Nonequilibrium Transport through a Kondo-dot in a Magnetic Field

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Abstract. Electron transport through a quantum-dot in the Coulomb blockade regime is modeled by a Kondo type hamiltonian describing spin-dependent tunneling and exchange interaction with the local spin. We consider the regime of large transport voltage V and magnetic field B with $\max(\mathrm{V,B}) \gg \mathrm{T_K}$, the Kondo temperature, and show that a renormalized perturbation theory can be formulated describing the local magnetization M and the differential conductance G quantitatively. Based on the structure of leading logarithmic corrections in bare perturbation theory we argue that the perturbative renormalization group has to be generalized to allow for frequency dependent coupling functions. We simplify the full RG equations in the spirit of poor man's scaling and calculate M and G in leading order of $1/\ell n[(V,B)/T_K]$.

The transport of electrons through a quantum dot in the limit of weak coupling to the leads is dominated by Coulomb interaction effects, forcing integral electron charge on the dot [1]. Adding or removing electrons requires a large energy, the charging energy E_c, and thus transport is blocked (Coulomb blockade). Only if energy levels of the addition spectrum of the dot are in the range between the chemical potentials of the two leads transport is possible. The differential conductance as a function of a gate voltage shifting the energy levels of the dot is then characterized by maxima separated by Coulomb blockade valleys. In the case that the total spin of the dot is nonzero, however, the antiferromagnetic exchange interaction of this local spin with the conduction electron spins in the leads gives rise to a Kondo resonance in the local density of states at the Fermi level. Electron transport may then take place via resonance tunneling, thus removing the Coulomb blockade [2,3]. The dramatic effect of the formation of the Kondo resonance at low temperatures in filling up the Coulomb blockade valley in the differential conductance, eventually restoring maximum conductance, has been seen in a number of experiments [4–8].

The Kondo resonance appears as a reaction of the conduction electron system to a continuous sequence of spin flip processes involving the local spin. From the point of view of the conduction electron system a local spin flip amounts to a sudden change of the impurity potential formed by the local spin. Since the scattering states of the conduction electron system before and after the switching are orthogonal, the system takes infinite time to adjust to the new situation (orthogonality catastrophe, [9]). In the course of this adjustment the system forms particle-hole excitations in numbers increasing with an inverse power of excitation energy. Since this process is interrupted when the next spin flip occurs, the infrared singularity in the spectrum associated with the orthogonality catastrophe is cut-off and a smooth resonance peak at the Fermi level is formed instead.

Powerful methods have been developed to describe this strong coupling problem, among them methods based on the renormalization group idea, in a simplified form [10] ("poor man's scaling") and in a full numerical implementation by Wilson [11], exact methods for calculating the thermodynamic quantities using the Bethe ansatz [12] and conformal field theory [13], and resummations of perturbation theory using auxiliary particle representations [14] to access dynamical quantities.

Necessary prerequisites for the existence of the Kondo effect are (1) the spin degeneracy of the ground state of the quantum dot, and (2) quantum coherence over sufficiently long time periods. The characteristic energy scale against which both of these perturbations have to be measured is the Kondo temperature $T_K \simeq D\sqrt{JN_0}\exp(-\frac{1}{2N_0J})$, where J is the exchange coupling constant, N_0 is the conduction electron density of states per spin at the Fermi level and D is the conduction band width.

A minimal spin exchange interaction model of a quantum dot features three exchange coupling constants, $J_{LL}(J_{RR})$ coupling the local spin to the left (right) lead and $J_{LR} = J_{RL}$ describing the spin-dependent tunneling through the dot. Within a standard Anderson impurity model, characterized by hopping amplitudes t_L, t_R from the dot to the left and right lead, respectively, one obtains by way of a Schrieffer-Wolff transformation $J_{\gamma\gamma'} = t_\gamma t_{\gamma'} / \mid \epsilon_d \mid, \gamma, \gamma' = L, R$ where ϵ_d is the energy of the dot level closest to the Fermi level. Within such a model one has the relation $J_{LR} = (J_{LL}J_{RR})^{1/2}$. More complex models of a quantum dot allowing for several orbitals and additional interactions may lead to less restrictive relations between the J's [16]. Of particular interest are quantum dots with $J_{LR} \ll J_{LL} = J_{RR}$, as in that case a new fixed point of two-channel character governs the behavior of the system over a wide range of energies [15], but probably not down to the lowest energies [17,18], depending on the applied bias voltage V.

