

## Effects of Fermi Liquid Interactions on the Shot Noise of an $SU(N)$ Kondo Quantum Dot

P. Vitushinsky,<sup>1</sup> A. A. Clerk,<sup>1</sup> and K. Le Hur<sup>2</sup>

<sup>1</sup>*Department of Physics, McGill University, Montréal, Québec, Canada H3A 2T8*

<sup>2</sup>*Department of Physics, Yale University, New Haven, Connecticut 06520, USA*

(Received 8 August 2007; published 24 January 2008)

We study shot noise in the current of quantum dots whose low-energy behavior corresponds to an  $SU(N)$  Kondo model, focusing on the case  $N = 4$  relevant to carbon nanotube dots. For general  $N$ , two-particle Fermi-liquid interactions have two distinct effects: they can enhance the noise via backscattering processes with an  $N$ -dependent effective charge, and can also modify the coherent partition noise already present without interactions. For  $N = 4$ , in contrast with the  $SU(2)$  case, interactions enhance shot noise solely through an enhancement of partition noise. This leads to a nontrivial prediction for experiment.

DOI: [10.1103/PhysRevLett.100.036603](https://doi.org/10.1103/PhysRevLett.100.036603)

PACS numbers: 72.70.+m, 71.10.Ay, 72.15.Qm, 73.63.Fg

Despite being over 40 years old, the Kondo effect remains among the most studied effects in condensed matter physics. It is a canonical example of how strong interparticle interactions can lead to nontrivial behavior; it is also attractive as it may be controllably studied in quantum dots [1]. In spite of ultimately being the result of strong interactions, much of the low-temperature phenomenology of the standard  $SU(2)$  Kondo effect in quantum dots relies on an effective noninteracting model, having a resonant level (the Kondo resonance) sitting at the Fermi energy. This, however, is not the full story: low-energy properties are properly described by a Fermi-liquid theory having weak interparticle interactions [2,3]. While the effects of these interactions in quantum dot systems can be subtle, it was shown in Ref. [4] that they play a large role in determining the size of current fluctuations. Fermi-liquid interactions lead to two-particle scattering and a resulting universal effective charge  $e^*$  (defined via the magnitude of the backscattering current noise) that is larger than 1:  $e^* = (5/3)e$ . This result is a clear sign of the interacting nature of the low-temperature Kondo state in a quantum dot.

In this Letter, we now turn to more exotic quantum dots that realize the higher symmetry  $SU(N)$  Kondo model ( $N \geq 2$ ), and which are described by an  $SU(N)$ -symmetric Fermi-liquid fixed point. Such Kondo effects involve the screening of an effective pseudospin that incorporates both orbital and spin degrees of freedom; various realizations of the  $N = 4$  case have been proposed [5–7] and even measured in carbon nanotube [8] and vertical [9] dots. Here, we will examine how Fermi-liquid interactions in these systems affect the low-voltage current noise  $S_I$ . Noise in carbon nanotubes is already under experimental study [10]; also, the first observations of noise through a  $SU(2)$  Kondo quantum dot have recently been made [11]. The picture we obtain will have relevance to the general question of how two-particle interactions affect coherent shot noise in a mesoscopic conductor. Formulating the Fermi-liquid theory in terms of scattering states and using the Keldysh technique, we find that the situation is markedly different from the  $SU(2)$  case. For general  $N$ , interactions

have *two* distinct effects on shot noise. The first (the only effect for  $N = 2$ ) is that Fermi-liquid interactions lead to new two-particle scattering processes; these *enhance* the shot noise in a Poissonian fashion (i.e., via uncorrelated tunneling events). Such two-particle scattering can result in the backscattering of either one or two particles, with respective rates  $\Gamma_1$  and  $\Gamma_2$ . We find that the effective charges associated with these processes are functions of  $N$ :

$$e_1^*/e = 2T_0 - 1 = 2\sin^2(\pi/N) - 1, \quad (1a)$$

$$e_2^* = 2e_1^*. \quad (1b)$$

Here,  $T_0$  is the transmission coefficient through the dot at zero temperature and voltage. Equations (1) reduce to the results of Ref. [4] in the  $SU(2)$  case, while in the  $SU(4)$  case, both these effective charges vanish.

