Magnetic field effects on the finite-frequency noise and ac conductance of a Kondo quantum dot out of equilibrium

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We present analytic results for the finite-frequency current noise and the nonequilibrium ac conductance for a Kondo quantum dot in presence of a magnetic field. We determine the line shape close to resonances and show that while all resonances in the ac conductance are broadened by the transverse spin relaxation rate; the noise at finite field additionally involves the longitudinal rate as well as sharp kinks resulting in singular derivatives. Our results provide a consistent theoretical description of recent experimental data [Phys. Rev. Lett. **108**, 046802 (2012)] for the emission noise at zero magnetic field, and we propose the extension to finite field for which we present a detailed prediction.

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Introduction.—The understanding of quantum many-body effects and their characteristic signatures in transport properties represents a fundamental topic in mesoscopic physics. Beside the average current, its fluctuations described by the current noise contain additional information on the interplay of strong correlations and quantum fluctuations. In particular, the finite-frequency noise reveals the characteristic time scales of the system and provides information about the dynamics of excitations. This triggered numerous experimental works on the noise in various systems ranging from Josephson junctions to single-electron transistors [1].

It is by now well established that strong correlations play a crucial role for the transport properties of quantum dots. For example, for quantum dots in the so-called Kondo regime the transport is dominated by spin fluctuations leading, at sufficiently low energies, to a universal conductance of $G = 2e^2/h$ due to resonant tunneling processes [2]. Recently it has also become possible to measure the current noise in such Kondo quantum dots realized in carbon-nanotube devices [3, 4]. In particular, Basset *et al.* [4] measured the finite-frequency emission noise and observed resonances when the external frequency equaled the applied bias voltage.

The nonequilibrium finite-frequency noise in quantum dots has theoretically been studied for the Anderson model, resonant level models, and spin valve systems [5]. For quantum dots in the Kondo regime previous studies focused either on the shot noise (zero frequency) [6] or on the exactly solvable Toulouse limit [7], while the finite-frequency noise has only very recently started to attract attention [4, 8]. Of particular interest in this context is the nontrivial interplay of the different energy scales, which manifests itself in the appearance of characteristic resonances whose line shapes contain information about the underlying microscopic relaxation mechanisms. For quantum dots in the Kondo regime these are the transverse and longitudinal relaxation of the dot spin, which are identical at zero magnetic field, but acquire different values when the rotational symmetry is broken.

In this Letter we provide an analytic analysis of the finite-frequency current noise and the ac conductance in the nonequilibrium Kondo model. We particularly focus on the effects of a finite magnetic field and show that it leads to (i) characteristic resonances as a function of the frequency and bias voltage, and (ii) the appearance of both the longitudinal and transverse spin relaxation rates in the broadening of these resonances as well as sharp kinks in the noise. We find excellent agreement with existing experimental data [4] for the emission noise at zero magnetic field, and propose the measurement at finite field for which we present a detailed analysis (see Fig. 3).

Model.—We consider a Kondo quantum dot consisting of a spin-1/2 $\boldsymbol{S}$ subject to a local magnetic field h_0, which is coupled to two noninteracting electronic leads via an isotropic exchange interaction (see inset in Fig. 1),

$$H = H_{\rm res} + h_0 S^z + \frac{1}{2} \sum_{\alpha\alpha' kk' \sigma\sigma'} J_{\alpha\alpha'} a^\dagger_{\alpha k\sigma} \boldsymbol{S}\cdot\boldsymbol{\sigma}_{\sigma\sigma'} a_{\alpha' k'\sigma'}. \quad (1)$$

Here $a^\dagger_{\alpha k\sigma}$ and $a_{\alpha k\sigma}$ create and annihilate electrons with momentum k and spin $\sigma = \uparrow, \downarrow$ in lead $\alpha = L, R$, and $\boldsymbol{\sigma}$ are the Pauli matrices. The leads are described by $H_{\rm res} = \sum_{\alpha k\sigma} \varepsilon_k a^\dagger_{\alpha k\sigma} a_{\alpha k\sigma}$, with a flat density of states in a band of width $2D$, and chemical potentials $\mu_{L/R} = \pm V/2$. The exchange interaction is assumed to be derived from an Anderson impurity model via the Schrieffer-Wolff transformation and thus satisfies $J^2_{\rm nd} = J_L J_R$, where $J_{\rm nd} = J_{RL} = J_{LR}$ and $J_\alpha = J_{\alpha\alpha}$. We use the parametrization $J_{L/R} = 2x_{L/R} J_0$ with $x_L + x_R = 1$. The system is at zero temperature and we use units such that $e = \hbar = k_B = 2\mu_B = 1$.

