

Inhomogeneous Floquet thermalization

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How a closed system thermalizes, especially in the absence of global conservation laws but in the presence of disorder and interactions, is one of the central questions in non-equilibrium statistical mechanics. We explore this for a disordered, periodically driven Ising chain. Our numerical results reveal inhomogeneous thermalization leading to a distribution of thermalization timescales within a single disordered sample, which we encode via a distribution of effective local temperatures. Using this, we find an excellent collapse *without any fitting parameters* of the local relaxation dynamics for the entire range of disorder values in the ergodic regime when adapting the disorder-averaged diagonal entanglement entropy as internal ‘time’ of the system. This approach evidences a remarkably uniform parametrization of the dynamical many-body evolution of local temperature within the otherwise highly heterogeneous ergodic regime, independent of the strength of the disorder.

Introduction: Spatially heterogeneous relaxation dynamics towards equilibrium is a hallmark of nonergodicity, being found in paradigmatic settings of glasses and jammed systems [1]. Such dynamical heterogeneity, e.g., evidenced in the coexistence of different relaxation timescales, can arise from spatial variations associated with the presence of metastable states. This can lead to global nonexponential decay of correlation functions in time despite local exponential decay rates. Alternatively, relaxation processes can be inherently complex also, contributing to local nonexponential decay [2]. In ergodic systems, such as supercooled liquids, even regions of slower relaxation eventually thermalize [3].

In quantum systems undergoing unitary dynamics, the nature of thermalization, or its absence, has been a focal point of research in recent years [4–6]. Thermalization in ergodic systems occurs, loosely speaking, as sub-system density matrices evolve to a thermal state, with the remainder of the system effectively acting as its bath. The presence of disorder may impede thermalization, for instance, by the emergence of a set of quasi-local integrals of motion, resulting in emergent effective integrability and non-ergodicity, a phenomenon dubbed many-body localization (MBL) [4, 7–12]. In this parameter regime, spatial and dynamical heterogeneity has been observed in local entanglement measures [13, 14] as demonstrated in models such as the XXZ spin chain. Indeed, entanglement can exhibit substantial spatial heterogeneity in this regime, as indicated by the subvolume scaling of the standard deviation of the cut-to-cut entanglement entropy [15]. Because of the difficulty of studying the real-time dynamics of disordered interacting quantum systems, a definite consensus on the nature of this regime has been slow to emerge, e.g., Refs. [16–29].

Here, we show that even the weakly disordered *ergodic* regime can exhibit considerable spatial structure, which we investigate in detail. We focus on a nonintegrable driven disordered Ising chain without global conservation laws. While even weak disorder tends to slow

down relaxation, eventual thermalization can be remarkably robust [16]. We analyze thermalization after a sudden quench via the time evolution of the local subsystem (inverse) temperature, $\beta_j(n)$, where j denotes a bond involving two sites after n the time steps. The $\beta_j(n)$ evolution reveals apparent glassiness in the sense of a distribution of local relaxation timescales within a single disorder configuration. Indeed, the time evolution of the spatial and disorder averaged $\beta_j(n)$ shows well-developed nonexponential behavior in finite-size numerical simulation. With increasing disorder, we further observe mixed dynamics, i.e., locally nonexponential decay of the $\beta_j(n)$ accompanied by regions in space with exponential relaxation in time towards ‘infinite temperature’.

The broad distribution of thermalization times, even at a single disorder value, suggests that relying on a single time scale may not be sufficient. Relatedly, it is challenging to compare the time evolution between different disorder realizations and, further, strengths. Here, we propose a unified description employing the ensemble-averaged diagonal entanglement entropy $S_d(t)$ to monitor the dynamical evolution of $\beta_j(n)$ across the entire ergodic regime. This extends the framework introduced by Evers *et al.* [28] for non-driven systems, which used the average entanglement entropy to track particle density imbalance decay. We find a straightforward data collapse of the evolution of average inverse temperature for several disorder values, offering a homogeneous perspective on thermalization dynamics despite its inherent heterogeneity. While previous work [28] required extra fitting parameters to obtain a scaling collapse, our intrinsic parametrization of time obviates this need.

