Kondo Effect in Quantum Dots Coupled to Ferromagnetic Leads

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We study the Kondo effect in a quantum dot coupled to ferromagnetic leads and analyze its properties as a function of the spin polarization of the leads. Based on a scaling approach, we predict that for parallel alignment of the magnetizations in the leads the strong-coupling limit of the Kondo effect is reached at a finite value of the magnetic field. Using an equation of motion technique, we study nonlinear transport through the dot. For parallel alignment, the zero-bias anomaly may be split even in the absence of an external magnetic field. For antiparallel spin alignment and symmetric coupling, the peak is split only in the presence of a magnetic field, but shows a characteristic asymmetry in amplitude and position.

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The Kondo effect [1] in electron transport through a quantum dot (QD) with an odd number of electrons is experimentally well established [2,3]. Screening of the dot spin due to the exchange coupling with lead electrons yields, at low temperatures, a Kondo resonance. The main goal of the present work is to investigate how ferromagnetic leads influence the Kondo effect. In the extreme case of half-metallic leads, minority-spin electrons are completely absent; i.e., the screening of the dot spin is not possible, and no Kondo-correlated state can form. What happens, however, for the generic case of partially spin-polarized leads? How does the spin-asymmetry affect the Kondo effect? Is there still a strong-coupling limit, and how are transport properties modified?

Based on a renormalization group scaling analysis, we first show that the strong-coupling limit can still be reached in this case if an external magnetic field is applied. This is familiar from the Kondo effect in QDs with an even number of electrons [4–7], which occurs at finite magnetic fields, although the physical mechanism is different in the present case. In the second part of the Letter we analyze within an equation of motion (EOM) approach the nonlinear transport through the QD. We find that for parallel alignment of the lead magnetizations the zero-bias anomaly is split. This splitting can be removed by appropriately tuning the strength of an external magnetic field B. In the antiparallel configuration of the lead magnetizations no splitting occurs at zero field.

The Anderson Hamiltonian for a QD with a single level at energy ϵ_0 coupled to ferromagnetic leads is

$$H = \sum_{rk\sigma}\varepsilon_{rk\sigma}c^\dagger_{rk\sigma}c_{rk\sigma} + \epsilon_0\sum_{\sigma}d^\dagger_\sigma d_\sigma + U d^\dagger_\uparrow d_\uparrow d^\dagger_\downarrow d_\downarrow + \sum_{rk\sigma}(V_{rk}d^\dagger_\sigma c_{rk\sigma} + V^*_{rk}c^\dagger_{k\sigma}d_\sigma) + g\mu_B B S_z, \qquad (1)$$

where $c_{rk\sigma}$ and d_σ are the Fermi operators for electrons with wave vector k and spin σ in the leads, $r = L, R$, and in the QD, V_{rk} is the tunneling amplitude, $S_z = (d^\dagger_\uparrow d_\uparrow - d^\dagger_\downarrow d_\downarrow)/2$, and the last term is the Zeeman energy of the dot. (Stray fields from the leads are neglected.) We assume identical leads and symmetric coupling, $V_{Lk} = V_{Rk}$. The ferromagnetism of the leads is accounted for by different densities of states (DOS) $\nu_{r\uparrow}(\omega)$ and $\nu_{r\downarrow}(\omega)$ for up- and down-spin electrons.

In the following we study the two cases of parallel (P) and antiparallel (AP) alignment of the leads' magnetic moments. For the AP configuration and zero magnetic field and bias voltage, the model is equivalent (by canonical transformation [8]) to a QD coupled to a single lead with DOS $\nu_{L\uparrow} + \nu_{R\uparrow} = \nu_{L\downarrow} + \nu_{R\downarrow}$. In this case, the usual Kondo resonance forms, which is the same as for nonmagnetic electrodes [1].