The nonequilibrium dc current through Kondo quantum dots has been intensively studied [9–11] in the past. Here we investigate the zero-temperature fluctuations of the current in the stationary state, which are captured by the symmetric $S^+(\Omega)$ and antisymmetric $S^-(\Omega)$ current noise

$$S^{\pm}(\Omega) = \frac{1}{2}\int_{-\infty}^{\infty} \mathrm{d}t\, e^{i\Omega t}\langle [I(t)-\langle I\rangle, I(0)-\langle I\rangle]_{\pm}\rangle, \quad (2)$$

where $I = -\dot{N}_L$ denotes the current operator in the left lead defined using the corresponding particle number N_L [12]. In equilibrium, $S^+(\Omega)$ and $S^-(\Omega)$ are related by the fluctuation dissipation theorem (FDT), which at $T = 0$ reads $S^-(\Omega) = \mathrm{sgn}(\Omega)S^+(\Omega)$. In addition, we study the nonequilibrium ac conductance $G(\Omega)$ exploiting the generalization of the Kubo formula to nonequilibrium distributions [13].

Method.—We calculate the current noise using the real-time renormalization group (RTRG) approach [14, 15] which is based on a derivation of the effective dot Liouvillian $L_{\mathrm{D}}^{\mathrm{eff}}(z)$ depending on a Laplace variable z. $L_{\mathrm{D}}^{\mathrm{eff}}(z)$ governs the time evolution of the reduced density matrix of the dot $\rho_{\mathrm{D}}(t) = \mathrm{Tr}_{\mathrm{res}}\rho(t)$ obtained by tracing out the lead degrees of freedom from the full density matrix of the system. The effective Liouvillian incorporates all information about the relaxation dynamics of the spin on the dot encoded in the renormalized magnetic field h and the longitudinal and transverse spin relaxation rates Γ_1 and Γ_2.

In general, the noise (2) can be rewritten as the real part of the auxiliary function $C_{II}^{\pm}(\Omega) = \int_{-\infty}^{0} \mathrm{d}t\, e^{-i\Omega t}\langle [I(0), I(t)]_{\pm}\rangle$, which itself can be expressed [15] as $C_{II}^{\pm}(\Omega) = -i\mathrm{Tr}_{\mathrm{D}}\left[\Sigma_{II}^{\pm}(\Omega, i0^+)\rho_{\mathrm{D}}^{\mathrm{st}}\right] - i\mathrm{Tr}_{\mathrm{D}}\left[\Sigma_I(\Omega)\frac{1}{\Omega - L_{\mathrm{D}}^{\mathrm{eff}}(\Omega)}\Sigma_I^{\pm}(\Omega, i0^+)\rho_{\mathrm{D}}^{\mathrm{st}}\right]$. Here $\rho_{\mathrm{D}}^{\mathrm{st}}$ denotes the stationary reduced density matrix obtained from $L_{\mathrm{D}}^{\mathrm{eff}}(i0^+)\,\rho_{\mathrm{D}}^{\mathrm{st}} = 0$, and the kernels $\Sigma_{II}^{\pm}(\Omega, i0^+)$, $\Sigma_I(\Omega)$ and $\Sigma_I^{\pm}(\Omega, i0^+)$ obey renormalization-group equations similar to that of the Liouvillian [16].

The RTRG weak-coupling analysis is based on a systematic expansion in the renormalized exchange couplings around the poor man's scaling solution $J(\Lambda)$ given by $J(\Lambda) = [2\ln(\Lambda/T_K)]^{-1}$. Here Λ denotes the flow parameter, and the Kondo temperature is defined by $T_K = De^{-1/2J_0}$. Before Λ reaches T_K in the flow from high to low energy scales, that is in the range $\Lambda \geq \Lambda_c \gg T_K$, where $\Lambda_c = \sqrt{\Omega^2 + V^2 + h^2}$, we can carry out an expansion of the noise in a power series of $J(\Lambda_c)$. In doing so, we are able to identify which resonant features in the noise get broadened by relaxation rates and which remain sharp. The latter effect, in particular, happens in the nonequilibrium setup at finite magnetic field and $\Omega = \pm V$. Technically, this is seen in the RTRG equations as the influence of the resolvent projection $P_0\frac{1}{\Omega - L_{\mathrm{D}}^{\mathrm{eff}}(\Omega)}$ onto the zero eigenvalue subspace of the Liouvillian. This represents a nontrivial feature of the two-point functions (2), in contrast to one-point functions which receive no contribution from the zero eigenvalue subspace [14].