Driven Ising Model: The time-dependent Hamiltonian of this periodically driven system is defined as,

$$H(t) = \begin{cases} 2H_x & \text{if } 0 < t < \frac{T}{2} \\ 2H_z & \text{if } \frac{T}{2} < t < T, \end{cases}$$

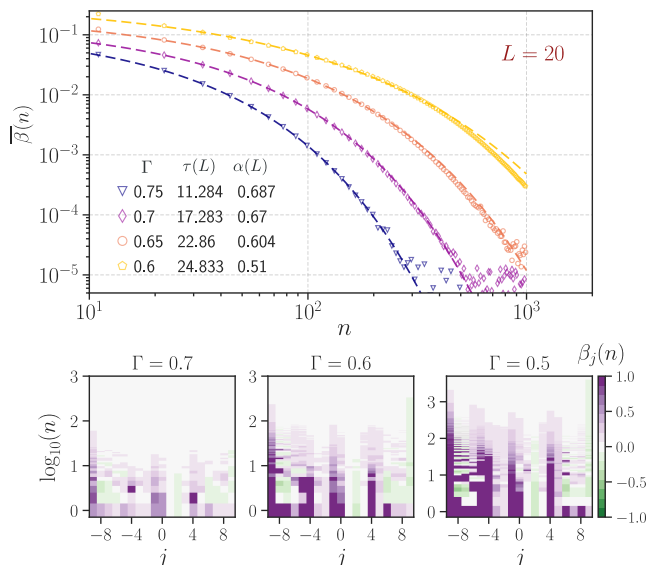


FIG. 1. Stretched exponential decay of the average local temperature $\bar{\beta}(n)$ with Floquet time step n . The line indicates a stretched exponent fit: $A \exp[-(t/\tau)^\alpha]$. The corresponding values of the thermalization time $\tau(L)$ and the stretched exponent $\alpha(L)$ are provided in the legend for $L = 20$. Lower panel: Spatiotemporal inhomogeneity of the evolution of local bond temperature $\beta_j(n)$ as defined in Eq. (2) for different values of $\Gamma = 0.7, 0.6, 0.5$ as a function of $\log(n)$ for a typical disorder configuration. The white coloring represents infinite temperature ($\beta_j=0$).

$$H_x = \sum_{i=1}^L g \Gamma \sigma_i^x,$$

$$H_z = J \sum_{i=1}^{L-1} \sigma_i^z \sigma_{i+1}^z + \sum_{i=1}^L (h + g \sqrt{1 - \Gamma^2 G_i}) \sigma_i^z, \quad (1)$$

σ_i^x and σ_i^z are the Pauli matrices on site i . We follow the standard parametrization of the model, see Refs. [30–32]. The interaction strength $J = 1$ and g, h and T are 0.9045, 0.8090 and 0.8, respectively. Such a parametrization is motivated by the clean static model, where strong thermalization is obtained with these parameter values for system sizes readily accessible in exact diagonalization studies [33]. The longitudinal field is disordered and chosen from a Gaussian distributed random variable G_i with zero mean and unit variance. The Γ controls the disorder strength, and the model is believed to have an MBL transition at $\Gamma \simeq 0.3$ [30, 31]. The stroboscopic time evolution is performed with Floquet operator $U_F(T) = e^{-iH_x T/2} e^{-iH_z T/2}$ using a Hadamard transformation, see Ref. [31] for further details. The initial state for these calculations is the Néel state $|1010\dots\rangle$.

Local temperature: The local temperature $\beta_j(n)$ is defined for each bond by minimizing the Frobenius norm distance between the canonical density matrix $e^{-\beta_j \mathcal{H}_b}$,

and the sub-system density matrix ρ_j^b ,

$$\min_{\beta_j} |e^{-\beta_j \mathcal{H}_b} - \rho_j^b(n)|, \quad (2)$$

with $\mathcal{H}_b(j) = \sigma_j^z \sigma_{j+1}^z + h_j/2 \sigma_j^z \otimes \mathbb{1}_{2 \times 2} + h_{j+1}/2 \mathbb{1}_{2 \times 2} \otimes \sigma_j^z + \{z \leftrightarrow x, h_j \leftrightarrow g_j\}$ is the local bond Hamiltonian with h_j, g_j being the local fields, and $\rho_j^b(n) = \text{Tr}_{L-\{j, j+1\}} |\psi(n)\rangle \langle \psi(n)|$ is the reduced density matrix for that bond. The norm is defined as $|A| \equiv \text{tr}(\sqrt{A^\dagger A})$. We expand the definition of β from Ref. [34] to the time domain. This approach includes the time-evolved wavefunction's structure through the sub-system density matrix. It is important to note that the precise value of $\beta_j(n)$ is contingent upon the chosen definition of norm, as examined in detail in Ref. [34].

