Weak Localization in the Conductance Peaks of Coulomb Blockade Quantum Dots

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Abstract

We derive closed expressions for the universal weak localization peak of the average conductance peak heights in Coulomb blockade quantum dots in the crossover from orthogonal to unitary symmetry. The scale for the crossover is independent of the number of channels in each lead, in contrast with the case of open dots. The functional form of the weak localization peak is independent of temperature. We also derive analytically the variance of the conductance peak heights as a function of the crossover parameter.

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Quantum dots are formed in the interface between layers of semiconductors by the electrostatic confinement of electrons in a two-dimensional electron gas. The transport properties (conductance) of such quantum dots can be measured by connecting them to external leads. In ballistic dots, the transport is dominated by scattering from the boundaries. Since dots with several hundred electrons are usually characterized by irregular shape, the scattering from the boundaries results in chaotic motion of the electrons. The conductance in such dots displays universal statistical fluctuations that are a signature of quantum chaos and that can be described by random matrix theory (RMT) [1]. Earlier studies of quantum dots focused mostly on open dots, where the movement of the electron into the dot is allowed classically [2]. As the point contacts are pinched off, the dot becomes closed, i.e., effective tunnel barriers are formed in the interface between the dot and the leads. In such dots the conductance is dominated by resonant tunneling through electron resonances that are narrow compared with their average spacing. This leads to a series of narrow peaks in the conductance as a function of gate voltage, known as Coulomb blockade peaks. A statistical theory based on RMT for the fluctuations of the conductance peak heights was developed in Ref. [3] and was recently confirmed experimentally [4,5].

The statistics of the peak heights are universal but sensitive to the underlying symmetries. For conserved time-reversal symmetry (no magnetic field) the appropriate ensemble is the Gaussian orthogonal ensemble (GOE), while for broken time-reversal symmetry the statistics are described by the Gaussian unitary ensemble (GUE). The crossover from preserved to fully broken time-reversal symmetry is also universal. The distribution of the wavefunction intensity at a fixed spatial point was calculated by the supersymmetry method in Ref. [6]. However the fluctuations of the conductance in the crossover regime depend on the wavefunction fluctuations at several spatial points. This statistics and the corresponding distributions of the peak heights were derived as a function of a dimensionless crossover parameter in Ref. [7]. This crossover parameter is linear in the magnetic field that induces the symmetry breaking.

One of the interesting phenomena associated with the onset of broken time-reversal sym-
metry in open dots is the suppression of weak localization. Weak localization in the absence of magnetic field originates from the constructive interference of pairs of time-reversed classical trajectories, which enhances the return probability. This results in a smaller average conductance. When a magnetic field is applied, the constructive interference is destroyed and the weak localization is suppressed. In open dots, the semiclassical theory yields a Lorentzian shape for the weak localization peak. An exact expression for open dots was derived from RMT using the supersymmetry method, and approaches a Lorentzian for a large number of open channels.

Weak localization also occurs in almost closed dots, where the average conductance peak height was found to be smaller in the GOE case than in the GUE case. On the other hand, no weak localization effect is found in the average of the conductance minima. In this paper we derive closed expressions for the weak localization peak of the conductance maxima in closed dots as a function of the crossover parameter and for any number of equivalent channels in the leads. We also study the temperature dependence of the weak localization peak and find that, when measured in units of the average GUE conductance, the peak is temperature-independent (assuming no phase breaking). We find that the width of the peak is only weakly dependent on the number of channels in the leads, in contrast to the case of open dots, where the width of the peak is proportional to the square root of the number of open channels. We also derive closed expressions for the variance of the conductance peak as a function of the crossover parameter.

For almost closed dots, the resonances in the dots are isolated and their average width is small compared with the mean level spacing \( \Delta \). At temperatures that are low compared with the mean level spacing \( (kT \ll \Delta) \), only one resonance contributes to a given Coulomb blockade peak, and the conductance peak amplitude \( G \) is

\[
G = \frac{e^2}{\hbar} \frac{\pi \Gamma}{4kT} g
\]

where

\[
g = \frac{1}{\bar{\Gamma}} \frac{\Gamma_l \Gamma_r}{\Gamma_l + \Gamma_r}.
\]

