Universal zero-bias conductance through a quantum wire side-coupled to a quantum dot

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A numerical renormalization-group study of the conductance through a quantum wire side-coupled to a quantum dot is reported. The temperature and the dot-energy dependence of the conductance are examined in the light of a recently derived linear mapping between the Kondo-regime temperature-dependent conductance and the universal function describing the conductance for the symmetric Anderson model of a quantum wire with an embedded quantum dot. Two conduction paths, one traversing the wire, the other a bypass through the quantum dot, are identified. A gate potential applied to the quantum wire is shown to control the flow through the bypass. When the potential favors transport through the wire, the conductance in the Kondo regime rises from nearly zero at low temperatures to nearly ballistic at high temperatures. When it favors the dot, the pattern is reversed: the conductance decays from nearly ballistic to nearly zero. When the fluxes through the two paths are comparable, the conductance is nearly temperature-independent in the Kondo regime, and a Fano antiresonance in the fixed-temperature plot of the conductance as a function of the dot energy signals interference. Throughout the Kondo regime and, at low temperatures, even in the mixed-valence regime, the numerical data are in excellent agreement with the universal mapping.

I. INTRODUCTION

The transport properties of nanostructured devices constitute a complex subject. Amongst its numerous facets, interference is one of the shiniest, because it challenges classical intuition, adds colorful features to experimental plots, and can be controlled. Such characteristics have motivated numerous studies, in either the Aharonov-Bohm or the T-shaped interferometric arrangement.[1,2,3,4,5,6,7,8,9,10,11,12,13,14,15,22] The minimal number of components in and the rich variety of effects stemming from the latter setup make it particularly interesting.

The T-shaped device comprises a quantum-dot side-coupled to a quantum wire, through which current flows in response to a small bias. Its dynamics parallels that of the single-electron transistor, the alternative arrangement that embeds the quantum dot in the conduction path. The same Hamiltonian models the two devices, the thermodynamical properties are hence identical. Nonetheless, the conductances in the two geometries are remarkably distinct.

In either geometry, if a gate potential $V_g$ applied to the quantum dot tunes its occupation $n_d$ to an odd integer, the resultant magnetic moment interacts antiferromagnetically with the wire electrons and induces a Kondo cloud. Thin at first, the cloud gradually shrouts the quantum dot as the temperature is reduced. Below the Kondo temperature $T_K$, the cloud and dot moment lock into a singlet. In energy space, the Kondo resonance in the low-energy spectrum of the model Hamiltonian identifies the consolidation of the Kondo singlet.[2,3]

The conductances of the single-electron transistor and of the T-shaped devices respond differently to the Kondo resonance. In the embedded arrangement, the resultant strong coupling between the conduction and the dot electrons breaks through the Coulomb blockade to allow conduction. This remarkable low-temperature property of the single-electron transistor found precise mathematical expression in the early numerical renormalization-group (NRG) study by Costi, Hewson, and Zlatic,[17] who described the temperature-dependent conductance for the spin-degenerate Anderson model[18] and showed that the conductance for the particle-hole symmetric model is a monotonically decreasing universal function $G^p(T/T_K)$ of the temperature scaled by the Kondo temperature.

In the side-coupled geometry, instead of enhancing the conductance, the Kondo droplet obstructs transport through the wire. While the thin high-temperature cloud offers little resistance to electronic flow, low temperatures raise a Kondo blockade in the wire sector closest to the dot. In experimental arrangements allowing conduction only through the wire, the conductance through the T-shaped device therefore rises with temperature. For the symmetric Anderson Hamiltonian, in particular, the conductance $G$ is complementary to the universal function $G^p$: $G(T/T_K) = G_2 - G^p(T/T_K)$, where $G_2 \equiv 2e^2/h$ is the quantum conductance.[15][19] To emphasize the significance of this theoretical finding, we define the complementary symmetric conductance $G^s_2(T/T_K) = G_2 - G^s(T_K)$.

More elaborate T-shaped setups accommodate an alternative conduction path through the quantum dot, a bypass around the blocked wire segment.[21] Considered on its own, the bypass is equivalent to a single-electron transistor: while a Coulomb blockade bars electronic transfer through the dot at high temperatures, the Kondo resonance allows conduction at low $T$.

Most frequently, the two conduction paths coexist. Three conduction patterns then stem from the combination of the currents through them.[22][23][24] If the experimental conditions favor the bypass (the path traversing the wire), the conductance will decay (rise) with temperature, following the pattern set by the universal (the complementary-conductance) function $G^s(T/T_K) [G^p(T/T_K)]$.

The third pattern arises when the amplitudes for electronic transfer along the two paths are comparable. The conductance is then a relatively flat function of temperature, but the low-temperature conductance profile (i. e., $G$ as a function of the
recently, the present authors showed that the temperature-device. Important results have been derived. T

ing is linear, its coe

\[ G \approx \frac{1}{N} \sum_k a_k, \] (1)

where \( N \) is the number of conduction states in the wire.

In standard notation, the Anderson Hamiltonian capturing the physics of this setup is

\[ H_A = \sum_k \epsilon_k a_k^\dagger a_k + W f_0^\dagger f_0 + V(f_0 c + V/H.c.) + H_d, \] (2)

where the states \( a_k \) form a structureless half-filled conduction band of halfwidth \( D \) and density of states \( \rho = N/2D \), the scattering potential \( W \) is controlled by the gate voltage \( V_w \), and the tunneling amplitude \( V \) couples the spin-degenerate dot level \( c_d \) to the \( f_0 \) orbital.

This paper closes a quartet dedicated to the conductance for the Anderson model of a wire coupled to a quantum dot. References [19] and [24] discussed the embedded geometry, the former paper having derived the mapping between the conductance and the universal function \( G^u(T/T_K) \), and the latter having presented NRG data to illustrate and to probe the accuracy of that mapping at the limits of the Kondo regime and beyond them. Reference [23] derived the mapping for the \( T \)-shaped device, and this report compares it with an NRG survey of the side-coupled geometry.

Our presentation is divided in 5 sections. Section [II] defines the model Hamiltonian and recalls basic concepts associated with it. The mapping is discussed in Section [III]. The comprehensive overview of the numerical results and comparison with the mapping to the universal function follow, in

Section [IV]. Finally, Section [V] summarizes our conclusions.