(Non-)Exponential Heating: An interacting driven system heats to an infinite temperature featureless state [35, 36]. The heating rate τ generally depends exponentially on the drive frequency, $\tau \propto \exp(\omega/J)$ for $\omega/J \gg 1$, where J is a microscopic energy scale [37–40]. In the presence of disorder, it has been seen that the (average) heating slows down considerably [31, 41–44]. As shown in Fig. 1, the time evolution of the disorder and spatially averaged (indicated by the overline) $\bar{\beta}(n)$ indeed exhibits a stretched exponential decay, with the decay exponent inversely correlated to the strength of the disorder, similar to the decay of the correlation function as reported earlier in Refs. [27, 31, 45]. A pronounced finite size effect is also observed and shown in the Appendix.

The lower panel of Fig. 1 highlights the spatially inhomogeneous evolution of $\beta_j(n)$ for a given disorder configuration for several Γ values. The upper panel of Fig. 2 shows $\beta_j(n)$ at different locations for a single disorder configuration. For weak disorder, $\Gamma = 0.7$, thermalization is exponential everywhere, e^{-n/τ_j} , however with spatially varying τ_j . The $|\beta_j(n)|$ decays to $\sim 10^{-3}$ for $L = 20$, and is expected to vanish in the thermodynamic limit. With increasing disorder, the variation in τ_j increases, and for even stronger disorder $\Gamma = 0.45$, the heating becomes slow, possibly as a stretched exponential (see Fig. 2(c)). Even there, a few subsystems still exhibit fast thermalization, i.e., an exponential decay of $|\beta_j(n)|$ in time, n . This distribution of thermalization time scale is reflected in a stretched exponential decay, $e^{-(t/\tau)^\alpha} = \int du P(u) e^{-(t/u)}$, of the average $\bar{\beta}(n)$ as highlighted in Fig. 1 with a timescale that increases, and an exponent α which decreases, with increasing disorder.

Thermalization time: Fitting each $\beta_j(n)$ trace is impractical because of fluctuations in the data and uncertainties associated with the fit. We instead extract a local decay time τ_j via

$$\tau_j := \frac{\int_0^T n |\beta_j(n)| dn}{\int_0^T |\beta_j(n)| dn}. \quad (3)$$

