Quantum interference correction to the shot-noise power in nonideal chaotic cavities

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We present analytical results for the leading (weak-localization) quantum interference correction to the shot-noise power in a ballistic chaotic cavity coupled nonideally to two electron reservoirs via leads with an arbitrary number of open scattering channels. The calculations were performed using two independent methods: *S*-matrix diagrammatic perturbation theory and quantum circuit theory. We obtained an unexpected amplification-suppression transition as a function of both the number of open channels and the barriers' transparencies. The effect results from a subtle combination of temporal and spatial quantum coherence in the electron propagation through the system.

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I. INTRODUCTION

Quantum phase coherence has dramatic manifestations in the transport properties of mesoscopic devices such as quantum dots and quantum wires. The most common signatures of quantum interference phenomena resulting from multiple coherent wave scattering in the device are the weak-localization (WL) correction and mesoscopic fluctuations. Both effects can in principle be observed in any transport observable, but the conductance and the shot-noise power are the ones that have received most of the attention from both the experimental and the theoretical literatures.

Being the leading quantum interference correction to the average semiclassical value of a transport observable, the WL term is a valuable source of information of the underlying mechanisms of the quantum coherent dynamics in the system. Its dependence on experimentally controllable parameters such as temperature, magnetic field, and barriers' transparencies has been the subject of much current interest.[1] The WL correction to the sample's conductance is by far the most studied and understood mesoscopic phenomena. Several fundamental time scales, such as the Ehrenfest time, the phase-coherence time, and the dwell time emerge naturally from theoretical models of dephasing.[2–4] Other important time scales appear in crossover regimes[5] between universality classes and depend on the relevant scattering processes, such as the spin-orbit time and the magnetic time.

More recently, the WL correction to the shot-noise power in a chaotic cavity with ideal contacts has been studied in the crossover regime.[6,7] In Ref. 6 a semiclassical trajectory approach is used to describe the orthogonal-unitary crossover due to the breaking of time-reversal symmetry by an external magnetic field. Employing random-matrix theory Béri and Cserti[7] extended this result to a very general class of crossover that includes simultaneously spin-orbit coupling and an external magnetic field. A general consequence of the expressions for the weak-localization correction to the shot-noise power obtained by these authors is a change in sign or a suppression-amplification transition as a function of the crossover parameters.

For quantum dots in semiconductor heterostructures it is also possible to experimentally control the transparencies of the barriers by manipulating gate voltages. It is therefore also important to understand how the WL signal depends on the value of these contact parameters. In this paper we provide such an analysis for the shot-noise power in systems belonging to two pure Wigner-Dyson universality classes (orthogonal and symplectic). We find an unexpected change in sign in the WL correction to the shot-noise power as a function of both the barriers' transparencies and the number of open scattering channels. Surprisingly, the effect is completely induced by the barriers and does not depend on the presence of spin-orbit coupling inside the dot. Since the shot-noise power is the second cumulant of the charge counting statistics[8] (thus reflecting the time coherence of transmission events) and the WL correction is a wave interference phenomenon associated with spatial coherence, our effect is a consequence of a subtle combination of time and spatial coherence of electron propagation through the device.

Changes in sign in the WL correction in the absence of spin-orbit scattering is a very unusual quantum interference effect. Recently such an effect was predicted for the conductance of a multiterminal network of quasi-one-dimensional diffusive wires.[9] The effect was shown to have a geometrical origin. More specifically, if $T_{\alpha,\beta}$ is the transmission probability to go from contact β to contact α, the WL correction was shown to be proportional to a weighted average of the Cooperon P_c over each bond $(\mu\nu)$ of the network $\delta T_{\alpha\beta} \propto \sum_{\mu\nu} \frac{\partial T^{\rm cl}_{\alpha\beta}}{\partial l_{\mu\nu}} \int_{(\mu\nu)} dx P_c(x,x)$, where $l_{\mu\nu}$ is the length of the bond $(\mu\nu)$ and $T^{\rm cl}_{\alpha\beta}$ is the classical transmission probability. For a system with N wires plunged in the middle of a quasi-one-dimensional conductor the WL correction to the transmission across the conductor reads $\delta T = \frac{1}{3}(-1+N/4)$. A geometrically induced change in sign occurs at $N=4$, indicating a depletion-enhancement transition in the conductance caused by multiple coherent backscattering in the network.