II. MODEL

The overview of the numerical results in Sec. [IV] will divide the parametrical space of the model in a number of regimes, easily identified by their nearly uniform conductances. Before examining the NRG results, it is therefore convenient to discuss the characteristic energies marking the boundaries of those regimes. This section defines the model Hamiltonian, relates the conductance to its eigenvalues and eigenvectors, and refers to the pioneer investigations that identified the characteristic scales in its spectrum.[23-27]

A. Hamiltonian

The quantum dot in Fig. 1 is side-coupled to a quantum wire. Current flows in response to a bias voltage applied to the wire. Gate potentials \( V_d \) and \( V_w \) control the dot occupation and the occupation of the Wannier orbital

![Diagram of a quantum dot side-coupled to a quantum wire](image-url)
It proves convenient to write the dot Hamiltonian $H_d$ in Eq. (2) in the form

$$H_d = (e_d + U/2) n_d - U/2 n_d \eta n_d,$$

(3)

which defines the antisymmetric component $e_d \equiv e_d + U/2$ of the dot energy, i.e., the component that changes sign under the particle-hole transformation $c_d \rightarrow -c_d$, $a_k \rightarrow a_k^\dagger$. The second term on the right-hand side of the Eq. (3), by contrast, remains invariant under the same transformation. For $W = e_d = 0$, only the second term survives, and the Anderson Hamiltonian reduces to its symmetric form.

To highlight interference, many a study of transport through a nanodevice has explicitly introduced two conduction paths in the model Hamiltonian. The ratio between the amplitudes for charge transfer along the two routes, i.e., the Fano factor $q$, is then a model parameter.

An alternative approach is possible. A sufficiently comprehensive model Hamiltonian may implicitly include two conduction routes and hence define an effective Fano factor $q_{\text{eff}}$. Fano’s pioneer analysis of self-ionization in the He atom provided the first example of this approach, which we adopt. A simple extension of Eq. (2) would provide another example. A momentum-dependent coupling $\sum_k (V d_k^\dagger c_d + H. c.)$ between the quantum wire and the quantum dot would introduce a conduction path alternative to transport straight through the quantum wire and define an effective Fano factor proportional to the Fermi-level value of the derivative $dV_k/dk$.

The momentum-independent coupling in Eq. (2) rules out this alternative route. Section [11] will show, however, that the wire gate potential $W$ opens a bypass through the quantum dot, so that an effective Fano factor $q_{\text{eff}}$ arises, which grows with $W$. As long as $W = 0$, Eq. (2) comprises two conduction paths, through which currents can flow and interfere.

To reproduce the model Hamiltonian in Ref. [23], a momentum-dependent coupling could have been substituted for the third term on the right-hand side of Eq. (2). Any of a variety of artificial couplings between the dot and the wire orbitals orthogonal to $f_0$ might equally well have been introduced, to generate another bypass and originate a different Fano parameter. In any case, the results would be equivalent to those in Sec. [11] because the universal mapping between the temperature-dependent conductance and the function $G^2(T/T_K)$ derived in Ref. [23] depends only on the ground-state phase shift $\delta$. Since $\delta$ decreases from $\pi/2$ to zero as the wire potential $W$ grows, our survey of the conductance for the structure in Fig. 1 samples well the dependence of the conductance on Fano factors. Alternative Hamiltonians would yield the same information.

**B. Conductance**

Reference [19] derived an expression for the temperature-dependent conductance $G(T)$ in the embedded configuration. The same analysis leads to an analogous equality for the conductance in the side-coupled configuration:

$$G(T) = G^2 \frac{\beta \pi T}{Z} \sum_{m>n} |\langle m | f_0 | n \rangle|^2 e^\beta E_n + e^\beta E_m,$$

(4)

where $Z$ is the partition function, $|m\rangle$ ($|n\rangle$) is an eigenstate of $H_A$, with eigenvalue $E_m$ ($E_n$), and $\Gamma \equiv \pi \rho V^2$ is the dot-level width, due to its coupling to the quantum wire.

To describe the temperature and dot-energy dependence of the conductance, it is sufficient to consider non-negative wire gate potentials. Since the model Hamiltonian and the conductance remain invariant under the modified particle-hole transformation

$$a_k \rightarrow a_k^\dagger, \\ c_d \rightarrow -c_d, \\ e_{\alpha} \rightarrow -e_{\alpha},$$

and

$$W \rightarrow -W,$$

given two energies $\alpha < 0$ and $\beta$, the conductance for $W = \alpha$ and $e_{\alpha} = \beta$ is given by the right-hand side of Eq. (4) for $W = -\alpha > 0$ and $e_{\alpha} = -\beta$. The conductance for a negative wire gate potential can therefore be obtained from the results for the symmetric (positive) potential.

**C. Characteristic energy scales**

Figure 3 in Ref. [19] depicts the five characteristic energies in the spectrum of the Hamiltonian [2]. The first (second), due to the dot Hamiltonian $[2]$, is the excitation energy $\Delta_0 = -e_{\alpha}$ ($\Delta_2 = U + e_d$) necessary to remove an electron from (add an electron to) the singly-occupied level. The only scale in the conduction band Hamiltonian is trivial: the halfwidth $D$.

The coupling to the dot, on the other hand, makes two important contributions: the level width $\Gamma$, and the Kondo energy $k_B T_K$. The scattering potential $W$ reduces the former to

$$\Gamma_w = \frac{\pi \rho V^2}{1 + (\pi \rho W)^2}.$$

(6)

The Kondo regime is defined by the condition $\max(k_B T, \Gamma_w) \ll \varepsilon_f$, where the dominant characteristic energy is $\varepsilon_f = \min(D, \Delta_0^*, \Delta_2^*)$. Here, $n_d = 0$ and $n_d = 2$ dot configurations are then energetically inaccessible, and a dot magnetic moment arises. As the temperature is reduced below the band halfwidth $D$, the antiferromagnetic interaction between the conduction electrons and the dot electron progressively screens the moment. The energy of the resulting low-temperature singlet defines the fifth characteristic scale in our set: the Kondo energy $k_B T_K$.

A side effect of the coupling between the dot and the conduction electrons, one that emerges in the spectrum of particle-hole asymmetric Hamiltonians, is the renormalization $\Delta_0 \rightarrow \Delta_0^*$, or $\Delta_2 \rightarrow \Delta_2^*$, of the lowest dot-excitation energy [25-26]. Although of no other practical consequence, this energy shift is important because it displaces both the Kondo and the mixed-valence regimes [25].
The latter regime comprises the two regions in the parametric space of the model Hamiltonian that enclose the Kondo regime. One of them is defined by the inequality $\Gamma_w \gtrsim \Delta_2^*$, the other, by $\Gamma_w \gtrsim \Delta_0^*$. In the mixed-valence regime the dominant characteristic energy is $\varepsilon_{c_0}^{\text{mix}} = 1$, and at high temperatures, the dot moment is only partially formed. As the thermal energy is reduced past $\Gamma_w$, the dot level $c_0$ couples strongly to the surrounding conduction electrons, and all physical properties approach their low-temperature limit.

