Mesoscopic features in the transport properties of a Kondo-correlated quantum dot in a magnetic field

Alberto Camjayi and Liliana Arrachea

Departamento de Física, FCEyN and IFIBA, Universidad de Buenos Aires, Pabellón 1, Ciudad Universitaria, 1428 Buenos Aires, Argentina

E-mail: alberto@df.uba.ar and lili@df.uba.ar

Received 2 October 2013, revised 20 November 2013
Accepted for publication 26 November 2013
Published 18 December 2013

Abstract
We study the transport behavior induced by a small bias voltage through a quantum dot connected to one-channel finite-size wires. We describe the quantum dot using the Hubbard–Anderson impurity model and we obtain solutions by means of a quantum Monte Carlo method. We investigate the effect of a magnetic field applied at the quantum dot in the Kondo regime. We identify mesoscopic oscillations in the conductance, which are introduced by the magnetic field. This behavior is analogous to that observed as a function of the temperature.

Keywords: Kondo, quantum dot, transport, Monte Carlo

(Some figures may appear in colour only in the online journal)

1. Introduction
The consequences of the Kondo effect for the transport properties of mesoscopic systems and nanostructures have attracted significant interest during the past two decades. After signatures of this effect were identified in quantum dots fabricated in semiconducting systems [1], it was later investigated in several other devices containing carbon nanotubes, and molecules attached to metallic electrodes [2–4].

All these systems contain a central piece, the quantum dot, which is connected to wires of non-interacting electrons. The Kondo effect arises as a consequence of the Coulomb repulsion, which gives rise to an effective coupling between the spin of a localized electron in the quantum dot and the spins of the electrons of the wires. For temperatures lower than the Kondo temperature $T_K$, these electrons form a singlet giving rise to a resonant state, which manifests itself as the opening of a transport channel for each spin component, corresponding to a conductance $G_0 = \frac{2e^2}{h}$ [5]. The electrons of the wires that intervene in the formation of these singlets define the so called ‘screening cloud’. The latter extends up to a length that is related to the Kondo temperature via $\xi_K \sim \frac{\hbar v_F}{k_B T_K}$, with $v_F$ being the Fermi velocity of the electrons of the wire. The nature of the wires is known to play a crucial role in the development of the Kondo effect and the concomitant behavior of $G$. In mesoscopic systems and nanostructures, where the wires have a finite length $L$ and a typical level spacing $\Delta$, the ratio $\xi_K / L$ (or, equivalently $\Delta / T_K$) defines different regimes for the behavior of the Kondo effect and $G$. This issue has been analyzed in detail for the case of clean wires in [6, 7]. When the mesoscopic wires are dirty, there is an interesting interplay between the length of the Kondo cloud and the localization length introduced by disorder, which has been recently analyzed [8].

Another interesting aspect is the influence of other magnetic phenomena in the formation of the Kondo state. This issue has been analyzed in the context of a Kondo impurity connected to a ferromagnetic environment [9], and for a quantum dot in an external magnetic field [10–15]. The formation of singlets leading to the development of Kondo resonance is clearly affected by the presence of an external magnetic field. In particular, the Zeeman splitting of the energy levels of the quantum dot destroys...
the Kondo resonance for strong enough magnetic fields. Interestingly, the conductance preserves the universality under the influence of the magnetic field, which means that the Kondo temperature $T_K$ sets the energy scale for the behavior of this quantity [5]. This property has been recently investigated experimentally [11] and theoretically [12]. All these analyses have been performed for the case of quantum dots with a perfect matching to the wires. They find that the thermal and the magnetic energy, scaled by the Kondo energy $k_B T_K$, play similar roles in the behavior of the conductance. The aim of the present work is to analyze whether this is also the case for a quantum dot embedded in mesoscopic wires with a finite length.

We model the dot as an Anderson impurity. An important aspect in the study of the Kondo effect in this model is tackling it with a reliable many-body technique. In this sense, treatments based on Bethe ansatz [16], slave-boson [17], perturbation theory in different limits [18], renormalization group [13, 14, 19], exact diagonalization [20] and quantum Monte Carlo methods [8, 21–26] have been widely used in the literature. Here we introduce one of the newly developed continuous time quantum Monte Carlo methods [26], in the context of transport properties in mesoscopic systems. In particular, we recover all the results presented for the conductance of a Kondo-correlated quantum dot embedded in mesoscopic wires [6, 7] as well as those for a magnetic field and wires with perfect matching [11, 12]. We then present results for finite-size wires and magnetic field and show that, in the case of perfect matching, the magnetic energy plays a similar role to the thermal energy in the behavior of the conductance. The paper is organized as follows. In section 2 we present the model, the expression for the conductance in terms of Green functions and the quantum Monte Carlo procedure followed to evaluate these functions. In section 3 we present and discuss the results without and with a magnetic field. Section 4 is devoted to a summary and conclusions.