Here \( \Gamma_l^{(r)} \) is the width of a resonance to decay into the left (right) lead and \( \bar{\Gamma} \) is the average width per channel. Note that the quantity \( g \) is dimensionless and temperature-independent.
In general we assume that the left (right) lead has $\Lambda^l$ open channels such that $\Gamma^l = \sum_c |\gamma^l_c|^2$, where $\gamma^l_c$ is the partial amplitude to decay into channel $c$ on the left (right). In $R$-matrix theory, the partial width amplitude $\gamma_c$ can be expressed as a scalar product (defined over the dot-lead interface) of the resonance wavefunction $\psi$ and the channel wavefunction $\phi_c$ through $\gamma_c = \langle \phi_c | \psi \rangle \propto \int dS \phi_c^* (r) \psi (r)$. We assume for simplicity that the $\Lambda$ channels are uncorrelated and equivalent, i.e. $\bar{\gamma}_c \bar{\gamma}_{c'} = \bar{\Gamma} \delta_{cc'}$. All average partial widths are assumed to be independent of the magnetic field. We also assume that the dot has symmetric leads, so that the average partial widths are the same in the left and right leads.

For a ballistic dot with chaotic dynamics of the electrons, we model the Hamiltonian $H$ of the dot by an ensemble of random matrices. When we apply a magnetic field to the dot, the ensemble describing the crossover from GOE to GUE is $[13,14]$

$$H = S + i\alpha A, \quad (2)$$

where $S$ and $A$ are uncorrelated symmetric and antisymmetric real matrices, respectively, chosen from Gaussian ensembles of the same variance, and $\alpha$ is a real parameter. The transition parameter describing the crossover is given by the ratio of the root-mean-square (rms) of a typical symmetry-breaking matrix element to the mean-level spacing $[13] \lambda = (H^2_{\text{break}})^{1/2}/\Delta = \alpha \sqrt{N}/\pi$. Alternatively $2\pi \lambda = \sqrt{\tau_H/\tau_{\text{mix}}}$, where $\tau_H = h/\Delta$ is the Heisenberg time and $\tau_{\text{mix}}$ is the mixing time defined by the spreading width $h/\tau_{\text{mix}} = 2\pi H^2_{\text{break}}/\Delta$ due to the interaction that breaks time-reversal symmetry $[8]$. The spectral properties of the ensemble (2) make the complete crossover for $\lambda \sim 1$. The average dimensionless conductance peak height $\bar{g}(\lambda)$ can be calculated by averaging (1) over the ensemble (2).

The statistics of the partial width amplitudes $\gamma$ were derived in Ref. $[7]$. The components $\psi_\mu$ of the resonance wavefunction (in a fixed basis) are decomposed into their real and imaginary parts $\psi_\mu = \psi_{\mu R} + i \psi_{\mu I}$ in a principal frame in the complex plane (determined by $\sum_{\mu=1}^N \psi_{\mu R} \psi_{\mu I} = 0$), and a parameter $t$ is defined by $t^2 \equiv \sum_{\mu} \psi_{\mu I}^2 / \sum_{\mu} \psi_{\mu R}^2$. At a fixed value of $t$, the real and imaginary parts of a finite number $\Lambda$ of components ($\Lambda \ll N$) are independent Gaussian variables. Except in the GOE and GUE limits (where $t = 0$ and
\[ t = 1, \text{ respectively}, \] \( t \) is not sharp but fluctuates according to a known distribution \( 3 \quad 4 \quad 10 \)
\[
P_{\lambda}(t) = \frac{\pi^2}{t^3} \lambda^2 e^{-\frac{\pi^2}{2} \lambda^2 (t-1/t)^2} \left\{ \phi_1(\lambda) + \left[ \frac{1}{4} \left( t + \frac{1}{t} \right)^2 - \frac{1}{2 \pi^2 \lambda^2} \right] [1 - \phi_1(\lambda)] \right\}, \quad (3)
\]
where \( \phi_1(\lambda) = \int_{0}^{\infty} e^{-2\pi^2 \lambda^2 (1-y^2)} \, dy \). These fluctuations of \( t \) are responsible for the long-range correlations of wavefunction intensities in the crossover from GOE to GUE \( 17 \).