**III. UNIVERSAL MAPPING**

The derivation of the universal expression mapping the conductance in Eq. (4) to the universal function $G^S(T/T_K)$ computed by Costi, Hewson, and Zlatidi\(^{1,2}\) has been summarized elsewhere.\(^{23}\) A detailed presentation of the analogous derivation for the single-electron transistor being more-over available, only recapitulation of the central result is required:

\[
G\left(\frac{T}{T_K}\right) = \frac{G_2}{2} = \left(\frac{G^S}{T_K} - \frac{G_2}{2}\right) \cos(2\delta) \quad \text{Eq. (7)}
\]

Here, $\delta$ is the ground-state phase shift of the conduction electrons in the wire.

To make the physical content of Eq. (7) more visible, we rewrite it in the equivalent form

\[
G\left(\frac{T}{T_K}\right) = G^S\left(\frac{T}{T_K}\right) \cos^2 \delta + G_0^S\left(\frac{T}{T_K}\right) \sin^2 \delta \quad \text{Eq. (8)}
\]

which shows that the conductance is a linear combination of the universal function $G^S(T/T_K)$ with its complement, $G_0^S(T/T_K)$. While the universal, monotonically decreasing function $G^S(T/T_K)$ describes the conductance through a particle-hole symmetric quantum dot in the embedded configuration, $G_0^S(T/T_K)$ is the monotonically increasing function that describes the conductance through a particle-hole symmetric quantum dot in the side-coupled configuration. For particle-hole asymmetric T-shaped devices, the first term on the right-hand side of Eq. (8) describes the current through the bypass, while the second term accounts for the current traversing the wire. If the two conduction paths were independent, the resultant conductance would be the sum $G_{\text{b}}(T/T_K) + G_{\text{t}}(T/T_K) = G_2$; the coefficients $\cos^2 \delta$ and $\sin^2 \delta$ thus account for the interference between the two fluxes.

**Special limits**

Four limits of Eq. (8) merit special attention. For $\delta = \pi/2$, the mapping reduces to

\[
G\left(\frac{T}{T_K}\right) = G_2 - G^S\left(\frac{T}{T_K}\right) \equiv G_0^S\left(\frac{T}{T_K}\right) \quad (\delta = \pi/2) \quad \text{Eq. (9a)}
\]

The equality, which can be derived from a diagrammatic expansion for the symmetric Anderson Hamiltonian\(^{23,24}\) shows that, with $\delta = \pi/2$, the conductances for the side-coupled and the embedded geometries are complementary. This result is particularly important because, in the absence of a wire gate potential, the Friedel sum rule\(^{23}\) pins the ground-state phase shift to the vicinity of $\pi/2$. For $W = 0$, the conductance through the $T$-shaped device is therefore expected to decay from nearly ballistic to nearly zero as the temperature rises.

In the second special limit, $\delta = 0$, Eq. (7) yields

\[
G\left(\frac{T}{T_K}\right) = G_2 \cos^2 \delta \quad (\delta = 0) \quad \text{Eq. (9b)}
\]

In the absence of a wire gate potential, the ground-state phase shift is close to $\pi/2$; a strong wire potential $W$ must be applied to make $\delta = 0$. The potential affects the conductance in two different ways: it induces a charge that blocks transport through the central region of the wire; at the same time, it opens a bypass through the quantum dot. At high temperatures, the Coulomb blockade against conduction through the dot makes the bypass ineffective, so that the conductance of the device vanishes. Upon cooling, however, the Kondo effect raises the blockade and the conductance rises, lifted by the very same mechanism that allows conduction through a single-electron transistor.\(^{23,24}\) In brief, Eq. (9b) tells us that a strong wire potential diverts the current to the bypass and makes the $T$-shaped device emulate a single-electron transistor.

The third simple limit of Eq. (7) is $G(T = 0)$. At low temperatures, the universal curve $G^S$ approaches $G_2$. From Eq. (7), it follows that

\[
G(T) = G_2 \cos^2 \delta \quad (T \ll T_K) \quad \text{Eq. (9c)}
\]

a result that can be derived from an extension of Langreth’s argument\(^{23}\) relating the low-energy dot-level spectral density to the ground-state phase shift.

Finally, at high temperatures, the universal function $G^S$ vanishes, and Eq. (7) yields

\[
G(T) = G_2 \sin^2 \delta \quad (T \gg T_K) \quad \text{Eq. (9d)}
\]

Analogous to Eq. (9c), this result also stems from Langreth’s expression. In the temperature range $k_B T_K \ll k_B T \ll \varepsilon_c^k$, within the Kondo regime, the Anderson Hamiltonian is in the local-moment regime, equivalent to a phase shifted conduction band weakly interacting with the dot moment.\(^{23}\) The interaction neglected, the Hamiltonian reduces to a free conduction band, to which Langreth’s reasoning applies.\(^{23}\) The phase shift $\delta_{LM}$ of the electron wires can be obtained from the Friedel sum rule, which associates a phase-shift difference $\pi/2$ with the Kondo cloud. We hence have that $\delta = \delta_{LM} + \pi/2$, and so $\delta_{LM} = \delta - \pi/2$. Substitution of $\delta_{LM}$ for $\delta$ in Eq. (9d) leads to Eq. (9c).

**IV. NUMERICAL RESULTS**

We are now ready to discuss the NRG results and compare them with Eq. (7). A previous report detailed the numerical procedure yielding the ground-state phase shift $\delta$, and
the temperature-dependent conductance $G(T)$ and tabulated the parameters controlling the accuracy of the computation. Here, to follow the structure of that paper, we will discuss $\delta$ before examining $G(T)$.

A. Phase shifts

Figure 2 shows the ground-state phase shift, extracted from the low-energy spectrum of the Hamiltonian $H_A$, as a function of the dot-energy $\varepsilon_d$, in the interval $|\varepsilon_d + U/2| < 3.5D$, i.e., for $|\varepsilon_d| < 3.5D$. It follows from the Friedel sum rule that the ordinate $2\delta/\pi$ is the extra charge $n_w$ piled up at the wire. That accumulation is the sum of two charges: the charge induced by the wire potential, the absolute value of which grow with $W$, and the Kondo screening charge, which is equal to the dot occupation $n_d$ and hence close to unity in the Kondo regime.