This changes for the P configuration. In this case, there is an overall asymmetry for up and down spins, say, $\nu_{L\uparrow} + \nu_{R\uparrow} > \nu_{L\downarrow} + \nu_{R\downarrow}$. To understand how this asymmetry affects the Kondo physics we apply the scaling technique [9], performed in two stages [10]. In the first stage, charge fluctuations dominate and lead to a renormalization of the QD's levels. Since the renormalization for the spin-down level is stronger than for the spin-up level, a level splitting between the two spin orientations is generated. This is one of the key mechanisms for all the effects discussed below. To reach the strong-coupling limit, it is, therefore, essential to apply an external magnetic field to compensate for the generated spin splitting. In the second stage, the resulting model is mapped onto a Kondo Hamiltonian, and the degrees of freedom involving spin fluctuations are integrated out. For simplicity we assume for the scaling analysis flat DOSs $\nu_{r\sigma}$ and neglect the k dependence of the tunnel amplitudes $V_{rk} = V$.

First we reduce the cutoff D from D_0, which is the smaller value of the bandwidth and the on-site repulsion U [10]. Charge fluctuations lead to the scaling equations

$$\frac{d\epsilon_\sigma}{d\ln(D_0/D)} = \frac{\Gamma_{\bar\sigma}}{2\pi}, \tag{2}$$

where we defined $\Gamma_\sigma = 2\pi|V|^2\sum_r \nu_{r\sigma}$, and $\bar\sigma$ is opposite to σ. This yields the solution $\Delta\tilde\epsilon = \tilde\epsilon_\uparrow - \tilde\epsilon_\downarrow = -(1/\pi)P\Gamma\ln(D_0/D) + \Delta\epsilon_0$, where $P = (\nu_{r\uparrow} - \nu_{r\downarrow})/(\nu_{r\uparrow} + \nu_{r\downarrow})$ measures the spin polarization in the leads, $\Gamma = (\Gamma_\uparrow + \Gamma_\downarrow)/2$,and $\Delta\epsilon_0 = g\mu_B B$ is the Zeeman splitting. The empty-dot state $|0\rangle$ hybridizes with states where the dot is singly occupied $|1\sigma\rangle$ with either spin-up or spin-down, while the singly occupied state $|1\sigma\rangle$ hybridizes only with the empty-dot state $|0\rangle$ (for $U \gg |\epsilon|$—asymmetric Anderson model). Because of the spin-dependent DOS in the leads the hybridization is spin dependent, which is the physical origin of the generated $\Delta\tilde\epsilon$.

To describe Kondo physics (for $\tilde\epsilon < 0$) we terminate [10] the scaling of Eq. (2) at $\tilde D \sim -\tilde\epsilon$, and perform a Schrieffer-Wolff transformation. Using the renormalized parameters $\tilde D$ and $\tilde\epsilon$, we get the effective Kondo Hamiltonian

$$H_{\mathrm{Kondo}} = \sum_{rr'kk'}\{J_+S^+c^\dagger_{rk\downarrow}c_{r'k'\uparrow} + J_-S^-c^\dagger_{rk\uparrow}c_{r'k'\downarrow} + S_z(J_{z\uparrow}c^\dagger_{rk\uparrow}c_{r'k'\uparrow} - J_{z\downarrow}c^\dagger_{rk\downarrow}c_{r'k'\downarrow})\}, \tag{3}$$

plus the term $-2S_z\tilde D(\nu_\uparrow J_{z\uparrow} - \nu_\downarrow J_{z\downarrow})$ and a potential scattering term. The initial values for the coupling constants are $J_+ = J_- = J_{z\uparrow} = J_{z\downarrow} = |V|^2/|\tilde\epsilon| \equiv J_0$. To reach the strong-coupling limit we tune the external magnetic field B such that the total effective Zeeman splitting vanishes, $\Delta\tilde\epsilon = 0$ (the field B will also slightly modify the DOS in the leads [6]). During the second stage of scaling, spin fluctuations will renormalize the three coupling constants $J_+ = J_- \equiv J_\pm$, $J_{z\uparrow}$, and $J_{z\downarrow}$ differently. The scaling equations are