Applying a similar decomposition for the partial width amplitudes \( \gamma_c = \gamma_{cR} + i \gamma_{cI} = \langle \phi_c | \psi_R \rangle + i \langle \phi_c | \psi_I \rangle \) in the principal frame of \( \psi \), it is found that at fixed \( t \), \( \gamma_R \) and \( \gamma_I \) are Gaussian statistically independent variables with \( \overline{\gamma^2_{cR}} = t^2 \gamma^2_{cR} \). The joint distribution of the partial widths is obtained by averaging the conditional distribution \( P(\gamma|t) \) over the \( t \) distribution, \( P_{\lambda}(\gamma) = \int_{0}^{1} P_{\lambda}(t) P(\gamma|t) \, dt \).

Using Eq. (1) and the statistical independence of the partial widths in the left and right leads at fixed \( t \) (i.e., \( P(\Gamma^l, \Gamma^r|t) = P(\Gamma^l|t) P(\Gamma^r|t) \) where \( P(\ldots|t) \) are conditional probabilities), the average conductance peak height can be written as
\[
\bar{g}(\lambda) = \left\langle \int_{0}^{\infty} d\Gamma Q^{(1)}(\Gamma, t)/\Gamma \right\rangle \quad (4a)
\]
\[
Q^{(1)}(\Gamma, t) \equiv \int_{0}^{\Gamma} d\Gamma' \Gamma' P(\Gamma'|t)(\Gamma - \Gamma') P(\Gamma - \Gamma'|t), \quad (4b)
\]
where \( \langle \ldots \rangle \) denotes here and in the following an average over the \( t \) distribution \( P_{\lambda}(t) \) given by Eq. (3). The function \( Q^{(1)}(\Gamma, t) \) is the convolution of \( \Gamma^l P(\Gamma^l) \) with itself, and its Laplace transform is \( Q^{(1)}(s, t) = (\partial P/\partial s)^2 \) where \( P(s, t) \equiv \int_{0}^{\infty} d\Gamma e^{-\Gamma s} P(\Gamma|t) \) is the Laplace transform of the width distribution. \( P(s, t) \) is easily calculated using the Gaussian nature of the partial width amplitudes at fixed \( t \) and \( \Gamma = \sum_e (|\gamma_c|^2 + |\gamma_I|^2) \). Assuming one-channel symmetric leads (i.e., \( \overline{\Gamma^l} = \overline{\Gamma^r} = \overline{\Gamma} \)) and measuring \( \Gamma \) in units of \( \overline{\Gamma} \), we find \( P_1(s, t) = (1 + s^2/x^2 + 2s)^{-1/2} \) where \( x \equiv (t^{-1} + t)/2 \), and thus \( Q^{(1)}_1(s, x) = x^2(s+x^2)^2(s^2 + 2sx + x^2)^{-3} \). The inverse Laplace transform \( Q^{(1)}_1(\Gamma, t) \) is calculated by the residue theorem using the poles (of degree three) of \( Q_1(s, x) \) at \( s = t \) and \( s = 1/t \). After performing the integral in Eq. (4a), we find the average conductance as a function of \( \lambda \) to be
\[
\bar{g}(\lambda) = \frac{1}{4} + \left\langle \left( \frac{t}{1 - t^2} \right)^2 \left( \frac{2t^2}{1 - t^2} \ln t + \frac{1}{2} \right) \right\rangle. \quad (5)
\]
The GOE and GUE limits are obtained for $t \to 0$ and $t \to 1$, respectively, where $\bar{g}_{\text{GOE}} = 1/4$ and $\bar{g}_{\text{GUE}} = 1/3$.

The above method can be generalized for any number $\Lambda$ of equivalent and uncorrelated channels (in each of the symmetric leads). Measuring $\Gamma$ in units of the average resonance width per channel, we find

$$P_\Lambda(s,t) = (1 + s^2/x^2 + 2s)^{-\Lambda/2}$$

and

$$Q_\Lambda^{(1)}(s,t) = \Lambda^2 x^{2\Lambda}(s + x^2)^2(s^2 + 2x^2s + x^2)^{-\Lambda - 2}.$$  

