Electron transport in quantum dots (QD) is affected by electron-electron interactions, the profound example being Coulomb blockade (CB) \(^1\) in QDs connected to leads by contacts of small conductance \(G \ll e^2/\pi \hbar \equiv G_q\). Due to charge quantization, at temperatures smaller than the charging energy \(E_c \gg \Delta\) (\(\Delta\) being the one particle level spacing), associated with the addition of one electron to the QD, transport through the dot is diminished except the mesoscopic fluctuations of the conductance are de-tering is expected to be strong, similar to antilocalization in the bulk systems. For the valleys, we consider the elastic cotunneling contribution to the conductance and calculate its moments at the crossover between ensembles of various symmetries.

In this Letter, we find the statistics of the peak heights for strong spin-orbit scattering in the presence and absence of time reversal symmetry. For the valleys we calculate the magnetic correlation for conductance, even in the presence and absence of time reversal symmetry. We find that the effect of spin-orbit scattering on the statistics of the conductance of a quantum dot with identical symmetries.

The results were extended to the crossover between ensemble s of various symmetries. In the case of quantum dots based on a 2D electron gas, the statistics of one particle energy levels and eigenfunctions are described by universal ensembles sensitive only to the underlying symmetries. The statistical distribution of the peak heights has been calculated based on RMT for the Gaussian orthogonal (GOE) and unitary (GUE) ensembles \(^3\). The results were extended to the crossover between those ensembles \(^4\), using the statistics of wavefunctions derived in Ref. \(^5\). The average peak height was observed to increase with the magnetic field, reminiscent of weak localization in bulk systems \(^6\). The statistics of conductance in the valleys was studied theoretically \(^7\) and experimentally \(^8\) for the same crossover. However, there is a gap in the theoretical literature regarding the effect of spin-orbit scattering on these statistics. The effect of breaking of time reversal symmetry on the CB peak heights in quantum dots with strong spin-orbit scattering is expected to be strong, similar to antilocalization in bulk systems \(^9\) or open dots \(^10\).

In this Letter, we find the statistics of the peak heights for strong spin-orbit scattering in the presence and absence of time reversal symmetry. For the valleys we calculate the magnetic correlation for conductance, even in the crossover between ensembles of various symmetries.

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Antilocalization in Coulomb Blockade

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We study the effect of spin-orbit scattering on the statistics of the conductance of a quantum dot for Coulomb blockade peaks and valleys. We find the distribution function of the peak heights exhibit at discrete values of the gate voltage corresponding to the QD, transport through the dot is diminished except the mesoscopic fluctuations of the conductance are de-

The QD attached to two leads, 1 and 2, is described by the Hamiltonian

\[
H = H_L + H_D + H_{LD},
\]

where

\[
H_L = v_F \sum_{\alpha=1,2} \sum_{\sigma} \int \frac{dk}{2\pi} k c_{\alpha,\sigma}^\dagger(k) c_{\alpha,\sigma}(k),
\]

corresponds to the leads, with \(v_F\) being the Fermi velocity, and \(\sigma = \uparrow, \downarrow\) labels the spin. We consider only single channel contacts. The dot Hamiltonian is

\[
H_D = \sum_{m,n=1}^{M} \sum_{\sigma_1,\sigma_2} c_{m,\sigma_1}^\dagger H_{mn}^{\sigma_1 \sigma_2} c_{n,\sigma_2} + E_c (\hat{n} - N)^2,
\]

where the first term describes the non-interacting dynamics of the closed dot, and the second term, with \(\hat{n} = \sum_m c_{m,\uparrow}^\dagger c_{m,\uparrow} + c_{m,\downarrow}^\dagger c_{m,\downarrow}\), corresponds to charging energy. Dimensionless parameter \(N\) is a linear function of the gate voltage. The matrix \(H\) belongs to an RMT ensemble. In the case of quantum dots based on a 2D electron gas, the spin-orbit interaction is atypical and is characterized by two parameters \(c_{\sigma}^{\parallel}\) and \(c_{\sigma}^{\perp}\). In the presence of a magnetic field with both perpendicular and parallel components to the gas plane (let the third axis be normal to the plane), \(H\) can be written as