A chaotic quantum dot is an open electron cavity with the form of a chaotic billiard. It is probably the simplest mesoscopic device that exhibits nontrivial quantum transport phenomena, such as weak-localization and universal mesoscopic fluctuations. In spite of its simplicity, the theoretical models that interrelate the observable quantum phenomena in this system are far from trivial. The main technical challenge is the development of an efficient scheme to remove the enormous dynamic redundancy caused by the chaotic motion inside the dot and in this way derive an optimal theory, i.e., the simplest mathematical representation of the relevant physical information. In the universal regime, defined by a mean dwell time much greater than both the ergodic time and the Ehrenfest time, and neglecting interaction effects, the random scattering matrix theory offers an essentially complete description since it incorporates nonperturbatively both temporal and spatial coherences. The description is however nonoptimal due to the large number of random variables. An optimal and quite powerful description is provided by Nazarov's quantum circuit theory.[10] The theory can describe perturbative effects such as the WL correction and universal mesoscopic fluctuations, as was recently shown by Campagnano and Nazarov.[11] An independent attempt to obtain an optimal description intimately connected to quantum circuit theory and including nonperturbative quantum fluctuations has been pursued in Refs. 12–15. In Sec. II, we describe the problem of full counting statistics (FCS) for charge transfer through a quantum dot in the scattering matrix formalism. The average shot-noise power, i.e., the second FCS cumulant, is calculated using a diagrammatic perturbation theory in inverse powers of the contact conductances between the dot and the leads. Large values of contact conductances are assumed so that Coulomb blockade effects are suppressed. In Sec. III, we rephrase the problem in the language of the supersymmetric nonlinear sigma model and quantum circuit theory. This offers an independent perturbative scheme, which we use to recalculate the dominant quantum correction (weak-localization term) of the average shot-noise power. The final expression, obtained only for equivalent channels, agrees with the general formula obtained in Sec. II. A brief discussion of the physical implications of our formula and conclusions is presented in Sec. IV.

II. SCATTERING MATRIX FORMALISM

Following Ref. 15, our starting point is the random scattering matrix description of the cumulant generating function of the FCS for charge transfer through a double-barrier chaotic quantum dot coupled to two leads, labeled as 1 and 2, with N_1 and N_2 open scattering channels, respectively. The FCS generating function is directly related to the following integral transform:

$$\Psi^{(\beta)}(\vec{\phi}) = \int dS\Omega^{(\beta)}(\vec{\phi},S)P^{(\beta)}(S), \tag{1}$$

where dS is the Haar measure over the appropriate unitary group, $P^{\beta}(S) \propto |\det(1-\bar{S}^{\dagger}S)|^{-\beta(N_1+N_2-1+2/\beta)}$ is the Poisson kernel and $\beta \in \{1,2,4\}$ is Dyson's symmetry index. The scattering matrix and its average read as

$$S = \begin{pmatrix} r & t' \\ t & r' \end{pmatrix}, \quad \bar{S} = \begin{pmatrix} r_1 & 0 \\ 0 & r_2 \end{pmatrix},$$

where r,r' and t,t' are the dot's reflection and transmission matrices, respectively, and $r_p = \mathrm{diag}(\sqrt{1-T_{p1}},\ldots,\sqrt{1-T_{pN_p}})$ is the reflection matrix of barrier p, which is fully characterized by its transmission coefficients T_{pn}, with $n=1,\ldots,N_p$.