$$\frac{d(\nu_\pm J_\pm)}{d\ln(\tilde D/D)} = \nu_\pm J_\pm(\nu_\uparrow J_{z\uparrow} + \nu_\downarrow J_{z\downarrow}), \tag{4}$$

$$\frac{d(\nu_\sigma J_{z\sigma})}{d\ln(\tilde D/D)} = 2(\nu_\pm J_\pm)^2 \tag{5}$$

with $\nu_\pm = \sqrt{\nu_\uparrow\nu_\downarrow}$, $\nu_\sigma \equiv \sum_r \nu_{r\sigma}$ [11]. To solve these equations we observe that $(\nu_\pm J_\pm)^2 - (\nu_\uparrow J_{z\uparrow})(\nu_\downarrow J_{z\downarrow}) = 0$ and $\nu_\uparrow J_{z\uparrow} - \nu_\downarrow J_{z\downarrow} = J_0 P(\nu_\uparrow + \nu_\downarrow)$ is constant as well. That is, there is only one independent scaling equation. All coupling constants reach the stable strong-coupling fixed point $J_\pm = J_{z\uparrow} = J_{z\downarrow} = \infty$ at the Kondo energy scale, $D \sim k_B T_K$. For the P configuration, the Kondo temperature in leading order,

$$T_{\mathrm{K}}(P) \approx \tilde D \exp\left\{-\frac{1}{(\nu_\uparrow + \nu_\downarrow)J_0}\frac{\operatorname{arctanh}(P)}{P}\right\}, \tag{6}$$

depends on the polarization P in the leads. It is a maximum for nonmagnetic leads, $P = 0$, and vanishes for $P \to 1$.

Finally, we point out an interesting consequence of the spin polarization in the leads. With nonmagnetic leads, the Kondo Hamiltonian couples the spin of the QD to the spin of the leads only, but not to its charge. To analyze the analogous situation in our case, we introduce the (pseudo) spin $\vec\sigma = (1/2)\sum_{kk'rr'\sigma\sigma'}c^\dagger_{kr\sigma}\sigma_{\sigma\sigma'}c_{k'r'\sigma'}/(2\tilde D\sqrt{\nu_\sigma\nu_{\sigma'}})$, where the spin-dependent normalization factor is crucial to ensure the proper spin commutation relations, and the (pseudo) charge $en = e\sum_{kk'rr'\sigma}c^\dagger_{kr\sigma}c_{k'r'\sigma}/(2\tilde D\nu_\sigma)$. The last term in Eq. (3) can, then, be written as $2\tilde D(\nu_\uparrow J_{z\uparrow} + \nu_\downarrow J_{z\downarrow})\sigma_z S_z$ plus $\tilde D(\nu_\uparrow J_{z\uparrow} - \nu_\downarrow J_{z\downarrow})nS_z$. The first term is analogous to the Kondo model with nonmagnetic leads, while the second term couples spin to charge. The latter does not scale up and the associated additional renormalization of the Zeeman splitting, $-(1/\pi)P\Gamma$, is negligible as compared to $\Delta\tilde\epsilon$ in the limit $D_0 \gg |\epsilon|$.

The unitary limit for the P configuration can be achieved by tuning the magnetic field appropriately, as discussed above. In this case, the maximum conductance through the QD is $G^{\mathrm{P}}_{\max,\sigma} = e^2/h$ per spin, i.e., the same as for nonmagnetic leads. This yields that the amplitude of the Kondo resonance for up- and down-spins at the Fermi level are different, since $G^{\mathrm{P}}_{\max,\sigma} \sim \Gamma_\sigma(E_{\mathrm{F}})\rho_\sigma(E_{\mathrm{F}})$ and, therefore, $\rho_\uparrow(E_{\mathrm{F}})/\rho_\downarrow(E_{\mathrm{F}}) = (1-P)/(1+P)$. For the AP configuration, the maximal conductance is reduced, $G^{\mathrm{AP}}_{\max,\sigma} = (1-P^2)e^2/h$, and vanishes for $P \to 1$.