In this paper we consider the effect of a large difference in the chemical potentials μ_L, μ_R in the two leads, generated by a bias voltage $V = \mu_L - \mu_R$. We model the leads as non-interacting electron systems. The voltage is driving a steady state–current I. We assume the temperature T to be uniform over the sample. At sufficiently low temperatures, $T \leq T_K$, the effect of the spin flip processes will be to generate resonance structures in the local density of states at both Fermi edges, μ_L and μ_R. At the same time, at finite bias voltage V, tunneling of electrons through the dot creates a momentary nonequilibrium state in the leads which is equilibrated by dissipative processes. In this way the levels in the dot are broadened by inelastic processes even at $T = 0$, leading to a suppression of the Kondo effect. We will show that at least for sufficiently large J_{LR}, say $J_{LR} = J_{LL} = J_{RR}$, the system remains in the regime where renormalized perturbation theory is applicable.

Transport through a Kondo system has been considered by several authors. Early on, Appelbaum [19] calculated the differential conductance $G(V,B)$ for large bias voltage V in a finite magnetic field B, assuming erroneously, however, that the local spin remains in thermal equilibrium. König et al. [20] calculated $G(V,B)$ for an Anderson model, employing a certain resummation of perturbation theory. Brief accounts of the work to be presented below have been published elsewhere [15,21].

1 Magnetization and conductance in lowest order perturbation theory

We study the following Kondo-type model Hamiltonian of a quantum dot:

$$
\begin{aligned}
H = & \sum_{\mathbf{k},\sigma,\alpha=L,R} (\varepsilon_{\mathbf{k}} - \mu_\alpha) c^\dagger_{\alpha\mathbf{k}\sigma} c_{\alpha\mathbf{k}\sigma} - BS_z \\
& + \frac{1}{2} \sum_{\mathbf{k},\mathbf{k}',\sigma,\sigma',\alpha,\alpha'=L,R} J_{\alpha'\alpha}\, \mathbf{S} \cdot (c^\dagger_{\alpha'\mathbf{k}'\sigma'} \tau_{\sigma'\sigma} c_{\alpha\mathbf{k}\sigma}),
\end{aligned}
\tag{1}
$$

where $\mathbf{S}$ is the local spin operator on the dot (assumed to be $S = \frac{1}{2}$) and τ is the vector of Pauli matrices. The $c^+_{\alpha\mathbf{k}\sigma}$ create conduction electrons with momentum $\mathbf{k}$ and spin σ in lead α. The chemical potential shifts induced by the bias voltage are given by $\mu_{L,R} = \pm V/2$ for the left (L) and right (R) lead. The exchange energies $J_{\gamma\gamma'}$, $\gamma = L, R$, will be assumed to have the symmetry $J_{RR} = J_{LL}$ and $J_{LR} = J_{RL}$. We shall use the dimensionless coupling constants $g_d = N_0 J_{RR}$ and $g_{LR} = N_0 J_{LR}$, with N_0 the local density of states at the Fermi energy, assumed to be flat in the accessible regime $\mid \omega \mid \leq V, B$. A magnetic field is included, inducing a Zeeman splitting of the magnetic levels of the local spin $\omega_\gamma = -\gamma B/2$, $\gamma = \pm 1$. The effect of the magnetic field on the conduction electrons gives only rise to particle-hole asymmetry terms, which are usually small and may be neglected.