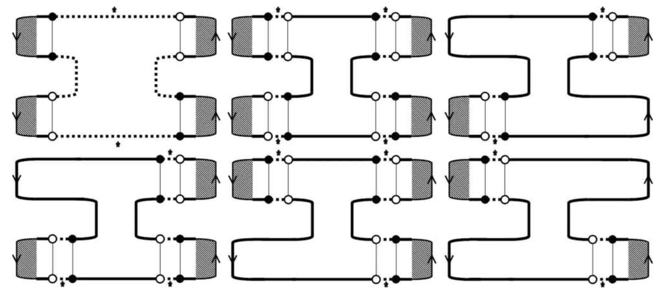

FIG. 1. Ladder diagrams contributing to the average shot-noise power. The maximally crossed diagrams that determine the weak-localization correction are obtained by crossing arms in each diagram.

The kernel of the integral transform is given by[16]

$$\Omega^{(\beta)}(\vec{\phi},S) = \prod_{\sigma=\pm} \det\left(\frac{1-\sin^2(\phi_{0\sigma}^{\beta}/2)tt^{\dagger}}{1+\sinh^2(\phi_{1\sigma}^{\beta}/2)tt^{\dagger}} \right), \tag{2}$$

where $\phi_{0\sigma}^{\beta} \equiv \phi_0 + \sigma\phi_{0\beta}$, with $\phi_{0\beta} = \phi_2\delta_{\beta,4}$, and $\phi_{1\sigma}^{\beta} = \phi_1 + \sigma\phi_{1\beta}$, with $\phi_{1\beta} = \phi_3\delta_{\beta,1}$. We have also defined the vector $\vec{\phi} \equiv (\phi_0,\phi_{0\beta},\phi_1,\phi_{1\beta})$. The FCS cumulants are given by

$$q_{l+1} = \lim_{\varepsilon\to 1}\left(\frac{\varepsilon}{2}\frac{d}{d\varepsilon}\right)^l \left.\frac{\varepsilon^2 I^{(\beta)}(\phi)}{\sin\phi}\right|_{\cos(\phi/2)=\varepsilon}, \quad l=0,1,\ldots,$$

where $I^{(\beta)}(\phi)$ is a quantity that plays the role of a pseudocurrent in quantum circuit theory. It is defined by

$$I^{(\beta)}(\phi) = -2^{m_\beta} \left.\frac{\partial\Psi^{(\beta)}(\vec{\phi})}{\partial\phi_{0+}^{\beta}}\right|_{\vec{\phi}=(\phi,\phi,i\phi,i\phi)}, \tag{3}$$

in which $m_\beta = 1+\delta_{\beta,4}$. From a mathematical point of view the problem is completely well defined. The average dimensionless conductance and the average dimensionless shot-noise power, for instance, are given by the first and second cumulant of FCS, respectively, i.e., $\langle g\rangle = q_1$ and $\langle p\rangle = q_2$. The calculation of the integral over the Haar measure, however, is a task that is far from trivial. Our calculation of $\langle p\rangle$ is based on a perturbative diagrammatic scheme for performing this integral devised by Brouwer and Beenakker.[17]

We start by writing the shot-noise power as

$$p = g - h, \tag{4}$$

where the conductance, $g = \mathrm{Tr}[tt^{\dagger}]$, and $h \equiv \mathrm{Tr}[(tt^{\dagger})^2]$ are linear statistics of transmission eigenvalues. Since the average conductance can be found in Ref. 17, we only need to compute $\langle h\rangle$ to obtain the average shot-noise power. It proves convenient to write h in terms of the scattering matrix,

$$h = \mathrm{Tr}[(C_1SC_2S^{\dagger})^2], \tag{5}$$

where C_1 and C_2 are projection matrices onto leads 1 and 2, with N_1 and N_2 open channels, respectively.[17]

Applying the diagrammatic rules to the average of Eq. (5) in the semiclassical limit of large contact conductances, we obtain a representation in terms of six distinct types of ladder diagrams, associated with diffusion modes in the system (see Fig. 1). We remark that symmetry considerations on the topology of the diagrams impose a counting factor of 2 when summing the diagrams of types 3 and 4. The sum of these six diagrams yields the dominant contribution of the average of h in the semiclassical limit,

$$\langle h\rangle = \frac{g_1^4\bar{g}_2 + 2g_1^3\bar{g}_1^2 + 2g_1^2\bar{g}_1^3 + g_2\bar{g}_1^4}{(g_1+\bar{g}_1)^4} + \mathcal{O}(1), \tag{6}$$

where $g_m = \Sigma_{n=1}^{N_1}(T_{1n})^m$ and $\bar{g}_m = \Sigma_{n=1}^{N_2}(T_{2n})^m$ are the parameters describing the contacts of the quantum dot to leads 1 and 2, respectively. The contact conductances, for instance, are given by g_1 and $\bar{g}_1$.