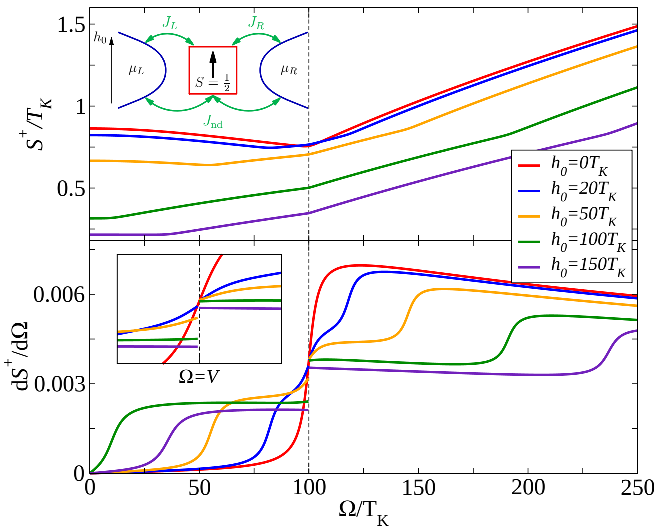

FIG. 1: (Color online) Upper panel: Symmetric noise $S^+(\Omega)$ for $V = 100\,T_K$, $r = 1$, and different magnetic fields. The inset shows a sketch of the Kondo dot. Lower panel: Derivative $\mathrm{d}S^+/\mathrm{d}\Omega$ showing a discontinuous jump at $\Omega = V$ for finite magnetic fields (zoom in the inset).

Finite-frequency noise.—We have analytically derived $S^{\pm}(\Omega)$ up to second order in the poor man's scaling solution $J = J(\Lambda_c)$. In the scaling limit ($D \to \infty$, $J_0 \to 0$ at fixed T_K) and for $\Omega \gg T_K$ they read

$$S^+(\Omega) = \pi J_{\mathrm{nd}}^2 Mh + \frac{\pi}{8}J_{\mathrm{nd}}^2\sum_{\sigma,\eta=\pm}|\Omega + \sigma V + \eta h|_2 + \frac{\pi}{2}J_{\mathrm{nd}}^2\sum_{\sigma=\pm}\left[M^2|\Omega + \sigma V| - \left(M^2 - \frac{1}{4}\right)|\Omega + \sigma V|_1\right] \quad (3a)$$

$$S^-(\Omega) = \frac{3\pi}{4}J_{\mathrm{nd}}^2\Omega + \frac{\pi}{4}J_{\mathrm{nd}}^2 M\sum_{\sigma,\eta=\pm}\eta|\Omega + \sigma V + \eta h|_2, \quad (3b)$$

where $|x|_i = (2x/\pi)\arctan(x/\Gamma_i)$ is the absolute value function smeared on the scale Γ_i. The longitudinal and transverse relaxation rates are given by $\Gamma_1 = \pi(J_L^2 + J_R^2)|h|/2 + \pi J_{\mathrm{nd}}^2\max\{|V|, |h|\}$ and $\Gamma_2 = \pi J_{\mathrm{nd}}^2|V|/2 + \Gamma_1/2$ respectively [9]. The dot magnetization is $M = -(1+r)^2 h/[2(1+r^2)|h| + 4r\max\{|V|, |h|\}]$, with the renormalized magnetic field $h = (1-J)h_0$ and the asymmetry $r = x_L/x_R$. One easily verifies $S^{\pm}(\Omega) = \pm S^{\pm}(-\Omega)$. In equilibrium, these $S^{\pm}$ also satisfy the FDT.

We stress that the RTRG method provides a consistent derivation of the relaxation rates appearing in (3) via the smeared absolute value function. In particular, at $\Omega = \pm V$ we find a contribution to S^+ which is not broadened by any microscopic decay rate. As shown in Fig. 1 the sharp kink in S^+ leads to a *discontinuity* in its derivative with the jump given by $\Delta = \pi J_{\mathrm{nd}}^2 M^2$. We attribute this behavior to the fact that at large fields, $h > V$, the spin on the dot is fixed to its ground state.