2. The theoretical approach

2.1. The model

The full system is sketched in figure 1. It consists of a quantum dot modeled as a Hubbard–Anderson impurity connected to one-dimensional left (L) and right (R) wires. A magnetic field is applied at the dot. The ensuing Hamiltonian is

$$H = H_d + \sum_{\alpha=L,R} H_\alpha + H_{\text{cont.}}$$

(1)

where the Hamiltonian for the dot includes the effect of the Coulomb repulsion $U$, the voltage gate $V_g$ and the Zeeman splitting $\Delta_z = g \mu_B B$ due to the external magnetic field $B$; we assume $B$ to be applied along the $z$-direction, $\mu_B$ is the Bohr magneton and $g$ is the gyromagnetic factor. It reads

$$H_d = \sum_{\sigma=\uparrow,\downarrow} \epsilon_{d,\downarrow} n_{d,\sigma} + U n_{d,\uparrow} n_{d,\downarrow},$$

(2)

where $\epsilon_{d,\downarrow} = V_g \pm \Delta_z/2$. We distinguish a finite region within the wires with $N_a$ sites, between the dot and the macroscopic reservoirs. Following previous works [6, 7], we model this piece using a tight-binding Hamiltonian with nearest neighbor hopping element $t$. We also consider a gate voltage represented by a uniform local energy $\varepsilon_W$.

The reservoirs are represented by semi-infinite tight-binding chains with hopping elements $t$. We have

$$H_{\alpha} = \sum_{l=1}^{N_a-1} \left[ \varepsilon_W c_{l,\uparrow,\alpha}^{\dagger} c_{l,\uparrow,\alpha} - t (c_{l+1,\uparrow,\alpha}^{\dagger} c_{l,\uparrow,\alpha} + \text{H.c.}) \right]$$

$$- t \sum_{l=N_a+1}^{\infty} \left[ c_{l,\uparrow,\alpha}^{\dagger} c_{l+1,\uparrow,\alpha} + \text{H.c.} \right].$$

(3)

The contacts between the wires and the quantum dot are described by the Hamiltonian

$$H_{\text{cont}} = - t \sum_{\alpha=L,R} \left[ c_{1,\uparrow,\alpha}^{\dagger} d_{\alpha} + \text{H.c.} \right].$$

(4)

In the limit of vanishing magnetic field, $\Delta_z = 0$, the model reduces to the Kondo impurity with an effective exchange constant $J = 2t^2/(1 + V_g) - 1/V_g$. The Kondo temperature is defined from the condition [5]

$$\frac{2}{J} = \int_{-\infty}^{+\infty} \frac{d\omega}{\omega} \tanh \left( \frac{\omega}{2T_K} \right) \rho_0(\omega),$$

(5)

where $\rho_0(\omega) = -\text{Im} \sum_{\alpha} \sigma^{\dagger} \rho_{\sigma}^{\alpha}(\omega)$ is the density of states of the wires. The Kondo resonance is a consequence of the formation of a spin singlet state, between the impurity and the conduction electrons at the wires, which extends within a length $\xi_K = \hbar v_F/ (k_B T_K)$. A mismatching in the contact between the wires and the reservoirs, $t_C \neq t$, introduces finite-size features in $\rho_0(\omega)$, thus affecting the value of the Kondo temperature and the Kondo cloud, as discussed in the value of the Kondo temperature and the Kondo cloud, as discussed in [6, 7] and section 2.2.

2.2. Conductance

The conductance through the dot for an infinitesimal applied voltage $V$, such that the chemical potentials of the L, R
where \( g \) is connected to the wires. The Fermi function \( G \) depends on the coupling \( J \) and the local density of states of the wire 

\[
G = \frac{e^2}{\hbar} \int d\omega \frac{\partial f(\omega)}{\partial \omega} \frac{\Gamma_{L}(\omega) \Gamma_{R}(\omega)}{\Gamma_{L}(\omega) + \Gamma_{R}(\omega)} \rho_{\sigma}(\omega). \tag{6}
\]

