Crossover Between Universality Classes in the Statistics of Rare Events in Disordered Conductors

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The crossover from orthogonal to the unitary universality classes in the distribution of the anomalously localized states (ALS) in two-dimensional disordered conductors is traced as a function of magnetic field. We demonstrate that the microscopic origin of the crossover is the change in the *symmetry* of the underlying disorder configurations that are responsible for ALS. These disorder configurations are of *weak* magnitude (compared to the Fermi energy) and of *small* size (compared to the mean free path). We find their shape explicitly by means of the direct optimal fluctuation method.

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Statistical properties of the wave functions, $\psi(\boldsymbol{r})$, in disordered conductors have been the subject of an intensive theoretical study during the last decade [1–4]. Considerable attention was devoted to the so-called anomalously localized states (ALS). These states constitute the large-$|\psi|^2$ "tail" of the wave function distribution, where this distribution deviates strongly from the prediction of the random-matrix theory. Analytical results obtained to date are summarized in the comprehensive review [5]. Calculations of the ALS density, based on the nonlinear σ model (NLSM) [1–4], yield exponentially small tails with different numerical factors in the exponent, depending on the universality class. A drawback of the NLSM approach is that it fails to pinpoint the actual disorder realizations responsible for ALS. In an attempt to overcome this deficiency, an alternative approach, the direct optimal fluctuation (DOF) method, was proposed [6]. In two dimensions the optimal disorder configuration obtained in [6] had a shape of a compact potential well with a radius of the order of the Fermi wavelength, supplemented by a periodic modulation at large distances. It was observed in [6] that within a numerical factor in the exponent, the NLSM-based results can be reproduced by considering a simplest configuration yielding the ALS, namely, a circular (in two dimensions) region of space surrounded by a high and thick potential barrier. Then the anomalously large values of the wave function, ψ, can be characterized by a single, intuitively transparent, parameter

$$\mathcal{T} = |\psi_{\text{in}}|^2/|\psi_{\text{out}}|^2, \tag{1}$$

where "in" and "out" refer to the inner and outer boundaries of the barrier. As a direct consequence of the compactness of the fluctuations, the probability of their formation is insensitive to weak magnetic fields, in apparent contradiction to the σ-model predictions. On the basis of this contradiction the authors [6] questioned the ability of the NLSM to handle the ALS.

In the present paper we demonstrate that an appropriate application of the DOF method yields the ALS density in the large-$\mathcal{T}$ domain that is sensitive to the universality class. In this domain the DOF method allows one to trace the crossover between the orthogonal and unitary classes with increasing magnetic field. Consideration is restricted to the two-dimensional case. We start with the observation that the parameter $\mathcal{T}$ can be interpreted as a dimensionless time (quality factor) during which an electron remains trapped within a given disorder configuration. This allows one to study the areal density of potential fluctuations with a given *trapping time* in an *infinite* sample rather than the distribution of the $\psi(\boldsymbol{r})$ in a finite-size sample. Later we establish the relation between the two distributions.

The sensitivity to a weak magnetic field within the direct optimal fluctuation approach is restored in two steps. First, instead of circle-shape fluctuations [6] we consider fluctuations of a ring shape, schematically depicted in Fig. 1. This step is in accord with the paper by Karpov [7], who has demonstrated that, in three dimensions, fluctuations of a toroidal shape, providing a given value of $\mathcal{T}$, are more probable than spherical fluctuations [6]. Ring-shape fluctuations, having a much bigger area, are sensitive to the magnetic field. However, this class of fluctuations *alone* does not exhibit a crossover between the orthogonal and unitary classes. To reveal this crossover, as a second step, we note that, on a ring, the states with angular momenta m and $-m$ are *degenerate* in the absence of magnetic flux. This suggests that the fluctuations lifting this degeneracy cause the *increase* of $\mathcal{P}(E_F, \mathcal{T})$, which is the probability to find a state with energy E_F and a given value of $\mathcal{T}$. The underlying reason is completely analogous to the Jahn-Teller effect [8]. Kusmartsev and Rashba [9] employed this argument to demonstrate that the fluctuations responsible for the *tail* states of a fourfold degenerate semiconductor valence band are of ellipsoidal rather than spherical shape.

