I. INTRODUCTION

Statistical properties of the electron eigenfunctions in disordered quantum dots have recently become a subject attracting considerable theoretical and experimental interest [1] - [8]. One of the reasons is that the problem of particle motion in a bounded disordered potential comprises a particular case of general chaotic systems, such as quantum billiards [9]. On the other hand, the development of the “microwave-cavity” technique [1] - [3] has opened up a unique possibility for experimental studies of wave-function statistics. Furthermore, the mesoscopic conductance fluctuations in a quantum dot in the Coulomb blockade regime contain fingerprints of the eigenfunctions of confined electrons [2], [10] - [13]. This also provides experimental access [14,15] to the wave-function statistics.

In order to characterize the eigenfunction statistics, we introduce two distribution functions. The first refers to the local amplitudes of electron eigenfunctions $\psi_n$ inside a dot of volume $V$,

$$P(v) = \langle \delta (v - V |\psi_n(r)|^2) \rangle,$$

while the second distribution function

$$Q(s) = \langle \delta (s - \sin \varphi_n(r)) \rangle$$

is associated with the eigenfunction phases, $\varphi_n = \arg \psi_n$. Angular brackets $\langle \cdots \rangle$ denote averaging over bulk impurities inside a quantum dot.

When all electronic states are extended (metallic regime), the random matrix theory (RMT) [6] predicts the Porter-Thomas distribution of eigenfunction intensity that only depends on the fundamental symmetry of the quantum dot: $P(v) = \frac{\beta}{\pi^2} \left( \frac{\beta v}{\pi} \right)^{\beta/2-1} \exp \left( -\beta v/2 \right)$. The parameter $\beta = 1$ for a system having time-reversal symmetry (orthogonal ensemble), whereas $\beta = 2$ when time-reversal symmetry is broken (unitary ensemble), and $\beta = 4$ for a system having time-reversal symmetry and strong spin-orbit interactions (symplectic ensemble).

The RMT predictions are known to be valid in the limit of large conductance $g = E_c/\Delta \gg 1$ (here, $E_c = hD/L^2$ is the Thouless energy, $D$ is the classical diffusion constant, $L$ is the system size, and $\Delta$ is the mean level spacing) providing the system size $L$ is much larger than the electron mean free path $l$. Corrections to the Porter-Thomas distribution calculated in the framework of the $\sigma$-model formalism [11] are of order $g^{-1}$ in the weak-localization domain.

The limiting case of orthogonal symmetry corresponds to systems with pure potential electron-impurity scattering, while the case of unitary symmetry applies to systems in a strong magnetic field that breaks the time-reversal symmetry completely. For intermediate magnetic fields, a crossover occurs between pure orthogonal and pure unitary symmetry classes. This crossover in the distribution function $P(v)$ was previously studied within the framework of supersymmetry techniques [12] and within an approach [13] exploiting the analogy with the statistics of radiation in the regime of the crossover between ballistic and diffusive transport.

We consider below the problem of eigenfunction statistics of chaotic electrons in a quantum dot in an arbitrary magnetic field in the framework of RMT. Our treatment is related to the case of the metallic regime, and describes statistical properties of eigenfunctions amplitudes and phases in the regime of the orthogonal-unitary crossover. The results are applied to describe distributions of level widths and conductance peaks for weakly disordered quantum dots in the Coulomb blockade regime in the presence of an arbitrary magnetic field.
II. DISTRIBUTION OF LOCAL AMPLITUDES AND PHASES OF EIGENFUNCTIONS