Using (4) we find that $\bar{g}_\Lambda$ changes from $\Lambda^2/2(\Lambda + 1)$ in the GOE limit ($x \to \infty$) to $\Lambda^2/(2\Lambda + 1)$ in the GUE limit ($x \to 1$). In the crossover regime we find that $\bar{g}_\Lambda(\lambda)$ has the form

$$\bar{g}_\Lambda(\lambda) = \frac{\Lambda^2}{2(\Lambda + 1)} + (-)^{\Lambda+1} \Lambda \left( \frac{2\Lambda}{\Lambda - 1} \right) \left\langle \left( \frac{t}{1 - t^2} \right)^{2\Lambda} \left[ \frac{2t^2}{1 - t^4} \ln t - \sum_{n=0}^{\Lambda-1} a_n \left( \frac{1 - t^2}{t} \right)^{2n} \right] \right\rangle. \quad (6)$$

The coefficients $a_n$ in Eq. (6) can be determined from the known GUE limit ($t \to 1$) of $\bar{g}_\Lambda$. For the limit on the r.h.s. of (6) to exist when $t \to 1$, the coefficients $a_n$ for $n = 0, 1, \ldots, \Lambda-1$ must be equal to the corresponding expansion coefficients of the function $2t^2/(1 - t^4) \ln t$ in powers of $[(1 - t^2)/t]^2$, and are therefore independent of $\Lambda$. The value of $\bar{g}_\Lambda$ in the GUE limit also determines the coefficient $a_\Lambda$ to be $a_\Lambda = (-)^{\Lambda+1}(\Lambda!)^2/2((2\Lambda + 1)!)$. Since the coefficients $a_n$ are independent of $\Lambda$, we conclude

$$a_n = (-)^{n+1} \frac{(n!)^2}{2(2n + 1)!}. \quad (7)$$

The function $\delta g_\Lambda(\lambda) \equiv \bar{g}_\Lambda^{\text{GUE}} - \bar{g}_\Lambda(\lambda)$, which is peaked at $\lambda = 0$ (GOE limit) and approaches zero for $\lambda \to \pm \infty$, is known as the weak localization peak. Its analytic form is plotted in Fig. 1 (in units of the average GUE conductance peak) for $\Lambda = 1, 2, 3$ and 5. The size of the peak is $\delta g_\Lambda(0) = \bar{g}_\Lambda^{\text{GUE}} - \bar{g}_\Lambda^{\text{GOE}} = \Lambda^2/[2(\Lambda + 1)(2\Lambda + 1)]$. For open dots with $\Lambda$ equivalent channels (with transmission coefficient 1) in each lead the size of the weak localization peak is $\Lambda/[2(2\Lambda + 1)]$, so that the weak localization peak in closed dots (measured in units of $\frac{2e^2}{h \pi kT}$) is smaller by a factor of $\Lambda/(\Lambda + 1)$ than the one in open dots (measured in units of $e^2/h$). In units of the average GUE conductance height, we find $\delta g_\Lambda(0)/\bar{g}_\Lambda^{\text{GUE}} = 1/[2(\Lambda + 1)]$ for closed dots, compared with $1/(2\Lambda + 1)$ for open dots.
It is interesting to compare the dependence of the width of the weak localization peak on the number of channels in open versus closed dots. The right inset of Fig. 1 shows the full width at half maximum (FWHM) of the weak localization peak as a function of the number of channels for both closed (solid circles) and open (open circles) dots. The FWHM in closed dots is approximately independent of the number of channels, unlike the case of open dots where the FWHM increases as $\sqrt{\Lambda}$ (for large $\Lambda$). In open dots, the crossover scale in the conductance is determined by the competition between the mixing time $\hbar/\tau_{\text{mix}} = 2\pi \Delta \lambda^2$ (associated with the breaking of time-reversal symmetry) and the decay time of resonances in the dot $\hbar/\tau_{\text{dec}} = (\Delta/2\pi)\Lambda$. The crossover in the conductance occurs when $\tau_{\text{dec}}/\tau_{\text{mix}} = 4\pi^2 \lambda^2 / \Lambda \sim 1$, i.e., for $\lambda_{\text{cr}}^{\text{open}} \sim \sqrt{\Lambda} / 2\pi$. The decay time in an open dot is shorter than the Heisenberg time by a factor of $\Lambda$, and for one open channel these two times are equal. For closed dots on the other hand, the decay time is much longer than the Heisenberg time (since the average level width is much smaller than $\Delta$), irrespective of the number of channels. In such a case the relevant time scale that competes with the mixing time is the Heisenberg time, and the crossover occurs on a scale of $\lambda_{\text{cr}}^{\text{closed}} \sim 1$, independent of the number of channels and comparable to the value of $\lambda_{\text{cr}}$ for an open dot with $\Lambda = 1$. This is confirmed by the exact RMT results of Fig. 1, where the full width at half maximum (FWHM) of the weak localization peak is $\Delta \lambda \approx 0.4$, irrespective of the number of channels in each lead.