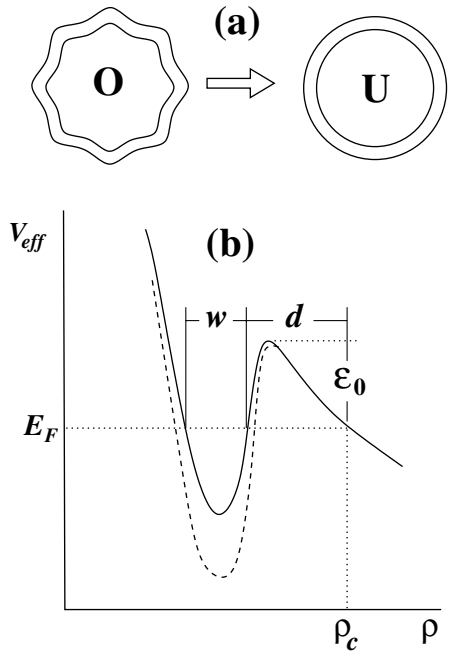

FIG. 1. (a) Illustration of the crossover between universality classes. Equipotential lines of the warped and ring-shaped fluctuations are shown schematically. (b) Effective potential for the angular harmonics m. Optimal $V_0(\rho)$ at $\delta = 0$ and at critical $\delta \approx 1/3$ are plotted with solid and dashed lines, respectively.

Summarizing, in a zero magnetic field $\mathcal{P}(E_F, \mathcal{T})$ is determined by the *warped* ring-shape fluctuations of the random potential. In a finite magnetic field the degeneracy $|m\rangle, |-m\rangle$ is lifted; the warp does not lead to an increase of $\mathcal{P}$ anymore. Then the fluctuations contributing to $\mathcal{P}$ are close to rings. Thus, in the statistics of rare events, the crossover between the orthogonal and unitary classes can be viewed as switching from *warped* to *almost perfect* rings. This is illustrated in Fig. 1(a).

The above discussion suggests the following analytical steps. First, the fluctuation $V(\boldsymbol{r})$ of the white-noise random potential is presented as a sum of the angular harmonics $V(\boldsymbol{r}) = \sum_m V_m(\rho) e^{im\phi}$, where $\rho = |\boldsymbol{r}|$. Second, in the probability $\mathcal{P}\{V(\boldsymbol{r})\}$ of the fluctuation $V(\boldsymbol{r})$

$$|\ln \mathcal{P}| = \frac{\pi \nu \tau}{2\hbar} \int dr V^2(\boldsymbol{r}) = \frac{\pi^2 \nu \tau}{\hbar} \sum_{m=-\infty}^{\infty} \int d\rho\, \rho |V_m(\rho)|^2, \tag{2}$$

only the angular harmonics $V_0(\rho)$ and $V_{\pm 2m}(\rho)$ are retained. This is because the $|m\rangle, |-m\rangle$ coupling is provided by the harmonics $\pm 2m$ of the random potential. In Eq. (2) τ and ν stand for the scattering time and the density of states, respectively. Then the system of equations relating the components χ_m and χ_{-m} of the wave function $\rho^{1/2}\psi(\rho, \phi)$ takes the form

$$\hat{H}_0 \chi = -\frac{\hbar^2}{2M} \frac{d^2\chi}{d\rho^2} + (V_{\text{eff}}^{(m)}(\rho) + \hat{V}_2 \hat{\sigma}_x)\chi = E_F \chi, \tag{3}$$

where M is the electron mass, $\chi = (\chi_m, \chi_{-m})$ is a two-component wave function, the matrix $\hat{V}_2 = \text{diag}(V_{2m}, V_{-2m})$ is diagonal, and $\hat{\sigma}_i$ are the Pauli matrices. The effective potential in (3) is a sum of a centrifugal potential $\frac{\hbar^2(m^2-1/4)}{2M\rho^2}$ and $V_0(\rho)$. The turning point, ρ_c, is determined by $V_{\text{eff}}^{(m)}(\rho_c) = E_F$ [see Fig. 1(b)]. As we will see below, the characteristic width w of the fluctuation $V_0(\rho)$ and the barrier thickness d [see Fig. 1(b)] are related as $w \ll d \ll \rho_c$. Similarly to [7], this separation of spatial scales allows a number of crucial simplifications: (i) since $w \ll d$, it follows that, within the barrier, $V_{\text{eff}}^{(m)}(\rho)$ is dominated by the centrifugal term; (ii) since $d \ll \rho_c$, this term can be linearized within the barrier region as $V_{\text{eff}}^{(m)}(\rho) = E_F + \varepsilon_0(\rho_c - \rho)d^{-1}$, where $\varepsilon_0 \approx \hbar^2 m^2 d / M\rho_c^3$ is the barrier height. Using the above simplifications the ratio $\mathcal{T} = |\psi_E(\rho_c)|^2 / |\psi_E(\rho_c - d)|^2$ can be readily calculated $\ln \mathcal{T} = 2\frac{\sqrt{2M}}{\hbar} \int_{\rho_c - d}^{\rho_c} d\rho [V_{\text{eff}}(\rho) - E_F]^{1/2}$, yielding

$$\ln \mathcal{T} = (2m/3)(2d/\rho_c)^{3/2} = (2m/3)(\varepsilon_0/E_F)^{3/2}. \tag{4}$$