\[
H = \frac{\Delta}{2\pi} \left[ \mathcal{H}_0 + i \mathcal{X}_e(x_1 + a_{\perp} \sigma^3) + i a (A_1 \sigma^1 + A_2 \sigma^2) + b_{\perp} B_h \sigma^3 \right] - \epsilon \vec{l} \cdot \vec{\sigma}/2,
\]

where \(\mathcal{H}_0\) and \(B_h\) are real symmetric \(M \times M\) matrices, with \(\langle \text{Tr} \mathcal{H}_0^2 \rangle = M^3\), \(\langle \text{Tr} B_h^2 \rangle = M^2\), and \(\mathcal{X}_e \) and \(A_i\) are real antisymmetric matrices with \(\langle \text{Tr} \mathcal{X}_e^T \rangle = M^2\) and \(\langle \text{Tr} A_i A_i^T \rangle = M^2 \delta_{ij}\) (the limit \(M \to \infty\) is taken eventually). Here \(a_{\perp} = \pi e c_{\sigma}^{\perp}/\Delta\), \(a_2 = \pi e c_{\sigma}^{\parallel}/\Delta\), \(\epsilon\) is the Zeeman splitting energy due to the parallel magnetic field in the direction given by the unit vector \(\vec{l} = (l_1, l_2, 0)\), and \(b_{\perp} = \pi e c_{\sigma}^{\perp}/\Delta\) where \(c_{\sigma}^{\perp}\) describes the combined effect of Zeeman splitting and spin-orbit scattering. The orbital effect of the magnetic field is characterized by the energy

\[
E_o = \frac{\Delta}{2}\left( x_1^2 + x_2^2 \right) + \frac{1}{2} (l_1^2 + l_2^2).
\]
\[ \epsilon_B = x^2 \Delta / \pi = \kappa E_{\text{T}L}(\Phi / \Phi_0)^2 \], where \( \kappa \) is a coefficient dependent on the shape of QD, \( \Phi \) is the flux of the magnetic field through the dot, and \( \Phi_0 \) is the flux quantum.

The tunneling Hamiltonian \( H_{LD} \) couples the states of the leads \( a = 1, 2 \) to orbital states in the QD:

\[ H_{LD} = \sqrt{\frac{M \Delta}{\pi \nu}} \sum_{a,n,k} \delta_{\alpha n} t_n c^\dagger_{\alpha}(k)c_n + \text{h.c.}, \] (4)

Here \( \nu = 1/(2\pi v_F) \) is the leads density of states per spin.

First we consider the statistics of the peaks. For the temperature range \( T \ll \Delta \ll E_c \) only the last occupied, possibly degenerate level participates in the electron transport close to the peaks. Using the Golden Rule, the escape rates of the level \( \alpha \) into the first and second leads are \( \Gamma_{1,2}^\alpha = \frac{\Delta}{2\pi i} \phi_{1,2} \), where \( \phi_i = 4|t_i|^2 y_i \) is the dimensionless conductance of the \( i \)-th contact. Here \( y_i = \frac{g_i}{(g_i)} = M \psi^\dagger_{\alpha}(i) \psi_{\alpha}(i), \quad (i = 1, 2) \), (5)
is a fluctuating quantity, and \( \psi_{\alpha}(i) \) are spinors with components \( (\downarrow, \downarrow) \). For the peaks we will confine ourselves to strong spin-orbit scattering \( \epsilon_\|^{\alpha}, \epsilon_\|^{\omega} \gg \Delta \). In this case, and for zero magnetic field, the random matrix \( H \) belongs to a Gaussian symplectic ensemble (GSE): it is a Hermitian matrix with quaternionic entries. In the notation \( 1^{\alpha} \), this corresponds to the symmetry class \( (\beta = 4, \Sigma = 1, s = 2) \). The number \( s = 2 \) signifies Kramer’s degeneracy of energy levels. In the presence of a strong enough magnetic field such that either \( \epsilon_\|^{\alpha} \gg \Delta \) or \( \epsilon_\|^{\omega} \gg \Delta \), the Hamiltonian \( H \) belongs to a GUE corresponding to \( (\beta = 2, \Sigma = 2, s = 1) \), and the symmetry group of the RMT ensemble is \( U(2M) \). In this case, the levels are not Kramer’s degenerate. In both cases, the random variables \( y_i \) in Eq. (5) are statistically independent, and their distribution is given by

\[ p(y) = 4ye^{-2y}. \] (6)