We now turn to the weak-localization correction. In the diagrammatic analysis, it contains two types of contributions, $h^{\mathrm{WL}} = \delta h_1 + \delta h_2$. The first one, denoted by δh_1, comes from the U-cycles of Fig. 1 (see Ref. 17 for a definition) and can be written as

$$\begin{aligned}\delta h_1 = &-2(g_1+\bar{g}_1)^{-5}(g_1^4\bar{g}_2 + 2g_1^3\bar{g}_1^2 - g_1^3\bar{g}_1\bar{g}_2 - g_1g_2\bar{g}_1^3 + 2g_1^2\bar{g}_1^3 \\ &+ g_2\bar{g}_1^4) + \mathcal{O}[(g_1+\bar{g}_1)^{-1}].\end{aligned} \tag{7}$$

The second contribution, denoted by δh_2, comes from maximally crossed or Cooperon's diagrams. They are systematically obtained by crossing the arms of the six diagrams of Fig. 1. Summing over the entire set of maximally crossed diagrams we obtain

$$\begin{aligned}\delta h_2 = &-2(g_1+\bar{g}_1)^{-6}(-g_1^2\bar{g}_1^3\bar{g}_2 + 4g_1^3\bar{g}_1^2\bar{g}_2 + 4g_1^2\bar{g}_1^3g_2 + 3g_1^4\bar{g}_1\bar{g}_2 \\ &- g_1^3\bar{g}_1^2g_2 - 2g_2\bar{g}_1^5 + 2g_1^5\bar{g}_3 + 2\bar{g}_1^5g_3 - 4g_1^4\bar{g}_2^2 + 2g_1\bar{g}_1^3\bar{g}_2g_2 \\ &+ 2g_1^3\bar{g}_1g_2\bar{g}_2 - 4g_1^2\bar{g}_1^2g_2\bar{g}_2 - 2g_1^5\bar{g}_2 - 4g_2^2\bar{g}_1^4 + 3g_1\bar{g}_1^4g_2 \\ &+ 2\bar{g}_1^4g_3g_1 + 2g_1^4\bar{g}_3\bar{g}_1 - 2g_1^4\bar{g}_1^2 - 4g_1^3\bar{g}_1^3 - 2g_1^2\bar{g}_1^4) \\ &+ \mathcal{O}[(g_1+\bar{g}_1)^{-1}].\end{aligned} \tag{8}$$

We are now in position to state our main result. Combining Eqs. (6)–(8) and using the results of Ref. 17 for the average conductance, we obtain the following expression for the average shot-noise power:

$$\begin{aligned}\langle p\rangle = &(g_1+\bar{g}_1)^{-4}(g_1^4\bar{g}_1 + g_1^3\bar{g}_1^2 - g_1^4\bar{g}_2 - \bar{g}_1^4g_2 + g_1^2\bar{g}_1^3 + g_1\bar{g}_1^4) \\ &+ \left(\frac{2}{\beta}-1\right)(g_1+\bar{g}_1)^{-6}(-3g_2\bar{g}_1^5 + 4g_1^5\bar{g}_3 + 4\bar{g}_1^5g_3 - 8g_1^4\bar{g}_2^2 \\ &- 3g_1^5\bar{g}_2 - 8g_2^2\bar{g}_1^4 + 4g_1^3\bar{g}_1g_2\bar{g}_2 + 4g_1\bar{g}_1^3\bar{g}_2g_2 - 8g_1^2\bar{g}_1^2g_2\bar{g}_2 \\ &- 3g_1^2\bar{g}_1^3\bar{g}_2 + 3g_1^3\bar{g}_1^2\bar{g}_2 + 3g_1^2\bar{g}_1^3g_2 + 3g_1^4\bar{g}_1\bar{g}_2 - 3g_1^3\bar{g}_1^2g_2 \\ &+ 3g_1\bar{g}_1^4g_2 + 4\bar{g}_1^4g_3g_1 + 4g_1^4\bar{g}_3\bar{g}_1) + \mathcal{O}[(g_1+\bar{g}_1)^{-1}].\end{aligned} \tag{9}$$