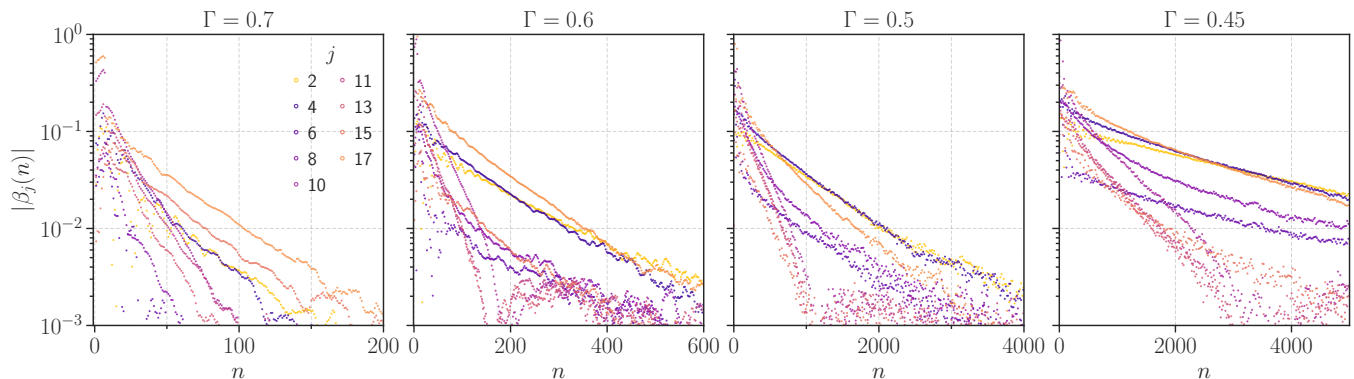


FIG. 2. Time evolution of local inverse temperature $|\beta_j(n)|$ at different locations, j , for several disorder values $\Gamma = 0.6, 0.5, 0.4, 0.45$ for $L = 20$. Log-linear scale highlights exponential decay at large Γ for a typical disorder configuration. With decreasing Γ , the thermalization time increases. For the small value of $\Gamma = 0.45$, the curves appear to fall into two groups characterized by rapid and slow decay rates.

Concretely, the pure exponential $|\beta_j(n)| \propto e^{-n/\tau'_j}$ yields $\tau_j = \tau'_j + T/(1 + e^{T/\tau'_j}) \rightarrow \tau'_j$ for sufficiently large number of Floquet time steps T . Indeed, for the larger $\Gamma = 0.7$, the individual samples show exponential decay within the simulation time window. With decreasing $\Gamma < 0.5$, not all the traces of $\beta_j(n)$ show pure exponential decay; instead, there is mixing of both exponential and stretched exponential decays (Fig. 2, right panel). This affects the calculation of the decay time as defined in Eq. (3), and we refrain from doing this analysis for smaller values of the Γ , i.e., larger disorder values close to the putative MBL transition. In this regime, the simple way of describing the exponential heating dynamics using the Fermi-Golden rule [40, 46, 47] is probably inapplicable, and one might need to go beyond this perturbative treatment.

Figure 3 shows the τ_j distribution in the ergodic phase for $\Gamma = 0.7, 0.6, 0.5$. A pronounced exponential tail is observed for all the disorder values, with a plateau forming at smaller Γ , revealing a broad distribution of timescales. At $\Gamma = 0.5$, significant finite size and time effects are evident (see also Appendix), suggesting the need for larger

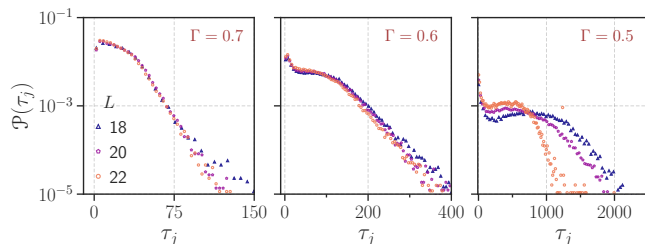


FIG. 3. Probability distribution of the local thermalization time τ_j for system sizes $L = 18, 20, 22$, and disorder values $\Gamma = 0.7, 0.6, 0.5$. The distribution is broad with an exponential tail. For smaller $\Gamma = 0.5$, the distribution shows strong finite-size effects. Typically, $\sim 2.5 \times 10^3$ disorder configurations are used for these distributions.

system sizes $L \gtrsim 22$ to observe thermalization in all parts of the sample.

For the largest system $L = 22$, the $\mathcal{P}(\tau_j)$ shifts towards the origin compared to smaller system sizes, indicating a slow flow of the full distribution towards shorter thermalization time.

Density matrix evolution: Having established a relatively broad distribution of τ_j , resulting in a stretched exponential decay of the spatially averaged $\beta_j(n)$ (Fig. 1), we now provide an entanglement perspective to the thermalization process. Figure 4 shows a typical evolution of the half-system density matrix $|\rho_{nn'}|/\max(\rho_{nn'})$ for different evolution times, n , for $\Gamma = 0.7$, which becomes thermal at these times. This behavior is representative of all samples reaching thermal equilibrium.

