Conductance fluctuations of open quantum dots under microwave radiation

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We develop a time-dependent random-matrix theory describing the influence of a time-dependent perturbation on mesoscopic conductance fluctuations in open quantum dots. The effect of external field is taken into account to all orders of perturbation theory, and our results are applicable to both weak and strong fields. We obtain temperature and magnetic field dependences of conductance fluctuations. The amplitude of conductance fluctuations is determined by electron temperature in the leads rather than by the width of electron-distribution function in the dot. The asymmetry of conductance with respect to inversion of applied magnetic field is the main feature allowing to distinguish the effect of direct suppression of quantum interference from the simple heating if the frequency of external radiation is larger than the temperature of the leads $\hbar \omega \gg T$.

INTRODUCTION

Transport coefficients of disordered and chaotic electron systems fluctuate from sample to sample. These fluctuations are commonly called mesoscopic fluctuations. Mesoscopic fluctuations of conductance of noninteracting systems are universal. The universality means that the variance of the conductance $\langle \delta g^2 \rangle$ is of the order $G_0^2$, where $G_0 = e^2/\pi \hbar$ is the quantum of conductance and is weakly dependent on the sample geometry, see Refs. 4 and 5.

The fluctuations of transport properties of electron systems is a quantum-mechanical phenomenon based on the interference of quantum states. As any other interference phenomena, conductance fluctuations are very sensitive to inelastic processes, commonly referred to as dephasing. The dephasing processes in open quantum dots were considered on purely phenomenological basis. First microscopic consideration of the microwave radiation on the weak localization in quantum dots was performed in Ref. 10. In this reference, the concept of the time-dependent random-matrix theory was used.

The purpose of the present paper is to extend the results of Ref. 10 to describe the effect of the external microwave radiation on the mesoscopic fluctuations of the conductance. The ultimate goal is to identify observable features that allow one to distinguish the effect of the external field to the dot from simple heating.

However, there is a significant difference in the calculation of the mesoscopic conductance fluctuations and the averaged conductance. Because, the dot is subjected to the external classical radiation that produces nonequilibrium in the dot, the dc current $I_0$ through the dot is finite (though randomly changing from one configuration to another) even if the dc voltage $V = V_l - V_r$ across the dot is zero, see Fig. 1. This current $I_0$ is due either to the photovoltaic effect or to rectification of ac bias across the dot, see Refs. 13 and 14. We are interested in the linear response to the applied dc voltage across the dot,

$$I_{dc} = I_0 + gV + O(V^2).$$

In principal, the linear in $V$ contribution to the current comes from two sources: (i) the nonequilibrium of the distribution functions in the leads, (ii) change in the photovoltaic current, correspondent to a different realization of the dot due to the finite bias. Nevertheless, we will show that due to the electroneutrality condition the nonequilibrium current prevails.

Closing the introductory part, we note that the present paper has a certain overlap with the recent preprint by Wang and Kravtsov, where the conductance fluctuations were calculated for open quantum dots subjected to a periodic ac pumping. Our treatment is different in several aspects. Firstly, our results are applicable for the frequencies of the external radiation $\omega$ smaller than the Thouless energy of the dot $E_T$, whereas treatment of Ref. 16 is valid in the opposite regime. Secondly, unlike Ref. 16, we will restrict ourselves to the case of the monochromatic radiation acting on the dot. Finally, we will highlight the role of the electroneutrality requirement in a separability of the photovoltaic effect and the mesoscopic conductance fluctuations, which was not done in Ref. 16.

MODEL

We apply the random matrix theory (RMT) to study the conductance of open quantum dots, see Ref. 4. All correc-
tions to the RMT are governed by a small parameter \(N_{ch}/g_{dot}\), where \(g_{dot} = E_d/\delta_1\) and \(\delta_1\) is the mean level spacing, see Ref. 17 for the detailed discussion. We consider the conductance fluctuations of quantum dots with a large number of open channels \(N_{ch} \gg 1\). In this approximation, we neglect the effects of the electron-electron interaction on the conductance that are as small as \(1/N_{ch}^2\), while the conductance fluctuations are proportional to \(1/N_{ch}\). The same condition allows us to use a conventional diagrammatic technique\(^1\) to take the ensemble average. External microwave radiation is modeled as a time-dependent random part of the Hamiltonian of the dot.