The hybridization function with the wire 

\[
\Gamma_{\alpha}(\omega) = \Gamma \rho_{\alpha}^{R}(\omega), \quad \rho_{\sigma}(\omega) = -2\text{Im}[G_{\alpha,\sigma}^{R}(\omega)]. \tag{7}
\]

depends on the coupling \( J \) and the density of states of the wire \( \rho_{\alpha}^{R}(\omega) = -2\text{Im}[g_{\alpha}^{R}(\omega)] \), where \( g_{\alpha}^{R}(\omega) \) is the retarded Green function of the wire \( \alpha \) connected to the corresponding reservoir but uncoupled from the dot. \( G_{\alpha,\sigma}^{R}(\omega) \) is the retarded Green function of the quantum dot connected to the wires. The Fermi function \( f(\omega) = 1/(1 + e^{(\omega-\mu)/T}) \) depends on the temperature \( T = 1/(\beta \hbar) \), as well as on the mean chemical potential of the reservoirs \( \mu \). At \( T = 0 \), the conductance simply results as

\[
G_{\sigma}(T = 0) = \frac{e^2}{\hbar} \frac{\Gamma_{L,\sigma}(\mu) \Gamma_{R,\sigma}(\mu)}{\Gamma_{L,\sigma}(\mu) + \Gamma_{R,\sigma}(\mu)} \rho_{\sigma}(\mu). \tag{8}
\]

While the calculation of the non-interacting Green functions entering \( G_{\sigma}(\omega) \) is a very simple task (see appendix), the calculation of the Green function \( G_{\alpha,\sigma}^{R}(\omega) \) implies using a suitable many-body method to treat the Coulomb interaction. In the next subsection, we present a procedure based on a quantum Monte Carlo approach.

2.3. The Monte Carlo method

In the study of local properties of the Anderson impurity model, like the Green function \( G_{\alpha,\sigma}^{R}(\omega) \), the quantum Monte Carlo (QMC) approach is one of the most efficient methods. It allows for the numerical evaluation of the Matsubara Green function \( G_{\alpha,\sigma}^{R}(i\omega_n) \), where \( \omega_n = (2n + 1)\pi/\beta \), with high precision, since it is free from the so-called ‘sign’ problem which usually plagues the QMC method in other fermionic models. In fact, this method has been widely used in dynamical mean field theory calculations [28], which involve a self-consistent evaluation of the Green function of the Anderson impurity model. In the context of quantum transport, this method has also been extended to treat vibrational degrees of freedom [24].

From the first implementation of QMC due to Hirsch and Fye [29], a new class of impurity solvers has been developed, the so-called continuous time quantum Monte Carlo (CTQMC) [26, 30] ones, which set the basis for band structure calculations by using for example LDA + DMFT [31]. In a previous work [8], we have used the CTQMC method based on hybridization expansion [32] to investigate the conductance of a quantum dot in the Kondo regime connected to disordered wires and we follow a similar strategy here.

The non-interacting Green functions \( g_{\alpha}(i\omega_n) \) enter as input for the CTQMC method for obtaining \( G_{\alpha,\sigma}^{R}(i\omega_n) \). In order to evaluate \( G_{\sigma} \) from (6) we also need the corresponding retarded Green functions \( g_{\alpha}^{R}(\omega) \). The functions \( g_{\alpha}(i\omega_n) \) and \( g_{\alpha}^{R}(\omega) \) are evaluated by using the recursive procedure of

the appendix. The delicate step consists in the analytical continuation of the output functions of the Monte Carlo method, \( G_{d,\sigma}(i\omega_n) \), to the real frequency axis, in order to obtain \( G_{d,\sigma}^{R}(\omega) \). A possible route to follow is a maximum entropy method [33]. These methods proved to be efficient for capturing relevant features of the spectral functions, like the number of peaks, as well as their positions, irrespectively of the magnitude of the energies involved. The shortcoming is that they fail to accurately predict the corresponding width of the peaks. In the present work, we follow an alternative route, which consists in a low frequency polynomial fit of the Green function on the imaginary frequency axis. This procedure is motivated in the fact that in equation (6) the derivative of the Fermi function defines for low temperatures a very narrow integration window of energies around the mean chemical potential of the wires \( \mu \). For practical purposes, we set \( \mu = 0 \) and change the voltage gate \( V_g \). Thus, in order to compute the conductance, all we need is a reliable analytical continuation within this narrow interval of energies centered at \( \omega = 0 \). From the CTQMC simulation we can obtain not only the Green function but also the self-energy on the Matsubara axis \( \Sigma(i\omega_n) = G_{d,\sigma}^{R}(i\omega_n) - G_{d,\sigma}^{R}(i\omega_n)^{-1} \).

In the metallic regime, within which the Kondo effect arises, this function is known to behave as a Fermi liquid, for which \( \text{Im}[\Sigma(\omega)] \propto \omega \) close \( \omega = 0 \). It is, then, natural to expect a behavior for \( \Sigma(i\omega_n) \) that can be fitted by a quadratic function of \( \omega_n \). In our case, we keep the discrete CTQMC values for \( \Sigma(i\omega_n) \) and fit the functions \( \Sigma^{(\nu)}(i\omega_n) \) with functions of the form \( \gamma(\omega) = a^{(\nu)} + b^{(\nu)}(\omega) + c^{(\nu)}(\omega^2) \). This procedure can be carried out with a very high accuracy and the analytical continuation can be easily done by simply substituting \( i\omega \to \omega + i\eta \), with \( \eta = 0^+ \) in the function \( \gamma(\omega) \).

