be combined into a single effective lead. In contrast, the 2CK fixed point is reached when g_{LR} remains small under RG while g_{LL} grows (and flows to intermediate coupling). Here, the two leads provide independent screening channels. This 2CK fixed point is infrared stable for $K<1/2$ (assuming $g_{\mathrm{LL}} = g_{\mathrm{RR}}$).[14,16]

For $K=1$ (free electron leads), where 1CK physics is realized,[8,9] the differential conductance $G(T) \equiv dI(T)/dV|_{V\to 0}$ [with $I(T)$ being the equilibrium current at finite temperature and V being the source-drain bias voltage] reaches the unitary limit at low temperatures, $G(T) = 2G_0\{1 - \mathcal{O}[(T/T_{\mathrm{K}})^2]\}$, where $G_0 = e^2/h$ is the conductance quantum. Here, $T_{\mathrm{K}} \equiv D_0 e^{-1/(g^0_{\mathrm{LR}} + g^0_{\mathrm{LL}})}$ is the Kondo temperature; whereas g^0_{LR} and g^0_{LL} are the bare values of the coupling constants.

For $T \gg T_K$, from $G(T) \propto g_{LR}(T)^2$, one finds $G(T) \sim 1/\ln^2(T/T_K)$.

For all $K<1$, g_{LR} grows slower under RG than $g_{LL(RR)}$. For $K \ll 1$, one can solve the RG equations analytically for large T. We may neglect g_{LR} in the RG equation for $g_{LL/RR}$ to obtain the approximate solution $g_{LL/RR}(T) \approx 1/\ln(T/T_K^*)$ with the shorthand $T_K^* \equiv D_0 e^{-1/g_{LL}^0}$. The coupling $g_{LR}(T)$ is found by substituting the approximate solution for $g_{LL/RR}(T)$ in the RG equation for the coupling g_{LR}. We evaluate

$$g_{LR}(T) \approx \frac{(T/D^*)^{1/2(1/K-1)}}{\ln^2(T/T_K^*)}, \qquad (2)$$

where $D^* = \frac{D_0}{[g_{LR}^0 \ln^2(D_0/T_K^*)]^s}$ with $1/s=(1/K-1)/2$. We deduce that, for $T \gg T_K^*$, the conductance $G(T) \propto g_{LR}^2(T)$ follows $T^{1/K-1}/\ln^4(T/T_K^*)$. Here, the power-law behavior is reminiscent of Luttinger physics whereas the logarithmic contribution is typical of Kondo correlations. The same power law (without logarithmic corrections) has been found in Ref. 15. Importantly, the conductance is *not* a universal function of (T/T_K^*) because transport arises from the subleading coupling g_{LR}.

For $1/2<K<1$, the low-temperature physics is governed by two scales, $T_{1CK}<T_{2CK}$ with 1CK behavior for $T \ll T_{1CK}$ and 2CK behavior for $T_{1CK} \ll T \ll T_{2CK}$. In the limits $K \to 1$ and $K \to 0$, we have $T_{1CK}=T_K$ and $T_{2CK}=T_K^*$, respectively, with T_K and T_K^* defined above. In general, $T_{1CK} \ll T_K$ due to interactions in the leads; also, $T_{2CK}>T_K^*$. In the presence of particle-hole symmetry, the conductance for $1/2<K<1$ reaches the unitary limit as $T \to 0$. However, potential scattering is a relevant perturbation with a scaling dimension $(1+K)/2<1$ and causes the conductance to decrease as $T^{1/K-1}$ as $T \to 0$; the leading irrelevant operator corresponds to the hopping (g_{LR}) term between the two leads with scaling dimension $(1/K+1)/2$. Similarly, near the 2CK fixed point reached for $K<1/2$ and $g_{LL}=g_{RR}$, the leading irrelevant operator (g_{LR} term) has dimension $1/2K$,[16] and therefore one expects $G(T) \propto T^{1/K-2}$ as $T \to 0$.

