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Asymptotic localization of stationary states in the nonlinear Schrödinger equation

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The mapping of the Nonlinear Schrödinger Equation with a random potential on the Fokker-Planck equation is used to calculate the localization length of its stationary states. The asymptotic growth rates of the moments of the wave function and its derivative for the linear Schrödinger Equation in a random potential are computed analytically and resummation is used to obtain the corresponding growth rate for the nonlinear Schrödinger equation and the localization length of the stationary states.

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I. INTRODUCTION

In this work we consider the problem of the one-dimensional Anderson localization [1, 2] for the nonlinear Schrödinger equation (NLSE) [3, 4, 5]. In spite of the extensive research, many fundamental problems are still open, and, in particular, it is not clear whether in one dimension (1D) Anderson localization can survive the effects of nonlinearities. This problem is relevant for experiments in nonlinear optics, for example disordered photonic lattices [6, 7], where Anderson localization was found in presence of nonlinear effects as well as experiments on Bose-Einstein Condensates (BEC) in disordered optical lattices [8, 9, 10, 11, 12]. The interplay between disorder and nonlinear effects leads to new interesting physics [9, 10, 13, 14, 15, 16]. In particular, the problem of spreading of wave packets and transmission are not simply related [17, 18], in contrast with the linear case.

We consider one-dimensional localization of stationary solutions of the nonlinear Schrödinger equation (NLSE) in a random δ -correlated potential $V(x)$ with a Gaussian distribution (white noise), of zero mean and variance $2D$, namely, $\langle V(x)V(x') \rangle = 2D\delta(x-x')$. The linear version of this model was studied extensively in the past [19]. Following our previous analysis [20] we study Anderson localization of stationary solutions with energies ω in the framework of the stationary NLSE

$$\omega\phi(x) = -\partial_x^2\phi(x) + \beta\phi^3(x) + V(x)\phi(x), \quad (1)$$

where $\phi(x)$ is real, the variables are chosen in dimensionless units and the Planck constant is $\hbar = 1$. For the lattice version of the model, it was established rigorously [3, 21, 22] that the stationary solutions of this equation are exponentially localized for a wide range of conditions. It was also argued that the rate of growth of moments for the stationary NLSE (1) coincides exactly with the linear case [20] and determines the localization length.

We will specifically calculate $\langle\phi^2(x)\rangle$ of solutions of Eq. (1) that are found for a certain ω , with given boundary conditions at some point, for example $\phi(x=0)$ and $\phi'(x=0)$, where prime means the derivative with respect to x . This will be done with the help of the analogy with the Langevin equation [19, 20, 23]. In particular, we will calculate the growth rate of the second moment,

$$2\gamma = \lim_{x \rightarrow \infty} \frac{\ln\langle\phi^2(x)\rangle}{x} > 0, \quad \xi = \frac{1}{\gamma}, \quad (2)$$

that will turn out to be independent of β , where ξ is the localization length. Note, it is different from the usually studied (in the linear case) self averaging quantity $\gamma_s = \frac{1}{2} \frac{d}{dx} \langle \ln \phi^2(x) \rangle = \frac{1}{2} \lim_{x \rightarrow \infty} \frac{\langle \ln \phi^2(x) \rangle}{x}$, and γ is a smooth function of energy. Since the distribution of the random potentials is translationally invariant, it is independent of the choice of the initial point as $x = 0$. Like in the linear case, starting from a specific initial condition, $\phi(x)$ will typically grow. For specific values of ω at some point this function will start to decay, so that a normalized eigenfunction is found. This is the approach *à la* Borland [24, 25] that was made rigorous for the linear case in [26, 27]. Here, following [20] we extend this approach, in a heuristic form, to the nonlinear case. The envelope of the wave function will grow exponentially if we start either from the right or from the left. The value of ω results from the matching condition, so that an eigenfunction has some maximum and decays in both directions as required by the normalization condition. The exponential decay is an asymptotic property, while the matching is determined by the potential in the vicinity of

the maximum. This observation is crucial for the validity of this approach and enables us to determine the exponential decay rate of states from the solution of the initial value problem (1). In [20] a linear equation for the moments of $\phi(x)$ of Eq. (1) and its Langevin analog was derived and it was shown that the exponents that control the growth of the moments are identical to the ones of the linear system ($\beta = 0$). In the present paper it will be shown that this is also correct for the asymptotic behavior of the moments.