In addition to the scattering-induced enhancement of  $S_I$ , a second effect of Fermi-liquid interactions is to modify the partition noise already present without Fermi-liquid corrections. Recall that in a phase-coherent conductor with transmission coefficient  $T_0$ , partitioning of electrons into transmitted and reflected streams results in  $S_I \propto T_0(1 - T_0)$ . We find that two-particle Fermi-liquid interactions enhance this contribution, in keeping with more phenomenological studies of shot noise [12]. This effect is absent in the ideal  $SU(2)$  case, as  $T_0 = 1$  [13]. In the especially relevant case  $N = 4$ , this partition noise modification is the *only* effect of interactions; this leads to a parameter-free prediction for the interaction enhancement of noise [cf. Eq. (17)]. Note that Fermi-liquid corrections affect only the nonlinear in voltage parts of the current and current noise; we thus will focus on these quantities. There will also be a linear-in-voltage contribution to the noise whose value follows immediately from the value of the Fermi-energy phase shift and the asymmetry of coupling to the leads.

*Model.*—For concreteness we consider the case of  $SU(4)$  Kondo physics in a carbon nanotube dot [6]. A small nanotube island (i.e., the dot) is attached via tunnel junctions to two nanotube leads. In addition to spin, electrons

here have an orbital degree of freedom  $m$  that can take one of two values ( $m = +, -$ );  $m$  labels two degenerate bands [6]. We take  $m$  to be conserved in lead-dot tunneling as is consistent with experiment [8]. Because of strong on-site Coulomb interactions, the dot can be tuned to a regime where it is singly occupied. It then acts as an effective 4-component pseudospin, characterized by both its physical spin and its orbital index  $m$ . Further, dot-lead tunneling becomes an effective exchange-like interaction between the dot and conduction electron pseudospin. The resulting model is the SU(4) Kondo model. The ground state corresponds to a singlet formed by the impurity pseudospin and a screening cloud of electrons in the leads. This model could be easily generalized to an SU( $N$ ) ( $N > 4$ ) Kondo model by allowing the orbital index  $m$  to take on  $N/2$  different values. In what follows, we will consider the general SU( $N$ ) case, paying special attention to  $N = 4$ .

*Fermi-liquid theory using scattering states.*—We will be interested in the behavior of our system at energies much smaller than  $T_K$ , the Kondo temperature. In this regime, the SU( $N$ ) Kondo effect is described by a local Fermi-liquid theory [3,15], which describes both elastic scattering off the Kondo singlet, as well as interparticle interactions induced by virtual excitations of the singlet. Both effects are described by the phase shift  $\delta_{\text{tot},\sigma} = \delta_0 + \delta_\sigma(\varepsilon)$  for  $s$ -wave scattering of an electron with pseudospin  $\sigma$  off the impurity.  $\delta_0$  is the phase shift at the Fermi energy; its value is determined by the Friedel sum rule to be  $\delta_0 = \pi/N$ .  $\delta_\sigma(\varepsilon)$  accounts both for the energy dependence of elastic scattering and for inelastic processes:

$$\delta_\sigma(\varepsilon) = \frac{\alpha\varepsilon}{k_B T_K} + \alpha' \left( \frac{\varepsilon}{k_B T_K} \right)^2 - \frac{\beta \sum_{\sigma' \neq \sigma} n_{\sigma'}}{\nu k_B T_K}. \quad (2)$$