In order to study the statistical properties of electron eigenfunctions within the RMT approach, we replace the microscopic Hamiltonian $H$ of an electron confined in a dot by the $N \times N$ random matrix $H = S_β \bar{\varepsilon} S_β^{-1}$ that exactly reproduces the energy levels $\varepsilon_n$ of the electron in a dot for a given impurity configuration. Here, $\bar{\varepsilon} = \text{diag}(\varepsilon_1, ..., \varepsilon_N)$, $S_β$ is a matrix that diagonalizes matrix $H$ (parameter $β$ reflects the system symmetry), and it is implied that $N \to \infty$. An ensemble of such random matrices reproduces the electron eigenvalues for microscopically different but macroscopically identical realizations of the random potential, and therefore it should describe the level statistics of electrons inside the dot. As long as we consider the metallic regime, the relevant ensemble of random matrices is known to belong to the Wigner-Dyson class \cite{16,17}. Correspondingly, averaging over randomness of the system is replaced by averaging over the distribution function $P[H] \propto \exp\{-\text{tr}V[H]\}$ of the random-matrix elements, where $V(\varepsilon)$ is a so-called “confinement potential” that grows at least as fast as $|\varepsilon|$ at infinity \cite{18,19}. In such a treatment, the columns of a diagonalizing matrix $S_β$ contain eigenvectors of the matrix $H$, that is $\phi_j(r_i) = (N/V)^{1/2}(S_β)_{ij}$, provided the space inside a dot is divided onto $N$ boxes with radius vectors $r_i$ enumerated from 1 to $N$ (the coefficient $(N/V)^{1/2}$ is fixed by the normalization condition $\int d\varepsilon |\phi_j(r)|^2 = 1$).

For the case of pure orthogonal symmetry, the eigenfunctions are real, being the columns of the orthogonal matrix $S_1$ (up to a normalization constant $(N/V)^{1/2}$). When time-reversal symmetry is completely broken, the eigenfunctions are complex and may be thought of as elements of the unitary matrix $S_2$. We note that in this case, the real and imaginary parts of the eigenfunction are statistically independent, and $\langle|\text{Re}\phi_j(r)|^2\rangle = \langle|\text{Im}\phi_j(r)|^2\rangle$. It is natural to assume that in the transition region between orthogonal and unitary symmetry classes, the eigenfunctions of electrons can be constructed as a sum of two independent (real and imaginary) parts with different weights, such as $\langle|\text{Im}\phi_j(r)|^2\rangle / \langle|\text{Re}\phi_j(r)|^2\rangle = \gamma^2$, where the transition parameter $\gamma$ accounts for the strength of the symmetry breaking. This means that in the crossover region, the eigenfunctions $(V/N)^{1/2}\phi_j(r_i)$ can be imagined as columns of the matrix $S$,

$$S = \frac{1}{\sqrt{1+\gamma^2}} (O + i\gamma \bar{O}), \quad OO^T = 1, \quad \bar{O}O^T = 1,$$

composed of two independent orthogonal matrices $O$ and $\bar{O}$. The parameter $\gamma$ in the parametrization Eq. (3) governs the crossover between pure orthogonal symmetry ($\gamma = 0$) and pure unitary symmetry ($\gamma = 1$). The values $0 < \gamma < 1$ correspond to the transition region between the two symmetry classes. From a microscopic point of view, this parameter is connected to an external magnetic field. (This issue will be discussed below.)

From Eq. (3), we obtain the moments

$$\mu_p(\gamma, N) = \langle|S_{ij}|^{2p}\rangle_S = \frac{1}{(1+\gamma^2)^p} \times \sum_{q=0}^{p} \binom{p}{q} \gamma^{2(p-q)} \langle|O_{ij}|^{2q}\rangle_O \langle|\bar{O}_{ij}|^{2(p-q)}\rangle_{\bar{O}}.$$

Here $\langle...\rangle_O$ stands for integration over the orthogonal group \cite{20}. The summation in Eq. (4) can be carried out using the large-$N$ formula

$$\langle|O_{ij}|^{2p}\rangle_O = \left(\frac{2}{N}\right)^p \Gamma\left(p + \frac{1}{2}\right) \sqrt{\pi},$$

leading to

$$\mu_p(\gamma, N) = \frac{p!}{N^p} \frac{P_p(X)}{X^p}, \quad X = \frac{1}{2} (\gamma + \gamma^{-1}),$$

which is valid over the whole transition region $0 \leq \gamma \leq 1$ in the thermodynamic limit $N \to \infty$. (Here $P_p(X)$ denotes the Legendre polynomial). One can easily see from the properties of Legendre polynomials that the ansatz introduced by Eq. (3) leads to the correct moments in the both limiting cases $\gamma = 0$ [$\mu_p(0, N) = (2p-1)!!/N^p$] and $\gamma = 1$ [$\mu_p(1, N) = p!/N^p$].