The first term, denoted by p^{CL}, is the dominant semiclassical contribution and coincides with Whitney's result[3] in the limit of vanishing Ehrenfest time. The second term, denoted by p^{WL}, is the weak-localization correction and constitutes the central result of this paper. We remark that it applies to all Wigner-Dyson ensembles where the symmetry index satisfies $\beta \in \{1,2,4\}$. In order to gain some insight into its physical meaning, we shall derive the particular case of equivalent channels, where $T_{pn} = T_p$ for all n, in Sec. IV by applying quantum circuit theory.

III. QUANTUM CIRCUIT THEORY

There is an alternative, much more intuitive, formulation of the problem that is based on the following remarkable exact transformation of color-flavor type:[18]

$$\int dS\Omega^{(\beta)}(\vec{\phi},S)P^{(\beta)}(S) = \int dQ e^{-S^{(\beta)}(Q,Q_{\hat{\phi}})}, \tag{10}$$

in which the "gluon field" represented by the unitary matrix S is replaced by a "meson field" represented by a supermatrix Q, satisfying the constraint $Q^2=1$. The action in flavor space reads as $S^{(\beta)}(Q,Q_{\hat{\phi}}) = m_\beta[S_1(Q,Q_{\hat{\phi}}) + S_2(Q,Q_0)]$, where

$$S_p(Q,Q') = \frac{1}{4}\sum_{n=1}^{N_p} \mathrm{Str}\,\ln\left(1 - \frac{T_{pn}}{4}(Q-Q')^2\right), \tag{11}$$

in which Str denotes as the supertrace operation. The source field $Q_{\hat{\phi}}$ is chosen to have the form,

$$Q_{\hat{\phi}} = \begin{pmatrix} 0 & e^{-i\hat{\phi}} \\ e^{i\hat{\phi}} & 0 \end{pmatrix}, \tag{12}$$

where $\hat{\phi} = \mathrm{diag}(i\phi_B, \phi_F)$, and

$$\phi_B \equiv \begin{pmatrix} \phi_1 & \phi_{1\beta} \\ \phi_{1\beta} & \phi_1 \end{pmatrix}, \quad \phi_F \equiv \begin{pmatrix} \phi_0 & \phi_{0\beta} \\ \phi_{0\beta} & \phi_0 \end{pmatrix}. \tag{13}$$

We have also defined $Q_0 = Q_{\hat{\phi}=0}$. The advantage of the supermatrix formulation is the natural reorganization of the semiclassical expansion in terms of a well-controlled expansion around the saddle point. The dominant contribution comes from the saddle-point equation,[15] which reproduces the pseudocurrent conservation law of quantum circuit theory,[10] $I^{(\beta)}(\phi) \simeq I_{sp}(\phi) = I_1(\phi-\chi) = I_2(\chi)$, where

$$I_p(\phi) = \sum_{n=1}^{N_p} \frac{2T_{pn}\tan(\phi/2)}{1+(1-T_{pn})\tan^2(\phi/2)}, \quad p=1,2. \tag{14}$$