III. NONEQUILIBRIUM PROPERTIES

We study the low-temperature conductance in the high-bias regime $V \gg T_K$, T_K^* where the nonequilibrium RG method can be applied.[2] (In the opposite limit, $V \ll T_K^*$, we expect the equilibrium results quoted above to be valid after replacing $T \to V$.) After bosonization and refermionization, we can rewrite the Kondo model coupled to Luttinger leads as an effective noisy Kondo Hamiltonian involving free-electron leads.[12] This mapping is useful to write down the nonequilibrium RG equations in the ultraviolet limit. The effective Kondo model reads: $H=H_{lead}+H_K+H_{\phi_b}$, where $H_{lead}=\sum_{k,\alpha=L,R,\sigma=\uparrow,\downarrow}(\epsilon_k-\mu_\alpha)c_{k\alpha\sigma}^\dagger c_{k\alpha\sigma}$ corresponds to the noninteracting Fermi-liquid leads with bias voltages $\mu_{L/R} = \pm V/2$ and $H_K=\sum_{k,\alpha,\beta=L,R} J_{\alpha\beta} e^{i\phi_{\alpha\beta}} \vec{S}_{\alpha\beta}^e \cdot \vec{S}_d$ contains the phase-dependent Kondo coupling terms, $\vec{S}_{\alpha\beta}^e = c_{k\alpha\gamma}^\dagger \vec{\sigma}_{\gamma\gamma'} c_{k\beta\gamma'}$ and $\vec{S}_d = d_{\alpha\gamma}^\dagger \vec{\sigma}_{\gamma\gamma'} d_{\beta\gamma'}$ being the spin operators of the electrons for the effective noninteracting leads and of the dot, respectively. Here, $\phi_{\alpha\alpha}=0$, $\phi_{LR}=-\phi_{RL}=\phi_b$, and H_{ϕ_b} describes an

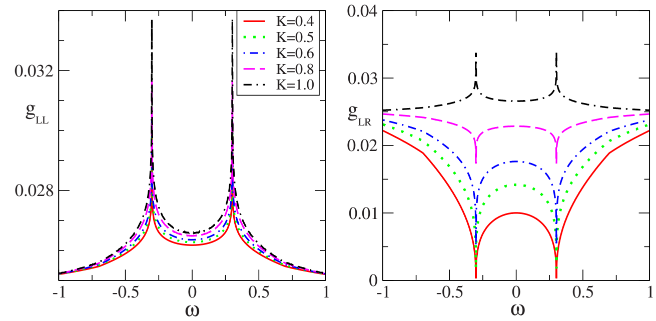

FIG. 1. (Color online) $g_{LL}(\omega)$ and $g_{LR}(\omega)$ for various Luttinger parameters K. The bare couplings are $g_{\alpha\alpha'}^0=g_{\alpha\alpha}^0=0.025$, resulting in $T_K^* \sim 4.2 \times 10^{-18}$ and $T_K \sim 2 \times 10^{-9}$ (in units of $D_0=1$). The bias voltage is $V=0.6 \gg T_K$, T_K^*.

effective Ohmic boson bath with the Green's function $G_{\phi_b}(i\omega) = \langle \phi_b(i\omega)\phi_b(-i\omega)\rangle = \pi(1/K-1)\frac{1}{|\omega|}$.[12] Following Refs. 2 and 12, we generalize the RG equations for the $J_{LL/RR/LR}$ terms to the nonequilibrium situation. These take the form

$$\frac{\partial g_{LL}(\omega)}{\partial \ln D} = -\sum_{\beta=-1,1}\left[g_{L\beta}\left(\frac{\beta V}{2}\right)\right]^2 \Theta_{\omega+\beta V/2},$$

$$\frac{\partial g_{LR}(\omega)}{\partial \ln D} = -\sum_{\beta=-1,1}\frac{1}{4}\left[1-\frac{1}{K}\right] g_{LR}\left(\frac{\beta V}{2}\right)\Theta_{\omega+\beta V/2} - g_{L\beta}\left(\frac{\beta V}{2}\right) g_{\beta R}\left(\frac{\beta V}{2}\right)\Theta_{\omega+\beta V/2}, \qquad (3)$$