Diagonal entanglement entropy: The entanglement entropy is defined as, $\mathcal{S}_E = -\text{Tr}(\rho^A \log \rho^A)$, where ρ^A is the reduced half-system density matrix. At long time, for the ergodic system, \mathcal{S}_E reaches the Page value i.e., $\mathcal{S}_E = L/2 \log(2) - 1/2$ [48]. In this regime, the diagonal elements ρ_{kk}^A dominate; these scale as $\rho_{kk}^A \propto 1/\sqrt{\mathcal{D}}$, while the off-diagonal terms are suppressed as $\sqrt{\mathcal{D}^{1/2}}$ [49], where $\mathcal{D} = 2^L$ is the Fock space dimension. The diagonal entropy in this basis is

$$\mathcal{S}_d = - \sum_k \rho_{kk}^A \log \rho_{kk}^A. \quad (4)$$

For a Haar random state, $\rho_{kk}^A \propto 2^{-L/2}$, therefore $\mathcal{S}_d \propto (L/2) \log(2)$, the volume law scaling for ergodic systems. Further, the diagonal entropy can be expressed in terms of the participation entropy $\mathcal{S}_P = 2\mathcal{S}_d$ (see Appendix for details), where $\mathcal{S}_P(n) = -\sum_{j=1}^{\mathcal{D}} p_j(n) \log p_j(n)$, with $p_j(n) = |\langle j | \psi(n) \rangle|^2$ is the probability of occupation of each spin basis state $|j\rangle$, and $|\psi(n)\rangle$ is the time evolved wavefunction.

Therefore, \mathcal{S}_d is an alternative measure of delocalization in configuration space. Indeed, Ref. [50] showed that

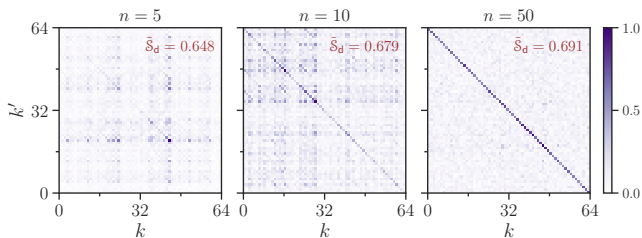


FIG. 4. Time evolution of the half-system density matrix $\rho_{kk'}^A$ (normalized such that its maximum value is unity) for $\Gamma = 0.7$, $L = 12$ and for a typical disorder configuration. As expected for a featureless random state, the diagonal elements dominate over time as the system heats up. Additionally, the diagonal entropy reaches maximum $\bar{S}_d = S_d/(L/2) \simeq \log(2)$, as indicated in the figure.

for a pure state with few random non-zero elements relative to the dimension of the space, the scaling of S_d with subsystem volume $L/2$ is exactly given by that of its participation entropy S_P .

Synchronized dynamics: The upper panel of Fig. 5 shows the time evolution of the disorder averaged $x(n) = -\log \beta_{\text{typ}}(n)$ with Floquet cycle n for several values of disorder $\Gamma = 0.7 \dots 0.3$. The typ. denotes the median value of $\beta_j(n)$ across the bonds j . For intermediate disorder values, $x(n)$ shows logarithmically slow propagation (see, e.g., $\Gamma = 0.45$); however, with increasing time, it accelerates as seen by the leftward bending of the curve to eventual thermalization. Such bending happens at progressively higher n with increasing disorder, and to observe this at an even larger disorder, larger L is necessary.

Most strikingly, when S_d is adapted as an ensemble average internal time of the system, we observe an excellent collapse of the mean $x(n)$ for several disorder values as seen in Fig. 5. This collapse, requiring no fitting parameters, implies that the diagonal entropy S_d faithfully describes central aspects of the thermalization of the closed system, such as the time to heat up to infinite temperature at finite disorder values. Concretely, disorder slows down entropy production, thus delaying thermalization. Once the simulation time is parametrized by the entropy itself, the universal nature of the dynamics is revealed.