Going beyond the saddle-point contribution one obtains quantum interference corrections so that the full expansion in inverse powers of the classical conductance reads as $I^{(\beta)}(\phi) = I_{sp}(\phi) + I^{\mathrm{WL}}(\phi) + \ldots$, where $I^{\mathrm{WL}}(\phi)$ denotes as the weak-localization correction. For systems with unitary symmetry ($\beta=2$) and ideal contacts ($T_{pn}=1$) the entire expansion can be explicitly calculated.[15] In a recent insightful paper, Nazarov and Campagnano[11] derived $I^{\mathrm{WL}}(\phi)$ directly from quantum circuit theory. Their expression is remarkably general since it contains all types of crossovers between the Wigner-Dyson symmetry classes. For the pure classes their result is expected to coincide with the above supermatrix formulation. In our notation it can be written as

$$I^{\mathrm{WL}}(\phi) = \frac{2-\beta}{\beta}\frac{d}{d\phi}\ln\frac{M_+(\phi)}{M_-(\phi)}, \tag{15}$$

where $M_{\pm}(\phi)$ are the eigenvalues of a finite element version of the Cooperon and are given by

$$M_{+}(\phi) = 2I_{sp}(\phi)\left[\cot\left(\chi(\phi)-\frac{\phi}{2}\right)-\cot\left(\chi(\phi)+\frac{\phi}{2}\right)\right] - 2\left[\frac{I'_{sp}(\phi)}{\chi'(\phi)+1/2} - I_{sp}(\phi)\cot\left(\chi(\phi)+\frac{\phi}{2}\right)\right] + 2\left[\frac{I'_{sp}(\phi)}{\chi'(\phi)-1/2} - I_{sp}(\phi)\cot\left(\chi(\phi)-\frac{\phi}{2}\right)\right],$$

$$M_{-}(\phi) = 2I_{sp}(\phi)\left[\cot\left(\chi(\phi)-\frac{\phi}{2}\right)-\cot\left(\chi(\phi)+\frac{\phi}{2}\right)\right],$$

where $I'_{sp}(\phi) \equiv dI_{sp}(\phi)/d\phi$ and the intermediate phase function $\chi(\phi)$ is obtained from the saddle-point equation written in the form, $I_{sp}(\phi)=I_1(\phi/2+\chi)=I_2(\phi/2-\chi)$. The main technical difficulty in obtaining explicit expressions for observables from Eq. (15) is the analytical control over the physical root of the saddle-point equation, which for the case of equivalent channels $T_{pn}=T_p$ can be transformed into a polynomial equation of fourth degree. In this case, the weak-localization correction to the shot-noise power can be obtained and for $\beta=1$; it reads as

$$p^{\mathrm{WL}} = \frac{G_1G_2(G_1-G_2)(G_1T_2+G_2T_1)[3(G_2^2-G_1^2)+4(G_1^2T_2-G_2^2T_1)]}{(G_1+G_2)^6}, \tag{16}$$

which coincides with the second term of Eq. (9) if we set $\beta=1$, $g_p=N_1T_1^p$, $\bar{g}_p=N_2T_2^p$, and $G_i=N_iT_i$, where $i=1,2$. Note that we recover the well-known result $p^{\mathrm{WL}}=N_1N_2(N_1-N_2)^2/(N_1+N_2)^4$ obtained by Beenakker[19] in the case of ideal contacts. In this case, the quantum correction is always non-negative and vanishes for $N_1=N_2$. Our more general result [Eq. (9)] predicts a much richer behavior as we describe in Sec. IV.