It is seen from Eq. (4) that the condition $\rho_c \gg d$ can be rewritten as $m \gg \ln \mathcal{T}$. Equation (4) relates $\mathcal{T}$ to the energy level position ε_0 in the potential well centered around $\rho_c - d - w/2$ [see Fig. 1(b)]. Thus the problem of finding the distribution of $\mathcal{T}$ reduces to the conventional problem [10,11] of finding the most probable configuration of random potential which confines the energy level of a depth $-\varepsilon_0$. In our case $-\varepsilon_0$ is the eigenvalue of the *matrix* Hamiltonian $\hat{H} = -\frac{\hbar^2}{2M}\frac{d^2}{dx^2} + V_0(x) + \hat{V}_2 \hat{\sigma}_x$, where $x = \rho - \rho_c + d + w/2$. The form of the Hamiltonian follows from the system Eq. (3) and the condition $w \ll d$ which allows one to neglect the change of the centrifugal potential within the potential well [see Fig. 1(b)].

The straightforward generalization of the approach [10,11] to the matrix Hamiltonian $\hat{H}$ implies minimizing the functional $F\{V_0, V_{\pm 2m}\} = |\ln \mathcal{P}| - \lambda(\chi^+ \hat{H} \chi)$. With an appropriate choice of the Lagrange multiplier, λ, we obtain the expressions for V_0 and $V_{\pm 2m}$ in terms of the two-component wave function, χ. Namely, $V_0 = \chi^+ \chi$, $V_{\pm 2m} = \chi^+ \hat{\sigma}_{\pm} \chi$. At this point it is convenient to introduce dimensionless variables as $z = x(2M\varepsilon_0)^{1/2}/\hbar$ and $\tilde{\chi} = \varepsilon_0^{-1/2}\chi$. With new notations, the matrix Schrödinger equation, $\hat{H}\chi = -\varepsilon_0 \chi$, reduces to

$$\tilde{\chi}'' + \frac{3}{2}(\tilde{\chi}^+ \tilde{\chi})\tilde{\chi} - \frac{1}{2}(\tilde{\chi}^+ \hat{\sigma}_z \tilde{\chi})\hat{\sigma}_z \tilde{\chi} - \tilde{\chi} = 0. \tag{5}$$

The probability Eq. (2) of the fluctuation, providing a given value of $\mathcal{T}$, takes the form

$$|\ln \mathcal{P}_m| = \pi^2 \frac{\nu \tau \rho_c}{\hbar} \int dx [|V_0(x)|^2 + 2|V_{2m}(x)|^2] = (3/8)\pi k_F l C_m \ln \mathcal{T}, \tag{6}$$

where $k_F = (2ME_F)^{1/2}/\hbar$ is the Fermi momentum, $l = \hbar k_F \tau / M$ is the mean free path; dimensionless factor C_m

in Eq. (6) is defined as

$$C_m = \frac{1}{2}\int dz[3(\tilde{\chi}^+\tilde{\chi})^2 - (\tilde{\chi}^+\hat{\sigma}_z\tilde{\chi})^2]. \quad (7)$$

Equations (5)–(7) express analytically the Jahn-Teller physics discussed in the introduction. Indeed, this system has a conventional instanton solution $\tilde{\chi}_{-m} = 0$ and $\tilde{\chi}_m = 2^{1/2}\cosh^{-1}z$, corresponding to the absence of warp ($V_{\pm 2m} \equiv 0$). This solution yields $C_m = 16/3$. On the other hand, the Jahn-Teller–type solution $\tilde{\chi}_m = \tilde{\chi}_{-m} = (2/3)^{1/2}\cosh^{-1}z$ leads to a *smaller* value $C_m = 32/9$, which corresponds to *exponentially* higher probability $\mathcal{P}_m$. Thus, the final result for the orthogonal case reads