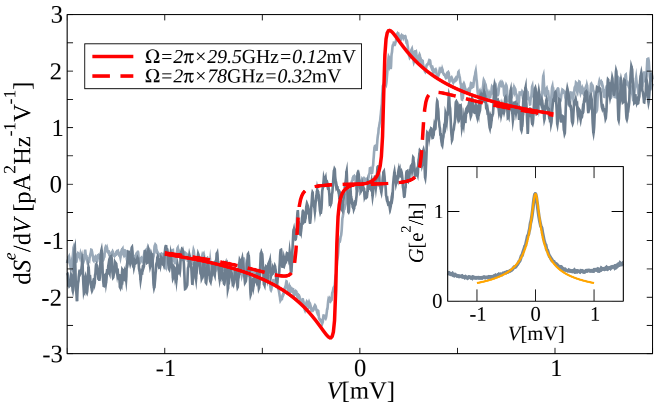

FIG. 2: (Color online) Comparison of Eq. (5) to the experimental data of Ref. [4] for the derivative of the emission noise $\mathrm{d}S^e/\mathrm{d}V$ at $h_0 = 0$. We stress that Eq. (5) does not contain any free parameter. Inset: Fit of $G(V)$ to the theoretical result [11].

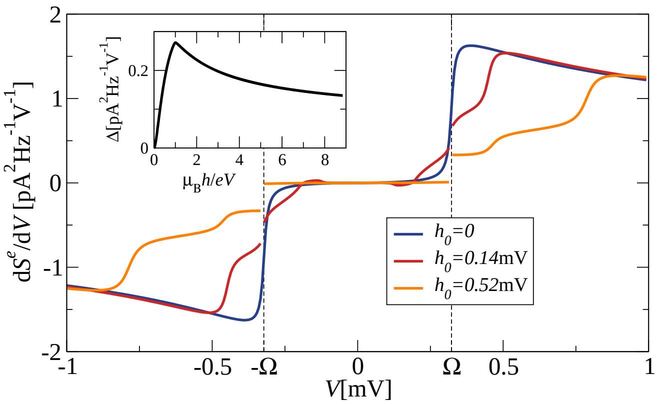

FIG. 3: (Color online) Voltage derivative of the emission noise for the experimental parameters of Ref. [4] with $\Omega/2\pi = 78\,\mathrm{GHz}$ (blue line). For comparison we show the result for small (red line) and large (orange line) magnetic fields. For finite fields we observe discontinuous jumps at $V = \pm\Omega$. The relation between h_0 and the applied field is given by $h_0 = g^* h_{\mathrm{app}}/2$ with the material specific effective g-factor g^*. Inset: Magnetic-field dependence of the jump height.

Processes at external frequencies $\Omega = V$ probe the charge transfer between the leads, which are not broadened due to the sharpness of the Fermi edges at zero temperature. At smaller fields, $0 < h < V$, the spin becomes dynamical, and virtual processes involving longitudinal spin fluctuations give an additional, *continuous*, contribution broadened by Γ_1. In turn, processes involving a spin flip on the dot, which appear at $\Omega = |V \pm h|$, are broadened by Γ_2. The latter behavior is also found for all resonances appearing in the current [9]. Thus the noise offers a way to study richer relaxation phenomena than those present in the current. We note that a discontinuity in the derivative of the noise was also found [7] in the strong-coupling regime of the Kondo model at the Toulouse point; we therefore expect it to be a generic feature of the finite-frequency noise in Kondo quantum dots. On the other hand, $S^-(\Omega)$ contains only terms broadened by Γ_2. This behavior is reflected in the ac conductance and will be discussed below.

For $\Omega \gg T_K$ the irreducible contribution to $C^{\pm}_{II}(\Omega)$ (given by its first term) is dominant, while the reducible one (the second term) is subleading $\sim O(J^4)$. However, for $\Omega \ll T_K$ the reducible term contributes in order J^2 supplementing Eq. (3a) in the limit $\Omega \to 0$ by

$$-\frac{\pi^2 J^4_{\mathrm{nd}}}{2\Gamma_1}\left[2VM + \left(M^2 + \frac{1}{4}\right) m(V,h)\right] m(V,h), \qquad (4)$$

where $m(V,h) = |V+h|_2 - |V-h|_2$. In total, this result generalizes the nonequilibrium shot noise [6] of a Kondo quantum dot to the case of finite magnetic fields.

Emission noise.—From the expressions (3) we determine $S^e(\Omega) = S^+(\Omega) - S^-(\Omega)$ describing the noise induced by photon emission [17]. It obviously inherits the features of the symmetric noise discussed above, which can be probed in the measurements of $\mathrm{d}S^e(\Omega)/\mathrm{d}\Omega$ or $\mathrm{d}S^e(\Omega)/\mathrm{d}V$.