With no wire potential, the charges $n_w$ and $n_d$ coincide. If $\varepsilon_a = 0$, then $H_A$ reduces to the symmetric Hamiltonian, the dot occupation is unitary, and $\delta = \pi/2$. If, on the other hand, $\varepsilon_a \neq 0$, then a particle-hole transformation, which reverses the sign of $\varepsilon_a$, transmutes $n_d$ into $2 - n_d$ and hence changes the sign of $\delta - \pi/2$. The curve through the circles in Fig. 2 thus remains invariant under a rotation of $180^\circ$ around the point $\varepsilon_a = 0$, $\delta = \pi/2$.

The solid line through the circles marks the Kondo regime, throughout which $\delta \approx \pi/2$. Comparison with the solid lines through the squares and triangles shows that the positive gate potential applied to the wires displaces the Kondo regime to higher dot energies.

The negative charge induced by the potential pushes down the phase shifts and spoils the symmetry about $\varepsilon_a = 0$, $\delta = \pi/2$. In the region $\varepsilon_a > 0$, the wire charge $n_w$ can now become negative. While the circles cross the $\delta = \pi/2$ line at a saddle point, the slope of the curves through the squares and triangles is markedly negative. Thus, in contrast with the circles, only in narrow ranges of dot energies do the squares and triangles dwell near $\delta = \pi/2$ or $\delta = 0$. As the data moreover show, except for $W \approx 0$, the range in which $\delta \approx 0$ ($\delta \approx \pi/2$) is pinned to the vicinity of the mixed-valence region $\varepsilon_a \approx U/2$ ($\varepsilon_a \approx -U/2$).

Combined with Eq. (7), the information in Fig. 2 determines the thermal dependence of the conductance through the device in Fig. [7]. Detailed comparison between the resulting curves and the NRG conductances will be presented in Secs. IV A-V D. Before that, to provide a more visual description of the numerical data, we present bird’s-eye views of the conductance as a function of the temperature and dot-level energy.

B. Conductance for $W = 0$

Figure 3 is a landscape consolidating 71 $G(T)$ curves calculated for $W = 0$, $U = 5D$, $T_w = 0.15D$, and dot energies in the range displayed in Fig. 2. The inset is a reversed-perspective view of the same plot offering an unobstructed depiction of the low-temperature region.

Eqs. (9a)-(9d) explain the salient features of the landscape. The central portion of the plot encompasses the Kondo regime. Here, as the circles in Fig. 2 show, the ground-state phase shift is close to $\pi/2$. At high temperatures, Eq. (9d) yields the nearly ballistic conductance depicted by the hood-like central plateau visible in the main plot and inset. At low temperatures, Eq. (9b) brings the conductance down to nearly zero; the resulting Kondo valley can be seen in the inset.

We defer to Sec. IV E 1 the quantitative discussion of the crossover between the plateau and the valley. Here, we note that the Kondo temperature rises with $|\varepsilon_d|$. For $|\varepsilon_d| \approx U/2$, in particular, i.e., for $\Delta_0 \equiv -\varepsilon_d \leq \Gamma$ or $\Delta_2 \equiv \varepsilon_d + U \leq \Gamma$, the Kondo temperature becomes comparable with $\Gamma$, an indication that the Hamiltonian has transposed the limits of the Kondo domain to enter the mixed-valence regime.

At high temperatures, the mixed-valence regime is identified by the two antiresonances bringing the conductance to $G/2$ on the frontal plane in the main plot, or the rear plane in the inset. As the temperature is reduced, the conductance descends into the Kondo valley. As Sec. IV D will show, the decay can still be mapped onto the universal function $G^0(T/T_K)$. Given the high crossover temperature, however, only the low-temperature tail of the numerical results agrees quantitatively with $G^0$.

Beyond the mixed-valence regime, as $|\varepsilon_d|$ grows past $U/2$, the dot occupation approaches an even integer, $n_d = 0$ or $n_d = 2$, the coupling to the quantum wire becomes ineffective,
and the electrons flow ballistically across the wire.

C. Conductance for weak scattering potentials

Even moderate gate potentials applied to the wire change qualitatively the conductance landscape. Figure 3 shows plots analogous to Fig. 3 calculated for the same Coulomb repulsion $U = 5D$ and effective dot-level width $\Gamma_w = 0.15D$, and for three weak potentials: (a) $\rho W = 0.1D$; (b) $\rho W = 0.2D$; and (c) $\rho W = 0.3D$. Compared to the symmetric landscape in Fig. 3 the new plots show distinctions that evolve rapidly under the rising gate potential.

The Kondo-valley conductance rise with $\rho W$, while the high-temperature plateau diminishes. For the highest potential shown, $\rho W = 0.3$, the Kondo-regime conductance is nearly uniform. Outside the Kondo regime, in the even-$n_d$ regions, the charge $2\delta_n/e$ induced by the wire potential lowers the ground-state phase shift to $-\delta_n (\pi - \delta_n)$ in the $n_d = 0$ ($n_d = 2$) domain. The right-hand side of Eq. (9c) is no longer ballistic: $G(T = 0) = G_0^2 \cos^2 \delta_n$.

In the mixed-valence regime, the transition from the Kondo regime to the $n_d = 0$ domain is now markedly different from the transition to the $n_d = 2$ domain. The former encompasses the narrow $\delta = 0$ dot-energy range in Fig. 3 for which the right-hand side of Eq. (9c) yields $G(T = 0) \approx G_0^2$; while the latter traverses the equally narrow $\delta = \pi/2$ range, for which Eq. (9c) yields $G(T = 0) = 0$. An insulating ravine develops in the conductance landscape, pinned near $\varepsilon_d = -U$, while a ballistic ridge arises near the $\varepsilon_d = 0$ plane. Plotted at fixed temperature $k_B T \ll \Gamma_w$ against the dot energy,
the conductance displays the trough-crest pair that defines a Fano profile. As evidenced by Fig. 4(c), this fiducial mark of interference between currents flowing along parallel paths becomes most pronounced for intermediate wire potentials.

A brief glance at Fig. 1 identifies the interfering fluxes. One of them, indicated by the white arrows, is subject to obstruction by the charge accumulated in the central segment of the wire. The second conduction path runs through the quantum dot. To provide an effective bypass, it must avoid the Wannier orbital \( f_0 \). At first sight, given that Eq. (2) couples the quantum dot to \( f_0 \), this may seem impossible, and in fact it is, for \( W = 0 \). The wire potential nonetheless spreads the coupling over wire states beyond the central orbital closest to the dot.