The outline of the paper is as follows. The analogy with the Langevin equation and dynamics of the moments is outlined in Section II. The generalized Lyapunov exponents, that are the eigenvalues which determine the growth of the various moments, are presented in Section II and their asymptotic behavior is derived analytically for the first time in Appendix A (they were found numerically in [28]). The resulting asymptotic expansion for the growth of the moments is presented in Section III. The results are summarized in Section IV.

II. THE FOKKER-PLANCK EQUATION AND LYAPUNOV EXPONENTS

Following Ref. [20], we perform the calculation of $\langle \phi^2(x) \rangle$ by using the analogy with the classical Langevin equation [19, 23]. Therefore, considering the x -coordinate as the formal time variable on the half axis $x \equiv \tau \in [0, \infty)$, Eq. (1) reduces to the Langevin equation

$$\ddot{\phi} + \omega\phi - \beta\phi^3 - V(\tau)\phi = 0 \quad (3)$$

with the δ correlated Gaussian noise $V(\tau)$. Now we introduce new variables $u = \phi$ and $v = \dot{\phi} \equiv \frac{d\phi}{d\tau}$ (that play the role of position and velocity in the Langevin equation) and a distribution function of these new variables is $P = P(u, v, \tau)$. The dynamical process in the presence of the Gaussian δ -correlated noise is described by the distribution function that satisfies the Fokker-Planck equation: (FPE) [20, 30]

$$\partial_\tau P - [\omega u - \beta u^3] \partial_v P + v \partial_u P - 2Du^2 \partial_v^2 P = 0, \quad (4)$$

where $2Du^2$ is the only nonzero component of the diffusion tensor.

A. Equation for moments

We are interested in the average quantum probability density $\langle \phi^2(x) \rangle \equiv \langle u^2(\tau) \rangle$, where

$$\langle u^2(\tau) \rangle = \int u^2 P(u, v, \tau) du dv.$$

From the FPE we obtain a system of equations for the moments

$$M_{k,l} = \langle u^k v^l \rangle, \quad (5)$$

where $k, l = 0, 1, 2, \dots$. Substituting $u^k v^l$ in the FPE and integrating over u and v , one obtains the following relation for $M_{k,l}$

$$\dot{M}_{k,l} = -l\omega M_{k+1,l-1} + kM_{k-1,l+1} + l(l-1)2DM_{k+2,l-2} + \beta l M_{k+3,l-1}, \quad (6)$$

where $M_{k,l}$ with negative indexes are assumed to vanish. We note that only terms with the same parity of $k+l$ are coupled. Since we are interested in $M_{2,0} = \langle u^2 \rangle$, we study only the case when this parity is even, namely $k+l = 2n$ with $n = 1, 2, \dots$. The sum of the indexes of the moments is $2n$, except the last term $\beta l M_{k+3,l-1}$, where the sum is $2(n+1)$. This leads to the infinite system of linear equations that can be written in the form

$$\dot{\mathbf{M}} = \mathcal{W}\mathbf{M}, \quad (7)$$

where the column vector is $\mathbf{M} = (M_{2,0}, M_{1,1}, M_{0,2}, M_{4,0}, M_{3,1}, \dots)$ and \mathcal{W} is the corresponding matrix. The matrix elements $\mathcal{W}_{k,l}$ are determined by Eq. (6). The solutions of the system of linear equations (7) are linear combinations of the eigenfunctions at time τ

$$\mathbf{M}_\lambda(\tau) = \exp(\lambda\tau)\mathbf{U}_\lambda, \quad (8)$$

where $\mathbf{U}_\lambda = \mathbf{M}_\lambda(t=0)$ is the eigenvector of \mathcal{W} corresponding to the eigenvalue λ found from the equation

$$\mathcal{W}\mathbf{U}_m = \lambda_m \mathbf{U}_m, \quad \mathbf{U}_m \equiv \mathbf{U}_{\lambda_m}. \quad (9)$$