Assuming \( \Delta > T \gg \Gamma^\alpha = (g_1 + g_2)\Delta/2\pi \hbar \), one can use the rate equations \( 2, 13 \). In the GSE case, we will label the Kramer’s degenerate levels by \( \psi_{\alpha,s} \), with the \( s = \pm \) states related by the operation of time reversal: \( \psi_{\alpha,-} = \text{i} \sigma^y \psi_{\alpha,+} \). Due to this relation, \( (\psi^\dagger_{\alpha,+} \psi_{\alpha,+}) = (\psi^\dagger_{\alpha,-} \psi_{\alpha,-}) \), and the escape rates are therefore equal for the \( s = \pm \) states (see Eq. (10)). In the absence of interactions, the level \( \alpha \) can be in four states: empty, doubly occupied, or singly occupied with \( s = \pm \). The picture changes when we take the interaction into account: due to the large charging energy, only three states with the number of electrons on the dot differing by 1 can be in resonance. For the peaks corresponding to \( N^* = 2j + 1/2, (j \in \mathbb{Z}) \), the doubly occupied state has an extra energy \( E_c \), and does not participate in transport (the case \( N^* = 2j - 1/2 \) gives the same result for the peak height). In the stationary state the rate equations yield

\[ G(N) = \left. \frac{I}{V} \right|_{V \to 0} = -G_\delta \frac{g_1 g_2}{T} \frac{\partial f_F/\partial x}{1 + f_F(x)}, \] (7)

where \( f_F(x) = 1/(1 + \exp x) \), and \( x = 2E_c(N - N^*)/T \). At the maximum which slightly deviates from \( x = 0 \)

\[ G_{\text{peak}}^{\beta=4,s=2}/G_\delta = \frac{3 - 2^{1/2}}{T} \frac{g_1 g_2}{g_1 + g_2} \] (8)

The only difference in the \( (\beta = 2, \Sigma = 2, s = 1) \) case is that now there is no Kramers’ degeneracy and the resonant level can only be in two states: empty or occupied. Similarly to Eq. (7), we find

\[ G(N) = -G_\delta \frac{\partial f_F/\partial x}{2T} \frac{g_1 g_2}{g_1 + g_2} \] (9)

The maximum occurs at \( x = 0 \), and we find

\[ G_{\text{peak}}^{\beta=2,s=1}/G_\delta = \frac{3 - 2^{3/2}}{8T} \frac{g_1 g_2}{g_1 + g_2} \] (10)

Since the \( y_i \)'s have the same distribution in both ensembles we see that \( G_{\text{peak}} \) has the same distribution in both cases except for a scaling. Using the probability distribution Eq. (6), we obtain \( G_{\text{peak}} \) in terms of the average conductances \( (g_1,2) \) and a single random variable \( \alpha \),

\[ G_{\text{peak}}^{\beta,s}/G_\delta = \frac{2(g_1/g_2)}{(g_1/2 + (g_2/1/2)^2) \chi_{\beta,s},} \] (11)

where \( \chi_{\beta=4,s=2} = 3 - 2^{3/2} \), and \( \chi_{\beta=2,s=1} = 1/8 \). The probability distribution for \( \alpha \) is given by (see Fig. 1)

\[ W(\alpha) = 16\alpha^3(1 - \alpha) e^{-2\alpha(1 + \alpha)} \{ \frac{1 + b^2}{2} K_0[2\alpha(1 - a)] + \left[ b + \frac{b^2 - 1/2}{2\alpha(1 - a)} \right] K_1[2\alpha(1 - a)] \}. \] (12)

Here \( K_0(x) \) and \( K_1(x) \) are MacDonald functions and

\[ a = \left( \frac{g_1^{1/2} - (g_2^{1/2})^2}{g_1^{1/2} + (g_2^{1/2})^2} \right)^2, \quad b = \frac{1 + a}{1 - a} \] (13)

For the average we obtain

\[ \langle \alpha \rangle = 8(1 - a)^4 \left\{ \frac{2(1 + b^2)}{105} F(5,1/2,11/2;a) + \frac{1 + a}{21} F(6,3/2,11/2;a) + \frac{b^2 - 1/2}{35} F(5,3/2,9/2;a) \right\}, \] (14)

with \( F(\ldots) \) being the hypergeometric functions. Equation (13) is well approximated by \( \langle \alpha \rangle \simeq (8 - 3a)/10 \).