To be more specific, let \( V \rightarrow 0 \) and consider the resulting wire Hamiltonian,

\[
H_w = \sum_k \varepsilon_k c_k^\dagger c_k - \frac{W}{N} \sum_{k,k'} \varepsilon_k^f c_k^f c_{k'}^\dagger.
\]

The diagonalization of this quadratic form yields

\[
H_w = \sum_t \varepsilon_t g_t^\dagger g_t,
\]

where

\[
g_t = \sum_k \alpha_{t,k} c_k,
\]

with coefficients \( \alpha_{t,k} \) that depend on \( W \). The eigenvalues \( \varepsilon_t \) are phase-shifted with respect to the \( \varepsilon_k \). In the vicinity of the Fermi level, in particular, the phase shifts \( \delta_w \) are uniform.\(^{17,26}\)

\[
\tan \delta_w = -\pi \rho W.
\]

In analogy with Eq. (11), we can therefore define

\[
\phi_0 \equiv \frac{1}{\sqrt{N}} \sum_t g_t.
\]

A straightforward calculation shows that \( \{ \phi_0, f_0 \} = \cos \delta_w \)\(^ {19}\). For \( W = 0 \), in particular, \( \phi_0 = f_0 \). It is also easy to check that the ground-state occupation \( n_0 \) of the \( \phi_0 \) orbital is \( 1 + 2\delta_w/\pi \), so that as required by the Friedel sum rule, the gate potential brings in an extra charge \( 2\delta_w/\pi \).

The operator \( \phi_0 \) plays the role that belonged to \( f_0 \) when \( W = 0 \). In the Kondo regime, the antiferromagnetic interaction between the dot moment and the moment of the electrons occupying the \( \phi_0 \) orbital is now responsible for the Kondo crossover\(^ {19,25}\) and at low temperatures, the \( \phi_0 \) occupation, equal to the charge in the Kondo cloud plus that induced by the wire potential, obstructs the transport along the white arrows in Fig. 1. At the same time, the overlap \( \sin \delta_w \) between \( f_0 \) and the wire states orthogonal to \( \phi_0 \) offers the dot bypass, through which the current runs to interfere with the flow across \( \phi_0 \) and generate the askew profile most visible on the \( k_BT = 10^{-10} D \) plane (i.e., the rear plane) in Fig. 4(c).

The amplitude for tunneling through the dot is \( \sin \delta_w \). The amplitude for tunneling across the orbital \( \phi_0 \) is \( \cos \delta_w \). The ratio between the two amplitudes defines the effective Fano parameter

\[
q_{\text{eff}} \equiv \tan \delta_w = -\pi \rho W.
\]

For \( W = 0 \), the right-hand side vanishes. The resulting conductance curve, stamped on the \( k_BT = 10^{-10} D \) plane in Fig. 5, exhibits the symmetry of the \( q = 0 \) Fano profile.\(^ {29}\) For \( \rho W = 0.3 \), by contrast, \( q_{\text{eff}} \approx -1 \), and the conductance profile, Fig. 4(c), is close to the maximum Fano asymmetry.

D. Conductance for strong scattering potentials

Under stronger gate potentials, while the amplitude for conduction along the white arrows in Fig. 1 diminishes, the dot bypass becomes more effective. The absolute value of the right-hand side of Eq. (15) grows, and the conductance profile at low temperatures gradually acquires the symmetry of the large-\( q \) Fano profile.

Illustrative landscapes are displayed in Figs. 5(a), (b), and (c), computed for \( \rho W = 0.4, 0.6 \), and 1.0, respectively, and \( U = 5 D \) and \( \Gamma_w = 0.15 D \). The dominant features of Fig. 4 reappear in the three plots. In particular, the ballistic ridge and the insulating ravine are still clearly visible in the mixed-valence regime. Sharp in (a) and (b) the Fano antiresonance becomes dull in Fig. 5(c), because as shown by the triangles in Fig. 2 the wire potential pushes \( \delta(\varepsilon_d) \) so far down that the phase-shift curve crosses the horizontal lines at the \( \delta = \pi/2 \) and \( \delta = 0 \) with relatively small slopes.

In the Kondo regime, while the high-temperature conductance falls, the low-temperature conductances rises steadily with \( \rho W \). In the end, with \( \rho W = 1 \), the conductance is nearly ballistic in the low-temperature Kondo plateau, while at high temperatures the Coulomb blockade (the charge induced by the wire potential) impedes transport through the dot (through the wire) and reduces the conductance to nearly zero.

The trend in Fig. 5 indicates that, in the large \( \rho W \) limit, the conductance landscape is complementary to the plot in Fig. 5, i.e., \( G_{W \rightarrow 0}(T) + G_{W \rightarrow D}(T) \approx G_2 \). In view of Eq. (2), we expect \( G_{W \rightarrow 0}(T) \) to reproduce the thermal-dependence of the conductance for a single-electron transistor\(^ {19,23}\) Comparison with Fig. 4 in Ref. 24 which describes a single-electron transistor with the same model parameters, confirms that, with \( \rho W \rightarrow \infty \), the conductances in the side-coupled and the embedded configurations are identical.

For \( W \rightarrow 0 \), instead of equal, the conductances are complementary. While the conductance for the \( T \)-shaped device displays the Kondo-valley profile in Fig. 3, the conductance through the single-electron transistor, which is insensitive to changes in the wire gate potential, retains the Kondo-plateau pattern to which the trend in Figs. 5(a)-(c) points.

E. Thermal dependence of the conductance

This section compares the NRG results for the temperature dependence of the conductance with the mapping (7). We fix
the parameters $U = 5D$ and $\Gamma_w = 0.15D$, and consider the eight gate wire potentials $\rho W$ in Table 1. To sample the dot-energy dependence of the data, four plots will be discussed, corresponding to the four dot energies indexed by the vertical arrows in Fig. 2. For each run, the table shows the Kondo temperature resulting from the definition $G(T) \equiv G_{1/2}$. The Kondo temperature marked with asterisks belong to Hamiltonians in the mixed-valence regime and, as explained in the text, had to be obtained by matching the solid lines to conductances computed at temperatures below $T_K$.