The Hamiltonian of the system is\(^1\)

\[
\mathcal{H}(t) = \mathcal{H}_D(t) + \mathcal{H}_L + \mathcal{H}_{LD},
\]

where \(\mathcal{H}_D\) is the Hamiltonian of the electrons in the dot, which is determined by the \(M \times M\) matrix \(H_{nm}\),

\[
\mathcal{H}_D(t) = \sum_{n,m=1}^{M} \psi_n^d H_{nm}(t) \psi_m + E_c n^2,
\]

\(\psi_n\) corresponds to the states of the dot and the thermodynamic limit \(M \to \infty\) is assumed, \(E_c\) is the charging energy, and \(n = \sum_{m=1}^{M} \psi_n^d \psi_m^n\), the last term in Eq. (3) is the largest contribution to the interaction effects in quantum dot, see Ref. 17 for the discussion of the status of this approximation.

Matrix \(\hat{H}_D(t)\) is given by

\[
\hat{H}_D(t) = \hat{H} + \hat{V}\psi(t).
\]

Here the time-independent part of the Hamiltonian \(\hat{H}\) is a random realization of a \(M \times M\) matrix, which obeys the correlation function

\[
\langle H_{nm}(\Phi_1) H_{n'm'}(\Phi_2) \rangle = M \left( \frac{\delta_1}{\pi} \right)^2 \left( 1 - \frac{N_d}{4M} \right) \delta_{nm} \delta_{nn'} \delta_{mm'},
\]

\[
+ \left( 1 - \frac{N_d}{4M} \right) \delta_{nm'} \delta_{mm'},
\]

where \(\delta_1\) is the mean level spacing of the dot and parameters \(N_{d,c}\) describe the effect of the magnetic field on the dot.\(^1\) These parameters can be estimated as \(N_{d,c} = g_{dot} = \Phi_0(\Phi_1 - \Phi_2)/2\Phi_0^2\), where \(\Phi_1, \Phi_2\) is the magnetic flux through the dot and \(\Phi_0 = h/e\) is the flux quantum. The time-dependent perturbation is described by symmetric \(M \times M\) matrices \(V_{nm}\) and function \(\psi(t)\) of time. We assume that the perturbation is harmonic with single frequency \(\omega\), \(\psi(t) = \cos \omega t\), even though most of the considerations [up to Eq. (26)] is valid for an arbitrary function. The effect of the perturbation on the system is totally determined by two parameters:\(^2\)

\[
Z = \frac{1}{M} \text{Tr} \hat{V}, \quad C_0 = \frac{\pi}{M^2 \delta_1} \text{Tr} \hat{V}^2.
\]

Parameter \(Z\) has a meaning of the average velocity of the energy levels of the dot under the external perturbation and can be omitted from our consideration due to screening (see below). The parameter \(C_0\) characterizes its typical deviation.\(^2\) Since all the physical responses of the system are characterized by the same parameters, the value of \(C_0\) can be eliminated by an independent measurement.

The electron spectrum in the leads near Fermi surface can be linearized,

\[
\mathcal{H}_L = \hbar \nu \sum_{\alpha,k} k \psi_\alpha^d(k) \psi_\alpha(k),
\]

where \(\psi_\alpha(k)\) denotes different electron states in the leads, \(k\) labels the continuum of momentum states in each channel \(\alpha\), \(\hbar \nu = 1/2\pi \nu\) is the Fermi velocity and \(\nu\) is the density of states per channel at the Fermi surface. We put \(\hbar = 1\) in all intermediate formulas below.

The coupling between the dot and the leads is

\[
\mathcal{H}_{LD} = \sum_{\alpha,n,k} \left[ W_{\alpha n} \psi_\alpha^d(k) \psi_n + \text{H.c.} \right].
\]

For the reflectionless point contacts, the coupling constants, \(W_{\alpha n}\), in Eq. (8) are given by:\(^4,17\)

\[
W_{\alpha n} = \begin{cases} \sqrt{\frac{M \delta_1}{\pi^2 \nu}} & \text{if } n = \alpha \leq N_{ch} \\ 0 & \text{otherwise} \end{cases}
\]

For open dots with a large number of open channels \(N_{ch} \gg 1\) the interaction term can be treated within mean-field approximation, so that the Hamiltonian (3) takes the form

\[
\mathcal{H}_{D}^{mf}(t) = \sum_{n,m=1}^{M} \psi_n^d [H_{nm}(t) + eV_d(t) \delta_{nm}] \psi_m^n,
\]

\(eV_d = 2E_c \langle n \rangle_q\),

where \(\langle n \rangle_q\) stands for the quantum mechanical (but not ensemble) of the number of electrons in the dot. Corrections to mean-field treatment (10a) were calculated in Ref. 18 and shown to be small as \(1/N_{ch}^2\).