Conclusions: For the thermalization of a disordered interacting Floquet system, we analyze the dynamics of sub-system temperature $\beta_j(n)$ in the ergodic regime. Generically, some blocks heat up faster than others, but all blocks eventually thermalize, leading to a broad distribution of thermalization times. Blocks with long thermalization times are not particularly *rare* and may exhibit either long exponential decay time constants or even nonexponential decay. This distribution of time scales leads to a slower decay of the spatial and disordered average inverse temperature $\beta_j(n)$, resulting in nonexponential heating over time in the ergodic phase, resembling

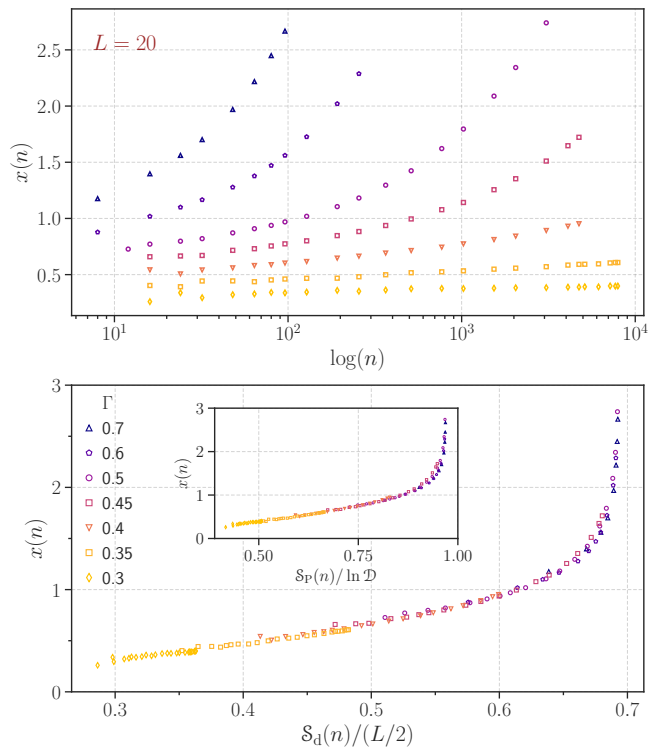


FIG. 5. Time evolution of the disorder averaged typical inverse local temperature $x(n) = -\log \beta_{\text{typ}}(n)$ for several disorder values until the putative transition point $\Gamma_c \simeq 0.3$ and for $L = 20$. The initial slow growth gives way to faster thermalization at later n . This feature has yet to manifest for smaller Γ . Below: The collapse of the evolution when diagonal entropy S_d is adapted as an implicit measure of time, n , for the above figure. Inset: The same data, plotted against the S_P , normalized with the Fock space dimension \mathcal{D} .

the relaxation dynamics of classical glassy liquids. Exploring the connection between inhomogeneous thermalization dynamics and the avalanche mechanism or many-body resonances [25–27, 51–56], which predict stretched exponential decay of correlation functions, is an obvious avenue for further study.

Identifying diagonal entropy S_d as an internal system time allows a collapse of the thermalization dynamics across all the disorder values in the ergodic phase, revealing a remarkable, albeit hidden, homogeneity. The thermalization process involves shrinking off-diagonal matrix elements of the density matrix. In this sense, the approximation of time with diagonal entropy S_d measures the heating rate along with the Fock space delocalization.

The prediction of any dynamical exponent is challenging due to limits imposed by finite time and system sizes, as copiously noted in Hamiltonian models [16, 17, 28, 57–60], disordered Floquet models [27, 31], and even in clean models, where $L \gtrsim 22$ is often necessary to observe heating towards infinite temperature [47]. Indeed, there is substantial variation in stretched exponents with increas-

ing system sizes, particularly evident in two-point correlators, and their fate even in the ergodic phase in the asymptotic limit [31, 45] is at this point unclear. We note that the data collapse we observe is largely independent of system size and thus appears less afflicted by finite-size effects.

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APPENDIX

Finite L dependence

Figure 6 upper panel shows the system size dependence of the decay of the disorder averaged inverse local temperature $\bar{\beta}(n)$. With decreasing Γ , i.e., increasing the strength of the disorder, the L -dependence is more severe for finite time simulations. The system thermalizes quickly for larger $\Gamma = 0.7$, and the data shows less finite-size corrections.

Estimation of decay time τ_j

Figure 6 lower panel shows the decay of $\beta_j(n)$ for two typical disorder configurations for different values of disorder strength, and $L = 20$. The dashed line represents the estimated curve using the decay time τ_j defined in Eq. (3). The estimate of τ_j reasonably reproduces the decay for pure exponential traces. When $\beta_j(n)$ shows non-exponential decay for a stronger disorder, the τ_j gives only a rough estimate as visible at smaller $\Gamma \gtrsim 0.5$.