where the dimensionless frequency-dependent Kondo couplings are given by $g_{\alpha\alpha}(\omega) \equiv J_{\alpha\alpha}(\omega)\rho_0$ with ρ_0 being the constant density of states of free-electron baths, $g_{\alpha\beta}(\omega) \equiv J_{\alpha\beta}(\omega)\rho_0(\frac{D}{D_0})^{1/2(1/K-1)}\Theta_{\omega-\mu_\beta}$, $\Theta_\omega = \Theta(D-|\omega+i\Gamma|)$, and $\beta=-1(+1)$ labels leads L(R). Here, we set $\hbar=k_B=e=1$. Note that within our one-loop RG scheme, the power-law tunneling density of states resulting from Luttinger liquid leads has been taken into account leading to the first term [linear in $g_{LR}(\omega)$] in the RG equation for $g_{LR}(\omega)$ [see Eq. (3)]. The symmetry of the RG equations gives $g_{LR}(\omega)=g_{RL}(\omega)$. Further, Γ is the decoherence (dephasing) rate at finite bias which cuts off the RG flow:[2]

$$\Gamma = \pi \sum_{\alpha\alpha'} \int d\omega f_\omega^\alpha (1-f_\omega^{\alpha'})[g_{\alpha\alpha'}(\omega)]^2, \qquad (4)$$

where the Fermi function obeys $f^\alpha(\omega)=1/(1+e^{(\omega-\mu_\alpha)/T})$. We note that there exists an additional contribution to Γ from electron dephasing caused by a finite potential drop in the Luttinger-liquid leads,[20] which will affect the subleading terms in Γ [given by $g_{LR}(V/2)$]. However, in the low-conductance regime of interest, this voltage drop is small and will be neglected henceforth. In general, the perturbative RG approach is valid for $V \gg T_{1CK}$, T_{2CK}. In the limit of $V \to 0$, Eq. (3) reduces to the equilibrium RG equations (with the flow cut off by temperature), and we recover Eq. (2).

The renormalized couplings are obtained by self-consistently solving Eqs. (3) and (4).[2] As shown in Fig. 1, $g_{LL(RR)}(\omega)$ exhibit peaks for all values of $K \leq 1$, indicating

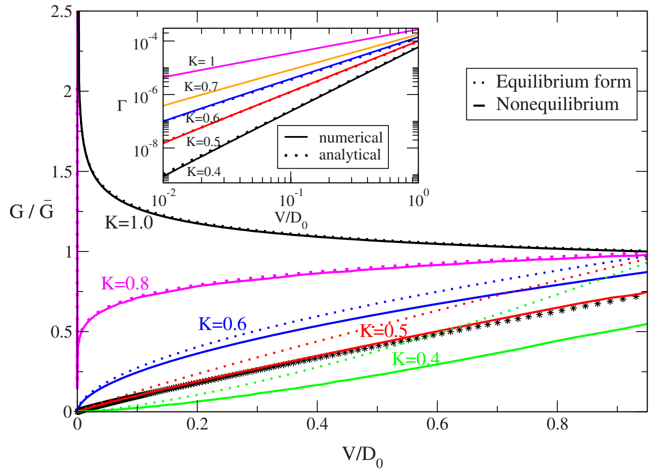

FIG. 2. (Color online) $G(V)$ normalized to $\bar{G}=3\pi(g^0_{\rm LR})^2/4$ for various K and Kondo couplings as in Fig. 1. Here, "equilibrium form" refers to the expression $G(V)=3\pi g_{\rm LR}(\omega=0)^2/4$. For $K=0.5$ we also show the analytical result from Eqs. (8) and (9) (stars). Inset: $\Gamma(\omega=0,V)$ for various K. The dashed lines are obtained from the analytical expression in Eq. (10). Here, $D_0=1$.