In the mean-field approximation (10a), one can introduce one-particle scattering matrix \(S_{\alpha \beta}(t_1, t_2)\) as

\[
\delta_{\alpha \beta}(t, t') = \delta_{\alpha \beta}(t - t') - 2\pi i \nu \text{Tr} \psi_\alpha^d W_{\alpha n} W_{n \beta}^{*} \psi_n^d t_1, t_2\),
\]

and the Green functions \(\hat{G}^{(R,A)}(t, t')\) are the solutions of

\[
\frac{\partial}{\partial t} \hat{H}_D(t) - eV_d(t) \pm i\pi \nu \hat{W} \hat{W}^d \hat{G}^{(R,A)}(t, t') = \delta(t - t'),
\]

where the matrices \(\hat{H}_D\) and \(\hat{W}\) are defined by Eqs. (4) and (9).

The dc current through the dot is given by, see Ref. 12:

\[
I_{dc} = e \int_0^{T_p} dt \int dt_1 dt_2 \text{Tr} \{ \hat{f}(t_1 - t_2) [\hat{S}(t_2, t) \hat{A}(t, t_1) - \hat{A}(t_2, t_1)] \},
\]

\(085115-2\)
where $T_p$ is the period of the external perturbation, $\hat{f}(t)$ is related to the Fourier transform of the electron-distribution function in the $\alpha$th channel as

$$f_{\alpha\beta}(t) = \delta_{\alpha\beta} \frac{iTe^{-iV_d t}}{\sinh \pi T t}$$

and

$$N_{\alpha\beta} = \delta_{\alpha\beta} \left\{ \begin{array}{ll} \frac{N_l}{N_{ch}} & \text{for } 1 \leq \alpha \leq N_l \\ -\frac{N_l}{N_{ch}} & \text{for } N_l+1 \leq \alpha \leq N_{ch}. \end{array} \right.$$  

The spin degeneracy is taken into account in Eq. (13). We assume that the degeneracy is not lifted by magnetic field.

To complete the theory one needs an equation for the averaged number of particles $\langle n(t) \rangle_q$, see Eq. (10b). It is found from the continuity relation as

$$\frac{d\langle n(t) \rangle_q}{dt} = -\int dt_1 dt_2 Tr[f(t_1-t_2)$$

$$\times \left\{ \hat{S}^\dagger(t_2,t)\hat{S}(t,t_1) - \delta(t_2-t_1) \right\}].$$

Equations (13)–(16) are similar to those used in Ref. 22 for studying the frequency dependence of the conductance of mesoscopic system.

**ENSEMBLE AVERAGING**

Our goal now is to perform calculations of the conductance correlation function

$$R(\Phi_1, \Phi_2) = \left\langle g(\Phi_1)g(\Phi_2) \right\rangle - \left\langle g(\Phi_1) \right\rangle \left\langle g(\Phi_2) \right\rangle,$$  

using the model outlined above.

We use the leading approximation in small parameter $1/N_{ch}$. The fluctuations of the conductance are smaller than its average value and we can use instead of sample specific Eq. (16) its ensemble-averaged counterpart,

$$\frac{d\langle n(t) \rangle_q}{dt} = -\frac{N_{ch} \delta_1}{2\pi} \langle n(t) \rangle_q + \frac{eN_l}{\pi} [V_l - V_d - Z\varphi(t)]$$

$$+ \frac{eN_r}{\pi} [V_r - V_d - Z\varphi(t)].$$  

Equation (17) is nothing but a discrete form of the diffusion equation for the bulk system and the last two terms correspond to the divergence of the drift current. Substituting Eq. (10b) into Eq. (17), solving the resulting differential equation, we find

$$eV_d(t) + Z\varphi(t) = \frac{4eE_c}{4E_c + \delta_1} \frac{N_lV_l + N_rV_r}{N_{ch}}$$

$$+ Z\frac{\delta_1 + (2\pi/N_{ch})\delta_1}{\delta_1 + 4E_c + (2\pi/N_{ch})\delta_1} \varphi(t).$$  