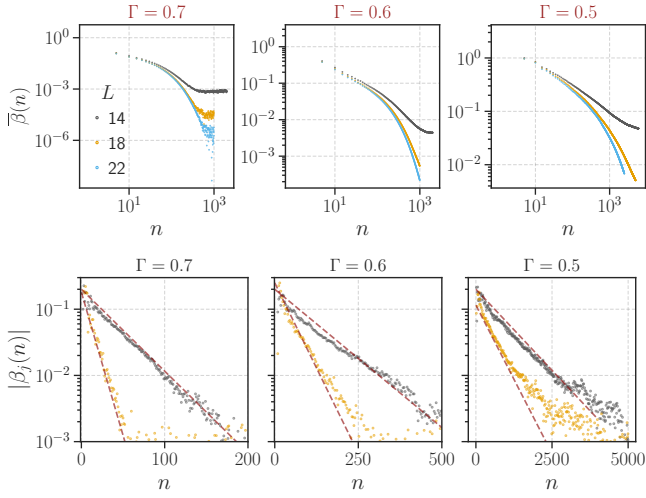


FIG. 6. Upper panel: L -dependence of the decay of $\beta(n)$ with n for different values of $\Gamma = 0.7, 0.6, 0.5$. Lower panel: Decay of $\beta_j(n)$ for two typical disorder configurations. The dashed lines show the curve corresponding to the estimated decay times defined in Eq. (3).

Proof of $\mathcal{S}_P = 2\mathcal{S}_d$

Here, we establish the relation between the Fock space (FS) participation entropy \mathcal{S}_P and the diagonal entropy in the FS basis defined in Eq. (4).

$$\mathcal{S}_P = - \sum_{i=1}^{2^L} |c_i(t)|^2 \log(|c_i(t)|^2)$$

from the definition of the participation entropy. Now, the probabilities, $|c_i(t)|^2$ can be written in terms of the diagonal elements of the density matrix of the system,

ρ in the following way. We note, $|c_i(t)|^2 = c_i(t)c_i^*(t) = \langle i|\psi\rangle\langle\psi|i\rangle = \langle i|\rho|i\rangle$, where $\rho = |\psi\rangle\langle\psi|$ is the system's density matrix.

$$\mathcal{S}_P = - \sum_{i=1}^{2^L} \rho_{ii} \log \rho_{ii}$$

Here, we can decompose the i -th diagonal element of the full density matrix, i.e., ρ_{ii} , as a product of diagonal elements of the sub-system density matrices.

$$\begin{aligned} \rho_{ii} &= \langle i|\rho_A \otimes \rho_B|i\rangle \\ &= \langle j|\langle k|\rho_A \otimes \rho_B|j\rangle|k\rangle \\ &= \rho_{jj}^A \rho_{kk}^B \end{aligned} \quad (5)$$

here, $|i\rangle = |j\rangle|k\rangle$. Therefore,

$$\begin{aligned} \mathcal{S}_P &= - \sum_{i=1}^{2^L} \rho_{ii} \log \rho_{ii} \\ &= \sum_{j=1}^{2^{L/2}} \sum_{k=1}^{2^{L/2}} \rho_{jj}^A \rho_{kk}^B \log(\rho_{jj}^A \rho_{kk}^B) \quad [\text{using Eq. 5}] \\ &= \sum_{j=1}^{2^{L/2}} \sum_{k=1}^{2^{L/2}} \rho_{jj}^A \rho_{kk}^B (\log(\rho_{jj}^A) + \log(\rho_{kk}^B)) \\ &= \sum_{k=1}^{2^{L/2}} \rho_{kk}^B \sum_{j=1}^{2^{L/2}} \rho_{jj}^A \log(\rho_{jj}^A) + \sum_{j=1}^{2^{L/2}} \rho_{jj}^A \sum_{k=1}^{2^{L/2}} \rho_{kk}^B \log(\rho_{kk}^B) \\ &= 2 \sum_{j=1}^{2^{L/2}} \rho_{jj}^A \log(\rho_{jj}^A) \\ &= 2\mathcal{S}_d \end{aligned} \quad (6)$$