3. Results

3.1. Without a magnetic field

A relevant test for the numerical procedure based on the quantum Monte Carlo approach is the analysis of the conductance \( G \) of the interacting quantum dot as a function of the gate voltage \( V_g \). The typical behavior for this quantity has been evaluated using several methods and published in several works [6, 7, 18, 20]. We have checked, as a benchmark, that our method is able to recover the main features discussed in the literature. Here and in the following sections we set the half-bandwidth of the reservoir \( W = 2\pi \) as the unit of energy. Typical plots are shown in figure 2 for three different values of the Coulomb interaction \( U \), \( T = 0.01 \) and perfect matching, i.e. \( l_c = \tau \). The values of the Coulomb repulsion shown in the figure each define different values of \( J \) which, in turn, imply different Kondo temperatures \( T_K \). The superscript on \( T_K^{(i)} \) indicates the Kondo temperature for ballistic wires. These can be evaluated from (5) and the results are shown in the inset of figure 2. Below the Kondo temperature the Kondo resonance is developed, opening a conduction channel. The conductance, thus, equals the conductance quantum \( G_0 \) for each spin degree of freedom, a situation which is referred to
as the ‘unitary limit’. Typically this arises for an interval of gate potentials $-U \leq V_g - U/2 \leq U$.

For $U = 2$, we show that the conductance corresponds to a situation between the so called mixed-valence regime and the Kondo regime. In this case, the unitary limit $G = G_0$ is achieved only for $V_g = U/2$, displaying a peak with a width $\sim U$. For $U = 4$, the temperature is $T = T_K^0$ for that value of $U$ and we clearly see that $G$ displays a behavior very close to that of the unitary limit, with a plateau with $G = G_0$ within $-U \leq V_g - U/2 \leq U$. Finally, we show the case of $U = 6$, which corresponds to a Kondo temperature $T_K^0 = 7 \times 10^{-4}$, which is lower than the temperature considered in our simulations. Comparison to the behavior of $G$ in the other cases reveals that a valley appears within the range of $V_g$ corresponding to the plateau in the Kondo regime, while two Coulomb blockade peaks centered at $V_g - U/2 = \pm U/2$ develop.

In order to verify that the quantum Monte Carlo method is also able to capture more subtle mesoscopic effects of the Kondo cloud, we analyze the behavior of the conductance in a wire with a mismatching in the contact modeled by $t_C < t$ and finite local energy within the finite-length wire $\epsilon_W \neq 0$. In this case, the wires and the dot define a finite-size box. The spectral density of the wires corresponds to a sequence of peaks at the energy of the internal levels.

We consider left and right chains of equal length, $N_L = N_R = N$. The corresponding mean level spacing of the $N$-site chain in this case is $\Delta \sim 4t/N$. In figure 3 we consider $\Delta = 0.4$, which satisfies $\Delta \gtrsim T_K^0$. We also consider $t_C = 0.8$, although the particular value of this parameter is not relevant provided that it satisfies $t_C < t$. In this regime, where the typical level spacing of the wires is larger than the Kondo temperature for ballistic wires, mesoscopic effects are expected to be relevant [6, 7]. An equivalent picture for this scenario corresponds to a Kondo cloud of length $\xi_K > L$, with $L = Na$, where $a$ is the lattice constant of the wire. In figure 3 we show the low temperature behavior of the conductance as a function of the gate voltage for three different values of $\epsilon_W$. Varying this parameter allows us to tune at the Fermi energy different regions of the energy spectrum of the wires. In one of the plots of the figure, we show the behavior of $G$ corresponding to parameters for which the Fermi energy of the infinite leads coincides with the energy of a resonant level of the finite wire (circles). In this case, the behavior of the conductance as a function of $V_g$ is basically the same as that of figure 2. A similar behavior is observed when the Fermi energy lies perfectly off-resonance (squares). This case corresponds to a valley between two consecutive peaks in the spectral function of the wire. The density of states at the Fermi energy $\rho_0(\mu)$ is, thus, much smaller than in the resonant case and equation (5) leads to a Kondo temperature $T_K^* < T_K^0$. The data set of figure 3 corresponds to a temperature $T < T_K^*$. For this reason, the unitary limit of the conductance $G \sim G_0^*/2$ can be observed for $V_g \sim 0$, although the conductance plateau as a function of $V_g$ is in this case narrower than for the resonant one. We also illustrate an intermediate situation (pentagons), which does not exactly correspond to resonance or to off-resonance, where anomalous asymmetric features are observed. For sufficiently large chains with $\Delta \ll T_K^0$, all these finite-size effects become irrelevant and the behavior of figure 2 is recovered. The picture for this case is a Kondo cloud fitting within the length of the wires.