that they grow under RG. For a given bias voltage, the Kondo coupling $g_{\rm LR}(\omega)$ shows a crossover from peak to dip structure as K decreases, traducing the fact that for a fixed bias voltage, $g_{\rm LR}(\omega)$ is either enhanced or decreased compared to its bare value $g^0_{\rm LR}$. Let us emphasize that for sufficiently large bias voltages, as soon as $K<1$ the coupling $g_{\rm LR}(\omega)$ exhibits a dip close to $\omega=\pm V/2$, signaling 2CK behavior for symmetric Kondo channels. The singular behavior at the peaks or dips is cut off by the decoherence rate [see Eq. (7)] while outside that regime the voltage serves to cut off the RG flow.

From the Keldysh calculation up to second order in the tunneling amplitudes, the current reads

$$I=\frac{3\pi}{4}\int d\omega\left[\sum_\sigma g_{\rm LR}(\omega)^2 f^{\rm L}_\omega(1-f^{\rm R}_\omega)\right]-({\rm L}\leftrightarrow{\rm R}). \quad (5)$$

For small bare couplings $g^0_{\alpha\alpha'}=g^0_{\alpha\alpha}$ this perturbative calculation remains valid for $V>T^*_{\rm K}$, implying at high bias voltage, $V\gg T$, contributions to the current from a frequency window $-V/2<\omega<V/2$. For $T^*_{\rm K}<V\ll D_0$ with decreasing V we find that the differential conductance $G(V)\equiv dI/dV$ approaches the equilibrium form of the conductance $G(T\to V)$ $\propto V^{1/K-1}/\ln^4(V/T^*_{\rm K})$ with $G(T)\propto g_{\rm LR}(T)^2$ [see Eq. (2) and Fig. 2].

In the remainder, we analyze $G(V)$ for larger bias voltages. For $K=1$, we checked that the nonlinear conductance satisfies $G(V)\propto 1/\ln^2(V/T_{\rm K})$ for $V\gg T_{\rm K}$.[23] Here, one can replace $g_{\rm LR}(\omega)$ by $g_{\rm LR}(\omega=0)\approx g_{\rm LR}(T\to V)$. When decreasing K, the double peak structure in $g_{\rm LR}(\omega)$ at $\omega=\pm V/2$ turns progressively into dips which acquire a complex shape as a result of the decoherence rate Γ and the electron-electron interaction which hinders the interlead electron tunneling. The effect becomes more pronounced for small K values associated with the 2CK fixed point, rendering the "flat" approximation $g_{\rm LR}(\omega)\approx g_{\rm LR}(\omega=0)$ not justified; see Fig. 2.

To gain an analytical understanding of the small-K nonequilibrium regime, we may treat $g_{\rm RL}(\omega)$ within the interval $-V/2<\omega<V/2$ as a semiellipse.[22] The current I reads

$$I\approx\frac{3\pi}{4}\left[\frac{\pi}{4}g_{\rm LR}(\omega=0)^2+\left(1-\frac{\pi}{4}\right)g_{\rm LR}(\omega=V/2)^2\right]. \quad (6)$$

For $K<1$, we manage to obtain an approximate analytical form for the couplings $g_{\rm LR}(\omega=V/2)$ and $g_{\rm LR}(\omega=0)$. Solving Eq. (3) in the limit $D\to 0$, we find

$$g_{\rm LR}(\omega=0)\approx g_{\rm LR}(T\to V)\mathcal{F}(K)$$

$$g_{\rm LR}\left(\frac{V}{2}\right)\approx\frac{4\left(\frac{\Gamma V}{D^{*2}}\right)^{1/4(1/K-1)}}{\ln^2\left(\frac{\Gamma V}{(T^*_{\rm K})^2}\right)}, \quad (7)$$

where $g_{\rm LR}(T\to V)$ is the equilibrium form of $g_{\rm LR}$ in Eq. (2) with T replaced by V, and we have defined $\mathcal{F}(K)$ $=2^{1+1/4(1-1/K)}-1$ with $\mathcal{F}(K=1)=1$. Using Eqs. (6) and (7) and $G=dI/dV$, we obtain a closed expression for the conductance