Equation (11) is our main result for the statistical distribution of the peak heights. We see that the application of the magnetic field causes the average conductance to drop by a factor \( 8(3 - 2^{3/2}) \simeq 1.37 \), similar to antilocalization for bulk systems. Surprisingly, the shape of the distribution remains the same. This is because at strong magnetic field the RMT ensemble crosses over to a unitary ensemble with two channel leads corresponding to
spin projections. The statistics of the eigenstates in the latter are the same as that of the symplectic ensemble, and hence the same distribution of peak heights. Though the drop in the average peak heights is a manifestation of the lifting of the Kramer’s degeneracy, the numerical factor for this drop is non-trivial. Had we ignored the charging energy, we would instead obtain a reduction by a factor of 2. The different numerical factor originates in the exclusion of the doubly occupied degenerate level from transport by the electron-electron interaction. We emphasize, however, that the distribution function changes for intermediate values of the perpendicular field.

Next we turn to the mesoscopic fluctuations of the valleys, $|N - N^*| > T/E_c$. Then, Eqs. 1−4 suggest an exponentially small conductance. However, this is incorrect, as the rate equations are based on processes of the first order in $H_{LD}$. For such processes a charge is transferred between a lead and QD, and the conductance in the valleys is small because of the charging gap in the final state. This is not the case for higher order processes that allow for both the initial and the final excitations to be in the leads, so there is no gap for the final state. Such processes are referred to as co-tunneling [2, 13], and are suppressed only algebraically, $G(N) \propto 1/(2E_c|N - N^*|)$.

We consider elastic co-tunneling which is the dominant process in the temperature range $T \ll (ΔE_c)^{1/2}$. In these processes a charge is transferred from lead to lead via a virtual transition to an excited state in the dot. The conductance calculated using the Golden rule is

$$G = 2\pi^2G_0\nu^2 \sum_{\sigma_1,\sigma_2=\pm 1/2} |A^{\sigma_1\sigma_2}_e + A^{\sigma_1\sigma_2}_h|^2,$$  \hspace{1cm} (15)

where the amplitudes $A^{\sigma_1\sigma_2}_e$ ($A^{\sigma_1\sigma_2}_h$) correspond to processes in which an electron(hole) in spin state $\sigma_1$ in the first lead is transferred to the second lead in spin state $\sigma_2$. The second order perturbation theory in $H_{LD}$ gives

$$\frac{A^{\sigma_1\sigma_2}_e(h)}{DE_{e(h)}} = \frac{\langle g_1 \rangle \langle g_2 \rangle}{4\pi^2 \nu} \sum_{\sigma_1,\sigma_2} \psi^{\sigma_2}_e(2)\psi^{\sigma_1}_e(1) \theta(\pm \varepsilon_\alpha),$$  \hspace{1cm} (16)

where $\alpha$ is summed over energy levels of the QD. Here

$$E_c = 2E_c(N^* - N), \quad E_h = 2E_c(N - N^* + 1),$$  \hspace{1cm} (17)

($N^* - 1 < N < N^*$), are the electrostatic part of the energy of the virtual state for the electron-like and hole-like processes respectively. The eigenenergies $\varepsilon_\alpha$ are measured from the last occupied level so that the step functions $\theta(\pm \varepsilon_\alpha)$ select empty(filled) states for electron(hole) like processes. Equation (15) can be expressed in terms of the Green functions (GF) for the closed dot:

$$G = G_0 \frac{\langle g_1 \rangle \langle g_2 \rangle}{4\pi^2} \frac{1}{\nu} \sum_{\sigma_1,\sigma_2=\pm 1/2} |F^{\sigma_1\sigma_2}_e + F^{\sigma_1\sigma_2}_h|^2,$$  \hspace{1cm} (18)

$$\frac{F^{\sigma_1\sigma_2}_e(h)}{M\Delta} = \int_{-\infty}^{\infty} d\epsilon \frac{\Delta^{\sigma_1\sigma_2}_{12}(\epsilon) - \Delta^{\sigma_1\sigma_2}_{21}(\epsilon)}{\epsilon \pm E_{e(h)}} \theta(\pm \epsilon).$$  \hspace{1cm} (19)