In figure 4 we observe another interesting effect of the finite-size wires, namely the oscillations of the conductance as a function of the energy $\epsilon_W$ of the wires and the change of the ensuing periodicity as the temperature grows above the Kondo temperature $T_K^*$. For these plots we keep the gate voltage of the dot fixed at $V_g = 0$. The low temperature regime can be understood from the behavior discussed in figure 3, where we found maxima of the conductance $G = G_0^*$ for both resonant and off-resonant configurations. Recalling that $\Delta$ is the mean level spacing of the wires, this corresponds to a periodicity equal to $\Delta/2$, consistent with a maximum of the conductance when the value of $\epsilon_W$ tunes a resonant as well as an off-resonant energy of the wire at the Fermi energy.

In the other limit with $T \gg T_K^*$, where the Kondo resonance is completely melted and the dot enters the Coulomb blockade regime, the conductance displays maxima only when $\epsilon_W$ tunes
Kondo resonance. As the magnetic field increases, the Kondo formation of the spin singlet, which is the origin of the spin of the electrons at the dot. This conspires against the field. The effect of the magnetic field is to polarize the conductance is a decreasing function of the magnetic

Figure 4. Conductance as a function of the energy shift $\varepsilon_W$ of the wires for finite wires with mismatching, $t_C = 0.8t$. Different plots correspond to different temperatures: $T < T_K^0$ (circles), $T \sim T_K^0$ and $T > T_K^0$ (pentagons and stars). Other parameters are $U = 4$, $V_g = U/2$, $N_L = N_R = 7$.

In figure 5 we illustrate the effect of the magnetic field on the conductance as a function of gate voltage for finite ballistic wires (without mismatching, $t_C = 1$ and two different values of the magnetic field. Other parameters are $U = 4$, $\varepsilon_W = 0$ and $T = 0.01$.

Figure 5. Conductance as a function of gate voltage for finite ballistic wires (without mismatching, $t_C = 1$) and two different sets of data in the figure are fitted by the same quadratic with coefficients $c_T$ and $c_B$ as in (9). Other parameters are $U = 3$, $V_g = U/2$.

we define the critical magnetic field $B_c$ as that for which $G(T, B_c) = G(T, 0)/2$.

The above figures show that, without finite-size wires, the magnetic field plays a role in the behavior of the conductance which is similar to that played by the temperature. In fact, in [12] a calculation based on a slave-boson treatment of the Kondo impurity and Fermi liquid theory shows the following universal law for the behavior of the ‘zero-bias’ conductance as a function of temperature and magnetic field:

$$G(T, \Delta_z)/G^0 = \left[1 - c_T \left(\frac{T}{T_K^0}\right)^2 - c_B \left(\frac{\Delta_z}{T_K^0}\right)^2\right].$$

where $[B]/T_K^0 < 1$, (9)

In figure 6 we show that the QMC results are consistent with the scaling behavior predicted by (9). In fact, we see that for low temperature and low magnetic field the QMC data can be fitted with a quadratic function of $T$ and $B$. Even at finite $U$, with a rescaling of the temperature by a factor close to $\pi$, both sets of data in the figure are fitted by the same quadratic function, which is consistent with the ratio (10).

Equation (9) and the behavior of figure 5 suggest that the magnetic field plays a similar role to the temperature as regards the behavior of the conductance. If we keep the temperature low, $T \ll T_K^0$, we can infer that the ratio $\Delta/B_c$ sets the scale for the mesoscopic effects playing a role in the behavior of the conductance. In particular, we can expect for $\Delta \ll B_c$, in analogy to the behavior for $B = 0$ and finite temperature analyzed in section 3.1, the conductance to be insensitive to the details of the wires, like the value of the local

Figure 6. Conductance as a function of the temperature squared $T^2$ and as a function of the Zeeman splitting squared $\Delta_z^2$ for a dot connected to ballistic wires. The symbols correspond to quantum Monte Carlo data, while the line corresponds to a quadratic fit with coefficients $c_T$ and $c_B$ as in (9). Other parameters are $U = 3$, $V_g = U/2$.

In figure 6 we illustrate the effect of the magnetic field on the behavior of the conductance as a function of the gate voltage at low temperature, $T \simeq T_K^0$. We see that the magnetic field tends to suppress the Kondo effect. In particular, the plateau $G \sim G^0$ is suppressed, giving rise to a valley between the two Coulomb blockade features.