Here,  $\nu$  is the density of states per electron flavor and the energy  $\varepsilon$  is measured from the Fermi level. Note that we retain the  $\varepsilon^2$  dependence of the elastic phase shift; while it plays no role in the SU(2) case, it will play a critical role for  $N = 4$ . Universality at the SU( $N$ ) Kondo fixed point ensures that the ratios between  $\alpha$ ,  $\alpha'$ , and  $\beta$  are fixed. The pinning of the energy of the Kondo resonance relative to the Fermi energy leads to the relation  $\alpha = (N - 1)\beta$  [15]. For  $N = 2$ , the fixed point has particle-hole symmetry, implying  $\alpha' = 0$ . For  $N > 2$ , the Kondo resonance is above the Fermi energy; consequently, the fixed point is not particle-hole symmetric. For  $N = 4$ , one finds  $\alpha' = \alpha^2$ . This follows from equating the elastic parts of the phase shift  $\delta_{\sigma,\text{tot}}$  with  $\arctan(\Gamma/(2(\varepsilon_d - \varepsilon)))$ , the phase shift of a resonant level (i.e., the Kondo resonance) having width  $\Gamma$  and energy  $\varepsilon_d$  [16]. For  $N = 4$ , expanding to order  $\varepsilon$  yields  $\varepsilon_d = \Gamma/2$  and  $\Gamma = T_K/\alpha$ ; expanding to  $\varepsilon^2$  yields  $\alpha^2 = \alpha'$ . These properties of the  $N = 4$  Kondo resonance are consistent with recent calculations [6].

Applying the Fermi-liquid picture to our quantum dot, we may replace the dot plus two lead system by a single one-dimensional channel described by a field operator  $\hat{\psi}_\sigma(x)$ : for  $x < 0$ , this describes the left lead, while for  $x >$

0, it describes the right lead. The scattering (both elastic and inelastic) due to the dot plus screening cloud may then be modeled as occurring locally at  $x = 0$ . It is convenient to work in a basis of scattering states that correspond to the Fermi-energy phase shift  $\delta_0$ . This can be converted to a transmission probability  $T_0$  via  $T_0 = \sin^2 \delta_0$ . We thus introduce scattering states  $\phi_L(x; k\sigma)$  and  $\phi_R(x; k\sigma)$ , which correspond (respectively) to waves incident on the dot from the left or from the right lead. For simplicity, we focus on the case where both leads are equally coupled to the dot; deviations will be considered in [14]. Inversion symmetry then implies that for  $x < 0$ :

$$\phi_L(x; k\sigma) = e^{i(k_F+k)x} - \sqrt{1-T_0} e^{-i(k_F+k)x}, \quad (3)$$

$$\phi_R(x; k\sigma) = -i\sqrt{T_0} e^{-i(k_F+k)x}. \quad (4)$$

Letting  $\hat{c}_{L,k\sigma}$ ,  $\hat{c}_{R,k\sigma}$  denote the annihilation operators for the scattering states, the field operator may be written as  $\hat{\psi}_\sigma(x) \propto \sum_k [\phi_L(x; k\sigma) \hat{c}_{L,k\sigma} + \phi_R(x; k\sigma) \hat{c}_{R,k\sigma}]$ . The current operator for  $x < 0$  then takes the form:

$$\hat{I}(x) = \hat{I}_D(x) + \hat{I}_{OD}(x), \quad (5)$$

with ( $R_0 \equiv 1 - T_0$ )

$$\hat{I}_D(x) = \frac{e}{\hbar\nu} \sum_{k,k';\sigma} [\hat{c}_{R,k'\sigma}^\dagger \hat{c}_{R,k\sigma} [-T_0 e^{-i(k-k')x}] + \hat{c}_{L,k'\sigma}^\dagger \hat{c}_{L,k\sigma} [e^{i(k-k')x} - R_0 e^{-i(k-k')x}]], \quad (6a)$$

$$\hat{I}_{OD}(x) = \frac{e}{\hbar\nu} \sum_{k,k';\sigma} [\hat{c}_{L,k'\sigma}^\dagger \hat{c}_{R,k\sigma} [-i\sqrt{T_0 R_0} e^{-i(k-k')x}] + \text{H.c.}]. \quad (6b)$$

Without Fermi-liquid corrections, the average current would be determined by  $\hat{I}_D$ , while the zero-temperature current noise would be determined by  $\hat{I}_{OD}$  [12].