Let us first consider vanishing magnetic field as in the recent experiments by Basset *et al.* [4] on the emission noise of a carbon nanotube quantum dot in the Kondo regime. For this case Eq. (3) yields

$$\frac{\mathrm{d}S^e(\Omega)}{\mathrm{d}V} = \frac{3}{4} J^2_{\mathrm{nd}} \left(\arctan\frac{\Omega+V}{\Gamma} - \arctan\frac{\Omega-V}{\Gamma}\right), \quad (5)$$

where $J^2_{\mathrm{nd}} = x_L(1-x_L)/\ln^2(\sqrt{\Omega^2+V^2}/T_K)$ and $\Gamma = \Gamma_1 = \Gamma_2 = \pi J^2_{\mathrm{nd}} V$. We emphasize that Eq. (5) contains only two unknown parameters, namely the Kondo temperature T_K and the asymmetry x_L, which are extracted from the differential conductance (see below). The previous theoretical analysis of the data in Ref. [4] used a frequency-dependent renormalization-group analysis which required, however, the fitting of the line shape close to the resonances with phenomenological relaxation rates. Here, in contrast, the rate Γ was derived consistently and does not contain free fit parameters.

In order to determine T_K and x_L we fit the measured differential conductance [4] to the theory [11] (see inset of Fig. 2). The asymmetry is extracted [18] from $G(V=0) = 1.194\, e^2/h$ and amounts to $x_L \approx 0.82$ (or 0.18), while the Kondo temperature is obtained from $G(V = T^*_K) = \frac{2}{3} G(V=0)$ and $T^*_K = 10.57\, T_K$ [19] and equals $T_K \approx 110\,\mathrm{mK} \approx 0.01\,\mathrm{mV}$. Using these parameters we plot Eq. (5) against the experimental results [4] in Fig. 2. In the range $|V| \lesssim 1\mathrm{mV}$ we find excellent agreement for both frequencies; for larger voltages charge fluctuations set in, and the Kondo model (1) is no longer adequate. For this reason our analysis of the features at $V = \pm\Omega$ is limited to $\Omega \lesssim 1\mathrm{mV}$. On the other hand, it is restricted by the weak-coupling condition $\Omega \gg T_K \approx 0.01\,\mathrm{mV}$, leaving two orders of magnitude in the window of admissible frequencies.

Here we propose to measure the emission noise of a quantum dot in the Kondo regime at finite magnetic field (see Fig. 3). As can be easily inferred from Eqs. (3) the energy scale h has two effects on $\mathrm{d}S^e(\Omega)/\mathrm{d}V$: (i) It introduces additional features $V = \pm|\Omega + h|$, $V = \pm|\Omega - h|$, and $V = \pm h$, which originate from the onset of additional transport processes as well as from the voltage dependence of the dot magnetization M. (ii) At large magnetic fields $h > \Omega$ the resonances at $V = \pm\Omega$ turn into discontinuous jumps. At smaller fields these jumps are superimposed with a contribution broadened by the longitudinal spin relaxation rate Γ_1, while all other resonances are broadened by the transverse rate Γ_2. For illustration we show $\mathrm{d}S^e(\Omega)/\mathrm{d}V$ for the parameters of Ref. [4] but finite magnetic fields in Fig. 3. In experiments the jumps at $V = \pm\Omega$ will be broadened by finite temperature T. This broadening is linear in T, in contrast to the other resonances which are broadened by $\Gamma_2 + \mathcal{O}(T)$. For this reason the features at $V = \pm\Omega$ stay much sharper than all other ones as long as $T \ll \Gamma_2 \sim T_K$.

AC conductance.—Finally we discuss the nonequilibrium ac conductance. We consider a setup [20] at finite dc bias V modulated by a small ac voltage δV, $V(t) = V + \delta V\, e^{-i\Omega t}$. This induces a frequency-dependent current $I(V, \delta V, \Omega)$ from which the nonequilibrium ac conductance can be extracted via $G(\Omega) = \lim_{\delta V \to 0} \frac{1}{\delta V}[I(V, \delta V, \Omega) - I(V)]$ with $I(V)$ denoting the stationary dc current. We stress that $G(\Omega)$ is the ac conductance in a nonequilibrium stationary state, i.e. in the presence of the finite dc bias V. By generalizing the Kubo formula to nonequilibrium distributions, $G(\Omega)$ can be related [13] to the auxiliary function $C^-_{II}(\Omega)$ introduced above via

$$G(\Omega) = \frac{1}{\Omega}\left[C^-_{II}(\Omega) - C^-_{II}(0)\right]. \qquad (6)$$