In lowest order perturbation theory the current through the dot is given by the Golden Rule expression

$$
\begin{aligned}
I = \frac{\pi e}{4\hbar} \int d\omega \sum_{\gamma=\pm 1} n_\gamma \Big[& g_1^2 f_{\omega-\mu_L} (1 - f_{\omega-\mu_R}) \\
& + 2 g_2^2 f_{\omega-\mu_L} (1 - f_{\omega-\mu_R-\gamma B}) \Big] - (L \leftrightarrow R)
\end{aligned}
\tag{2}
$$

where g_1 and g_2 are the dimensionless coupling constants for spin-nonflip and spin-flip interaction, $g_1 = g_2 = N_0 J_{LR}$, $f_{\omega-\mu_\alpha}$ is the Fermi function in the lead $\alpha = L, R$, and γB is the Zeeman energy transfer taking place in a spin flip process. At low temperatures, $T \ll V, B$, the product of Fermi functions limits the energy integration to the window $\mu_R < \omega < \mu_L$ and $\mu_R + \gamma B < \omega < \mu_L$, respectively. The current is seen to depend on the occupation $n_\gamma = \frac{1}{2}(1+\gamma M)$ of the local spin states, where M is the magnetization.

For sufficiently small voltage, $V \ll T$, the local spin system is in thermal equilibrium. In the opposite limit, $V \gg T$, the stationary current through the dot drives the system out of equilibrium. The occupation numbers n_γ are then determined by the rate equation $\frac{\partial}{\partial t} n_\gamma = C_\gamma\{n_{\gamma'}\} = 0$, where C is the collision integral. This leads to the condition (assuming $J_{LR} = J_{LL}$ here)

$$n_\uparrow \sum_{\alpha,\alpha'=L,R} \int d\omega g_2^2 f_{\omega-\mu_\alpha}\Big(1 - f_{\omega-\mu_{\alpha'}-B}\Big) = \\ = n_\downarrow \sum_{\alpha,\alpha'=L,R} \int d\omega g_2^2 f_{\omega-\mu_\alpha}\Big(1 - f_{\omega-\nu_{\alpha'}+B}\Big) \, . \tag{3}$$

The spin-flip coupling constant g_2 cancels out of this equation, yielding the nonequilibrium magnetization in the limit $J \to 0$,

$$M = 2/\Big[\coth\frac{B}{2T} + \frac{\sinh\frac{B}{T} - \frac{V}{B}\sinh\frac{V}{T}}{\cosh\frac{B}{T} - \cosh\frac{V}{T}}\Big] \, . \tag{4}$$

In the limit $V \to 0$, the equilibrium result $M = \tanh\frac{B}{2T}$ is recovered, whereas in the limit $V \gg B, T$, $M = \frac{2B}{V}$, independent of temperature. The result (4) has been obtained independently in [22,24].

2 Leading logarithmic corrections to M and I

In order to calculate higher order contributions in perturbation theory in J we switch now to a more formal description in terms of nonequilibrium Green's functions. We find it convenient to represent the local spin operator in pseudo fermion (PF) language [23],

$$\mathbf{S} = \frac{1}{2}\sum_{\gamma\gamma'} f_\gamma^+ \tau_{\gamma\gamma'} f_{\gamma'} \ , \tag{5}$$

where f_γ^+ creates a pseudofermion of spin $\gamma =\uparrow,\downarrow$ at the dot. The projection onto the physical sector of Hilbert space, with pseudofermion occupation number $Q = \sum_\gamma f_\gamma^+ f_\gamma = 1$, is done by adding a term λQ to the Hamiltonian and taking the limit $\lambda \to \infty$. This means that the PF system is taken in the low density limit. The Feynman diagrams of perturbation theory in the exchange interaction will therefore have one PF loop at most.