The vector \mathbf{M} consists of the block vectors \mathbf{M}_n : $\mathbf{M} = (\mathbf{M}_1, \mathbf{M}_2, \dots, \mathbf{M}_n, \dots)$, where \mathbf{M}_n is a vector of $2n+1$ elements defined in Eq. (5) and corresponds to the moments of the order of $2n$. Therefore, the initial vector at $\tau = 0$ is

$$(\mathbf{M}_1, \mathbf{M}_2, \dots, \mathbf{M}_n, \dots) = \sum_{n'} \sum_{l'} c_{n', l'} \mathbf{U}_{n', l'}. \quad (16)$$

Assume that at the initial point $\tau = 0$, the wave function and its derivative are small, of the order of ϵ in units of $\sqrt{|\omega/\beta|}$. Then, at that point the moments scale as $\mathbf{M}_n \sim \epsilon^{2n}$ with ϵ arbitrary small. As follows from Eq. (16), one finds $c_{n, l} U_{n, l} \sim \bar{c}_n \epsilon^{2n} + o(\epsilon^{2n})$ with bounded \bar{c}_n for any n , as demonstrated in Appendix B. In the linear case ($\beta = 0$) the growth rate of each moment of the order of $2n$ corresponding to the n -th block is determined by the eigenvalue with the largest real part, and we denote it by $\lambda_{\max}(n) \equiv \max(\text{Re } \lambda_{n, l})$, where the maximum is over the $2n+1$ eigenvalues corresponding to the n th block, indicated by l . As shown in Appendix A, the asymptotic behavior of the generalized Lyapunov exponent $\lambda_{\max}(n)$ when $n \rightarrow \infty$ is given by $\lambda_{\max}(n) \sim \mathcal{A} n^{4/3}$, where $\mathcal{A} = \frac{3}{4} 2^{4/3} D^{1/3}$ is a constant (see Eq. (12)). The leading contribution to the growth of $\mathbf{M}_n(\tau)$ in the nonlinear case ($\beta \neq 0$) is determined by the following sum

$$\tilde{\mathbf{M}}_n(\tau) = \sum_{m \geq n} \bar{c}_m \epsilon^{2m} e^{\mathcal{A} m^{4/3} \tau}, \quad (17)$$

as is clear from Eqs. (13)–(16).

B. Resummation

This series in Eq. (16) has a vanishing radius of convergence and probably has to be interpreted as an asymptotic series. It can be used to study the behavior of $\tilde{\mathbf{M}}_n(\tau)$ after being resummed. Such a resummation is done with the help of the identity

$$\exp(K^2) = \frac{1}{\sqrt{\pi}} \int_{-\infty}^{+\infty} du \exp(-u^2 + 2Ku), \quad (18)$$

known as the Hubbard-Stratonovich transformation. We can rewrite the above series as follows

$$\tilde{\mathbf{M}}_n(\tau) = \sum_{m \geq n} \bar{c}_m \epsilon^{2m} \exp(\mathcal{A} m^{4/3} \tau) = \frac{1}{\sqrt{\pi}} \int_{-\infty}^{+\infty} du \exp(-u^2) \Phi(2u\sqrt{\mathcal{A}\tau}), \quad (19)$$

where the function Φ is given by

$$\Phi(y) = \sum_{m \geq n} \bar{c}_m \epsilon^{2m} \exp(m^{2/3} y). \quad (20)$$

If the coefficients \bar{c}_m do not grow too fast (which we shall assume hereafter) the function $\Phi(y)$ is well defined at least for small values of y . Note that for the resummation of Eq. (19) it was not crucial that the power of m is $4/3$. Such a resummation can be performed for any power $\bar{\alpha} < 2$ replacing $4/3$.