For $V_g = 0$, and assuming a fixed temperature $T \ll T_K^0$, the conductance is a decreasing function of the magnetic field. The effect of the magnetic field is to polarize the spin of the electrons at the dot. This conspires against the formation of the spin singlet, which is the origin of the Kondo resonance. As the magnetic field increases, the Kondo resonance melts down. In order to quantify this mechanism,
energy of the wires \( \varepsilon_W \) relative to \( V_0 \). However, for \( \Delta_C \gtrsim B_c \), we expect to observe mesoscopic features.

We set again left and right chains of equal length and mismatching in the contact. In section 3.1, we choose \( N = 7 \) and \( l_C = 0.8 \). As mentioned before, any other value of \( l_C < 1 \) would yield the same features, while changing \( N \) implies a change in the mean level spacing \( \Delta = 4t/N \) of the wires. The particular \( N \) that we consider corresponds to \( \Delta \sim 0.57 \), which satisfies \( \Delta > B_c \), corresponding to the regime where the mesoscopic features are expected to be significant.

In figure 7 we show the conductance as a function of the local energy of the wires \( \varepsilon_W \) for different values of \( B \). For \( B \ll B_c \) we observe a periodicity of \( \Delta/2 \). This is consistent with perfect transmission when a level of the wires is resonant as well as when the Fermi energy is exactly off-resonant. This is precisely the behavior that we discussed for \( \varepsilon_W \), as well as when the Fermi energy is exactly off-resonant.

In figure 8 we show the conductance for each spin species corresponding to the same parameters as in figure 7. These plots satisfy \( G_{\uparrow} + G_{\downarrow} = G \), where \( G \) is the one shown in figure 7. For \( B = 0 \), the Kondo effect is robust and the two contributions to the conductance show the same pattern of oscillations with a period \( \Delta/2 \), and maxima when the levels of the wires are resonant or off-resonant with the Kondo peak. As \( B \) is switched on, a splitting of the Kondo peak is expected to take place. These changes reflect the changes in the behavior of the local density of states of the up and down spins at the dot close to the Fermi energy, \( \rho_{\uparrow,\downarrow}(\omega) \sim \mu \). For the largest value of \( B \) shown in figure 8 (see plots in stars), the Kondo effect is expected to be broken and the emerging picture for the behavior of the conductance is the following. For \( \varepsilon_W = 0 \), the density of states of the wires has a resonant peak at the Fermi energy, while the resonant peaks of \( \rho_{\uparrow,\downarrow}(\omega) \) are shifted to \( \omega = \pm \Delta_C \), respectively. The effect of increasing \( \varepsilon_W \) is a downwards rigid shift of the density of states of the wires. For \( \varepsilon_W \) increasing from zero, there is a resonant peak of the density of states of the wires, which was originally centered at \( \omega = 0 \) and tends to get aligned with the resonant peak of \( \rho_{\uparrow}(\omega) \) at the same time as it further departs from that of \( \rho_{\downarrow}(\omega) \). The result is an increasing (decreasing) conductance for the (up) spins. For even larger \( \varepsilon_W \sim \Delta/2 \) the density of states of the wires has a valley at the Fermi energy and, as the Kondo effect is not active, the conductance decreases for both spins. For larger values of \( \varepsilon_W \) within the range \( \Delta/2 < \varepsilon_W < \Delta \), a resonant peak of the wires which was originally at \( \omega = \Delta \) is shifted downwards and tends to be aligned with the peak at \( \omega = \Delta_C \) of \( \rho_{\uparrow}(\omega) \). Hence, the conductance of \( \uparrow \) spins increases, while that for \( \downarrow \) spins continues the decreasing behavior.

4. Summary and conclusions

We have studied the mesoscopic features in the conductance of a quantum dot with a magnetic field, modeled as a Hubbard–Anderson impurity connected to finite one-dimensional wires. We have focused on the linear regime in the bias voltage and evaluated the conductance by recourse to
a quantum Monte Carlo method. In order to benchmark this numerical technique, we first focused on a vanishing magnetic field and identified the change in the period of the mesoscopic conductance oscillations as the temperature grows above the Kondo temperature.

We then studied the effect of the magnetic field. In the absence of mismatching between the wires and the reservoirs, we have recovered universal relations in the behavior of the conductance as a function of the temperature and the magnetic field. These relations indicate that the magnetic field and the temperature play similar roles in the behavior of the conductance. In particular, both quantities have the effect of destroying the Kondo singlet, thus tending to decrease the conductance. Furthermore, the departure from the perfect conductance quantum is quadratic both in the temperature and in the magnetic field. It is, thus, not surprising that the magnetic field may introduce similar effects in the behavior of the mesoscopic features in the presence of finite-size wires. In fact, the conductance displays oscillations as a function of a gate voltage applied at the wires. As the magnetic field overcomes a critical value, for which the Kondo effect is sufficiently weak, the period of such oscillations changes to twice its original value, as happens at zero magnetic field, and the temperature increases over the Kondo temperature.