Equation (18) gives

$$V_d = \frac{N_lV_l + N_rV_r}{N_{ch}}$$  

We notice from Eq. (18) that the characteristic energy scale governing charge dynamics is $E_cN_{ch}/2\pi$. Usually, this scale is of the order of the Thouless energy, $E_T$. Because the random-matrix theory is capable to describe the energy scale smaller than $E_T$, we can consider only $\omega \ll E_T = E_cN_{ch}/2\pi$. Moreover, for the small quantum dot $E_c \gg \delta_1$, so that Eq. (18) gives

$$V_d = \frac{N_lV_l + N_rV_r}{N_{ch}}$$  

FIG. 2. Two diagrams, which contribute to the conductance correlation function $R$.
and the time-dependent perturbation can be considered as traceless, \( Z = 0 \). This constant electric potential of the dot can be removed from Eq. (12) for the Green function by the following gauge transformation:

\[
\hat{G}(t,t') = \hat{G}|_{V_d=0}(t,t') \exp[-ieV_d(t-t')].
\] (20)

Substituting Eq. (20) into Eq. (13) and expanding up to the first power in \( V = V_r - V_l \), we find

\[
g = \frac{\partial I_{dc}}{\partial V} = g_d + \frac{G_0}{\hbar} \int_0^\infty dt_1 dt_2 F(T(t_1-t_2))
\]

\[
\times \text{Tr}[\hat{S}(t_1)\hat{\Lambda} S^\dagger(t_2,t)\hat{\Lambda}],
\] (21)

The diffuson and the Cooperon are defined by the following equations, see Fig. 3:

\[
C(\tau_1,\tau_2,t) = \Theta(\tau_1 - \tau_2) \exp\left(-\frac{1}{2} \int_{\tau_2}^{\tau_1} d\tau K_c(\tau,t)d\tau\right),
\] (24a)

\[
D(t_1,t_2,\tau) = \Theta(t_1 - t_2) \exp\left(-\int_{t_2}^{t_1} d\tau K_d(\tau,t)d\tau\right),
\] (24b)

where

\[
K_{d,c} = \gamma_{d,c} + C_0 (\varphi(t + \tau/2) - \varphi(t - \tau/2))^2
\] (25a)

\[
\gamma_{d,c} = \frac{\delta_1}{2\pi}(N_{clh} + N_{d,c}),
\] (25b)

and parameters \( N_{d,c} \) describe the effect of the magnetic field, see Eq. (5).

Derivation of Eq. (23) deserves a little bit of additional discussion. One notices, that the diagrams of Fig. 2 do not contain any piece corresponding to the classical distribution function in the dot. We can trace it into the expression for conductance fluctuations (21) that contains traceless vertices \( \hat{\Lambda} \), which cannot be dressed by the dashed line. On the other hand, any vertex with finite trace corresponds to the modified-distribution function in the dot. We can trace it into the expression for conductance fluctuations, and the temperature dependence of the conductance fluctuations is uniquely determined by the electron temperature in the leads. That means that contrary to the common beliefs, e.g., Ref. 7, the amplitude of the mesoscopic fluctuations cannot be used to study the distribution function of electrons in the dot. From the theoretical side, it is important to emphasize, that the appearance of the traceless vertices is determined solely by the electroneutrality condition (19), any other choice of the dot bias would lead to the change in the photovoltaic current.2

IV. LIMITING CASES

Below we consider the limit of high, \( \hbar \omega \gg \max(C, \gamma_{d,c}) \), and low, \( \hbar \omega \ll \gamma_{d,c} \), frequencies. For the high frequencies, \( \omega \gg C \), we obtain

\[
R = \frac{\delta_1^2 \delta_{cl}^2}{4\pi^2} \left[ \frac{1}{\gamma_d^2} Q_d \left( C_0, T, \gamma_d \right) + \frac{1}{\gamma_c^2} Q_c \left( C_0, T, \hbar \omega, \gamma_c \right) \right],
\] (26)

where the dimensionless \( Q \) functions are given by

\[
Q_d(x,y) = \int_0^\infty d\tau F^2(y\tau) \int_0^\infty \frac{\exp[-\tau(1+2x\sin^2\xi/2)]}{1+2x\sin^2\xi/2} \frac{d\xi}{2\pi},
\] (27a)

\[
Q_c(x,y,z) = \int_0^\infty d\tau F^2(y\tau)
\]

\[
\times \int_0^{2\pi} \frac{\exp[-\tau(1+2x\sin^2\xi/2)]}{1+x(\sin^2\xi/2 + \sin^2(\xi/2 + z\tau/2))} \frac{d\xi}{2\pi}.
\] (27b)
where time-reversal symmetry in the system with time-dependent Sager relation is a simple manifestation of the lifting of the symmetry of the conductance.\textsuperscript{23} For the limit of high temperature, $T \gg 1$, we obtain

$$Q_c(x,y,z) \approx Q_d(x,y) = \frac{\pi^2}{1+2x}.$$  

(28)