Using this relation we can express all the moments of the conductance in terms of the average of a product of GFs. Furthermore the condition $\Delta \ll E_c, E_h$, allows us to use the diagrammatic technique in terms of diffusions and cooperons to calculate the latter. In this approximation the GFs become Gaussian variables with average $\langle \Delta^{\sigma_1\sigma_2}_{12}(\epsilon) \rangle \propto \delta_{nm}$, and variances given in terms of the Diffusion and Cooperon matrices [2, 7].

$$\text{Tr} \langle \sigma^\mu \Delta^{\sigma_1\sigma_2}_{12}(\epsilon; 1_1) \sigma^\nu \Delta^{\sigma_1\sigma_2}_{12}(\epsilon; 1_2) \rangle = \frac{4\pi}{M\Delta^2} \epsilon^{\mu\nu} D^{\delta\gamma}(\epsilon; B_1, B_2),$$  \hspace{1cm} (20a)

$$\text{Tr} \langle \sigma^\mu \Delta^{\sigma_1\sigma_2}_{12}(\epsilon; 1_1) \sigma^\nu \Delta^{\sigma_1\sigma_2}_{12}(\epsilon; 1_2) \rangle = \frac{4\pi}{M\Delta^2} \epsilon^{\mu\nu} C^{\delta\gamma}(\epsilon; B_1, B_2),$$  \hspace{1cm} (20b)

where $\omega = \epsilon_1 - \epsilon_2$, $\sigma^{\mu\nu} = \hat{1}$, and $\sigma^i$, $i = 1, 2, 3$ are Pauli matrices. Here we defined $\Delta^{\sigma_1\sigma_2}_{12}(\epsilon; B_1, B_2)$ for the Hamiltonian Eq. 13, the inverse of the diffusion and cooperon matrices are given by

$$D^{-1}(\omega; 1_1, B_2) = -i(\hat{\epsilon} + i\hat{\nu})^T \cdot \hat{S},$$

$$+ \left( \sqrt{\hat{\nu}^2 + \sqrt{\hat{\nu}^2} \hat{S}_3} \right)^2 + \epsilon_0^\sigma \left( \hat{S}_1^2 + \hat{S}_2^2 \right) + \frac{\hat{S}_3^2 \hat{S}_1^2}{\hat{S}_3^2},$$

$$C^{-1}(\omega; 1_1, B_2) = -i(\hat{\epsilon} + i\hat{\nu})^T \cdot \hat{S} + \left( \sqrt{\hat{\nu}^2 + \sqrt{\hat{\nu}^2} \hat{S}_3} \right)^2 + \epsilon_0^\sigma \left( \hat{S}_1^2 + \hat{S}_2^2 \right) + \left( \hat{1} - \hat{S}_3 \right)^2 \epsilon_0^\sigma,$$  \hspace{1cm} (21)

where $\hat{S}_1^j = -ie^{ij}k$, and $\hat{\nu}_{\mu\nu} = \nu_{\mu\nu} = \nu_{\mu\nu}$, with $\hat{l}$ the unit vector in the direction of parallel field, and

$$\epsilon_{B_1} = \kappa E_{Th} \left( \frac{B_1 - B_2}{2\Phi_0/A} \right)^2, \quad \epsilon_{B_1} = \kappa E_{Th} \left( \frac{B_1 + B_2}{2\Phi_0/A} \right)^2,$$  \hspace{1cm} (22)

where $A$ is the dot’s area.