The changes in the behavior of these mesoscopic oscillations of the conductance as a function of the temperature have been proposed as providing a way to detect the Kondo screening length, relative to the length of the mesoscopic features in the presence of finite-size wires. It is, however, much easier to control changes of the temperature within such a small scale. It is, however, much easier to control changes of the temperature within such a small scale. It is, however, much easier to control changes of the temperature within such a small scale. It is, however, much easier to control changes of the temperature within such a small scale.

Acknowledgments

This work was supported by CONICET, MINCyT and UBACYT, Argentina.

Appendix. Calculation of the non-interacting retarded Green functions

An exact procedure is to numerically evaluate the Matsubara and retarded Green functions $g_n(i\omega_n)$ and $g_R^n(i\omega)$ is the recursive solution of the Dyson equation $g_R^n(i\omega) = g_{1,n}(i\omega_n)$, with

$$g_{l,a}(i\omega_n)^{-1} = i\omega_n - i\omega_n - t^2 g_{l+1,a}(i\omega_n),$$

(A.1)

for $1 \leq l < N_{\alpha}$, and

$$g_{N_{\alpha},a}(i\omega_n)^{-1} = i\omega_n - i\omega_n - t^2 g_{N_{\alpha},a}(i\omega_n),$$

(A.2)

where $g_{N_{\alpha},a}(i\omega_n)$ is the Matsubara Green function corresponding to a semi-infinite tight-binding chain with hopping $t$. The corresponding expressions for the retarded functions can be easily obtained by making the substitution $i\omega_n \rightarrow \omega + i\eta$. 