A more explicit resummation procedure for the r.h.s. of equation (17) can be developed with the help of fractional derivatives. First, let us expand the exponential function

$$\exp(\mathcal{A}\tau m^{4/3}) = \sum_{k=0}^{\infty} \frac{(\mathcal{A}\tau m^{4/3})^k}{k!}. \quad (21)$$

Then, writing $\epsilon^{2m} = e^{m \ln \epsilon^2} \equiv e^{\zeta m}$, we obtain

$$\tilde{\mathbf{M}}_n(\tau) = \sum_{k=0}^{\infty} \frac{(\mathcal{A}\tau)^k}{k!} \sum_{m \geq n} \bar{c}_m m^{4k/3} e^{\zeta m}. \quad (22)$$

We now introduce the Weyl fractional derivative of order q of a function $f(z)$ by the Weyl integral, see e.g. [32],

$$\frac{d^q f(z)}{dz^q} \equiv \frac{1}{\Gamma(-q)} \int_{-\infty}^z \frac{f(y) dy}{(z-y)^{1+q}} \quad (23)$$

where for $q > 0$ the integral should be properly regularized [32, 33], and $\Gamma(-q)$ is the gamma function. For $f(y) = e^{\nu\zeta}$ it takes the form

$$\frac{d^q e^{\nu\zeta}}{d\zeta^q} = \nu^q e^{\nu\zeta}. \quad (24)$$

Substitution of Eq. (24) in Eq. (22) with $\nu = m$ yields

$$\tilde{\mathbf{M}}_n(\tau) = \sum_{k=0}^{\infty} \frac{(\mathcal{A}\tau)^k}{k!} \frac{d^{4k/3}}{d\zeta^{4k/3}} \sum_{m \geq n} \bar{c}_m \epsilon^{2m}. \quad (25)$$

If \bar{c}_m are bounded, as shown in Appendix B, by some \bar{C}_n , the sum in Eq. (25) is bounded by $\frac{\bar{C}_n \epsilon^{2n}}{1 - \epsilon^2}$, hence $\tilde{\mathbf{M}}_n$ is the fractional derivative of some well defined function presented in Eq. (23) with its regularization.

Therefore, Eqs. (25) and (19) describe the long time asymptotics of the moments. Therefore, Eq. (17) is an asymptotic expansion and is a good approximation as long as we sum decreasing terms. The condition is (for bounded \bar{c}_m)

$$\epsilon^{2m} \exp(\mathcal{A}m^{4/3}\tau) > \epsilon^{2(m+1)} \exp(\mathcal{A}(m+1)^{4/3}\tau). \quad (26)$$

For large m this inequality is $\ln \frac{1}{\epsilon} > \frac{2}{3}\mathcal{A}\tau m^{1/3}$ with $m \geq n$. Consequently, for the time of the order

$$\tau < \tau_0^{(n)} \equiv \frac{3 \ln(1/\epsilon)}{2 \mathcal{A}n^{1/3}} \quad (27)$$

the n -th moment will be dominated by the leading terms and will grow as in the linear case ($\beta = 0$).

C. The growth of the second moments

The second moments are of particular interest for the present work. Their growth for a time that is shorter than $\tau_0^{(1)}$ of Eq. (27) is dominated by the leading term, namely $\tilde{\mathbf{M}}_2 = \bar{c}_2 \epsilon^2 e^{\lambda_{\max}(1)\tau}$. Consequently, for $\tau < \tau_0^{(1)}$,

$$\langle u^2 \rangle = M_{2,0} = \bar{c}_2 \epsilon^2 e^{\lambda_{\max}(1)\tau}. \quad (28)$$

This result was verified numerically. Using the analogy between the stationary Schrödinger equation (1) and the Fokker-Planck equation (4) we identify the generalized Lyapunov exponent $\lambda_{\max}(1)$ with the growth rate (2) as

$$2\gamma = \lambda_{\max}(1) = \lim_{x \rightarrow \infty} \lim_{\epsilon \rightarrow 0} \frac{\ln \langle \phi^2(x)/\epsilon^2 \rangle}{x} > 0. \quad (29)$$

For $\tau > \tau_0^{(1)}$, non-linear saturation effects become relevant and the full nonlinear theory should be used.