\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|c|c|}
\hline
\multicolumn{2}{|c|}{Figure} & $\rho W$ & $\delta/\pi$ & $k_BT/D$ \\
\hline
6 & $\circ$ & 0.00 & 0.50 & $8.1 \times 10^{-7}$ & $\bullet$ & 0.10 & 0.40 & $8.6 \times 10^{-7}$ \\
6 & $\Box$ & 0.20 & 0.32 & $1.0 \times 10^{-6}$ & $\Box$ & 0.30 & 0.26 & $1.3 \times 10^{-6}$ \\
6 & $\diamond$ & 0.40 & 0.22 & $1.6 \times 10^{-6}$ & $\uparrow$ & 0.60 & 0.16 & $2.5 \times 10^{-6}$ \\
6 & $\Delta$ & 0.80 & 0.13 & $3.8 \times 10^{-6}$ & $\downarrow$ & 1.00 & 0.11 & $6.0 \times 10^{-6}$ \\
7 & $\circ$ & 0.00 & 0.51 & $4.4 \times 10^{-6}$ & $\uparrow$ & 0.10 & 0.41 & $5.9 \times 10^{-6}$ \\
7 & $\Box$ & 0.20 & 0.33 & $8.9 \times 10^{-6}$ & $\Box$ & 0.30 & 0.27 & $1.4 \times 10^{-5}$ \\
7 & $\diamond$ & 0.40 & 0.23 & $2.2 \times 10^{-5}$ & $\uparrow$ & 0.60 & 0.17 & $5.7 \times 10^{-5}$ \\
7 & $\Delta$ & 0.80 & 0.15 & $1.4 \times 10^{-4}$ & $\downarrow$ & 1.00 & 0.13 & $3.6 \times 10^{-4}$ \\
8 & $\circ$ & 0.00 & 0.52 & $8.8 \times 10^{-5}$ & $\uparrow$ & 0.10 & 0.43 & $1.4 \times 10^{-4}$ \\
8 & $\Box$ & 0.20 & 0.35 & $2.5 \times 10^{-4}$ & $\Box$ & 0.30 & 0.29 & $4.5 \times 10^{-4}$ \\
8 & $\diamond$ & 0.40 & 0.25 & $9.6 \times 10^{-4}$ & $\uparrow$ & 0.60 & 0.22 & $3.0 \times 10^{-3}$ \\
8 & $\Delta$ & 0.80 & 0.22 & $1.0 \times 10^{-2}$ & $\downarrow$ & 1.00 & 0.25 & $1.2 \times 10^{-2}$ \\
9 & $\circ$ & 0.00 & 0.42 & $8.0 \times 10^{-3}$ & $\uparrow$ & 0.10 & 0.34 & $4.5 \times 10^{-3}$ \\
9 & $\Box$ & 0.20 & 0.27 & $2.7 \times 10^{-3}$ & $\Box$ & 0.30 & 0.21 & $2.0 \times 10^{-3}$ \\
9 & $\diamond$ & 0.40 & 0.17 & $1.4 \times 10^{-3}$ & $\downarrow$ & 0.60 & 0.12 & $6.6 \times 10^{-4}$ \\
9 & $\Delta$ & 0.80 & 0.09 & $3.5 \times 10^{-4}$ & $\downarrow$ & 1.00 & 0.08 & $1.9 \times 10^{-4}$ \\
\hline
\end{tabular}
\caption{Phase shifts and Kondo temperatures for the 32 NRG runs depicted in Figs. 5. The ground-state phase shifts $\delta$ come from the low-energy spectrum of the model Hamiltonian, and the Kondo temperatures, from the definition $G(T = T_K) \equiv G_{1/2}$. The Kondo temperature marked with asterisks belong to Hamiltonians in the mixed-valence regime and, as explained in the text, had to be obtained by matching the solid lines to conductances computed at temperatures below $T_K$.}
\end{table}

1. Symmetric dot Hamiltonian

Figure 6 shows the thermal dependence of the conductance for $\varepsilon_d = 0$ and the eight gate potentials $\rho W$ in Table 1. As the potential grows from $\rho W = 0$ to $\rho W = 1$, the ground-state phase shift decreases from $\pi/2$ to nearly zero, and the conductance $G(T)$ evolves from monotonically increasing to monotonically decreasing. In all cases, the model Hamiltonian lying well within the Kondo regime, the agreement with the solid curves representing Eq. (7) is excellent.

2. Asymmetric models

Figure 7 displays temperature-dependent conductances for $\varepsilon_d = -3.4D$, equivalent to $\varepsilon_a = -0.9D$. The other model parameters coincide with those in Fig. 6. As Table 1 shows, the enhanced particle-hole asymmetry makes each ground-state phase shift somewhat larger than the corresponding $\delta$ in Fig. 6. For $W = 0$ (open circles) the phase shift in Fig. 6 was $\pi/2$, so that the factor $\cos(2\delta)$ on the right-hand side of Eq. (7) is now larger than it was. The more negative dot energy thus flattens the conductance curve.

For each nonzero potentials, the phase shift associated with Fig. 7 in Table 1 makes the factor $\cos(2\delta)$ smaller than the corresponding factor in Fig. 6. Equation (9c) therefore pushes down the low-temperature conductances, while Eq. (9b) pushes up the high-temperature conductances. The Kondo temperatures are now spread over a wider range.
Apart from such minor changes, however, Figs. 6 and 7 display the same picture. In particular, all Hamiltonians and temperatures being in the Kondo regime, the numerical conductances are in very good agreement with Eq. (7).

More negative dot energies can displace the Hamiltonian into the mixed-valence regime. As indicated by the solid lines in Fig. 6 for $\rho W = 0$ the Kondo domain extends from $\epsilon_d = -(U/2 - T)$ to $\epsilon_d = U/2 - T$; with $\epsilon_d = -4 D$, i.e., $\epsilon_d = -1.5 D$, the model is still within it. As $\rho W$ grows, however, the Kondo regime is uniformly shifted toward higher dot energies. It results that the open and filled triangles in Fig. 8 which describe conductances for $\rho W = 0.8$ and 1, respectively, represent Hamiltonians in the mixed-valence regime.

While the other curves, computed for smaller $\rho W$’s, agree with Eq. (7), the deviations separating the triangles from the solid lines at high temperatures, $k_B T \gtrsim 10^{-2} D$, are substantial. The discrepancies are reminders that, in the mixed-valence regime, the dominant characteristic energy is $E_{\text{C}}^0 = \Gamma_w$, which restricts the domain of Eq. (7) to $k_B T \ll \Gamma_w = 0.15 D$. Our numerical study of the embedded configuration reported similar deviations.

Since the conductance curves for large $\rho W$ cross the $G = G_2/2$ horizontal at relatively high temperatures, which fall in the nonuniversal thermal range $k_B T \lesssim \Gamma_w$, it would be inappropriate to extract Kondo temperatures from the identification $(G(T)/T_K) = G_2/2$. To evaluate each of the two numbers marked with asterisks in Table, we therefore adjusted $T_K$ so that the corresponding solid line ran through the triangle closest to $G = 0.7 G_2$.