In the remainder of this Letter we analyze the DOS of the QD and address nonequilibrium transport. For a qualitative discussion, we should employ the simplest technique which accounts for both the formation of Kondo resonances and the influence of the spin-dependent renormalization of the dot level on spin fluctuations. The EOM technique with the usual decoupling procedure [12,13] for higher-order Green functions satisfies the first requirement but not the second. We, therefore, extend this scheme by calculating the level splitting $\Delta\tilde\epsilon$ self-consistently. For a more quantitative analysis, one could include higher-order (than usual) Green's functions in the EOM approach or higher-order diagrams in the resonant-tunneling approximation [14], or use more advanced schemes such as real-time [15] or numerical renormalization group [16] methods. These techniques are, however, much more complex [17] and are not pursued here.

Within the Keldysh formalism, the transport current $I = \sum_\sigma I_\sigma$ through a QD for $\Gamma_{R\sigma}(\omega) = \lambda_\sigma\Gamma_{L\sigma}(\omega)$ is

$$I_\sigma = \frac{e}{\hbar}\int d\omega\frac{\Gamma_{L\sigma}(\omega)\Gamma_{R\sigma}(\omega)}{\Gamma_{L\sigma}(\omega) + \Gamma_{R\sigma}(\omega)}[f_L(\omega) - f_R(\omega)]\rho_\sigma(\omega), \tag{7}$$

where $\rho_\sigma(\omega) = -(1/\pi)\,\mathrm{Im}\,G^{\mathrm{ret}}_\sigma(\omega)$. For strong interaction ($U \to \infty$), the retarded Green's function is

$$G^{\mathrm{ret}}_\sigma(\omega) = \frac{1 - \langle n_{\bar\sigma}\rangle}{\omega - \epsilon_\sigma - \Sigma_{0\sigma}(\omega) - \Sigma_{1\sigma}(\omega) + i0^+}, \tag{8}$$

where $\Sigma_{0\sigma}(\omega) = \sum_{k\in L,R}|V_k|^2/(\omega - \varepsilon_{k\sigma})$ is the self-energy for a noninteracting QD, while

$$\Sigma_{1\sigma}(\omega, \Delta\tilde{\epsilon}) = \sum_{k\in L,R} \frac{|V_k|^2 f_{L/R}(\varepsilon_{k\bar{\sigma}})}{\omega - \sigma\Delta\tilde{\epsilon} - \varepsilon_{k\bar{\sigma}} + i\hbar/2\tau_{\bar{\sigma}}} \quad (9)$$

appears for interacting QDs only. The average occupation of the QD with spin σ is obtained from $\langle n_\sigma\rangle = -(i/2\pi)\int d\omega\, G^<_\sigma(\omega)$. Extending the standard derivation [12], we replaced on the right-hand side of Eq. (9) $\Delta\epsilon \rightarrow \Delta\tilde{\epsilon}$, where $\tilde{\epsilon}_\sigma$ is found self-consistently from the relation

$$\tilde{\epsilon}_\sigma = \epsilon_\sigma + \mathrm{Re}[\Sigma_{0\sigma}(\tilde{\epsilon}_\sigma) + \Sigma_{1\sigma}(\tilde{\epsilon}_\sigma, \Delta\tilde{\epsilon})], \quad (10)$$

which describes the renormalized dot-level energy, where the real part of the denominator of Eq. (8) vanishes [1]. We emphasize that without this self-consistency relation the Kondo resonances will, in general, appear at different positions, which disagree with the conclusions from the scaling analysis [19]. The procedure simulates higher-order contributions and the influence of the renormalization of the dot-level on spin fluctuations. Following Ref. [12] we introduce, in a heuristic way, a lifetime $\tau_\sigma(\mu_L, \mu_R, \tilde{\epsilon}_\uparrow, \tilde{\epsilon}_\downarrow)$ which describes decoherence due to a finite bias voltage V or level splitting $\Delta\tilde{\epsilon}$. It is obtained in second-order perturbation theory and depends on the electrochemical potentials in the leads, $\mu_{L(R)}$, and the level positions. Again, we replace the bare levels by the renormalized ones. In the numerical results presented below we use Lorentzian bands of width $D = 100\Gamma$.