References

[1] Goldhaber-Gordon D, Shtrikman H, Mahalu D, Abusch-Magder D, Meirav D and Kastner M A 1998 Nature 391 156
[2] Cronenwett S M, Oosterkamp T H and Kouwenhoven L P 1998 Science 281 540
[3] Nygard J, Cobden D H and Lindelof P E 2000 Nature 408 342
[4] Delatre T et al 2009 Nature Phys. 5 208
[5] Zarchin O, Zaffalon M, Heiblum M, Mahalu D and Umansky V 2008 Phys. Rev. B 77 241303
[6] For a review see Glazman L I and Pustilnik M 2005 Nanophysics: Coherence and Transport ed H Bouchiat et al (Amsterdam Elsevier) pp 427–78 (and references therein)
[7] Hewson A C 1993 The Kondo Problem to Heavy Fermions (Cambridge: Cambridge University Press)
[8] Affleck I and Simón P 2001 Phys. Rev. Lett. 86 2854
[9] Simon P and Affleck I 2002 Phys. Rev. Lett. 89 206602
[10] Cornaglia P S and Balseiro C A 2003 Phys. Rev. Lett. 90 216801
[11] Camjayi A and Arrachea L 2012 Phys. Rev. B 86 235143
[12] Cornaglia P S and Balseiro C A 2003 Phys. Rev. Lett. 90 216801
[13] Quay C H L, Cumings J, Gamble S J, de Picciotto R, Katamura H and Goldhaber-Gordon D 2007 Phys. Rev. B 76 245311
[14] Costi T A 2000 Phys. Rev. Lett. 85 1504
[15] Pustilnik M, Avishai Y and Kikoin K 2000 Phys. Rev. Lett. 84 1756
[16] Andrei N 1980 Phys. Rev. Lett. 45 379
[17] Andrei N, Furuya K and Lowenstein J H 1983 Rev. Mod. Phys. 55 331
[18] Vigman P B 1980 JETP Lett. 31 364
[19] Tselick M A and Wiegmann P B 1983 Adv. Phys. 32 453
[20] Moore J and Wen X 2000 Phys. Rev. Lett. 85 1722
[21] Coleman P 1984 Phys. Rev. B 30 3035
[22] Bulka R and Lipiński S 2003 Phys. Rev. B 67 024404
[23] López R and Sánchez D 2003 Phys. Rev. Lett. 90 116602
[24] Choi M, Sánchez D and López R 2004 Phys. Rev. Lett. 92 056601
[25] Costi T A 2000 Phys. Rev. Lett. 85 1504
[26] Basset J, Kasumov A Yu, Moca C P, Zaránd G, Simon P, Bouchiat H and Deckel R 2012 Phys. Rev. Lett. 108 046802
[27] Hewson A C 1993 The Kondo Problem to Heavy Fermions (Cambridge: Cambridge University Press)
[28] Affleck I and Simon P 2001 Phys. Rev. Lett. 86 2854
[29] Simon P and Affleck I 2002 Phys. Rev. Lett. 89 206602
[30] Cornaglia P S and Balseiro C A 2003 Phys. Rev. Lett. 90 216801
[31] Quay C H L, Cumings J, Gamble S J, de Picciotto R, Katamura H and Goldhaber-Gordon D 2007 Phys. Rev. B 76 245311
[32] Kretinin A V, Shirikian H, Goldhaber-Gordon D, Hanl M, Weichselbaum A, von Delft J, Costi T and Mahalu D 2001 Phys. Rev. B 84 245316
[33] Smirnov S and Grifoni M 2013 New J. Phys. 15 073047
[34] López R and Sánchez D 2003 Phys. Rev. Lett. 90 116602
[35] Choi M, Sánchez D and López R 2004 Phys. Rev. Lett. 92 056601
[36] Kretinin A V, Shirikian H, Goldhaber-Gordon D, Hanl M, Weichselbaum A, von Delft J, Costi T and Mahalu D 2001 Phys. Rev. B 84 245316
[37] Smirnov S and Grifoni M 2013 New J. Phys. 15 073047
[38] López R and Sánchez D 2003 Phys. Rev. Lett. 90 116602
[39] Choi M, Sánchez D and López R 2004 Phys. Rev. Lett. 92 056601
[40] Costi T A 2000 Phys. Rev. Lett. 85 1504
[41] Pustilnik M, Avishai Y and Kikoin K 2000 Phys. Rev. Lett. 84 1756
[42] Eto M and Nazarov Y V 2000 Phys. Rev. Lett. 85 1306
[43] Andrei N 1980 Phys. Rev. Lett. 45 379
[44] Andrei N, Furuya K and Lowenstein J H 1983 Rev. Mod. Phys. 55 331
[45] Vigman P B 1980 JETP Lett. 31 364
[46] Tselick M A and Wiegmann P B 1983 Adv. Phys. 32 453
[47] Moore J and Wen X 2000 Phys. Rev. Lett. 85 1722
[48] Coleman P 1984 Phys. Rev. B 29 3035
[49] Bulka R and Lipiński S 2003 Phys. Rev. B 67 024404
[50] López R and Sánchez D 2003 Phys. Rev. Lett. 90 116602
[51] Dong B and Lei X L 2001 Phys. Rev. B 63 235306
[52] Levy Yeyati A, Martín-Rodero A and Flores F 1993 Phys. Rev. Lett. 71 2991
[53] Oguri A 2001 Phys. Rev. B 63 115305
[54] Aligia A 2006 Phys. Rev. B 74 155125
[55] Zarea M, Ulloa S E and Sandler N 2012 Phys. Rev. Lett. 108 046601
[56] Wilson K G 1975 Rev. Mod. Phys. 47 773
[57] Krishna-murthy H R, Wilkins J W and Wilson K G 1980 Phys. Rev. B 21 1003
[58] Hofstetter W and Schoeller H 2001 Phys. Rev. Lett. 88 016803
[59] Izaume W, Sakai O and Tarucha S 2001 Phys. Rev. Lett. 87 216803
[60] Büsser C A, Anda E V, Lima A L, Davidovich M A and Chieppa G 2000 Phys. Rev. B 62 9907
Anda E V, Chiappe G, Büsser C A, Davidovich M A, Martins G B, Heidrich-Meisner F and Dagotto E 2008
Phys. Rev. B 78 085308
[21] Oguri A, Ishii H and Saso T 1995 Phys. Rev. B 51 4715
Oguri A 1997 Phys. Rev. B 56 13422
[22] Sakai O, Suzuki S, Izumida W and Oguri A 1999 J. Phys. Soc. Japan 68 1640
[23] Hirsch J E and Fye R M 1986 Phys. Rev. Lett. 56 2521
Hützen R, Weiss S, Thorwart M and Egger R 2012 Phys. Rev. B 85 121408
[24] Arrachea L and Rozenberg M J 2005 Phys. Rev. B 72 041301
[25] Hützen R, Weiss S, Thorwart M and Egger R 2012 Phys. Rev. B 85 121408
[26] Gull E, Millis A J, Lichtenstein A I, Rubtsov A N, Troyer M and Werner P 2011 Rev. Mod. Phys. 83 349
[27] Meir Y and Wingreen N S 1992 Phys. Rev. Lett. 68 2512
[28] Georges A, Kotliar G, Krauth W and Rozenberg M J 1996 Rev. Mod. Phys. 68 13
[29] Hirsch J E and Fye R M 1986 Phys. Rev. Lett. 56 2521
[30] Werner P, Comanac A, de’ Medici L, Troyer M and Millis A J 2006 Phys. Rev. Lett. 97 076405
[31] Kotliar G, Savrasov S Y, Haule K, Oudovenko V S, Parcollet O and Marianetti C A 2006 Rev. Mod. Phys. 78 865
[32] Haule K 2007 Phys. Rev. B 75 155113
[33] Jarrell M and Gubernatis J E 1996 Phys. Rep. 269 133