Using the RMT mapping described above, we can write the distribution function $P(v)$, Eq. (5), in the form

$$P(v) = \int \frac{d\omega}{2\pi} \exp\{-i\omega v\} \sum_{p=0}^{\infty} \frac{(i\omega N)^p}{\Gamma(p+1)} \langle|S_{rn}|^{2p}\rangle_S,$$

whence we get, with the help of Eq. (3)

$$P(v) = X \exp\{-vX^2\} I_0 \left(vX \sqrt{X^2 - 1}\right),$$

where $I_0$ is the modified Bessel function. Equation (8) gives the distribution function of the local amplitudes of electron eigenfunctions in a quantum dot [see Fig. 1]. It is easy to confirm that for pure orthogonal ($X \to \infty$) and pure unitary ($X = 1$) symmetries, Eq. (8) yields

$$P(v)\big|_{X \to \infty} = \frac{1}{\sqrt{2\pi v}} \exp\left(-\frac{v}{2}\right),$$

$$P(v)\big|_{X = 1} = \exp(-v).$$

2
A different approach to the issue of eigenfunction statistics was recently proposed in Ref. [3], whose authors obtained a single-integral representations for $P(v)$ [see their Eq. (17)]. It can be shown that our Eq. (8) coincides with Eq. (17) of the cited work provided the parameter $X$ is related to the parameter $Y$ appearing in Ref. [3] by $X = \sqrt{1+1/Y}$.

The advantage of the random-matrix approach is that it allows one to calculate in a rather simple way, the distribution functions after appropriate rescaling of our phenomenological parameter $X$ [see inset in Fig. 1, where the dis-

Using Eq. (3), we obtain after straightforward calculations

$$\sigma_{2p} (\gamma) = \frac{1}{\Gamma (p)} \int_0^\infty dx x^p \sum_{q=0}^\infty \frac{(-1)^q x^q}{\Gamma (q+1)}$$

Equations (10) and (14) yield the following formula for the phase distribution function in the crossover regime:

$$Q (s) = \frac{1}{2 \sqrt{1-s^2}} \frac{1}{\gamma^2 + s^2 (1-\gamma^2)}.$$  

As can be seen from Eq. (13), the limiting case of pure orthogonal symmetry is characterized by the $\delta$-functional phase distribution

$$\frac{1}{2} \sqrt{1-s^2} Q (s)|_{\gamma \rightarrow 0} = \frac{1}{2} \delta (s),$$

whereas the case of pure unitary symmetry is described by the uniform distribution

$$\frac{1}{2} \sqrt{1-s^2} Q (s)|_{\gamma \rightarrow 1} = \frac{1}{2 \pi}.$$
In accordance with Ref. [3], the microscopic parameter $X_m = \left| \phi / \phi_0 \right| (\alpha_g E_c / \Delta)^{1/2}$, where $\alpha_g$ is a factor depending on the sample geometry, $\phi$ is the magnetic field flux penetrating into the cross-sectional area of a sample, and $\phi_0$ is the flux quantum. It can be seen from Eq. [3] that for very weak breaking of time-reversal symmetry ($X_m \ll 1$), $X \approx 2 \sqrt{\pi} / 3 X_m$. In the opposite limit of weak deviation from unitary symmetry ($X_m \gg 1$), $X \approx 1 + 1 / 2 X_m^2$.

The connection between parameters $X$ and $X_m$ given by Eqs. [3] and [3] allows to use the simple random-matrix model proposed here for describing various phenomena occurring in quantum dots in arbitrary magnetic fields.