Similar to the  $N = 2$  case [1,17], the effects of  $\delta_\sigma(\varepsilon)$  may be incorporated via the Hamiltonian:

$$H \equiv \sum_{k,\sigma,j=L,R} \xi_k \hat{c}_{j,k\sigma}^\dagger \hat{c}_{j,k\sigma} + H_\alpha + H_{\alpha'} + H_\beta, \quad (7a)$$

$$H_\alpha = -\frac{\alpha}{2\pi\nu T_K} \sum_{k,k',\sigma} (\xi_k + \xi_{k'}) \hat{a}_{k\sigma}^\dagger \hat{a}_{k'\sigma}, \quad (7b)$$

$$H_{\alpha'} = -\frac{\alpha'}{2\pi\nu (T_K)^2} \sum_{k,k',\sigma} (\xi_k + \xi_{k'})^2 \hat{a}_{k\sigma}^\dagger \hat{a}_{k'\sigma}, \quad (7c)$$

$$H_\beta = \frac{\beta}{2\pi\nu^2 T_K} \sum_{k',q',k,q} \sum_{\sigma\sigma'} \hat{a}_{k\sigma}^\dagger \hat{a}_{k'\sigma} \hat{a}_{q'\sigma'}^\dagger \hat{a}_{q\sigma'}. \quad (7d)$$

Here,  $\xi_k = \hbar v_F k$  and  $\hat{a}_{k\sigma} = (\hat{c}_{L,k\sigma} + \hat{c}_{R,k\sigma})/\sqrt{2}$  (only the symmetric combination enters due to inversion symmetry).  $H_\alpha + H_{\alpha'}$  and  $H_\beta$  represent, respectively, the elastic scattering from the impurity [to order  $(\varepsilon/T_K)^2$ ] and the presence of two-particle scattering. This Hamiltonian explicitly reproduces the phase shift  $\delta_\sigma$  [Eq. (2)]; its form follows from the usual relation between scattering and potential matrices (see, e.g., [18]). The description

presented here is consistent with the results obtained in Ref. [7] for  $N = 4$  using conformal field theory. As we show, using Eqs. (7) and the current operator in Eq. (5) allows one to both recover results for the current obtained from  $\delta_{\text{tot},\sigma}$  and calculate current fluctuations.

*Effect of elastic scattering.*—We focus on the regime of zero temperature, and small voltage  $\mu_L - \mu_R = eV \ll T_K$ . As we treat the case of equal couplings to the left and right leads, we use a symmetric voltage bias:  $\mu_L = -\mu_R = eV/2$  [i.e., the Kondo resonance position is fixed relative to  $(\mu_L + \mu_R)/2$ ]. Consider first the effects of  $H_\alpha + H_{\alpha'}$ , which describe purely elastic scattering. Keeping only these terms, our system is equivalent to a noninteracting system having an energy-dependent transmission coefficient  $T(\varepsilon)$  given by

$$T(\varepsilon) = \sin^2 \left[ \delta_0 + \alpha \frac{\varepsilon}{k_B T_K} + \alpha' \left( \frac{\varepsilon}{k_B T_K} \right)^2 \right]. \quad (8)$$

We have verified through an explicit perturbative calculation in  $H_\alpha + H_{\alpha'}$  that to order  $(eV/T_K)^3$ , the elastic scattering contributions to the current and noise do indeed correspond to  $T(\varepsilon)$ . In particular, the elastic contributions to the average current are given by [19]

$$\begin{aligned} \langle I \rangle|_{\beta=0} &\equiv \frac{Ne^2}{h} \int_{\mu_R}^{\mu_L} d\varepsilon T(\varepsilon) \\ &= \frac{Ne^2 V}{h} \left[ T_0 + \frac{\alpha^2 \cos 2\delta_0 + \alpha' \sin 2\delta_0}{12} \left( \frac{eV}{k_B T_K} \right)^2 \right. \\ &\quad \left. + O((eV/k_B T_K)^4) \right]. \end{aligned} \quad (9)$$

$$\begin{aligned} &= \frac{Ne^2 V}{h} \left[ T_0 + \frac{\alpha^2 \cos 2\delta_0 + \alpha' \sin 2\delta_0}{12} \left( \frac{eV}{k_B T_K} \right)^2 \right. \\ &\quad \left. + O((eV/k_B T_K)^4) \right]. \end{aligned} \quad (10)$$

Note for  $N = 4$ , there is no  $\alpha^2$  contribution  $\langle I \rangle$ .