Hence we find $\mathrm{Re}\, G(\Omega) = S^-(\Omega)/\Omega$ and

$$\mathrm{Im}\, G(\Omega) = -\frac{M}{2\Omega} J^2_{\mathrm{nd}} \sum_{\sigma,\eta=\pm} \eta \Big[\mathcal{L}_2(\Omega + \sigma V + \eta h) - \mathcal{L}_2(\sigma V + \eta h)\Big], \qquad (7)$$

where $\mathcal{L}_2(x) = x\, \ln(\Lambda_c/\sqrt{x^2 + \Gamma_2^2})$. One can check that $\mathrm{Re}\, G(\Omega) = \mathrm{Re}\, G(-\Omega)$ and $\mathrm{Im}\, G(\Omega) = -\mathrm{Im}\, G(-\Omega)$ obey the Kramers-Kronig relations. Similarly to $S^+(\Omega)$, for small frequencies $G(\Omega)$ is supplemented by the contribution $\frac{\pi}{2} J^2_{\mathrm{nd}} m(V,h)\,(\partial M/\partial V)$ for $V > h$, where $\partial M/\partial V = -M\Gamma_1^{-1}(\partial\Gamma_1/\partial V)$, originating from the reducible part of $C^-_{II}(\Omega)$. Thus in the limit $\Omega \to 0$ one recovers the nonequilibrium dc conductance [9], while for $V \to 0$ one obtains the equilibrium ac conductance [21]. We note that in contrast to the symmetric noise no feature occurs at $\Omega = V$. Moreover, all resonances are broadened by the transverse rate Γ_2, thus the finite-frequency noise contains more information on the relaxation processes than the conductance.

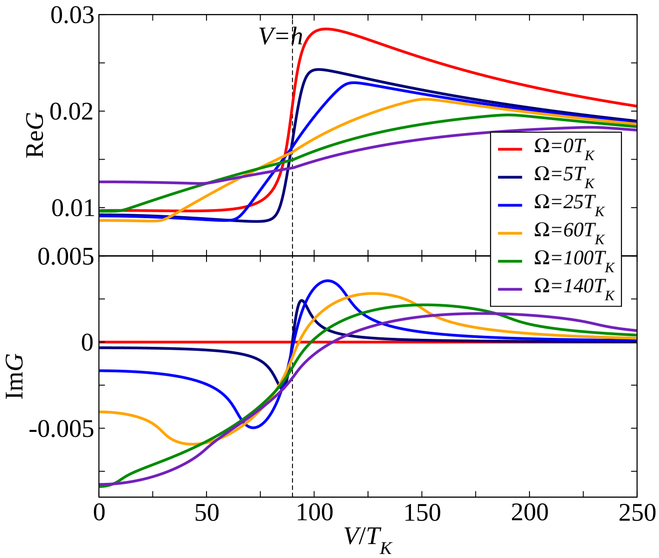

FIG. 4: (Color online) Nonequilibrium ac conductance $G(\Omega)$ [22] for $h_0 = 100\, T_K$, $r = 1$, and different frequencies Ω.

In Fig. 4 we plot the voltage dependence of $G(\Omega)$ at fixed frequency. The real part exhibits a characteristic enhancement at $V = h$ due to the onset of inelastic cotunneling processes. For finite frequencies this step-like enhancement is replaced by a continuous increase in the range $V = |h \pm \Omega|$ with reduced height. The slight change in the slope at $V = h$ is due to the voltage dependence of the dot magnetization. The dependence of $G(\Omega)$ as a function of the frequency shows a qualitatively similar behavior with features at $\Omega = |V \pm h|$.

Conclusion.—We have studied the effects of a finite magnetic field on the finite-frequency current noise and nonequilibrium ac conductance of a Kondo quantum dot. Due to the interplay of the different energy scales both observables exhibit various resonances, close to which the lineshapes are governed by self-consistently derived decay rates. In particular, at finite magnetic field the noise possesses a sharp feature at $\Omega = \pm V$ which is not broadened by any microscopic decay rate. We propose to measure the emission noise of a Kondo quantum dot at finite magnetic field, for which we have derived the full line shape including the characteristic resonances and a discontinuous jump in its derivative.

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