The local magnetization can be calculated in terms of the PF Green's function $G^<_\gamma(\omega)$ as $M = \sum_{\gamma=\pm 1} \gamma \int \frac{d\omega}{2\pi i} G^<_\gamma(\omega)$, where $G^<$ is found by solving a quantum kinetic equation. In steady state one finds

$$G^<_\gamma(\omega)\Gamma_\gamma(\omega) = \Sigma^<_\gamma(\omega) A_\gamma(\omega) \tag{6}$$

where $A_\gamma(\omega)$ is the PF spectral function, $\Sigma^<_\gamma(\omega)$ is the lesser component of the self-energy and $\Gamma_\gamma(\omega)$ its imaginary part. In general Eq. (6) is an integral equation for $G^<_\gamma(\omega) = in_\gamma(\omega)A_\gamma(\omega)$, which, however, can be approximated in a controlled way by making use of the sharply peaked form of $A_\gamma(\omega)$ to give an algebraic equation for $n_\gamma(0)$. In this form Eq. (6) corresponds to the

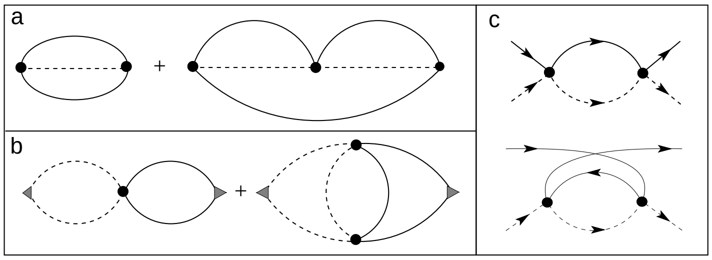

Fig. 1. Feynman diagrams for (a) PF self-energies, (b) current and (c) vertices entering the 1-loop RG equation. PF (electron-) propagators are displayed as dashed (full) lines.

rate equation discussed above, with higher order processes included in the transition amplitudes. The diagrams of the self-energy up to third order in J are depicted in Fig. 1a. Note that the lines represent Keldysh matrix Green's functions dressed with arrows in all topologically different ways. The result including leading logarithmic corrections is found as

$$M = \frac{Z}{\coth\frac{B}{2T}\left[\frac{Z}{2} + g_-^2 B(1+\mathcal{L}(B))\right] + g_{LR}^2(c(B)+c(-B))} \tag{7}$$

where

$$Z = 2g_+^2 B(1+\mathcal{L}(B)) + 2g_{LR}^2\Big[2B + (V+B)\mathcal{L}(V+B) - (V-B)\mathcal{L}(V-B)\Big] \tag{8}$$

and we use the abbreviations $g_\pm^2 = g_d^2 \pm g_{LR}^2$, $L(B) = 2g_d \ell n \frac{D}{|B|}$, and $c(B) = \coth\frac{V+B}{2T}[(V+B)(1+\mathcal{L}(V+B)) + V\mathcal{L}(V) + B\mathcal{L}(B)]$. We note that logarithmic corrections of order $g^3 \ell n D$ to the self-energies lead to corrections of order $g \ell n D$ to M. These corrections are larger than the leading logarithmic corrections in equilibrium ($V \to 0$), which are of order $g^2 \ell n D$. We have neglected these sub-leading corrections in (7) and note in passing that they can be obtained by including the energy shifts of the PF induced by ReΣ. On the other hand, the effects of ImΣ are important as will be discussed below.

The current through the dot may be expressed in terms of a spin-spin correlation function. Defining the conduction electron charge density operator in lead α at the dot, $\hat{n}_\alpha = e\sum_{\mathbf{k},\sigma} c^+_{\alpha\mathbf{k}\sigma} c_{\alpha\mathbf{k}\sigma}$, the current operator $\hat{I}$ can be expressed as

$$\hat{I} = -\frac{\partial}{\partial t}\hat{n} = \frac{ie}{\hbar} J_{LR} \mathbf{S} \cdot \sum_{\mathbf{k},\mathbf{k}',\sigma,\sigma'} \left(c^+_{L\mathbf{k}\sigma} \tau_{\sigma\sigma'} c_{R\mathbf{k}\sigma'} - c^+_{R\mathbf{k}\sigma} \tau_{\sigma\sigma'} c_{L\mathbf{k}'\sigma'} \right) . \tag{9}$$