$$\ln\mathcal{P}_m = -(4/3)\pi g \ln\mathcal{T}, \quad (8)$$

where $g = k_F l$ is the dimensionless conductance. Remarkably, $\ln\mathcal{P}_m$ *does not* depend on the value of the angular momentum m, provided that $m \gg \ln\mathcal{T}$. In fact, the condition $m \gg \ln\mathcal{T}$ justifies all the assumptions made in deriving the result Eq. (8). Indeed, making use of the relation $\varepsilon_0 = E_F(3\ln\mathcal{T}/2m)^{2/3}$, one can verify that this condition insures $V_0(x) \ll E_F$, i.e., the potential well is *shallow*. Other assumptions used in the calculation are less restrictive. Indeed, the typical width of the potential well is $w \sim \hbar/(M\varepsilon_0)^{1/2} \sim k_F^{-1}(m/\ln\mathcal{T})^{1/3}$. On the other hand, for barrier width d we have from Eq. (4) $d = (\rho_c/2)(3\ln\mathcal{T}/2m)^{2/3}$. Since $\rho_c = m/k_F$, we have $d = m^{1/3}(3\ln\mathcal{T}/2)^{2/3}/2k_F$, so that $w/d \sim \ln^{-1}\mathcal{T} \ll 1$. We have also assumed that the change of the centrifugal potential within the potential well is much smaller than ε_0. This change can be estimated as $\varepsilon_0\frac{w}{d} \sim \varepsilon_0/\ln\mathcal{T} \ll \varepsilon_0$.

We now turn to the case of a finite magnetic field, B. Because of the azimuthal symmetry of the problem the system of equations (3) relating the radial functions χ_m, χ_{-m} can be easily modified to

$$\hat{H}_0\chi + \frac{\hbar^2}{2M}\left(\frac{m}{l_B^2}\hat{\sigma}_z + \frac{\rho^2}{4l_B^4}\right)\chi = E_F\chi, \quad (9)$$

where $l_B = (c\hbar/eB)^{1/2}$ is the magnetic length. As can be seen from Eq. (9) there exists a following hierarchy of magnetic field strengths. When $l_B \gg m^{1/2}/k_F$, then the "confining" term ρ^2/l_B^4 can be neglected. The latter condition is equivalent to $R_L \gg \rho_c$, where $R_L = k_F l_B^2$ is the Larmor radius for an electron with energy E_F. On the other hand, it is obvious that the term $\pm\hbar^2 m/(2Ml_B^2)$, lifting the $|m\rangle, |-m\rangle$ degeneracy, becomes important when it is of the order of ε_0. This yields $R_L \sim \rho_c(\ln\mathcal{T}/m)^{-2/3}$. Since $\ln\mathcal{T} \ll m$, we conclude that there exists an interval of magnetic fields which affect the potential well but do not affect the barrier. Condition $\varepsilon_0 \sim \hbar^2 m/(Ml_B^2)$ determines the characteristic magnetic field for orthogonal-unitary crossover: $B_0^{(m)} = (\Phi_0/2\pi)k_F^2(3\ln\mathcal{T}/2)^{2/3}m^{-5/3}$, where Φ_0 is the flux quantum. The remaining task is to find the dependence $C_m(\delta)$, where $\delta = B/B_0^{(m)}$, which describes the change of C_m between "orthogonal," $C_m(0) = 32/9$, and "unitary," $C_m = 16/3$, values with increasing magnetic field, δ. At finite δ the system Eq. (5) takes the form

$$\tilde{\chi}'' + [3(\tilde{\chi}^+\tilde{\chi}) - (\tilde{\chi}^+\hat{\sigma}_z\tilde{\chi})\hat{\sigma}_z]\frac{\tilde{\chi}}{2} - \tilde{\chi} = \delta(1 - \hat{\sigma}_z)\tilde{\chi}, \quad (10)$$

whereas Eqs. (6) and (7) retain their form. At small $\delta \ll 1$ the dependence $C_m(\delta)$ can be found by solving Eq. (10) perturbatively and substituting $\tilde{\chi}(x)$ into Eq. (7). This program yields the following analytical result:

$$C_m(\delta) = C_m(0)(1 + 3\delta/2). \quad (11)$$

Our important observation is that perturbative expression Eq. (11) remains valid with high accuracy up to $\delta = 1/3$, when the unitary solution $\tilde{\chi}_{-m} = 0$, $\tilde{\chi}_m = 2^{1/2}\cosh^{-1}z$ takes over. The crossover behavior $C_m(\delta)$ is illustrated in Fig. 2. In the same figure we show the orthogonal solution of Eq. (10) at the crossover point.