According to Eq. 13, the amplitudes $F_e, F_h$ are linear in $\Delta^R$ and $\Delta^A$ and are also Gaussian variables with zero average, so that we can calculate the moments of the conductance, Eq. 13, using Wick’s theorem. From Eqs. 18 and 20 for $\nu = 1\mu$, we obtain the result

$$\langle G \rangle = G_0 \frac{\langle g_1 \rangle \langle g_2 \rangle}{4\pi^2} \left[ \frac{\Delta}{E_c} + \frac{\Delta}{E_h} \right],$$  \hspace{1cm} (22)
independent of all the crossover parameters. Using Eqs. (20)-(21) together with (18), we calculate the correlation of the conductances at different values of the perpendicular magnetic field. In the vicinity of the peaks corresponding to the condition $E_h \ll E_c$ or $E_c \ll E_h$, for $C_{12} = \langle \delta G(B_1) \delta G(B_2) \rangle / \langle G \rangle^2$ we obtain

$$C_{12} = \sum_{i=C,D} \sum_{\alpha, \beta=1, \ldots, 4} \Lambda \left( \frac{\lambda^i}{E} \right) \Lambda \left( \frac{\lambda^j}{E} \right) \frac{\delta_{\alpha\beta} + A^j_{\alpha} A^i_{\beta}}{8},$$

(23)

where $\Lambda(z) = \tilde{\Lambda}(z) - \tilde{\Lambda}(-z)$ and

$$\tilde{\Lambda}(z) = \frac{1}{\pi z} \left[ \ln |z| \ln (1 - iz) + \frac{1}{4} Li_2(-z^2) \right],$$

(24)

for complex $z$, with $Li_2(x)$ being the dilogarithm function. The asymptotic behavior is $\Lambda(z) = 1 + z \ln z / \pi$, for $|z| \ll 1$, and $\Lambda(z) = (\pi z)^{-1} \ln^2 z$, for Re $z \gg 1$. $E = \min(E_c, E_h)$, and $\lambda^i_D$ and $\lambda^j_C$ are the four eigenvalues of $D^{-1}(0; B_1, B_2)$ and $C^{-1}(0; B_1, B_2)$ matrices respectively. The $4 \times 4$ matrices $A_{D,C}$ are given by

$$A_D = U_D^{-1} L_D U_D, \quad A_C = U_C^{-1} L_C U_C,$$

(25)

where $U_D$ and $U_C$ are the matrices whose columns are the eigenvectors of $D^{-1}(0; B_1, B_2)$ and $C^{-1}(0; B_1, B_2)$ respectively, $L_D^{-1} = 2 \delta_{\mu\nu}^3 \delta^3 \nu$ and $L_C^{-1} = 2 \delta_{\mu\nu} - 2 \delta_{\nu\delta} \delta_\nu$. For zero parallel field and small perpendicular fields Zeeman splitting can be ignored and the eigenvalues of inverse cooperator and diffusion correspond to the usual singlet and triplet states, and we have (Fig. 2)

$$C_{12} = \sum_{i=D,C} \sum_{j=0,1} \sum_{m=-j}^j \Lambda \left( \frac{\lambda^i_{jm}}{E} \right)^2 \chi^{D(C)}_{jm} = \left( \sqrt{\epsilon_{B}} + m \sqrt{\epsilon_{\parallel}} \right)^2 \epsilon_{\perp} (j(j+1) - m^2).$$

Furthermore for strong spin-orbit scattering, $\epsilon_{\parallel}^0 \gg E$, we can calculate all the moments of the conductance $G$, and obtain the distribution functions for $\gamma = G / \langle G \rangle$:

$$P(\gamma) = \theta(\gamma) \frac{e^{-\gamma^g}}{\lambda^2} \left[ \frac{8 \gamma \cosh 4 \gamma - 1 - \lambda^2}{\lambda^2} \right],$$

(26)

where $\lambda = \Lambda (\Phi^2 / \Phi_c^2)$, $\gamma = \gamma / (1 - \lambda^2)$, $\Phi$ is the magnetic flux, and $\Phi_c = \Phi_0 / \sqrt{E / k E_{TH}}$.

In conclusion, we studied the statistics of the conductance for CB peaks and valleys in the presence of spin-orbit scattering. We calculated the distribution function of the peak heights for strong spin-orbit scattering in the presence and absence of time-reversal symmetry. We found that the average peak height is reduced in the latter case. For the valleys we calculated the average conductance and the correlation function of conductance for a 2DEG QD, as a function of perpendicular magnetic field.

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