If time-reversal symmetry is broken by a magnetic field, then $\lambda = \Phi/\Phi_{\text{cr}}$, where $\Phi_{\text{cr}}$ is the crossover flux (defined as the flux where $\lambda = 1$). It is possible to estimate $\Phi_{\text{cr}}$ semiclassically assuming single-particle chaotic motion [3,18]. The rms of the electron action difference between a pair of time-reverse orbits is given by $[(\Theta^2)^{1/2}/A](\Phi/\Phi_0)$ (in units of $\hbar$), where $(\Theta^2)^{1/2}$ is the rms area accumulated by the electron and $A$ is the area of the dot. Since the area accumulation in a chaotic dot is diffusive, we have for an open dot

$$(\Theta^2)^{1/2}/A = \alpha_g^{1/2} \sqrt{\tau_{\text{cr}}/\tau_{\text{dec}}},$$

where $\tau_{\text{cr}}$ is the time scale for the electron to cross the dot, and $\alpha_g$ is a geometrical factor that is dot specific [18]. For the breaking of time reversal symmetry, we require the rms action difference to be of order 1, leading to $\Phi_{\text{cr}}^{\text{open}}/\Phi_0 = \ldots$
\[ \pi \alpha_g^{-1/2} \sqrt{\Lambda \tau_{cr}/\tau_H} \]. For the closed dot, on the other hand, \( \tau_H \) replaces \( \tau_{dec} \) to give a \( \Lambda \)-independent crossover flux of \( \Phi_{cr}^{\text{closed}}/\Phi_0 = \pi \alpha_g^{-1/2} \sqrt{\tau_{cr}/\tau_H} \). A weak localization of the average conductance peak height was observed experimentally in semi-open dots \([11]\) with an FWHM of \( \approx 6.2 \text{ mT} \), corresponding to \( B_{cr} \approx 14.9 \text{ mT} \) or \( \Phi_{cr} \approx 2\Phi_0 \). The desymmetrized stadium billiard in a uniform magnetic field gives \( \Phi_{cr} \approx 2\Phi_0 \) \([18]\) with about the same number of electrons (\( \sim 1000 \)). Thus the experimental result is about three times larger than the billiard model estimate, similar to the discrepancy estimated for the correlation field \([7]\). The exact discrepancy is uncertain because of the unknown geometrical factor of the dot used in the experiment. The discrepancy does not indicate a breaking of RMT (which only predicts the universal form of the correlator \([19]\) and not its magnetic flux scale), but rather that effects beyond single-particle dynamics are important. Indeed, a model that takes into account electron-electron interactions in the dot gives a correlator whose form is consistent with the RMT correlator, but with a correlation field that is larger by about a factor of three as compared with a non-interacting model \([20]\).

It is also possible to calculate in closed form the variance of the conductance as a function of \( \lambda \). The second moment of the conductance is given by \( \bar{g}^2(\lambda) = \left\langle \int_0^\infty d\Gamma Q^{(2)}(\Gamma, t)/\Gamma^2 \right\rangle \), where \( Q^{(2)} \) is the convolution of \( (\Gamma^4)^2 P(\Gamma^4|t) \) with itself. The Laplace transform is then calculated from \( Q^{(2)}(s, t) = [d^2 P(s|t)/ds^2]^2 \). For one-channel leads we find

\[
\bar{g}^2(\lambda) = \frac{3}{16} + \frac{27}{2} \left\langle \left( \frac{t}{1 + t^2} \right)^2 \left( \frac{1 + t^2}{1 - t^2} \right)^4 \left[ \frac{1}{1 + t^2} \ln t + 1 + \frac{1}{12} \left( \frac{1 - t^2}{t} \right)^2 - \frac{2}{27} \left( \frac{1 - t^2}{t} \right)^4 \right] \right\rangle. \tag{8}
\]