Overall, the DOF approach yields in two dimensions $|\ln\mathcal{P}| \sim g\ln\mathcal{T}$ rather than $|\ln\mathcal{P}| \sim g\ln^2\mathcal{T}$ [6]. This is because the optimal fluctuations are close to rings rather than to circles. The ring area, $S_m \sim \rho_c w$, being much bigger than the area of a circle $\sim k_F^{-2}$ in [6] is the origin of a high sensitivity of rings to the magnetic field (in contrast to a conventional suppression of the Cooperon modes in the metallic regime). On the other hand, due to the same relation, $S_m \gg k_F^{-2}$, the rings are much more "vulnerable" to the perturbations caused by the harmonics with small angular momenta, for which the centrifugal barrier is low. To demonstrate this we write down the expression for the decay rate of a quasilocal state in a ring

$$\mathrm{Im}E = 2\pi\sum_\mu\left(\langle m|V_0|\mu\rangle^2 + \sum_{n\neq 0}\langle m|V_n|\mu\rangle^2\right)\delta(E - E_\mu) = \frac{\hbar}{\tau_1} + \frac{\hbar}{\tau_2}, \quad (12)$$

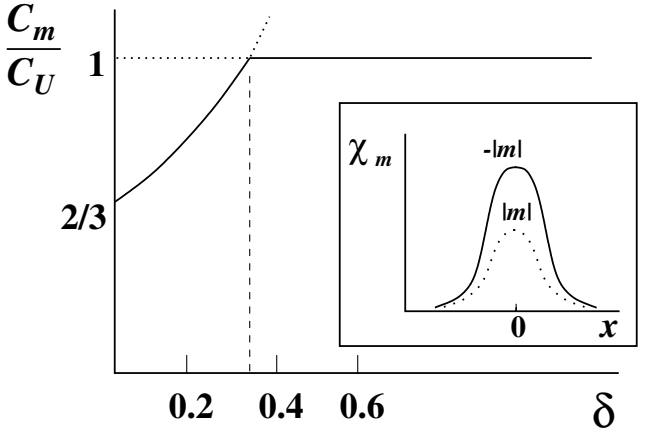

FIG. 2. Normalized logarithm of the ALS density is shown versus the dimensionless magnetic field, δ. The inset shows the "orthogonal" optimal wave function $(\tilde{\chi}_m, \tilde{\chi}_{-m})$ for crossover value $\delta \approx 1/3$.

where $\langle m|V_n|\mu\rangle$ are the matrix elements of the harmonics $V_n(\rho)$ [see Eq. (2)]; $|\mu\rangle$ is the state of the continuous spectrum in the absence of random potential (the corresponding energy is E_μ). Above we were looking for the fluctuations with anomalously large $\mathcal{T}$. Note that within a numerical factor $\mathcal{T} = \varepsilon_0\tau_1/\hbar$. It is seen from Eq. (12) that our consideration is justified only if $\tau_2 \gtrsim \tau_1$. Within the DOF approach, harmonics with small n, contributing to $\hbar/\tau_2$, do not affect the principal exponent. They enter at the stage of the prefactor calculation [12]. A *typical* value of τ_2 is the scattering time τ, which is much shorter than τ_1. The relation $\tau_2 \gtrsim \tau_1$ is satisfied only for sparse configurations, in which the harmonics V_n are *suppressed* within the ring. Therefore, our result for $\mathcal{P}_m$ should be multiplied by the probability, P, to find such a configuration. To calculate this probability we consider the distribution function of time τ_2:

$$P(\tau_2) = \left\langle \delta\left(\frac{1}{\tau_2} - \int d\boldsymbol{r}d\boldsymbol{r}' S(\boldsymbol{r},\boldsymbol{r}')V(r)V(r')\right)\right\rangle_{V(r)} \quad (13)$$

where the expression for $S(\boldsymbol{r},\boldsymbol{r}')$ follows from Eq. (12)

$$S(\boldsymbol{r},\boldsymbol{r}') = \frac{2\pi}{\hbar}\rho_c\chi_m(\rho)\chi_m^*(\rho')\sum_\mu \psi_\mu(\boldsymbol{r})\psi_\mu^*(\boldsymbol{r}')\delta(E - E_\mu).$$