IV. SUMMARY

In this work the asymptotic behavior of the generalized Lyapunov exponents of the linear Fokker-Planck equation (3) with $\beta = 0$, was found analytically and is given by Eq. (12). The resulting expression for the moments (17) is divergent, but it can be resummed in the form Eqs. (19) or (25). Therefore, for short time $\tau < \tau_0^{(n)}$, the first term in Eq. (17) provides a good approximation of the moments. In particular, for the second moment this result enabled to identify the generalized Lyapunov exponent of the linear system with the asymptotic growth rate for the nonlinear one, namely

$$2\gamma = \lambda_{\max}(1).$$

According to the implementation of the *à la* Borland method, as outlined in [20], this is the decay rate of the stationary states of the random nonlinear equation (1), showing that it is independent of the nonlinearity β .

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APPENDIX A: ASYMPTOTIC BEHAVIOR OF THE GENERALIZED LYAPUNOV EXPONENTS

In this appendix, $\lambda_{\max}(n)$, the maximal eigenvalue of the matrix (11) will be evaluated. Following [28] we call this quantity a *generalized Lyapunov exponent*. By elementary dimensional analysis [28], one finds that the following scaling is suggestive,

$$\lambda_{\max}(n) = D^{1/3} L_n \left(\frac{\omega}{D^{2/3}} \right). \quad (\text{A1})$$

In the long time limit the moments of order $2n$ grow as $e^{\lambda_{\max}(n)t}$. In [28], Zillmer and Pikovsky studied $\lambda_{\max}(n)$ for different values of n : they give exact expressions for $n = 2$, for $n \rightarrow 0$ (which corresponds to the usual Lyapunov exponent) and consider limiting cases for large values of the dimensionless parameter $\omega/D^{2/3}$. They also study numerically the behavior of $\lambda_{\max}(n)$ as $n \rightarrow \infty$, keeping the dimensionless parameter fixed. They found numerically the scaling law:

$$\lambda_{\max}(n) \propto n^\alpha \quad \text{with} \quad \alpha \simeq 1.4. \quad (\text{A2})$$

The fact that the scaling exponent α is different from 2, implies deviations from Gaussian behavior and breakdown of monoscaling; the consequences of this breakdown of single parameter scaling for conductance distribution were studied by Schomerus and Titov [29]. We also remark that related studies were also carried out in the context of the harmonic oscillator with random frequency [30, 31].

In this appendix, we study analytically the behavior of the generalized Lyapunov exponents in the limit $n \rightarrow \infty$. For the special case of $\omega = 0$, we prove the following asymptotic formula:

$$\lambda_{\max}(n) \simeq \frac{3}{4} D^{1/3} (2n)^{4/3}. \quad (\text{A3})$$

We shall then argue that this behavior remains valid for any finite value of ω .

We now outline the proof of the scaling equation (A3) for $\omega = 0$ by studying the large n behavior of the coefficients of the characteristic polynomial $P(X)$ of the matrix A_n . Recalling that if $\lambda_{n,1}, \dots, \lambda_{n,2n+1}$ are the eigenvalues of A_n we have

$$P(X) = \prod_{i=1}^{2n+1} (X - \lambda_{n,i}) = X^{2n+1} - \left(\sum_{i=1}^{2n+1} \lambda_{n,i} \right) X^{2n} + \left(\sum_{i \neq j} \lambda_{n,i} \lambda_{n,j} \right) X^{2n-1} + \dots + \prod_{i=1}^{2n+1} \lambda_{n,i}. \quad (\text{A4})$$