The average current is given in terms of the Keldysh component of the contour-ordered correlation function

$$D_{LR}(t,t') = -i\langle \hat{T}_c \Big\{ \sum_{\sigma,\sigma'} c^+_{L\mathbf{k}'\sigma'}(t) \tau_{\sigma'\sigma} c_{R\mathbf{k}\sigma}(t) \cdot \mathbf{S}(t') \Big\} \rangle \tag{10}$$

as

$$I = -\frac{e}{\hbar} J_{LR} \mathrm{Re} \Big\{ D^K_{LR}(t,t) \Big\} . \tag{11}$$

In Fig. 1b the Feynman diagrams for D^K_{LR} to order J^3 are shown. A long but straightforward calculation yields the current including leading logarithmic corrections:

$$\begin{aligned} I \; &= \frac{e\pi}{2\hbar} g^2_{LR} \Big[\frac{3}{2} V + V\mathcal{L}(V) + (V+B)\mathcal{L}(V+B) \\ &\qquad + (V-B)\mathcal{L}(V-B) - M(c(B) - c(-B)) \Big] \end{aligned} \tag{12}$$

In the limit of zero magnetic field, Eq. (12) reduces to $I = \frac{e\pi}{4\hbar} g^2_{LR} 3V(1 + 2L(V))$, in agreement with Ref. [25].

3 Resummation of perturbation theory in nonequilibrium: a poor man's scaling approach

Even for small coupling constants g and for sufficiently large V and B,such that $V, B \gg T_K$, but still in the scaling regime $V, B \ll D$, such that $g\ell n D/\mid V \mid \ll 1$, bare perturbation theory converges slowly. It is necessary to sum the leaving logarithmic contributions in all orders of PT. In the equilibrium state a powerful method is available to perform this resummation in a controlled way: the perturbative renormalization group method [10]. It makes use of the fundamental idea that a change of the cut-off D can be fully absorbed into a redefinition of the coupling constants g. As long as the running coupling constant $g(D)$ is small, the change of g under an infinitesimal change of D,

$\partial g/\partial \ell n D$, may be calculated in perturbation theory. It is well known that in the equilibrium Kondo problem the coupling constant is found to grow to infinity, thus leaving the perturbative regime beyond $g \sim 1$. A nonperturbative treatment is then necessary, as shown by Wilson [11] in his pioneering work on the numerical RG. In the nonequilibrium situation, the RG flow is cut-off by inelastic processes already within the weak coupling regime, such that perturbation theory is valid [25,15]. This is true at least for the case $J_{LR} = J_{LL}, J_{RR}$, considered here. While the discussion of the RG flow in [15] was a qualitative one, motivated by results obtained in the so-called "Non-crossing approximation" for the Anderson model, and applied only to the limit $B = 0$, here we consider the RG formulation on a more fundamental level. A different and considerably more involved real-time RG scheme has been developed by Schoeller and König [26].

First we observe that a straightforward renormalization of g like in the equilibrium situation is not possible for $V \gg T_K$. In Eq. (7) logarithmic corrections appear (even in the limit $B \to 0$) as $2\ell n(D/V)$ in the denominator, but as $\ell n(D/V) + \ell n(D/T)$ in the numerator. This is related to the fact that the energy integrals in the numerator are confined to the vicinity of the Fermi energy, whereas in the denominator the energy integral covers a finite range of width V. This suggests that the frequency dependence of the couplings becomes important. In order to understand how these frequency dependencies are generated on the lowest level, i.e. in one-loop order, it is sufficient to analyze the behavior of the vertex corrections shown in Fig. 1c under a change of cut-off. Logarithmic corrections in perturbation theory are generated by PF - conduction electron bubbles containing the product of a real part of a PF Green's function, $1/(\omega \pm B/2)$, and the Keldysh component of the local conduction electron line, $-2\pi i N_0 \tanh[(\omega - \mu_\alpha)/2T]$. If the energy of the PF - conduction electron intermediate state is within the interval $[-D, D]$ the process will contribute to the renormalization of the coupling function $g(\omega)$, otherwise it will not. In general the coupling functions depend on three frequencies (taking energy conservation into account). Using the fact that the spectral function of the PFs is sharply peaked at $\omega = \mp B/2$, we set two of the three frequencies to $\omega = \mp B/2$, keeping only one frequency variable. In addition, the frequency dependence may be neglected within the running band-width, $\mid \omega \mid < D$. The spin structure of the general coupling functions is given by two invariant amplitudes, $\tilde{g}_2$ for spin flip and $\tilde{g}_1$ for spin non-flip processes:

$$\begin{aligned} g^{\alpha\sigma,\omega;\alpha'\bar{\sigma},\omega-\gamma B}_{\gamma,-\gamma B/2;\bar{\gamma},B/2} &= (\tau^x_{\gamma\bar{\gamma}}\tau^x_{\sigma\bar{\sigma}} + \tau^y_{\gamma\bar{\gamma}}\tau^y_{\sigma\bar{\sigma}})\, \tilde{g}_2(\omega - \gamma B/2) \\ g^{\alpha\sigma,\omega;\alpha\sigma,\omega}_{\gamma,-\gamma B/2;\gamma,-\gamma B/2} &= \tau^z_{\alpha\alpha}\tau^z_{\gamma\gamma}\tilde{g}_1(\omega)\ , \end{aligned} \tag{13}$$

where $g^{a;a'}_{b;b'}$ denotes the coupling function for conduction electrons of spin σ and energy ω in lead $\alpha(a = (\alpha, \sigma, \omega))$ interacting with a pseudo fermion in state $b = (\gamma, \omega_f)$ and going into states a', b'. The two frequency-dependent running coupling functions $\tilde{g}_1(\omega)$ and $\tilde{g}_2(\omega)$ obey the following flow equations,

$$
\begin{aligned}
\frac{\partial \tilde{g}_1(\omega)}{\partial \ln D} &= -\frac{\tilde{g}_2(\frac{B+V}{2})^2}{2}\left(\Theta_{\omega+B+\frac{V}{2}} + \Theta_{\omega-B-\frac{V}{2}}\right) \\
&\quad -\frac{\tilde{g}_2(\frac{B-V}{2})^2}{2}\left(\Theta_{\omega-B+\frac{V}{2}} + \Theta_{\omega+B-\frac{V}{2}}\right) \\
\frac{\partial \tilde{g}_2(\omega)}{\partial \ln D} &= -\frac{\tilde{g}_1(\frac{V}{2})\tilde{g}_2(\frac{B+V}{2})}{2}\left(\Theta_{\omega+\frac{B+V}{2}} + \Theta_{\omega-\frac{B+V}{2}}\right) \\
&\quad -\frac{\tilde{g}_1(\frac{V}{2})\tilde{g}_2(\frac{B-V}{2})}{2}\left(\Theta_{\omega-\frac{B-V}{2}} + \Theta_{\omega+\frac{B-V}{2}}\right) ,
\end{aligned}
\tag{14}
$$

with initial condition $\tilde{g}_1(\omega) = \tilde{g}_2(\omega) = JN_0$. In the limit $V, B \to 0$, Eq. (14) reduces to the well-known scaling equations (8).

So far the flow equations (14) reflect the nonequilibrium physics in two ways: they describe the renormalization of the coupling constants at the (up to four) different resonant energies, and they account for the variation with frequency in the regions between and somewhat outside the Fermi edges. There is, however, still one additional effect missing: the finite relaxation of the spins even in the limit $T \to 0$. Within the PF representation this effect shows up as a finite imaginary part of the PF self-energy. The relaxation is modified by vertex corrections, but is certainly not canceled altogether [15]. The relaxation rate may be interpreted as due to spin relaxation processes, leading to inelastic broadening of the local spin levels, thus cutting off the RG flow. While different rates exist for the different coupling functions, in leading approximation we may neglect the differences between them and assume a single rate Γ, which we identify with the transverse spin relaxation rate $\Gamma = 1/T_2$.