$$G(V)\approx\frac{3\pi^2}{16}\left(\frac{V}{D^*}\right)^{1/K-1}\mathcal{R}(V)+12\pi\left(1-\frac{\pi}{4}\right)\mathcal{W}'(V), \quad (8)$$

where $\mathcal{W}'(V)=d\mathcal{W}/dV$ and

$$\mathcal{R}(V)=\mathcal{F}(K)^2\left[\frac{1/K}{\ln^4\left(\frac{V}{T^*_{\rm K}}\right)}-\frac{4}{\ln^5\left(\frac{V}{T^*_{\rm K}}\right)}\right],$$

$$\mathcal{W}(V)=\frac{V(\Gamma V/D^{*2})^{1/2(1/K-1)}}{\ln^4[\Gamma V/(T^*_{\rm K})^2]}. \quad (9)$$

For completeness, we have kept the less dominant contribution $\sim 1/\ln^5(V/T^*_{\rm K})$ in $\mathcal{R}(V)$. To rigorously define the function $\mathcal{W}(V)$, we need to provide an analytical expression for the decoherence rate in Eq. (4). Using an analogous reasoning as for the nonequilibrium current I [$\Gamma=\frac{4}{3}I$ at $T=0$ via Eqs. (4) and (5)], to second order in $g^0_{\rm LR}$, we extract [see Eqs. (6) and (7)]

$$\Gamma\approx\frac{\pi^2}{4}\mathcal{F}(K)^2\frac{V\left(\frac{V}{D^*}\right)^{1/K-1}}{\ln^4\left(\frac{V}{T^*_{\rm K}}\right)}. \quad (10)$$

Close to $K=1/2$, we can safely neglect contributions in $(g^0_{\rm LR})^{2+(1/K-1)}$ and therefore to second order in $g^0_{\rm LR}$, we find $\Gamma\approx(\pi^2/4)[g_{\rm LR}(\omega=0)]^2$. We have checked our analytical expression of Γ against a numerical treatment of Eqs. (3) and (4); see inset in Fig. 2. Notably, the decoherence rate contributes to a "distinct" power law $\sim V^{1/2(1/K^2-1)}$ in the nonequilibrium conductance $G(V)$, where $(1/K^2-1)/2>1/K$ -1 for $K<1$, rendering the second term in Eq. (8) to be subleading. The conductance becomes smaller than its equilibrium counterpart since $g_{\rm LR}(\omega=V/2)<g_{\rm LR}(\omega=0)$. A comparison between the analytical formula in Eq. (8) and the numerical integration of Eqs. (3)–(5) are shown in Fig. 2. As

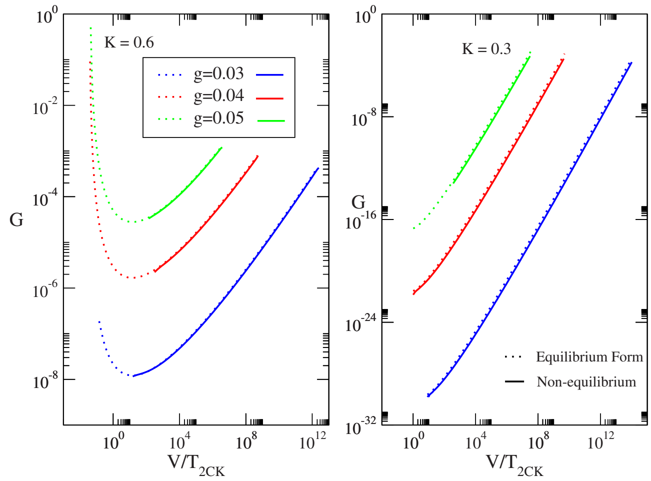

FIG. 3. (Color online) G versus V/T_{2CK} for K=0.6 (left panel) and K=0.3 (right panel) for various initial Kondo couplings $g^0_{\alpha\alpha'}=g^0_{\alpha\alpha}=g$ (in units of $D_0=1$).