**III. STATISTICS OF RESONANCE CONDUCTANCE OF A QUANTUM DOT**

The issue of eigenfunction statistics is closely connected to the problem of conductance of a quantum dot that is weakly coupled to external leads in the Coulomb blockade regime. At low temperatures, the conductance of a dot exhibits sharp peaks as a function of the external gate voltage. The height of conductance peaks strongly fluctuates since the coupling to the leads depends on the fluctuating magnitudes of electron eigenfunctions near the leads. Thus far both theoretical [10] - [13] and experimental [14,15] studies were restricted to the pure symmetry classes (with conserved or completely broken time-reversal symmetry) even in the simplest case of two pointlike leads. Here we present the analytical treatment of the problem for arbitrary magnetic fields.

Let us consider non-interacting electrons confined in a quantum dot of volume $V$ with weak volume disorder (metallic regime). The system is probed by two pointlike leads weakly coupled to the dot at the points $r_L$ and $r_R$. The Hamiltonian of the problem can be written in the form [11]

$$H = \frac{\hbar^2}{2m}(i \nabla + \frac{e}{c \hbar} A)^2 + U(r)$$

$$+ \frac{i}{\tau_H} V [\alpha_R \delta(r-r_R) + \alpha_L \delta(r-r_L)],$$

where $U$ consists of the confinement potential and the potential responsible for electron scattering by impurities, $\tau_H$ is the Heisenberg time, and $\alpha_{R(L)}$ is the dimensionless coupling parameter of the right (left) lead. We also suppose that the coupling between the leads and the dot is extremely weak, $\alpha_{R(L)} \ll 1$, so that the only mechanism for electron transmission through the dot is tunneling.

It can be shown that the heights of conductance peaks are entirely determined by the partial level widths

$$\gamma_{\nu R(L)} = \frac{2 \alpha_{R(L)} V}{\tau_H} |\psi_{\nu}(r_{R(L)})|^2$$

in two temperature regimes. When $T \ll \alpha_{R(L)} \Delta$, the heights $g_{\nu} = \frac{1}{2e^2} G_{\nu}$ are given by the Breit-Wigner formula

$$g_{\nu} = \frac{4 \gamma_{\nu R} \gamma_{\nu L}}{(\gamma_{\nu R} + \gamma_{\nu L})^2}.$$ (22)

At higher temperatures, $\alpha_{R(L)} \Delta \ll T \ll \Delta$, the heights are determined by the Hauser-Feshbach formula

$$g_{\nu} = \frac{\pi}{2 T} \gamma_{\nu R} \gamma_{\nu L}.$$ (23)

Using the distributions $P_X(\gamma_{\nu R(L)})$ of the partial level widths, Eqs. (21) and Eq. (8),

$$P_X(\gamma_{\nu R(L)}) = \frac{X \tau_H}{2 \alpha_{R(L)}^2} \exp \left( - \frac{X^2 \tau_H}{2 \alpha_{R(L)}^2} \gamma_{\nu R(L)} \right)$$

$$\times I_0 \left( \frac{X \sqrt{X^2 - 1} \tau_H}{2 \alpha_{R(L)}^2} \gamma_{\nu R(L)} \right),$$ (24)

and assuming that the eigenfunctions of electrons near the left and right contacts are uncorrelated, that is $|r_R - r_L| \gg \lambda$, we can derive analytical expressions for the distribution of the conductance peaks heights.

**A. Breit-Wigner regime**

From Eqs. (22), we conclude that the distribution of the conductance peaks is given by

$$R_X(g_{\nu}) = \int_0^\infty d^2 \gamma_{\nu R} \int_0^\infty d^2 \gamma_{\nu L} P_X(\gamma_{\nu R}) P_X(\gamma_{\nu L})$$

$$\times \delta \left( g_{\nu} - \frac{4 \gamma_{\nu R} \gamma_{\nu L}}{(\gamma_{\nu R} + \gamma_{\nu L})^2} \right).$$ (25)