IV. DISCUSSION AND CONCLUSIONS

For the sake of clarity let us concentrate on the case of equivalent channels. The results can be best understood through diagrams in the planes (T_1,T_2) for fixed $a\equiv N_2/N_1$ and (a,T_2) for fixed T_1. The main new physical effect is the existence of regions in parameter space, denoted by (−) and (+) in panels (2) and (4) of Fig. 2, where $p^{\mathrm{WL}}<0$ and $p^{\mathrm{WL}}>0$, respectively. Consequently, the system exhibits a barrier-induced suppression-amplification transition in the quantum correction to the shot-noise power. We can also define regions, denoted as (0), (I), and (II) in panels (1) and (3) of Fig. 2 where p^{WL} has, as a function of the remaining variable, zero (0), one (I), or two (II) sign changes. The shapes of the boundary lines separating these regions are independent of the remaining variable. The most interesting point in the diagrams is $p_c^{\mathrm{WL}}=(3/4,3/4)$. At p_c^{WL}, $p^{\mathrm{WL}}=0$ for all values of a. Therefore p_c^{WL} separates two a-independent regions in panel (1) of Fig. 2: one where $p^{\mathrm{WL}}<0$ for points on the line (T,T) with $T<3/4$ and another where $p^{\mathrm{WL}}>0$ for points on the line (T,T) with $T>3/4$. Consequently, the lower boundary lines between regions (+) and (−) shown in panel (2) of Fig. 2 (both the full and the dashed line) are located inside the region (II) of panel (1) of Fig. 2. Finally, the linear suppression of p^{WL} in the opaque limit, defined in Ref. 3 as $T_i\to 0$ with G_i fixed, is an interesting feature of our

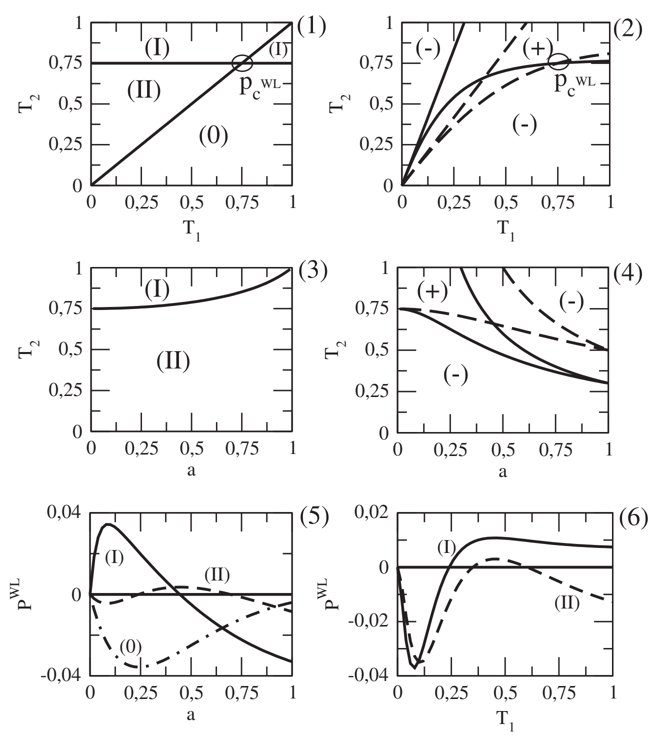

FIG. 2. Diagrams showing the transition. (1) In regions (0), (I), and (II) p^{WL} has, respectively, zero, one, and two sign changes as functions of a in the interval [0,1]; (2) Diagrams (T_1,T_2) separating positive (+) and negative (−) regions for $a=3/10$ (continuous lines) and $a=6/10$ (dashed lines); (3) Diagram (a,T_2) separating regions (I) and (II), where p^{WL} has one and two sign changes as functions of T_1, respectively; (4) Positive (+) and negative (−) regions in the (a,T_2) plane for $T_1=3/10$ (continuous lines) and $T_1=5/10$ (dashed lines); (5) Plots of p^{WL} showing the sign changes as a function of a for points at each region in diagram (1). (6) Plots of p^{WL} showing the sign changes as functions of T_1 for each region in diagram (3).

result. A similar effect for the weak-localization correction to the conductance has a nice physical explanation in the semiclassical approach.[3]

We believe that the transitions predicted in this work could become a useful tool to detect experimentally, in a controlled way, the weak-localization correction to shot-noise power in ballistic chaotic quantum dots. Experiments with tunable barriers, such as the ones discussed in Refs. 20 and 21, are already able to detect small variations in the electric current and in this way extract both the full counting statistics and cumulants up to the fifth order. From a conceptual point of view, our result provides strong links interconnecting quantum circuit theory, the supersymmetric nonlinear sigma model, the random scattering matrix approach, and the trajectory-based semiclassical theory.

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