Averaging over random potential $V(\boldsymbol{r})$ in Eq. (13) can be carried out explicitly. The result is expressed in terms of eigenvalues, λ_n, of the integral operator with the kernel $S(\boldsymbol{r},\boldsymbol{r}')$: $\int d\boldsymbol{r}' S(\boldsymbol{r},\boldsymbol{r}')\phi_n(\boldsymbol{r}') = \lambda_n\phi_n(\boldsymbol{r})$. With the accuracy of a numerical factor, the result can be obtained qualitatively. This is because $P(\tau_2)$ is determined by the first N_0 eigenvalues, which are almost equal to each other: $\lambda_n \approx \lambda_0$ for $n < N_0$. The value N_0 can be estimated as the number of squares with a side k_F^{-1} within the area of a ring S_m, i.e., $N_0 \sim \rho_c w k_F^2$. On the other hand, there is an exact sum rule $\sum_n \lambda_n = 2\pi\nu$ for the eigenvalues λ_n. Since $\sum_n \lambda_n \approx \lambda_0 N_0$, we can find λ_0. This leads to the following result for the distribution Eq. (13) $P(\tau_2) \propto (\nu\tau/\lambda_0\tau_2)^{N_0} \propto (\tau/\tau_2)^{N_0}$. The sought product $\tilde{\mathcal{P}}_m = \mathcal{P}_m P(\tau_2 \sim \tau_1)$, describing the probability to find ALS with "preexponential" accuracy, can be then presented as $\ln\tilde{\mathcal{P}}_m = \ln\mathcal{P}_m - c_2 N_0 \ln(\mathcal{T}/g)$, i.e.,

$$\ln\tilde{\mathcal{P}}_m - \ln\mathcal{P}_m = -c_2 m^{4/3}\ln^{-1/3}\mathcal{T}\ln(\mathcal{T}/g), \quad (14)$$

where $c_2 \sim 1$ is a numerical factor. The above expression for $\tilde{\mathcal{P}}_m$ allows one to find the optimal value of the angular momentum, m. Indeed, the prefactor *falls off* rapidly with m, while the m-dependent correction to the main term due to weak nonlinearity of the barrier potential is $\sim \ln\mathcal{P}d/\rho_c = -c_1 g(\ln\mathcal{T}/m)^{2/3}\ln\mathcal{T}$, *increases* with m. Here $c_1 \sim 1$ is another numerical coefficient. Probability $\tilde{\mathcal{P}}_m$ is maximal for $m = m_{\rm opt} = c_0(g\ln\mathcal{T})^{1/2}$, where $c_0 = (c_1/2c_2)^{1/2}$. The resulting expression for $\tilde{\mathcal{P}}_m$ with preexponential accuracy reads

$$|\ln\tilde{\mathcal{P}}_m| = \frac{4}{3}\pi g\ln\mathcal{T}\{1 + (9c_0^{4/3}/8)[\ln\mathcal{T}/g]^{1/3}\}. \quad (15)$$

It is seen from Eq. (15) that the above consideration is valid within the domain $1 \ll \ln\mathcal{T} \ll g$. For the optimal m, the ring radius is $\rho_c = c_0 k_F^{-1}(g\ln\mathcal{T})^{1/2} \ll l$.

Within a factor $\sim k_F^2$, Eq. (15) describes the density of traps with a trapping time $\hbar\mathcal{T}/\varepsilon_0$ in an *infinite* sample. To relate Eq. (15) to the statistics of the wave functions, consider a *closed* sample with a radius $R \gg l$. Then the DOF result Eq. (15) describes the probability to find a stationary eigenfunction for which ψ^2 inside the ring is by a factor $\mathcal{T}$ bigger than outside [see Eq. (1)], so that almost the entire weight of this eigenstate is concentrated inside the ring. On the other hand, the NLSM yields the ALS, in which ψ^2 is "amplified" by a factor of $\mathcal{T}$ from the periphery of the sample towards the center in a power-law fashion [5]. This ratio of ψ^2 at the center and at the periphery governs the trapping time (or, in other words, relaxation time [5]) after the sample is slightly opened. The NLSM result for the density of states with a given ratio $\mathcal{T}$ is $|\ln\mathcal{P}(\mathcal{T})| = \frac{1}{8}\pi g\ln^2\mathcal{T}/\ln(R/l)$, that is valid for $\mathcal{T} \gg (R/l)^4$ [5]. Comparison of this result to Eq. (15) suggests that NLSM captures anomalously high values of ψ^2 within the domain $(R/l)^4 \ll \mathcal{T} \ll (R/l)^{32/3}$. For larger $\mathcal{T}$ the ALS are governed by the warped rings.

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