Physica E 72 (2015) 140–148

Contents lists available at ScienceDirect

Physica E

journal homepage: www.elsevier.com/locate/physe

A new possible transition from two- to one-channel Kondo physics in mesoscopic transport

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HIGHLIGHTS

- Both DOS and conductance of the QD in the Kondo regime.
- DOS and differential conductance exhibit different behavior.
- Quantum phase transition in the tunneling characteristics from 2CK to 1CK physics.

ARTICLE INFO

Article history:
Received 16 December 2014
Received in revised form
18 April 2015
Accepted 21 April 2015
Available online 21 April 2015

Keywords:
Kondo effect
Quantum dots
Luttinger liquid
Joint effects

ABSTRACT

We study the joint effects of the Coulombic dot–lead interaction and intralead electron interaction on nonequilibrium transport through a quantum dot coupled to Luttinger liquid leads in the Kondo regime. By applying the nonperturbative canonical transformation technique and extended equation of motion method of nonequilibrium Green function approach, we find two types of zero-bias anomalies with an enhanced conductance for the renormalized exponent $Y < 1$ and a suppressed conductance for $Y > 1$. The transition between the two behaviors is found at $Y \approx 1$. For $Y > 1$ and $Y < 1$, tunneling between the dot and Luttinger liquid leads is suppressed and enhanced, respectively, and they reflect two different types of exciton-assist tunneling. These physics phenomena provide opportunities to easily study the transition of two-channel Kondo physics in the future experiments.

1. Introduction

As techniques in manufacturing devices of nanometer scale develop, research on the properties of strongly correlated systems has become a very active field in condensed matter physics. One of the most striking examples exhibiting many-body correlations in mesoscopic physics is the Kondo effect in a quantum dot (QD) coupled to leads. The Kondo effect was first studied in dilute magnetic alloys [1] and then was observed in QD systems [2,3], which lead to extensive theoretical and experimental investigations. Much attention has been paid to the interplay of the dot–lead tunneling and on-dot Coulomb repulsion interactions. However, the effect of a Coulombic interaction between the charges on the dot and leads has been studied little in the Kondo regime, even though it should be clearly present from the aspect of physical grounds.

Another important type of strongly correlated systems is one-dimensional electronic systems whose low energy dynamics is governed by the Luttinger liquid (LL) theory instead of Fermi liquid (FL) theory. When the leads connecting to the QD were taken to be FL, the interactions in the leads could be ignored. This has no problem if the leads are two or three dimension, where interactions affect the low energy properties only perturbatively. However, in a one-dimensional electron system, arbitrary weak interactions completely modify the ground state, and the low energy excitations are described by the LL theory which predicted a series of unique characters, such as the absence of Landau quasiparticles, spin-charge separation, suppression of the electron tunneling density of states (DOS), anomalous power laws for transport coefficients with interaction-dependent exponents, and so on [4,5].

The physics of the Coulomb dot–lead coupling in LL has received attention. Several theoretical techniques have been developed and applied to address the role of a Coulombic interaction between the charges on the dot and LL leads on transport through the system, such as the scattering Bethe Ansatz [6], perturbative and numerical renormalization group (NRG) methods [7], the Hershfield Y-operator [8], the time-dependent density matrix renormalization group method [9] sophisticated field theory

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http://dx.doi.org/10.1016/j.physe.2015.04.018
1386-9477/