For nonmagnetic leads, $P = 0$, and zero magnetic field, $B = 0$, the proposed approximation is identical to the standard EOM scheme [12]. For finite magnetic field, $B \neq 0$, the self-consistency condition yields a splitting of the Kondo resonances which is slightly smaller than $2g\mu_B B$, in agreement with both experimental [2] and theoretical findings [20,21]. For $B = 0$ and $P > 0$ in the parallel configuration, we obtain a value of the splitting $\Delta\tilde{\epsilon}$ comparable to the result from scaling, Eq. (2).

In Fig. 1 we plot the DOS of the QD for spins in AP and P configurations with spin polarization $P = 0.2$ in the leads. In the AP configuration there is one Kondo resonance [Fig. 1(a)] and the DOS is the same as for the case of nonmagnetic leads. For the P configuration, however, the Kondo resonance splits [Fig. 1(c)], which can be compensated by an external magnetic field B [Fig. 1(d)]. In the latter case, the amplitude of the Kondo resonance for spin-down significantly exceeds that for spin-up (as discussed above). A finite bias voltage, $\mu_R - \mu_L = eV > 0$, again leads to a splitting for both the AP and the P configurations [Figs. 1(b) and 1(e)]. In the AP configuration, the amplitudes of the upper and the lower Kondo peaks appear asymmetric [Fig. 1(b)].

In Fig. 2 we show the differential conductance as a function of the transport voltage. For nonmagnetic leads, there is a pronounced zero-bias maximum [Fig. 2(a)], which splits in the presence of a magnetic field [Fig. 2(b)]. For magnetic leads and parallel alignment,

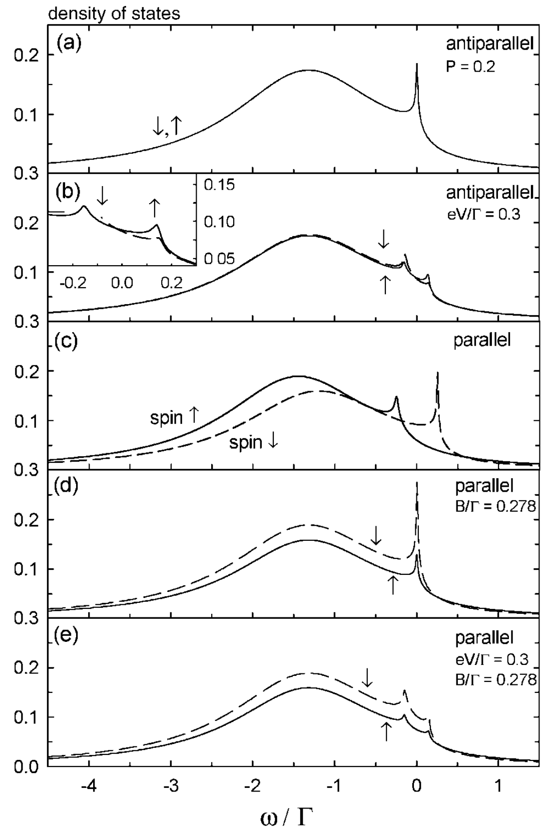

FIG. 1. Spin-dependent DOS for spin-up (solid line) and spin-down (dashed line), calculated for P and AP alignment (as indicated), for a spin polarization of the leads $P = 0.2$. The parts (d),(e) include the effect of an applied magnetic field B and (b),(e) of an applied bias voltage V. The other parameters are $T/\Gamma = 0.005$ and $\epsilon/\Gamma = -2$.