Similarly, the elastic scattering contributions to the zero frequency current noise  $S_I \equiv 2 \int dt \langle \delta I(t) \delta I(0) \rangle$  are given by the noninteracting formula:

$$\begin{aligned} S_I|_{\beta=0} &\equiv \frac{2Ne^3}{h} \int_{\mu_R}^{\mu_L} d\varepsilon T(\varepsilon) [1 - T(\varepsilon)] \\ &= \frac{2Ne^3 V}{h} \left[ T_0(1 - T_0) + \frac{\alpha^2 \cos(4\delta_0) + \alpha' \sin(4\delta_0)/2}{12} \right. \\ &\quad \left. \times \left( \frac{eV}{k_B T_K} \right)^2 + O((eV/k_B T_K)^4) \right]. \end{aligned} \quad (11)$$

Thus, the elastic part of  $\delta_\sigma(\varepsilon)$  yields  $V^3$  terms in both  $\langle I \rangle$  and  $S_I$ . In the case  $N = 4$ , their ratio is simply

$$\frac{1}{2e} \frac{d^3 S_I / dV^3}{d^3 \langle I \rangle / dV^3} \Big|_{\beta=0, V=0} = -\frac{\alpha^2}{\alpha'} = -1. \quad (13)$$

This is just the result for a noninteracting resonant level. As we will see, two-particle interactions modify both the sign

and the magnitude of this ratio; this relative enhancement of noise over a noninteracting model is our key prediction for experiment.

*Effect of two-particle scattering.*—We have calculated the effect of  $H_\beta$  [Eq. (7d)] on the average current and current noise perturbatively using the Keldysh approach. Heuristically,  $H_\beta$  induces two-particle scattering between the scattering states  $\phi_L, \phi_R$ . For  $\mu_L = \mu_R + eV$  ( $eV > 0$ ) and zero temperature, there is a range of energies  $\varepsilon$  where the  $\phi_L(\varepsilon)$  states are occupied, while the  $\phi_R(\varepsilon)$  are unoccupied. As a result, three types of inelastic processes are possible: (a)  $\phi_L \phi_L \rightarrow \phi_R \phi_L$ , (b)  $\phi_L \phi_R \rightarrow \phi_R \phi_R$ , (c)  $\phi_L \phi_L \rightarrow \phi_R \phi_R$ . Using Fermi's golden rule in  $H_\beta$ , one finds that for a fixed initial and final spin configuration, processes (a) and (b) occur at a rate  $\Gamma_1$ , while process (c) occurs at a rate  $\Gamma_2$ . These rates have no explicit- $N$  dependence, and are given by  $\Gamma_1 = \Gamma_2/8 = \beta^2 \frac{1}{24} \frac{eV}{h} \left( \frac{eV}{k_B T_K} \right)^2$ , as in Ref. [4].

The current associated with these scattering processes does not follow immediately from the rates, as unlike the SU(2) case,  $T_0 \neq 1$  and hence the scattering states  $\phi_L, \phi_R$  are not eigenstates of current. A long but straightforward calculation using the Keldysh technique yields the following contribution to the current to order  $\beta^2$ :

$$\delta \langle I \rangle|_\beta = -(N_a + N_b)e_1^* \Gamma_1 - N_c e_2^* \Gamma_2. \quad (14)$$

Here, the ‘‘effective charges’’  $e_1^*, e_2^*$  are given in Eq. (1). The factors  $N_a, N_b, N_c$  count the multiplicity of each of the three processes (i.e., the number of distinct initial and final states once pseudospin is included):  $N_a = N_b = 2N_c = N(N-1)$ . The contribution to the current given in Eq. (14) can be combined with that in Eq. (10) to yield the total low-voltage current to order  $(eV/T_K)^3$ .