The equality between functions $Q_c(x,y,z)$ and $Q_d(x,y)$ means the magnetic field symmetry of the conductance.\textsuperscript{23} Indeed, using Eqs. (25), (26), and (28) we observe, that

$$R(\Phi_1,\Phi_2) = R(\Phi_1, -\Phi_2).$$  

(29)

However, in low temperature limit $y \ll 1$, we obtain for $x \gg 1$

$$Q_d(x,0,z) = \frac{1+x}{(1+2x)^{3/2}} \approx \frac{1}{2 \sqrt{2x}}.$$  

(30a)

$$Q_c(x,0,z) = \frac{1}{2x}.$$  

(30b)

Comparison of Eqs. (26) and (30) reveals an important fingerprint of the dephasing by the external radiation—violation of the Onsager relation

$$\frac{R(\Phi_1, -\Phi_2)}{R(\Phi_1, \Phi_2)} = \sqrt{\frac{2 \tilde{\gamma}}{1/\gamma_c}},$$  

(31)

where $\tilde{\gamma} = \gamma_d(\Phi_1, \Phi_2) = \gamma_c(\Phi_1, -\Phi_2)$, provided that $\gamma_c(\Phi_1, \Phi_2) = \gamma_d(\Phi_2, -\Phi_1)$. This breakdown of the Onsager relation is a simple manifestation of the lifting of the time-reversal symmetry in the system with time-dependent Hamiltonian.

Figure 4 shows $Q_c(x,y,z)$ and $Q_d(x,y)$ as a function of temperature $y$. At low temperature, $Q_c(x,y,z)$ is smaller than $Q_d(x,y)$ and it approaches $Q_d(x,y)$ as temperature $y$ increases.

In the limit of low frequency $\hbar \omega \ll \gamma_d, c$, the contribution from the Cooperon and diffusion parts are described by the same function, so that the conductance correlation function can be represented in the form

$$R = \frac{\delta \xi^2}{4 \pi^2} \left[ \frac{1}{\gamma_d} Q(x,0, T) - \frac{1}{\gamma_c} Q(x,0, T) \right],$$  

(32)

so that the Onsager relation (29) holds. Here

$$Q(x,y) = \int_0^{\pi/2} \frac{d\xi d\tilde{\xi}}{4 \pi^2} \int_0^\infty F^2(y \tau) \times \exp \left\{ - (1 + 4x \sin^2 \tilde{\xi}/2 \sin^2 \xi/2) \tau \right\} d\tau.$$  

(33)

This expression in the limit of high temperature $T \gg \gamma_d, c$ has an asymptotic behavior

$$Q(x,y) \approx \frac{\pi}{3y} K(-4x).$$  

(34)

At zero temperature $Q(x,y)$ is given by the expression

$$Q(x,0) = \frac{1}{\pi} E(-4x) + (1+4x) K(-4x) \times \frac{1}{1+4x},$$  

(35)

where $K(x)$ and $E(x)$ are the elliptic integrals of the first and second kind, respectively.

$$K(x) = \int_0^{\pi/2} \frac{d\varphi}{\sqrt{1-x \sin^2 \varphi}};$$  

$$E(x) = \int_0^{\pi/2} \frac{d\varphi}{\sqrt{1-x \sin^2 \varphi}}.$$

We conclude that the conductance fluctuations are suppressed by external radiation even in the limit of the low frequency, see Eq. (32). Indeed, during one period of time, the system goes along a closed loop in the parameter space and the contribution to the dc conductance is effectively determined by the equilibrium conductance, correspondent to each point of the loop. The equilibrium conductance fluctuates along this loop. Thus, the observed dc conductance is already partially averaged over some realizations of the quantum dot and its fluctuations decrease. The perturbation strength is related to the length of the contour in the parameter space and effectively determines how many different dot’s configurations contribute to the dc conductance. Consequently, the stronger perturbation, over the larger number of the realizations the dc conductance is averaged and the smaller fluctuations of the dc conductance.