Since the wire potential displaces the Kondo regime to higher dot energies, for $\epsilon_d \to 0$ we expect the $\rho W = 0$ Hamiltonian to leave the Kondo regime before the $\rho W > 0$ Hamiltonians. Accordingly, Fig. 9 contrasts triangles and squares very well fitted by Eq. (7) with circles and diamonds that depart significantly from the solid lines representing the universal mapping for $k_B T > 3 \times 10^{-2} D$.

The ground-state phase shift $\delta$ decreases with $\epsilon_d$. Thus, while the phase shift for each curve in Figs. 7 and 8 exceeds the phase shift for the same $\rho W$ in Fig. 6 the phase shift for each curve in Fig. 9 is smaller than the corresponding phase shift in Fig. 6. In compliance with Eq. (9c), the low-temperature conductances in Fig. 9 are higher than the corresponding $G(T \to 0)$ in Fig. 6. At high temperatures, the analogous comparison is unproductive, because the relatively high Kondo temperatures place the high-temperature end of the axis in
A background current being detected, three parameters were involved in the fit to the first conductance profile: the background conductance, the phase shift, and the Kondo temperature $T_K$. The same first two parameters and a different Kondo temperature then fitted the second curve. In each case, within the dispersion of the experimental data, optimum agreement resulted.

Fig. 9 outside the reach of Eq. (9d).

**F. Comparison with experiment**

The numerical results in Figs. 3-5 can be compared to a number of experimental studies of transport in $T$-shaped devices. Particularly illustrative are Fig. 1(c) in Ref. 7 and Fig. 5 in Ref. 20. The former, which shows low-temperature conductance profiles for various level widths, exhibits valleys, plateaus, and antiresonances—all the conspicuous features in the conductance profiles at $kt^2 > 10^{-2} D$ in Figs. 3-5.

Even more instructive is the illustration in Ref. 20, which demonstrates that the wire gate potential controls the division of the current between the wire and the bypass. The evolution of the low-temperature conductance profile under an increasing wire gate potential is documented in the plot. The profile starts out as Kondo valley analogous to the low-temperature profile in Fig. 3. It then evolves to a curve marked by interference, analogous to the low-temperature profile in Fig. 4(c). Finally, it reaches a Kondo plateau, analogous to the low-temperature profile in Fig. 5(c). The NRG results and Eq. (7) are therefore in qualitative agreement with the experimental results.

What is more important, the mapping describes quantitatively the thermal dependence of measured conductances. A fit of the temperature-dependent conductances resulting from two gate potentials applied by Sato et al.24 to a $T$-shaped device was presented in Ref. 23. A background current being detected, three parameters were involved in the fit to the first conductance curve: the background conductance, the phase shift, and the Kondo temperature $T_K$. The same first two parameters and a different Kondo temperature then fitted the second curve. In each case, within the dispersion of the experimental data, optimum agreement resulted.

**V. CONCLUSIONS**

Equation (7) offers a unifying view of electrical conduction through a quantum wire side-coupled to a quantum dot. Valid over the entire Kondo regime, it captures with error $O(k_{B}T/D)$ the conductance crossover from the local-moment regime (i.e., high-temperatures, $T \gg T_K$), equivalent to a conduction band weakly interacting with the dot magnetic moment, to the low-temperature regime ($T \ll T_K$), in which the dot electron and the conduction electrons around it look into a singlet that reduces the model Hamiltonian to a phase-shifted conduction band.

The phase shift controls the thermal dependence of the conductance. With no wire gate potential, the Friedel makes $\delta \approx \pi/2$, and the Kondo cloud lowers the conductance from nearly ballistic at high $T$ to nearly zero at low $T$. For $W = 0$, the conductances for the $T$-shaped device and the single-electron transistor are complementary.

A (positive) gate potential applied to the wire in the $T$-shaped device induces electric charge that reduces the ground-state phase shift. As $\rho W$ pushes $\delta$ below $\pi/4$, the conductance curve first becomes flat and then reverses the pattern in Fig. 3. For large $\rho W$, as indicated by Fig. 5(c), the conductance through the $T$-shaped device approaches the conductance through a single-electron transistor.

The Friedel sum rule explains why the two devices respond differently to the wire gate potential. In the embedded configuration, since the dot is in the conduction path, its occupation $n_d$ controls the conductance. Although the wire gate potential $W$ determines the amount of charge in the wire, it has little influence upon $n_d$: according to the Friedel sum rule, $n_d = 2(\delta - \delta_w)/\pi$, so that the charge $2\delta_w/\pi$ induced by the gate potential $W$ is subtracted from the total charge $2\delta/\pi$ piled up in the quantum wire. We therefore expect the conductance $G$ through the single-electron transistor to be independent of $W$. The mapping between $G(T/T_K)$ and the universal function $G^2(T/T_K)$, which is parametrized by the phase-shift difference $\delta - \delta_w$, rationalizes this reasoning.

By contrast, the mapping between the conductance through the device in Fig. 4 and the universal function $G^2(T/T_K)$ is parametrized by the phase shift $\delta$. According to the Friedel sum rule, the wire charge is $2\delta/\pi$. Much as the dot charge $n_d = 2(\delta - \delta_w)/\pi$, which controls conduction in the embedded configuration, parametrizes the mapping to the conductance through the single-electron transistor, the wire charge $2\delta/\pi$ controls conduction in the side-coupled configuration and parametrizes the mapping. While the dependence on the dot charge practically shields the single-electron transistor from gate potentials applied to the wire, the dependence on the wire charge makes the conductance through the $T$-shaped device remarkably sensitive to such potentials.

The wire gate potential affects Fig. 4 in two ways. First, the total charge in the $\phi_0$ orbital, equal to the charge induced by the wire potential plus, at low temperatures, the charge in the Kondo droplet, obstructs transport through the central portion of the wire. Second, the potential couples $e_c$ to states orthogonal to $\phi_0$ and thus opens a bypass through the quantum dot. For very weak (strong) potentials, the path through $\phi_0$ ($e_c$)
is dominant; the landscape in Fig. 3 is thus complementary (very similar) to the landscape for a single-electron transistor.

For intermediate potentials [Fig. 4(c)], the interference between the currents along the two paths makes the landscape markedly asymmetric in the mixed-valence region. At the same time, in the Kondo regime, the two parallel conduction paths, one efficient at low temperatures, the other efficient at high temperatures, make the landscape remarkably flat.