We also have $\lambda(n) = \max_l (\text{Re } \lambda_{n,l})$. The coefficients of the characteristic polynomial $P(X)$ are symmetric functions of the eigenvalues $\lambda_{n,i}$. Thanks to the Newton formulae, all these coefficients can be written as linear combinations of traces of powers of A_n . Hence, we have

$$P(X) = X^{2n+1} - \text{Tr}(A_n) X^{2n} + \frac{1}{2} \left([\text{Tr}(A_n)]^2 - \text{Tr}(A_n^2) \right) X^{2n-1} + \dots \quad (\text{A5})$$

In principle, we can obtain the eigenvalues from the traces of the $2n + 1$ powers of A_n . In practice, since we are interested only in the asymptotic behavior of $\lambda_{\max}(n)$, it will be computed from the traces of high powers. From equation (11) we observe that $\text{Tr}(A_n) = 0$. In the case $\omega = 0$, we also have $\text{Tr}(A_n^2) = 0$. More generally, for $\omega = 0$, we can show that $\text{Tr}(A_n^k) \neq 0$ only when k is a multiple of 3. Indeed, writing

$$A_n = d + g \quad (\text{A6})$$

with

$$d = \begin{pmatrix} 0 & 2n & 0 & & & \\ 0 & 0 & 2n-1 & & & \\ & & & \ddots & & \\ & & & & 0 & 1 \\ & & & & & 0 \end{pmatrix} \quad \text{and} \quad g = D \begin{pmatrix} 0 & & & & & \\ 0 & 0 & & & & \\ 2 & 0 & 0 & & & \\ & 6 & 0 & 0 & & \\ & & 12 & 0 & 0 & \\ & & & \ddots & \ddots & \ddots \\ & & & & 2n(2n-1) & 0 & 0 \end{pmatrix} \quad (\text{A7})$$

we obtain

$$(d+g)^k = \sum_{k_1+k_2=k} W_{k_1,k_2} \quad (\text{A8})$$

where W_{k_1,k_2} is a product of k factors with k_1 factors equal to d and k_2 factors equal to g , with $k_1+k_2=k$ (there are 2^k such terms because the matrices d and g do not commute). We remark that d is an upper-diagonal band matrix and its non vanishing terms are all on the band located at level +1 above the diagonal. Similarly g is a lower-triangular band matrix and its non vanishing terms are all on the band located 2 levels below the diagonal. Therefore a product of k_1 matrices d and k_2 matrices g will have a non-zero diagonal term only if $k_1 = 2k_2$ i.e. if $k = 3k_2$; hence, k is a multiple of 3. For example, we have

$$\text{Tr}(A_n^3) = \text{Tr}(ddg + dgd + gdd) = 3\text{Tr}(d^2g) = 3D \sum_{l=1}^{2n-1} l(l+1)(2n-l+1)(2n-l) \simeq 3D(2n)^5 \int_0^1 x^2(1-x)^2 dx = D \frac{16n^5}{5}. \quad (\text{A9})$$

More generally, the terms that contribute to $\text{Tr}(A_n^{3k})$ are obtained by taking the product of $2k$ factors d and k factors g written in all possible orders (there are $\frac{(3k)!}{(2k)!k!}$ such terms):

$$\text{Tr}(A_n^{3k}) = \sum \text{Tr}(W_{2k,k}) = d^{2k}g^k + \text{'permuted terms'} \quad (\text{A10})$$

In particular, we have

$$\begin{aligned} \text{Tr}(d^{2k}g^k) &= D^k \sum_l l(l+1)(l+2)\dots(l+2k-1)(2n-l+1)(2n-l)\dots(2n-l-2k+2) \\ &\simeq D^k (2n)^{4k+1} \int_0^1 x^{2k}(1-x)^{2k} dx = D^k (2n)^{4k+1} \frac{(2k)!(2k)!}{(4k+1)!}. \end{aligned} \quad (\text{A11})$$