approaches [9,10], master equation approach [11,12], density matrix renormalization group, and classical Monte Carlo simulations [13–15], and equation of motion (EOM) of nonequilibrium Green function [16]. For example, with the help of the time-dependent density matrix renormalization group method [9] as well as sophisticated field theory approaches [9,10], Boulat and Saleur studied an impurity coupled to one-dimensional leads in terms of an anisotropic interacting resonant-level model and found that a strong impurity-wire interaction could enhance transport in the low-temperature regime. By use of the density matrix renormalization group and classical Monte Carlo simulations, Goldstein et al. [13–15] investigated the effect of a Coulombic dot–lead interaction on the dynamics of the population of a QD. Elste et al. studied the effect of the dot–lead coupling on the relaxational dynamics and transport of a QD coupled to LL leads [11,12] with the master equation approach. They found that this Coulomb dot–lead coupling enhanced the electron tunneling in the regime of weak tunneling. Yang et al. studied the joint effects of the Coulomb dot–lead interaction and intralead interaction on transport through a QD with the EOM method of the nonequilibrium Green functions [16] which allowed an expansion to all orders in the dot–lead hybridization. In spite of the theoretical researches mentioned above, the effect of the Coulombic dot–lead interaction on physics properties in the Kondo regime has been paid little attention so far.

Owing to the lack of satisfactory treatment of the Kondo problem, little effort has been made to investigate how the Kondo physics and the Coulomb dot–lead interaction developed and manifested themselves in the system of the strongly correlated QD coupled to the LL leads. Since the QD system has some tunable parameters in an experiment, various aspects of the Kondo effect can be controlled. Here we intend to apply an easier and straightforward approach, which will combine the improved nonperturbative canonical transformation treatment of the Coulomb dot–lead interaction with the extended EOM method [17–20] of the nonequilibrium Green functions [21]. This approach, although not rigorous in quantitative description, has the advantage of intuitiveness and can provide a semi-quantitative understanding of the Coulomb dot–lead interaction-assisted Kondo effect.

Using this approach, in the present work, we shall mainly focus on the joint effects of Coulomb dot–lead interaction and intralead interaction both on DOS and conductance of the QD in the Kondo regime, which can be probed by tunneling spectroscopies. DOS may reflect the details of the tunneling process [22,23].

We find that the DOS and differential conductance exhibit different behavior when the Coulomb dot–lead interaction and intralead interaction vary. The Kondo-satellite peaks or dips can appear at both sides of the main Kondo peak depending on the interaction strengths. These results clearly show the single-channel Kondo (1CK) effect and two-channel Kondo (2CK) effect, respectively, and the quantum phase transition in the tunneling characteristics from the 2CK physics to the 1CK physics can occur.

It was ever stressed that transition of the phonon-assisted satellite dips changing into the peaks could occur due to the electron–phonon interaction, but the quantum phase transition of the Kondo dip to the peak at the zero bias voltage has not been found [24].

The paper is organized as follows. In Section 2 we introduce the model and the methods. The canonical transformation is introduced to eliminate the Coulomb dot–lead interaction at the expense of adding operator content to the tunneling Hamiltonian. Section 3 presents the numerical calculations and analysis. Finally, we summarize our findings in Section 4.

2. Model Hamiltonian and method

The system studied is usually regarded as the Anderson model coupled to 1D interacting leads. The Hamiltonian reads

$$H = H_{leads} + H_{dot} + H_T + H_C. \tag{1}$$

The first-term $H_{leads} = H_L + H_R$ describes the left (L) and right (R) LL leads. Their bosonized Hamiltonian can be written as follows [25,26]:

$$H_\alpha = \hbar \int_0^\infty k(\upsilon_{c\alpha} a_{k\alpha}^\dagger a_{k\alpha} + \upsilon_{s\alpha} c_{k\alpha}^\dagger c_{k\alpha})\, dk, \quad \alpha = L, R, \tag{2}$$

where $a_{k\alpha}^\dagger (c_{k\alpha}^\dagger)$ and $a_{k\alpha} (c_{k\alpha})$ are the creation and annihilation bosonic operators describing charge (spin) density fluctuations with the renormalized velocity $\upsilon_{c\alpha}$ ($\upsilon_{s\alpha}$), and satisfy the canonical bosonic commutation relations. The detailed derivation of Eq. (2) can be found in Ref. [26].