we find a splitting of the peak in the absence of a magnetic field [Fig. 2(c)], which can be tuned away by an appropriate magnetic field [Fig. 2(d)]. In the AP configuration, the opposite happens, no splitting at $B = 0$ [Fig. 2(e)] but finite splitting at $B > 0$ [Fig. 2(f)] with an additional asymmetry in the peak amplitudes as a function of the bias voltage. This asymmetry is related to the asymmetry in the amplitude of DOS [Fig. 1(b)]. All these findings are in good agreement with our scaling analysis. In Fig. 2(g) we show the tunnel magnetoresistance (TMR) which can be much larger than for conventional TMR systems. Finally, we find that the positions of the peaks in the AP configuration in the presence of a magnetic field are slightly shifted as a function of the polarization P [Fig. 2(h)]. This can be explained in the similar way as in Ref. [7] by an additional level splitting $\delta\Delta\tilde{\epsilon} = (1/4\pi)P\Gamma/\tilde{\epsilon}\, eV$ at finite bias voltages due to spin accumulation in the QD.

We finally comment on the observability of our proposal and how one can attach ferromagnetic leads to a QD. A conceivable realization might be to put carbon nanotubes in contact with ferromagnetic leads [22]. The Kondo effect has been observed already for nonmagnetic

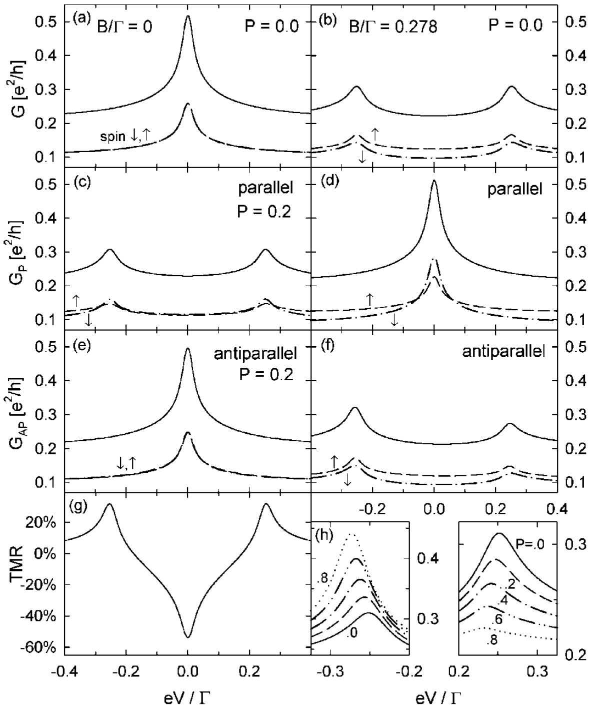

FIG. 2. Total differential conductance (solid lines) as well as the contributions for spin-up (dashed lines) and spin-down (dot-dashed lines) vs the applied bias voltage V at zero magnetic field $B = 0$ (a),(c),(e) and at finite magnetic field (b),(d),(f),(h) for normal (a),(b) and ferromagnetic leads with parallel (c),(d) and antiparallel (e),(f),(h) alignment of the lead magnetizations. (g) The tunnel magnetoresistance, $\mathrm{TMR} = (G_\mathrm{P} - G_\mathrm{AP})/G_\mathrm{AP}$, for the cases (c) and (e). (h) The conductance $G_\mathrm{AP}/(1 - P^2)$ for several values of P as indicated. Other parameters are as in Fig. 1

electrodes [5,23]. Alternatively, one might use magnetic tunnel junctions with magnetic impurities in the barrier, or spin-polarized STM [24,25].

In conclusion, we presented a qualitative study of the Kondo effect in QDs coupled to ferromagnetic leads. In particular, we found a splitting of the Kondo resonance for parallel alignment of the leads' magnetizations, even in the absence of a magnetic field. Our results are based on a scaling approach and an EOM technique. Further investigations on a more quantitative level using more advanced techniques are desirable.

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