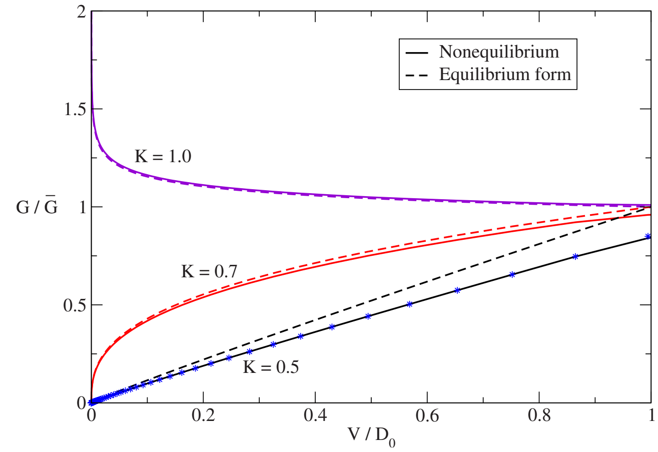

FIG. 4. (Color online) $G(V)$ normalized to $\bar{G}=3\pi(g^0_{LR})^2/4$ for various K with channel asymmetric bare Kondo couplings: $g^0_{LL}=0.02D_0$, $g^0_{RR}=0.01D_0$, and $g^0_{LR}=\sqrt{g^0_{LL}g^0_{RR}}$; $D_0=1$. Equilibrium form refers to that defined in Fig. 2; symbols (stars) are analytic nonequilibrium $G(V)$ at $K=1/2$.

our results are based on one-loop RG, we may expect both corrections to the power-law prefactors and further subleading terms upon including higher-loop contributions.

Our results also show that $G(V)$ for voltages $T_{2CK}\ll V\ll D_0$ is *not* an universal function of V/T_{2CK} (even for fixed K). Figure 3 displays G versus V/T_{2CK} for various initial Kondo couplings with T_{2CK} extracted from the RG flow. Note that due to the logarithmic scale for V/T_{2CK} shown in Fig. 3(left), the difference in equilibrium and nonequilibrium conductances is barely noticeable for K=0.6 though the distinction between them is clearly seen in the linear plot of $G(V)$ (see Fig. 2). This difference becomes noticeable for a smaller value of K=0.3 [Fig. 3(right)]. As it becomes also clear from Eqs. (8) and (9), the nonequilibrium conductance for $V\gg T_{2CK}$ is a function of both V/D_0 and V/T^*_K and hence is nonuniversal. This is again related to the fact that transport arises from the subleading term g_{LR}.

Finally, we have extended our results in the presence of a finite channel asymmetry $g^0_{LL}\neq g^0_{RR}$. We find that the nonequilibrium conductance follows the same deviation from its equilibrium counterpart (see Fig. 4). Results are perfectly consistent with the fact that the elliptic dip of $g_{LR}(\omega)$ is characteristic of large interactions in the leads. As shown in Fig. 4, for $K=1/2$ the analytical results for the nonequilibrium $G(V)$ (stars) which involves two distinct Kondo energy scales (associated with the two leads) agree well with the numerics (blue solid line).

IV. CONCLUSIONS

In conclusion, we have studied nonequilibrium transport through a Kondo dot coupled to Luttinger-liquid leads and calculated the conductance profile at bias voltages larger than the Kondo scales of the system. In particular, for symmetric couplings with the leads, the RG flow at large bias shows signatures of intermediate 2CK physics for all Luttinger parameters $K<1$. On the other hand, as the conductance G arises from the coupling g_{LR} which is subleading, $G(T\to 0)$ is not a universal function of V/T_{2CK} as it also depends on V/D_0. Our results push forward the knowledge of correlation effects in nanosystems far from equilibrium and should stimulate further experimental works on transport through dots coupled to quantum wires and carbon nanotubes.

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