Straightforward calculation lead to the following integral representation

$$R_X(g_{\nu}) = \frac{\Theta(1 - g_{\nu})}{\sqrt{1 - g_{\nu}}} \int_0^\infty d\mu \mu$$

$$\times \left\{ \prod_{k=\pm} \exp \left( -X S^{(k)} f^{(k)} \mu \right) I_0 \left( \frac{S^{(k)} f^{(k)} \mu \sqrt{X^2 - 1}}{f^{(k)} \mu} \right) \right\}$$

$$+ \prod_{k=\pm} \exp \left( -X S^{(k)} f^{(k)} \mu \right) I_0 \left( \frac{S^{(k)} f^{(k)} \mu \sqrt{X^2 - 1}}{f^{(k)} \mu} \right),$$ (26)
where
\[ S^{(\pm)} = 1 \pm \sqrt{1 - g_{\nu}}, \]
\[ f^{(\pm)} = a \pm \sqrt{a^2 - 1}, \]
\[ a = \frac{1}{2} \left( \sqrt{\frac{\alpha_R}{\alpha_L}} + \sqrt{\frac{\alpha_L}{\alpha_R}} \right). \]  
(27)

The integral in Eq. (26) can be calculated, yielding the distribution function for the conductance peaks
\[ R_X (g_{\nu}) = \frac{X \Theta (1 - g_{\nu})}{2\pi \sqrt{1 - g_{\nu}}} \]
\[ \times \sum_{i=1}^{2} \frac{E (k_i)}{\mathcal{M}_i \sqrt{\mathcal{M}^2_i + g_{\nu} (X^2 - 1)}}. \]  
(28)

where the following notation has been used:
\[ \mathcal{M}_i = a + (-1)^{i+1} \sqrt{1 - g_{\nu}} \sqrt{a^2 - 1}, \]
\[ k_i = \sqrt{X^2 - 1} \frac{\sqrt{g_{\nu}}}{\sqrt{\mathcal{M}^2_i + g_{\nu} (X^2 - 1)}}, \]  
(29)
and \( E (k) \) stands for the elliptic integral. Equation (28) is valid for arbitrary magnetic field flux, and it interpolates between the two distributions corresponding to the quantum dot with completely broken time-reversal symmetry
\[ R_{X=1} (g_{\nu}) = \frac{\Theta (1 - g_{\nu})}{2\sqrt{1 - g_{\nu}}} \left[ 1 + (2 - g_{\nu}) \left( a^2 - 1 \right) \right]^2 \]  
(30)

and with conserved time-reversal symmetry:
\[ R_{X \to \infty} (g_{\nu}) = \frac{\Theta (1 - g_{\nu})}{\pi \sqrt{g_{\nu} \sqrt{1 - g_{\nu}}}} \frac{a}{1 + g_{\nu} (a^2 - 1)}. \]  
(31)

Comparing Eqs. (28) and (31), one can see that the influence of the magnetic field is most drastic in the region of small heights of the conductance peaks, \( g_{\nu} \approx 0 \). [Fig. 3]. Equations (28) yields
\[ R_X (0) = \frac{X}{2} (2a^2 - 1). \]  
(32)

Thus, we conclude that the probability density of zero-height conductance peaks decreases from infinity to the value \( (2a^2 - 1) \sqrt{\pi}/3X_m \) at arbitrary small magnetic fluxes penetrating into the sample.

**B. Hauser-Feshbach regime**

When \( \alpha_{R(L)} \Delta \ll T \ll \Delta \), the distribution of the conductance peaks is given by
\[ F_X (g_{\nu}) = \int_0^\infty d\gamma_{\nu R} \int_0^\infty d\gamma_{\nu L} P_X (\gamma_{\nu R}) P_X (\gamma_{\nu L}) \]
\[ \times \delta \left( g_{\nu} - \frac{\pi}{2T} \frac{\gamma_{\nu R} \gamma_{\nu L}}{\gamma_{\nu R} + \gamma_{\nu L}} \right). \]  
(33)