Assuming the RG flow to be stopped at the scale Γ, we replace the step functions Θ_ω in Eq. (14) by $\Theta(D - \sqrt{\omega^2 + \Gamma^2})$. The decay rate Γ has to be calculated self-consistently with the solution of the flow equations (14). Starting from the golden rule expression for the transverse relaxation rate,

$$
\begin{aligned}
\Gamma = \frac{\pi}{4\hbar} \sum_{\alpha,\alpha'=L,R,\gamma=\uparrow,\downarrow} \int d\omega \Big[&\tilde{g}_1(\omega)^2 f_{\omega-\mu_\alpha}(1 - f_{\omega-\mu_{\alpha'}}) \\
&+ \tilde{g}_2(\omega - \gamma B/2)^2 f_{\omega-\mu_\alpha}(1 - f_{\omega-\mu_{\alpha'}-\gamma B})\Big]
\end{aligned}
\tag{15}
$$

the renormalized Γ is obtained by replacing $g_{1,2}$ by $\tilde{g}_{1,2}(\omega)$ as determined from the solution of (2). In Fig. 2 the running coupling functions $\tilde{g}_1(\omega)$ and $\tilde{g}_2(\omega)$ are shown for different values of D. One can see how a Kondo resonance structure begins to form at the two Fermi energies, each one split by twice the Zeeman energy B. The growth of $\tilde{g}_{1,2}$ stops, however, due to the relaxation effects embodied in Γ, such that $\tilde{g}_{1,2}(\omega) \ll 1$ for all ω and D.

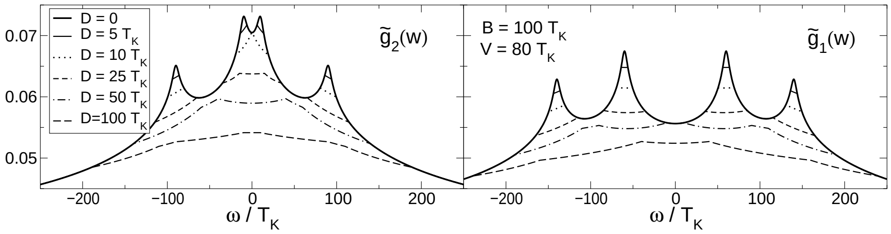

Fig. 2. Renormalized coupling constants $\tilde{g}_1(\omega)$ (right patel) and $\tilde{g}_2(\omega)$ (left panel) for $B = 100T_K$ and $V = 80T_K$ and for different values of the cut-off D.

We are now ready to calculate further physical quantities. The renormalized value of the magnetization is obtained by substituting $\tilde{g}_{1,2}(\omega)$ in place of $g_{1,2}$ in the Golden Rule expression (3). In Fig. 3b we show the fully renormalized result for M as a function of V/B for different values of B.

The charge current I is calculated from Eq. (2) inserting the renormalized coupling functions. Fig. 3a shows a comparison of the differential conductance $G(V) = dI/dV$ obtained in this way, with the bare result (2) and the PT result (12).

The peak structures in $G(V, B)$ appearing at $V = \pm B$ have been detected in experiment. In order to reach large values of B/T_K, it is necessary to have relatively low T_K. This happened to be the case in transport through metallic point contacts, containing a magnetic impurity [27]. In Fig. 3c the result of our theory using the value of $T_K \simeq 30mK$ given in [27] is compared to the experimental data, after subtracting a background contribution [21]. The agreement is seen to be excellent.