This should be contrasted with the suppression of the averaged magnetoresistance.\textsuperscript{8,10} There, the stationary field does not do anything because the result is already ensemble averaged. In order to suppress the average quantum correction, the field should change on the time scale of the order of $1/\gamma_{esc}$, where $\gamma_{esc} = \delta \xi N_{esc} / 2\pi$ is the escape rate from the dot. That is why the effect of the low-frequency radiation on...
conductance fluctuations and weak localization corrections are significantly different. At high frequency $\hbar \omega \gg \gamma_{esc}$ the dc conductance no longer can be represented in terms of the stationary conductance and the suppression of both the conductance fluctuations and the weak localization correction to the conductivity can be interpreted as dephasing.

**COMPARISON WITH EXPERIMENT**

Our results still contain an unknown parameter $C_0$ characterizing the strength of the perturbation. There is a way, however, to present the results in a form not depending on this parameter, thus eliminating a need for additional fitting. Following Ref. 7, we represent the parametric dependence of the weak localization correction $\delta g_{wl}$ vs $g$, where $\delta g_{wl}$ is given by:

$$\delta g_{wl} = \frac{-e^2}{\pi \hbar} \frac{N_fN_r}{(N_f+N_r)^2} P \left( \frac{C_0}{\gamma_{esc}}, \frac{\hbar \omega}{\gamma_{esc}} \right),$$

$$P(x,z) = \int_0^\infty e^{-\xi-x}P I_0[x\phi]d\xi, \phi = \xi - \frac{\sin z \xi}{z}.$$  \hfill (36)

The conductance variance is determined from Eq. (23) for broken time-reversal symmetry $\gamma_{c,d} \to \infty$. Figure 5 shows the parametric dependence for various values of the parameters $C_0$ and $\omega$ and $T=10 \gamma_{esc}$.

We observe that the shape of the curves depends on the frequency of external radiation. Particularly, in the limit of low frequency $\hbar \omega \ll \gamma_{eig}$ the weak localization correction is not changed by the radiation, while the conductance fluctuations may be significantly suppressed. At high frequency, $\hbar \omega \gg C_0$, $\gamma_{esc}$ the curves become nonsensitive to the radiation frequency.

The authors of Ref. 7 found that the radiation applied to their device produced curves in $\var g$ vs $\delta g_{wl}$ plane identical to the curve produced by increasing temperature of the device for a wide range of frequencies. This observation apparently demonstrates that the radiation produces the heating of electrons and the effect of dephasing without heating, see Ref. 24, is not observed in experiments. It was also suggested that the main mechanism is the increase of the temperature in the dot due to the Joule heat by induced ac source-drain bias.

Although the present data of Ref. 7 support the heating mechanism of suppression of the weak localization correction to the conductance and the conductance fluctuations, we believe that more detailed analysis has to be done. According to our theory, see Eq. (21) and the paragraph below Eq. (25b), (i) mesoscopic fluctuations are sensitive only to the temperature in the leads, and therefore, the notion of heating of electrons in the dot responsible for $1/T$ dependence of the mesoscopic fluctuations is not relevant: if there is a heating, it manifests itself only through the temperature of the leads; (ii) high-frequency curves of our theoretical Fig. 5 quantitatively agree with data on Fig. 3 of Ref. 7, for frequencies $f$ = 1, 10, and 25 GHz. An exception is the lowest-frequency curve ($f=100$ MHz) represented in this plot, for the dot with $\delta_1=2.4\mu$eV, $N_{ch}=2$ corresponds to $\hbar f/\gamma_{esc}=0.5$, so according to our Fig. 5 it should have observable deviations from the high-frequency curves, which is not seen. However, taking into account uncertainty in determination of the level spacing $\delta_1$ from the geometrical area of the dot, this does not unambiguously rule out the microwave dephasing mechanism.

We believe that the “smoking gun” evidence for the mechanism considered in the present paper is the violation of the Onsager relation (31) in high-frequency regime, $\hbar \omega \gg \gamma_{c,d}, \gamma$. The dependence of this violation on the amplitude of the field $C_0$ is the main prediction of the theory.

**SUMMARY**

We constructed the time-dependent random-matrix theory to describe the effect of the nonequilibrium external radiation on conductance fluctuations of an open quantum dot. The main experimental feature to reveal such a mechanism is the breakdown of the magnetic symmetry of the conductance by high-frequency radiation, $\hbar \omega \gg T$, see Eq. (31).

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CONDUCTANCE FLUCTUATIONS OF OPEN QUANTUM . . .

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20 Note that $C_0$ is defined in Ref. 10 with a different numerical factor.