Carrying out the double integration and using Eq. (24), we obtain the following formula for the function \( T_X (\xi) = 2g_{\nu} T \tau_H / \pi \alpha_{L} \alpha_{R} \):
\[ T_X (\xi) = (\frac{X}{2})^2 \xi \exp (-aX^2 \xi) \int_0^\infty d\mu \left( 1 + \frac{1}{\mu} \right)^2 \exp \left[ -\frac{X^2}{2} \xi \left( \mu f(+) + \frac{1}{\mu} f(+) \right) \right] \]
\[ \times I_0 \left( \frac{X \sqrt{X^2 - 1}}{2} (1 + \mu) f(+) \right) I_0 \left( \frac{X \sqrt{X^2 - 1}}{2} (1 + \frac{1}{\mu}) f(+) \right). \]  
(34)

This equation interpolates between two limiting distributions
\[ T_{X=1} (\xi) = \xi \exp (-a \xi) \left[ K_0 (\xi) + a K_1 (\xi) \right] \]  
(35)
and
\[ T_{X \to \infty} (\xi) = \sqrt{\frac{1 + a}{2}} \xi \exp \left( -\frac{a + 1}{2} \xi \right) \]  
(36)
for unitary and orthogonal symmetries, respectively.

The distribution function Eq. (34) can be represented as a series
\[ T_X (\xi) = (\frac{X}{2})^2 \xi \exp (-aX^2 \xi) \]
\[ \times \sum_{n=0}^\infty \sum_{m=0}^\infty \left( \frac{1}{n! m!} \right)^2 \left( \frac{X \sqrt{X^2 - 1}}{4} \xi \right)^{2(n+m)} \left( f(+) \right)^{2(n-m)} \]
\[ \times \sum_{s=1}^{s+1+2m} \sum_{s=-1-2n}^{2(n+m+1)} \left( \frac{2(n+m+1)}{2n+s+1} \right) \left( f(+) \right)^s K_s (X^2 \xi) \]  
(37)
that is convenient for the calculation of the distribution of conductance peaks in the case of weak deviations from unitary symmetry. As is the case of very low temperatures, the distribution function of the heights of the conductance peaks given by Eq. (38) is most affected by magnetic field in the region $g_\nu \approx 0$:

$$T_X (0) = aX,$$  \hspace{1cm} (38)

and therefore, the probability density of the conductance peaks of zero height immediately drops from infinity to the value $T_X (0) \approx 2a\sqrt{\pi}/3X_m$ when an arbitrary small magnetic field ($X_m \ll 1$) is applied to a system.

\section*{IV. CONCLUSIONS}

We have introduced a one-parameter random matrix model for describing the eigenfunction statistics of chaotic electrons in a weakly disordered quantum dot in the crossover regime between orthogonal and unitary symmetry classes. Our treatment applies equally to the statistics of local amplitudes and local phases of electron eigenfunctions inside a dot in the presence of an arbitrary magnetic field. The transition parameter $X$ entering our model is related to the microscopic parameters of the real physical problem, and therefore the distributions calculated within the proposed random-matrix formalism can be used for the interpretation of experiments.

This random matrix model has also been applied to describe the distribution function of the heights of conductance peaks for a quantum dot weakly coupled to external pointlike leads in the regime of Coulomb blockade in the case of crossover between orthogonal and unitary symmetry. We have shown that the magnetic field exerts a very significant influence on the distribution of the heights of the conductance peaks in the region of small heights. The effect of the magnetic field consists of reducing the probability density of zero-height conductance peaks from infinity to a finite value for arbitrary small magnetic flux.

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Fig. 1. Distribution function $\varphi (\tau)$ for different parameters $X$. Inset: the same distribution function calculated from Eq. (5) (lines) and from Eq. (13) in Ref. [2] (figures) for parameters $X$ and $X_m$ connected by Eq. (18). Solid line: $X = 2.56$, triangles $X_m = 0.5$; dashed line: $X = 1.52$, squares: $X_m = 1$; dot-dashed line: $X = 1.14$, circles: $X_m = 2$.

Fig. 2. Phase distribution $q(s)$. Parameters: $X = 1.5$ (curve 1), $X = 5$ (curve 2), $X = 10$ (curve 3).

Fig. 3. Distribution of the heights of the resonance conductance peaks in the Breit-Wigner regime, $a = 1$. Parameters: $X = 1.2$ (curve 1), $X = 2$ (curve 2), $X = 4$ (curve 3).

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