There is a simple heuristic way to understand the effective charges appearing in Eq. (14): they are proportional to the total change in the scattered-wave amplitude associated with a particular interaction-induced scattering process. An interaction-induced scattering event leads to a sudden change in the amplitude of the scattered waves associated with a particular pair of scattering states. For example, consider the  $\Gamma_2$  process, which takes a pair of  $\phi_L$  scattering states to a pair of  $\phi_R$  scattering states (i.e.,  $\phi_L \phi_L \rightarrow \phi_R \phi_R$ ); without loss of generality, consider the current in the left lead. Initially, we have two left scattering states; the associated scattered-wave amplitude [cf. Eq. (3)] is thus  $2(1 - T_0)$ . In the final state, we have instead two right scattering states; the associated scattered-wave amplitude [cf. Eq. (4)] is thus  $2T_0$ . The  $\Gamma_2$  process thus suddenly changes the scattered-wave amplitude by an amount  $2(2T_0 - 1) = e_2^*/e$ . A sudden change in scattered-wave amplitude implies a delta-function blip in the current of magnitude  $e_2^*$ :  $\delta I(t) \simeq e_2^* \delta(t)$ . Such ‘‘blips’’ occur at a rate  $\Gamma_2$ , hence the expression in Eq. (14). A similar analysis may be given of the  $\Gamma_1$  processes and the charge  $e_1^*$ . We stress that Eq. (14) follows from a rigorous perturbative calculation of the current.

The utility of the interpretation of Eq. (14) in terms of effective charges becomes apparent when we now turn to the interaction contribution to  $S_I$ . Based on this interpretation, we would expect a Poissonian contribution to  $S_I$ :

$$\delta S_I|_{\beta,D} = 2(N_a + N_b)(e_1^*)^2\Gamma_1 + 2N_c(e_2^*)^2\Gamma_2. \quad (15)$$

This is precisely what we find in our full calculation when we sum all contributions independent of  $\hat{I}_{OD}$  [cf. Eq. (6b)]; in the case where  $T_0 \rightarrow 1$  or  $T_0 \rightarrow 0$ , these are the only contributions, as  $\hat{I}_{OD} = 0$ . Thus, interaction-induced scattering events enhance the shot noise in a simple way, namely, via an uncorrelated Poisson process. For  $N = 2$ , these contributions reduce to what was found in Ref. [4]. For  $N = 4$ ,  $e_1^*$  and  $e_2^*$  are zero, and these scattering processes make no contribution.

We now turn to the second interaction-induced contribution to the noise, namely, the contribution from diagrams involving  $\hat{I}_{OD}$  [cf. Eq. (6b)]. Recall that  $\hat{I}_{OD}$  is responsible for the partition shot noise already present without Fermi-liquid corrections [i.e., first term in Eq. (12)]. We find that two-particle interactions *enhance* this noise; summing all diagrams involving  $\hat{I}_{OD}$  to order  $\beta^2$ :

$$\begin{aligned} \delta S_I|_{\beta,OD} &= \left(2(N-1) - \frac{8}{3}\right)N(N-1) \\ &\times \frac{e^3V}{h}T_0R_0\beta^2\left(\frac{eV}{k_B T_K}\right)^2. \end{aligned} \quad (16)$$

As will be discussed in Ref. [14], the dominant effect here is an enhancement of noise due to an interaction-induced enhancement of the particle-hole density of states. Note that in the  $N = 2$  case this contribution vanishes, as there is no partition noise in this case. Also note that it is impossible to express this contribution purely in terms of the scattering rates  $\Gamma_1$  and  $\Gamma_2$  introduced above.

Focusing now on the experimentally relevant case  $N = 4$ , we may again compare the nonlinear-in- $V$  dependencies of the current noise and current. Combining all contributions [Eqs. (10), (12), and (14)–(16)], we find

$$\frac{1}{2e} \frac{d^3 S_I / dV^3}{d^3 \langle I \rangle / dV^3} \Big|_{V=0} = \frac{-\alpha^2 + 15\beta^2}{\alpha'} = \frac{2}{3}. \quad (17)$$

This ratio is modified by a factor of  $-2/3$  compared to the noninteracting expectation [cf. Eq. (13)]. This universal interaction-induced noise enhancement is a central result of this Letter; it is a stringent test of the theory involving three Fermi-liquid parameters. We derived Eq. (17) for a symmetric dot-lead coupling; for a small coupling asymmetry  $x$  and small residual potential scattering (with phase shift  $\delta_r \ll 1$ ), the ratio in Eq. (17) becomes (to leading order)  $2/3 + c_1\delta_r + c_2x^2$ , where  $c_1, c_2$  are constants of order unity [14]. We stress again that in the SU(4) case, Fermi-liquid interactions enhance shot noise through a modification of the coherent partition noise. This is in

contrast to the SU(2) case, where the modification is due to interaction-induced scattering events.