The trace of any term $W_{2k,k}$ is given by the same expression at the leading order: indeed, the elements of the matrix dg are of order n^3 whereas those of the commutator $[d,g]$ are of order n^2 . The reason is that both dg and gd are triangular with all nonvanishing matrix elements one level below the diagonal, of the form $(dg)_{l,l-1}$ and $(gd)_{l,l-1}$. In the center of the matrix, $l \approx n$, a generic term is of the order of $n^3 + O(n^2)$. Therefore both dg and gd are dominated by n^3 , while $[d,g]$, given by the difference of two such terms is dominated by the $O(n^2)$ corrections. We know that any $W_{2k,k}$ differs from $d^{2k}g^k$ by a finite number of commutators. Therefore, $\text{Tr}(W_{2k,k}) = \text{Tr}(d^{2k}g^k) +$ subleading terms. We thus have

$$\text{Tr}(A_n^{3k}) \simeq \frac{D^k (2n)^{4k+1}}{(4k+1)} \frac{(3k)!(2k)!}{(4k)!k!}. \quad (\text{A12})$$

Rewriting this trace in terms of the eigenvalues of the matrix A_n , we obtain

$$\sum_{i=1}^{2n+1} \lambda_{n,i}^{3k} \simeq \frac{D^k (2n)^{4k+1}}{(4k+1)} \frac{(3k)!(2k)!}{(4k)!k!}. \quad (\text{A13})$$

Besides, from equation (A12), we deduce that

$$\frac{[\text{Tr}(A_n^3)]^k}{\text{Tr}(A_n^{3k})} \sim n^{k-1}. \quad (\text{A14})$$

$2n$	$\lambda_{\max}(n)$	$\frac{3D(2n)^{4/3}}{4\lambda_{\max}(n)}$
10	15	1.077
40	101	1.015
80	257	1.006
100	347	1.003
120	443	1.0021
140	544	1.0022
160	650	1.0018
180	761	1.0016
200	876	1.0013
300	1504.7	1.0010

TABLE I: Behavior of the dominant eigenvalue in the case $\omega = 0$, $D = 1$.

$2n$	$\lambda_{\max}(n)$	$\frac{3D(2n)^{4/3}}{4\lambda_{\max}(n)}$
10	13.2	1.223
40	95.9	1.072
80	248	1.043
100	336	1.035
120	431	1.031
140	530	1.028
160	635	1.026
180	745	1.023

TABLE II: Behavior of the dominant eigenvalue in the case $\omega = 1$, $D = 1$.

This equation shows that the trace is not dominated by the largest eigenvalue alone: otherwise this ‘participation’ ratio would be of order one. Rather, a finite fraction ρ of the eigenvalues has a scaling behavior similar to that of $\lambda_{\max}(n)$ and we can write: $\text{Tr}(A_n^{3k}) \sim \rho n \lambda_{\max}(n)^{3k}$. Thus, we deduce from equation (A13) the following behavior

$$\rho n \lambda_{\max}(n)^{3k} \simeq \frac{D^k (2n)^{4k+1} (3k)!(2k)!}{(4k+1) (4k)!k!}. \quad (\text{A15})$$

Finally, using Stirling formula, we obtain for large k , assuming ρn does not vary strongly with n

$$\lambda_{\max}(n) \simeq D^{1/3} (2n)^{4/3} \left(\frac{(3k)!(2k)!}{(4k)!k!} \right)^{1/3k} \simeq \frac{3}{4} D^{1/3} (2n)^{4/3}. \quad (\text{A16})$$

This ends the proof of equation (A3) or (12) in the case $\omega = 0$. In Table I the highest eigenvalue of the matrix A_n is computed numerically for various values of $2n$ up to $2n = 300$ (for $\omega = 0$ and $D = 1$). The asymptotic scaling given in equation (A3) is well satisfied.