The second-term H_{dot} in Eq. (1) describes the QD and has the form

$$H_{dot} = \sum_\sigma \varepsilon_d d_\sigma^\dagger d_\sigma + U d_\uparrow^\dagger d_\uparrow d_\downarrow^\dagger d_\downarrow, \tag{3}$$

where $d_\sigma^\dagger (d_\sigma)$ represents electron creation (annihilates) operators with spin-σ, ε_d denotes the discrete energy-levels of the QD which can be tuned by a gate voltage, and U is the on-site Coulomb repulsion, which is high enough to forbid double occupancy ($U \to \infty$).

The third-term in Eq. (1),

$$H_T = \sum_{\alpha\sigma} (t_\alpha d_\sigma^\dagger \psi_{\alpha\sigma} + \text{h. c. }), \tag{4}$$

describes the tunneling processes between the QD and electrodes, where t_α is the tunneling coupling between the dot and α lead. $\psi_{\alpha\sigma}^\dagger (\psi_{\alpha\sigma})$ creates (annihilates) an electron at the end point of the α lead. The electron field operator $\psi_{\alpha\sigma}$ at the boundary could be written in a "bosonized" form [25]

$$\psi_{\alpha\sigma} = \sqrt{\frac{2}{\pi\alpha'}} \exp\left[\int_0^\infty dk\, e^{-\alpha' k/2} \left(\frac{a_{k\alpha} - a_{k\alpha}^\dagger}{\sqrt{2kK_{c\alpha}}} + \sigma \frac{c_{k\alpha} - c_{k\alpha}^\dagger}{\sqrt{2kK_{s\alpha}}} \right) \right], \tag{5}$$

where α' is a short-distance cutoff being of the order of the reciprocal of the Fermi wave number k_F, and $K_{c\alpha}$ and $K_{s\alpha}$ are the interaction parameters for the charge and spin sectors respectively. Because the electron–electron interaction is repulsive, $K_{c\alpha}$ and $K_{s\alpha}$ are less than 1.

The last term of the Hamiltonian H_C describes the electrostatic Coulomb interaction between the dot and leads and is given by

$$H_c = \sum_{\alpha\sigma} \lambda_\alpha d_\sigma^\dagger d_\sigma \psi_{\alpha\sigma}^\dagger(0) \psi_{\alpha\sigma}(0), \tag{6}$$

which is also called excitonic coupling term. To our knowledge, as long as the coupling strength λ_α is finite, the Hamiltonian equation (1) cannot be solved exactly even in the case of $K_{c\alpha} = 1$ and $K_{s\alpha} = 1$. We here make use of the phase fields $\varphi_{\alpha\sigma}(x)$ which describe the slow varying spatial component of the electron density (plasmons) relevant to $\partial_x \varphi_{\alpha\sigma}/2\pi$. Then, in terms of $\varphi_{\alpha\sigma}(x)$, the electron annihilation field operator can be expressed as

$$\psi_{\alpha\sigma}(x) = \sqrt{\frac{2}{\pi\alpha'}} e^{-i\varphi_{\alpha\sigma}(x)}, \tag{7}$$

where $\varphi_{\alpha\sigma}(x) = i[\int_0^\infty dk(\exp(-\alpha' k/2)/\sqrt{2k})\{ch\beta_c (a_{k\alpha} e^{-ikx} - a_{k\alpha}^\dagger e^{ikx})$ $-sh\beta_c (a_{k\alpha} e^{ikx} - a_{k\alpha}^\dagger e^{-ikx}) + \sigma[ch\beta_s (c_{k\alpha} e^{-ikx} - c_{k\alpha}^\dagger e^{ikx})]$ $-[sh\beta_s (c_{k\alpha} e^{ikx} - c_{k\alpha}^\dagger e^{-ikx})]\}]$ with $ch\beta_{c/s} = (\sqrt{K_{c/s}} + 1/\sqrt{K_{c/s}})/2$ and $sh\beta_{c/s} = (\sqrt{K_{c/s}} - 1/\sqrt{K_{c/s}})/2$. The electron density is given by