4 Conclusion

We considered the transport of electrons through a Kondo dot in the regime of large transport voltage V and in the presence of a magnetic field B, such that $\max(V, B) \gg T_K$ is satisfied. A finite difference of the chemical potentials in the two leads, $\mu_L - \mu_R = V$, opens an energy window for inelastic processes even at $T = 0$, leading to a broadening Γ of the spin sublevels of the dot and providing a cut-off for the scaling towards the Kondo fixed point. A finite magnetic field induces a local magnetization M at the dot, which for $V \gg T_K, T$ is independent of temperature, being solely controlled by the voltage V. This can be seen on an elementary level by considering the rate equations for the spin occupation number. The fact that M is controlled by V has been overlooked in early work from the sixties on this problem [19]. The differential conductance is likewise affected. We calculated the leading logarithmic corrections to M and G and found that the structure of these log-terms changes in nonequilibrium. An analysis of the reasons for these changes lead us to conjecture a renormalization group formulation allowing

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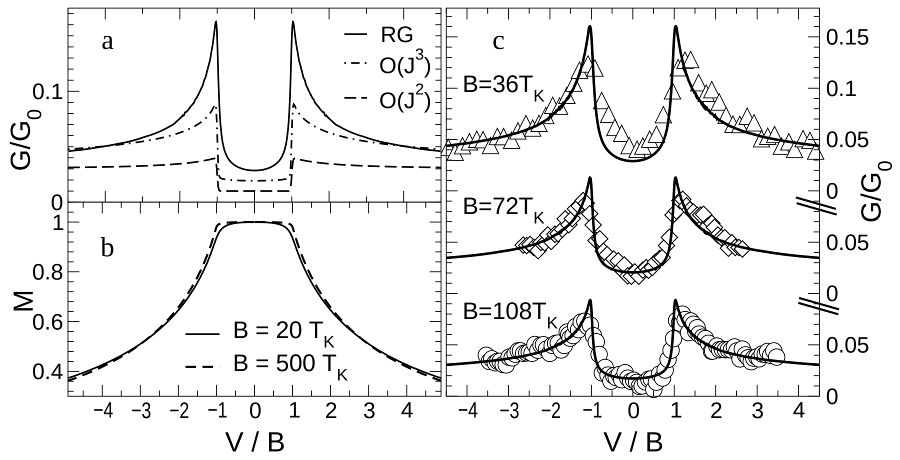

Fig. 3. a) Conductance $G(V/B)$ calculated in leading (dashed line) and next-to-leading (dot-dashed line) order PT compared to the result of perturbative RG (solid line) for $B = 100T_K$, $D = 10^4 T_K$, $T_K = D\sqrt{g}e^{-1/2g}$). b) Local magnetization $M(V/B, B/T_K)$ of a symmetric dot for fixed magnetic field $B = 20, 500\, T_K$. c) Experiments by Ralph and Buhrman [27] (symbols) on transport through metallic point contacts in magnetic fields $0.85, 1.7, 2.55\,\mathrm{T}$ ($B = 36, 72, 104\, T_K$ where $T_K \approx$ 30mK [27]). Assuming that the corresponding point contact is described by a single-channel ($J_{LR} = \sqrt{J_{LL}J_{RR}}$) Kondo model, $\frac{J_{RR}}{J_{LL}} \approx 4.2$ is determined from $G(V{=}0, B{=}0, T{=}50\mathrm{mK}) = \frac{4J_{LL}J_{RR}}{(J_{LL}+J_{RR})^2} G_{\mathrm{sym}}(T/T_K)$, where G_{sym} is known exactly from NRG calculations. This fixes *all* parameters for our RG calculation (solid lines) which uses a straightforward generalization of [21] for $J_{LL} \neq J_{RR}$. As the (B-dependent) background is not known experimentally, we subtract $\Delta G = G_B - 5.2 \cdot 10^{-5} G_0 \frac{V}{T_K}$, where G_B is fitted to our results at large V. Our calculations are at $T = 0$; the experimental $T = 50$mK leads to an extra small broadening at $V = B$.

for frequency dependent coupling functions [21]. We derived a set of RG equations within a poor man's scaling approximation, which for once reproduce the bare PT result, but then may be integrated to give the fully renormalized result. The integration is done including the self-consistently determined cut-off Γ. The results obtained in this way are valid up to corrections of order $1/\ell n[(V, B)/T_K]$. It is clear that the fully renormalized result features much more pronounced peaks than either Eq. (2) or (12), demonstrating the necessity of the RG treatment.

5 Acknowledgments

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