When ω is different from 0, we can still write A_n as a sum of two matrices as in equation (A6): the upper-diagonal matrix d remains the same as in equation (A7) but the lower-triangular part g' is now given by the previous g plus a band-diagonal matrix $g_2(\omega)$ (with a band-diagonal at level -1) which contains terms proportional to ω . However, in the large n limit (and keeping the value of ω fixed) the matrix elements of g are much bigger than those of $g_2(\omega)$: hence $g' \simeq g$ up to subdominant contributions and the large n scaling of the maximal eigenvalue is insensitive to ω at leading order, therefore equation (A3), derived for $\omega = 0$, remains true for finite values of ω . In Tables II and III we give numerical results $\omega = 1$ and $\omega = -1$ (taking $D = 1$ for both cases). The scaling behavior, proportional to $(2n)^{4/3}$, is well satisfied and it seems also that the prefactor $3/4$ remains correct.

$2n$	$\lambda_{\max}(n)$	$\frac{3D(2n)^{4/3}}{4\lambda_{\max}(n)}$
10	18	0.900
40	107.5	0.952
80	267	0.968
100	358	0.972
120	455	0.975
140	557.5	0.977
160	666	0.980
180	777	0.981

TABLE III: Behavior of the dominant eigenvalue in the case $\omega = -1$, $D = 1$.

APPENDIX B: BOUNDNESS OF THE \bar{c}_n

Let us assume an expansion (of the initial moments (16))

$$M_{n,l} = \sum_{m=0}^{\infty} b_{n,l}^{(n+m)} \epsilon^{2(n+m)} \quad (\text{B1})$$

for the l -th moment of order n as defined after the matrix (11). The b 's are bound (by construction). Let us expand the $c_{n,l}$ of Eq. (16) in powers of ϵ as

$$c_{n,l} = \sum_{m'} \bar{c}_{n,l}^{(m')} \epsilon^{2m'}. \quad (\text{B2})$$

Since only even moments are considered, the powers of ϵ are even. Now we write Eq. (16) in the form

$$M_{n,l} = \sum_{n',l'} c_{n',l'} U_{n',l'}^{(n,l)}, \quad (\text{B3})$$

where we equate the components of the vectors. now we use the property (14) of the eigenvectors $U_{n',l'}^{(n,l)}$, namely $U_{n',l'}^{(n,l)} = 0$ for $n > n'$, and solve the equation order by order in ϵ .

Take first $n = 1$, and leading order in ϵ

$$b_{1,l}^{(1)} = \sum_{l'} \bar{c}_{1,l'}^{(1)} U_{1,l'}^{(1,l)} \quad l = 0, 1, 2. \quad (\text{B4})$$

These are three linear equations for the $\bar{c}_{1,l'}$, since the eigenvectors $U_{1,l'}^{(1,l)}$ are given.

Next take $n = 2$,

$$b_{1,l}^{(2)} = \sum_{l'=0}^2 \bar{c}_{1,l'}^{(2)} U_{1,l'}^{(2,l)} + \sum_{l'=0}^4 \bar{c}_{2,l'}^{(2)} U_{2,l'}^{(2,l)} \quad (\text{B5})$$

$$b_{2,l}^{(2)} = \sum_{l'=0}^2 \bar{c}_{1,l'}^{(2)} U_{1,l'}^{(2,l)} + \sum_{l'=0}^4 \bar{c}_{2,l'}^{(2)} U_{2,l'}^{(2,l)}. \quad (\text{B6})$$

These are eight equations for the $\bar{c}_{2,l'}$, $\bar{c}_{1,l'}$.

This process can be continued to any order in ϵ . The independence of the eigenvectors implies that finite solutions for the $\bar{c}_{n,l}^{(m')}$ can be obtained by Kramer's rule. Note, that $b_{n,l}^{(m')} = 0$ for $n > m'$, therefore, also $\bar{c}_{n,l}^{(m')} = 0$ for $n > m'$. Therefore,

$$c_{n,l} \sim \bar{c}_n \epsilon^{2n} \quad (\text{B7})$$

for small ϵ .

This shows that the \bar{c}_n are bounded for any finite n , but it does not imply the existence of a uniform bound for all n and m .

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