The practical value of universality has been demonstrated in both experimental arrangements. In the embedded configuration, early in the history of the single-electron transistor, the universal function $G(T/T_K)$ guided the interpretation of conductance data. In the side-coupled arrangement, Eq. (7) was shown to fit conductance curves generated in the laboratory accurately enough to determine the Kondo temperature and ground-state phase shift within deviations set by the dispersion of the experimental data.

Neither the universal mapping (7), nor the corresponding expression for the embedded configuration determine explicitly any of the model parameters, let alone the physical constants emulated by the model Hamiltonian. They offer indirect information that may assist future ab initio descriptions of the physical properties of single-electron transistors or T-shaped devices. The two mappings of the Kondo-regime conductances to the universal function $G(T/T_K)$ redefine the ultimate target of such ab initio problems, from the computation of conductance curves to the computation of phase shifts and Kondo temperatures. In a less challenging arena, the mappings bring Kondo-regime conductance curves within the reach of the Bethe-ansatz approach. Finally, on the basis of the experience with thermodynamical properties, we expect the mappings to offer benchmarks to check the accuracy of numerically or analytically computed transport properties for model Hamiltonians describing quantum-dot arrays.

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References

1. Present address: Instituto de Física
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2. H. Aikawa, K. Kobayashi, A. Sano, S. Katsumoto, and Y. Iye, Phys. Rev. Lett. 92, 176802 (pages 4) (2004), URL http://link.aps.org/abstract/PRL/v92/e176802
3. A. Aharoni, O. Entin-Wohlman, T. Otsuka, S. Katsumoto, H. Aikawa, and K. Kobayashi, Phys. Rev. B 73, 195329 (pages 8) (2006).
4. A. Fuhrer, P. Brusheim, T. Ihn, M. Sigrist, K. Ensslin, W. Wegscheider, and M. Bichler, Phys. Rev. B 73, 205326 (pages 9) (2006).
5. L. G. G. V. Dias da Silva, N. P. Sandler, K. Ingersent, and S. E. Ulloa, Phys. Rev. Lett. 97, 096603 (pages 4) (2006).
6. I. Maruyama, N. Shibata, and K. Ueda, J. Phys. Soc. Japan 73, 3239 (2004).
7. M. Sato, H. Aikawa, K. Kobayashi, S. Katsumoto, and Y. Iye, Phys. Rev. Lett. 95, 066801 (2005).
8. K. Kobayashi, H. Aikawa, A. Sano, S. Katsumoto, and Y. Iye, Phys. Rev. B 70, 035319 (2004).
9. S. Katsumoto, J. Phys.: Condens. Matter 19, 233201 (2007).
10. B. R. Bulka and P. Stefanski, Phys. Rev. Lett. 86, 5128 (2001).
11. A. A. Clerk, X. Waintal, and P. W. Brouwer, Phys. Rev. Lett. 86, 4636 (2001).
12. M. E. Torio, K. Hallberg, A. H. Ceccatto, and C. R. Proetto, Phys. Rev. B 65, 085302 (2002).
13. R. Franco, M. S. Figueura, and E. V. Anda, Phys. Rev. B 67, 155301 (2003).
14. R. Zitko and J. Bonca, Phys. Rev. B 73, 035332 (2006).
15. I. Maruyama, N. Shibata, and K. Ueda, Physica B 378-380, 938 (2006).
16. A. C. Hewson, The Kondo Problem to Heavy Fermions (Cambridge University Press, Cambridge, 1993).
17. T. Costi, A. Hewson, and V. Zlatic, Journal of Physics–Condensed Matter 6, 2519 (1994).
18. P. W. Anderson, Phys. Rev. 124, 41 (1961).
19. M. Yoshida, A. C. Seridonio, and L. N. Oliveira, Universal zero-bias conductance for the single-electron transistor (2009), URL arXiv:0906.4063
20. S. Katsumoto, M. Sato, H. Aikawa, and Y. Iye, Physica E: Low-dimensional Systems and Nanostructures 34, 36 (2006).
21. T. Otsuka, E. Abe, S. Katsumoto, Y. Iye, G. L. Khym, and K. Kang, J. Phys. Soc. Japan 76, 084706 (2007).
22. L. G. G. V. Dias da Silva, K. Ingersent, N. Sandler, and S. E. Ulloa, Phys. Rev. B 78, 153304 (2008).
23. A. C. Seridonio, M. Yoshida, and L. N. Oliveira, Europhys. Lett. 86, 67006 (2009).
24. M. Yoshida, A. C. Seridonio, and L. N. Oliveira, Universal zero-bias conductance for the single-electron transistor: ii: Numerical results (2009), URL arXiv:0906.4228
25. F. D. M. Haldane, Phys. Rev. Lett. 40, 416 (1978).
26. H. R. Krishna-murthy, J. W. Wilkins, and K. G. Wilson, Phys. Rev. B 21, 1044 (1980).
27. H. R. Krishna-murthy, J. W. Wilkins, and K. G. Wilson, Phys. Rev. B 21, 1003 (1980).
28. U. Fano, Phys. Rev. 124, 1866 (1961).
29. B. Rulla, T. A. Costi, and T. Pruschke, Rev. Mod. Phys. 80, 395 (2008), ISSN 0034-6861.
30. D. C. Langreth, Phys. Rev. 150, 516 (1966).
31. D. Goldhaber-Gordon, H. Shtrikman, D. Mahalu, D. Abusch-Magder, U. Meirav, and M. A. Kastner, Nature 391, 156 (1998).
32. L. I. Glazman and M. E. Raikh, JETP Lett. 47, 452 (1987).
33. T. K. Ng and P. A. Lee, Phys. Rev. Lett. 61, 1768 (1988).
34. D. Goldhaber-Gordon, J. Göres, M. A. Kastner, H. Shtrikman, D. Mahalu, and U. Meirav, Phys. Rev. Lett. 81, 5225 (1998).
35. N. Andrei, K. Furuya, and J. H. Lowenstein, Rev. Mod. Phys. 55, 331 (1983).
36. A. M. Tsvelick and P. B. Wiegmann, Advances in Physics 32, 453 (1983).
37 J. B. Silva, W. L. C. Lima, W. C. Oliveira, J. L. N. Mello, L. N. Oliveira, and J. W. Wilkins, Phys. Rev. Lett. 76, 275 (1996).
38 V. L. Campo Jr and L. N. Oliveira, Phys. Rev. B 70, 153401 (2004).
39 Even if the Fano factor is defined with the model Hamiltonian, electronic correlations may renormalize it. An example was provided by L. N. Oliveira and J. W. Wilkins, Phys. Rev. B, 32, 696 (1985).