\( \bar{g}^2 \) increases from 3/16 (GOE) to 1/5 (GUE), while the variance \( \sigma^2(g) = \bar{g}^2 - \bar{g}^2 \) decreases from 1/8 to 4/45. Closed expressions for \( \bar{g}^2(\lambda) \) can be similarly derived for any number of equivalent channels \( \Lambda \). The left inset of Fig. \([4]\) shows the analytically calculated \( \sigma(g)/\bar{g} \) as a function of \( \lambda \) for \( \Lambda = 1, 2, 3 \) and 5 channels. We see that the distribution becomes sharper in the crossover from conserved to broken time-reversal symmetry. In the GOE and GUE limits we have

\[
\sigma(g)/\bar{g} = \begin{cases} 
\left[ (\Lambda^2 + 5\Lambda + 2) \Lambda^2(\Lambda + 3) \right]^{1/2} (\text{GOE}) \\
\left[ (2\Lambda^2 + 5\Lambda + 1) 2\Lambda^2(2\Lambda + 3) \right]^{1/2} (\text{GUE})
\end{cases}. \tag{9}
\]
For large $\Lambda$, this quantity decreases in the crossover from $1/\sqrt{\Lambda}$ (GOE) to $1/\sqrt{2}\Lambda$ (GUE), i.e., effectively doubling the number of channels.

At finite temperature, several levels $i$ contribute to the conductance peak with known thermal weights $w_i(T)$ \[12,21\] which are determined by the single-particle spectrum and the number of electrons in the dot. The dimensionless peak height is given by $g = \sum_i w_i(T) g_i$, where $g_i$ are the level conductances as in Eq. (1), but with $\Gamma^{l,r}$ depending now on the level $i$. In calculating the statistics of the peak heights at finite $T$, a good approximation is to neglect the fluctuations in the dot’s spectrum, and $w_i(T)$ are then fixed. Since $\bar{g}_i(\lambda)$ are independent of $i$, the weak localization peak factorizes $\delta g(\lambda, T) \approx [\sum_i w_i(T)] \delta g(\lambda)$, where $\delta g(\lambda)$ is the weak localization peak at temperatures much smaller than $\Delta$. In particular, since $\bar{g}_{\text{GUE}}(T) \approx [\sum_i w_i(T)] \bar{g}_{\text{GUE}}$, the scaled weak localization peak $\delta g/\bar{g}_{\text{GUE}}$ is temperature-independent (assuming no phase breaking).

Another quantity of interest is the size of the weak localization peak relative to the rms of the (GUE) peak-heights distribution. For dots with one-channel leads, we find $\delta g(0)/\sigma_{\text{GUE}}(g) \approx (\sqrt{5}/8)[(\sum w_i)^2/\sum w_i^2]^{1/2}$, which increases with temperature. At high temperatures, the number of levels that contribute to a given peak height is of the order of $T/\Delta$, and $\delta g(0)/\sigma_{\text{GUE}}(g) \propto (T/\Delta)^{1/2}$. At low temperatures it is difficult to measure the weak localization peak because its size is comparable to the conductance fluctuations. In open dots, the large number of open channels enhances the peak relative to the fluctuations by $\sqrt{\Lambda}$. In closed dots (with $\Lambda = 1$), the temperature can be used instead to enhance the peak relative to the background fluctuations.

In conclusion, we have derived in closed form the weak localization peak of Coulomb blockade peak heights in quantum dots. In the absence of phase breaking, this peak is temperature-independent up to an overall scale.

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FIG. 1. The weak localization peak in closed dots (measured in units of the average GUE conductance) $\delta g(\lambda)/\bar{g}_{\text{GUE}}$ versus the crossover parameter $\lambda$. Shown are the cases with $\Lambda = 1, 2, 3$ and 5 channels (peak is decreasing as $\Lambda$ increases). Right inset: The FWHM of the weak localization peak $\Delta \lambda$ as a function of the number of channels $\Lambda$ in each lead for closed and open dots (solid and open circles, respectively). Left inset: The ratio $\sigma(g)/\bar{g}$ between the rms and average values of the conductance peak as a function of the crossover parameter $\lambda$ for Coulomb blockade dots with $\Lambda = 1, 2, 3$ and 5 channels ($\sigma(g)/\bar{g}$ is smaller for larger $\Lambda$).