In conclusion, we have shown how Fermi-liquid interactions affect current noise through a SU( $N$ ) Kondo quantum dot, identifying two distinct physical mechanisms. In the experimentally relevant  $N = 4$  case, we show that two-particle interactions modify the shot noise in a nontrivial and universal manner; this serves as a robust prediction for experiment. Our approach in this Letter is also well suited to study the effect of interactions on current noise in a variety of mesoscopic conductors [20].

We thank T. Kontos for useful discussions and C. Mora for a very constructive conversation. A. C. and K. L. H. acknowledge support from NSERC and CIFAR.

- 
- [1] L.I. Glazman and M. Pustilnik, arXiv:cond-mat/0501007v2.
  - [2] P. Nozières, *J. Low Temp. Phys.* **17**, 31 (1974).
  - [3] P. Nozières, *J. Phys. (Paris)* **39**, 1117 (1978).
  - [4] E. Sela *et al.*, *Phys. Rev. Lett.* **97**, 086601 (2006).
  - [5] L. Borda *et al.*, *Phys. Rev. Lett.* **90**, 026602 (2003); G. Zaránd, A. Brataas, and D. Goldhaber-Gordon, *Solid State Commun.* **126**, 463 (2003); K. Le Hur and P. Simon, *Phys. Rev. B* **67**, 201308 (2003); K. Le Hur, P. Simon, and L. Borda, *Phys. Rev. B* **69**, 045326 (2004); R. Lopez *et al.*, *Phys. Rev. B* **71**, 115312 (2005).
  - [6] M.-S. Choi, R. Lopez, and R. Aguado, *Phys. Rev. Lett.* **95**, 067204 (2005); J. S. Lim *et al.*, *Phys. Rev. B* **74**, 205119 (2006).
  - [7] K. Le Hur, P. Simon, and D. Loss, *Phys. Rev. B* **75**, 035332 (2007).
  - [8] P. Jarillo-Herrero *et al.*, *Nature (London)* **434**, 484 (2005); A. Makarovski *et al.*, *Phys. Rev. B* **75**, 241407 (2007).
  - [9] S. Sasaki *et al.*, *Phys. Rev. Lett.* **93**, 017205 (2004).
  - [10] N. Y. Kim *et al.*, *Phys. Rev. Lett.* **99**, 036802 (2007).
  - [11] M. Heiblum (to be published).
  - [12] Y. M. Blanter and M. Büttiker, *Phys. Rep.* **336**, 1 (2000).
  - [13] The modification of partition noise due to interactions does play a role for an SU(2) dot that is asymmetrically coupled to the leads, as will be discussed in [14].
  - [14] P. Vitushinsky, A. A. Clerk, and K. Le Hur (to be published).
  - [15] P. Nozières and A. Blandin, *J. Phys. (Paris)* **41**, 193 (1980).
  - [16] This correspondence to a Lorentzian resonance is consistent with exact Fermi-liquid relations on the local impurity electron self-energy, as will be discussed in Ref. [14].
  - [17] I. Affleck, *Nucl. Phys.* **B336**, 517 (1990).
  - [18] A. A. Clerk, S. M. Girvin, and A. D. Stone, *Phys. Rev. B* **67**, 165324 (2003).
  - [19] For an asymmetric bias, we find for  $N = 4$  a  $V^2$  term in the current, exactly as found in Ref. [7]. In this case,  $\alpha'$  plays no role to leading order.
  - [20] M. C. Goorden and M. Büttiker, *Phys. Rev. Lett.* **99**, 146801 (2007).