

Fermionic transport through a driven quantum point contact: Breakdown of Floquet thermalization controlled by the driving frequency

Ivan V. Dudinets 

*Russian Quantum Center, 30 Bolshoy Boulevard, Building 1, Skolkovo Innovation Center Territory, Moscow 121205, Russia
and Moscow Institute of Physics and Technology, Institutskii per. 9, Dolgoprudnyi 141700, Russia*

Oleg Lychkovskiy 

*Skolkovo Institute of Science and Technology, the Territory of the Skolkovo Innovation Center, Bolshoy Boulevard, 30, p.1, Moscow 121205, Russia;
Steklov Mathematical Institute of Russian Academy of Sciences, Gubkina Street, 8, Moscow 119991, Russia;
and Russian Quantum Center, 30 Bolshoy Boulevard, Building 1, Skolkovo Innovation Center Territory, Moscow 121205, Russia*



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We study a quantum system that consists of two fermionic chains coupled by a driven quantum point contact (QPC). The QPC contains a bond with a periodically varying tunneling amplitude. Initially, the left chain is packed with fermions, while the right one is empty. We numerically track the evolution of the system and demonstrate that, at frequencies above a certain threshold, the current through the QPC halts, and the particle imbalance between the chains remains forever. This implies a spectacular breakdown of the Floquet version of the eigenstate thermalization hypothesis that predicts a homogeneous particle density profile at large times. We confirm the effect for various driving protocols and interparticle interactions.

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I. INTRODUCTION

Leveraging a rich phenomenology of quantum transport is essential for the advancement of science and technology at nano- and microscales. The unceasing theoretical efforts in understanding this phenomenology are of crucial importance, as they provide the foundation to interpret experimental results, predict unexplored phenomena, and guide the design of next-generation quantum and microelectronic devices [1–5].

Here, we contribute to these efforts by studying a system of two tight-binding fermionic chains connected by a quantum point contact (QPC) with a periodically driven tunneling amplitude (see Fig. 1). Initially, the left chain is filled by fermions, while the right one is empty.

This system has been studied previously in the case of noninteracting fermions [6]. It has been shown to exhibit a nonequilibrium phase transition under the variation of the driving frequency: While for low frequencies the QPC is conducting, for frequencies above the single-particle band gap the QPC turns to be insulating, and the initial imbalance of particle densities lasts forever [6,7].

At first sight, the latter behavior should become impossible when the interactions between the fermions come into play. This expectation is based on the Floquet version of the eigen-

state thermalization hypothesis (the Floquet ETH) [10,11], which is believed to hold for generic interacting periodically driven quantum many-body systems [10–16]. The Floquet ETH predicts that a state locally indistinguishable from the infinite-temperature thermal state will establish in the long run, which implies the homogeneous distribution of particles across both chains.

Here, we demonstrate that this expectation is wrong: For sufficiently high frequencies, the particle imbalance between the chains remains forever even when interparticle interactions are on. This implies the breakdown of the Floquet ETH for these frequencies.

Most of our results are obtained numerically by means of exact diagonalization. They are presented in Sec. III, preceded by Sec. II where the model and observables are introduced. We carefully perform a finite-size scaling to demonstrate the validity of our results in the thermodynamic limit. A special case where the Floquet Hamiltonian can be obtained explicitly is studied in Sec. IV. In the final section, we compare our findings to apparently similar phenomena, discuss potential origins of the observed Floquet ETH violation, and outline possible directions for future research.

II. MODEL

The system under study consists of two one-dimensional chains harboring spinless fermions (see Fig. 1). The lengths of the chains are $(L + 1)$ and L . They are chosen to be unequal to avoid degeneracies due to the reflection symmetry. The chains are connected by a driven QPC. The total Hamiltonian reads

$$H_t = H_L + H_R + V_{\text{int}} + V_t. \quad (1)$$

*Contact author: dudinets@phystech.edu

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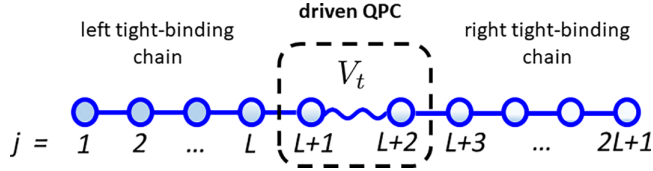


FIG. 1. The system under study consists of two fermionic tight-binding chains coupled by a bond with a periodically driven tunneling amplitude. The bond constitutes the time-dependent quantum point contact.

Here, H_L and H_R describe hopping along left and right chains, respectively,

$$H_L = -\frac{1}{2} \sum_{j=1}^L (c_j^\dagger c_{j+1} + c_{j+1}^\dagger c_j), \quad (2)$$

$$H_R = -\frac{1}{2} \sum_{j=L+2}^{2L} (c_j^\dagger c_{j+1} + c_{j+1}^\dagger c_j), \quad (3)$$

c_j are fermionic annihilation operators, V_{int} is the interparticle interaction term, and V_t is the time-periodic Hamiltonian of the QPC acting exclusively on the pair of boundary sites (see Fig. 1).

To illustrate our findings, in the present section we focus on the specific V_t and V_{int} . Namely, we choose

$$V_t = -\frac{1}{2} f_t (c_{L+1}^\dagger c_{L+2} + c_{L+2}^\dagger c_{L+1}), \quad (4)$$

$$f_t = \sin \omega t, \quad (5)$$

where the periodic function f_t is referred to as driving protocol, ω being its frequency, and

$$V_{\text{int}} = U \sum_{j=1}^{2L} n_j n_{j+1}, \quad (6)$$

where $n_j = c_j^\dagger c_j$ is the particle number operator at the j th site and U is the interaction strength. We take $U = 0.5$ unless specified otherwise.

Other types of driving protocols f_t , QPC terms V_t , and interaction terms V_{int} are considered in Appendixes B, C, and D, respectively. A noninteracting case of $U = 0$ was studied earlier in Ref. [6]. Additionally, in Sec. IV we will study a special instance of QPC where an explicit Floquet Hamiltonian can be found.

The total number of fermions is conserved throughout the evolution. Initially, the fermions fill the left chain, while the right chain is empty. In the main text, we focus on the initial state with the full filling of the left chain:

$$\Psi_0 = \left(\prod_{j=1}^{L+1} c_j^\dagger \right) |\text{vac}\rangle. \quad (7)$$

States away from the full filling are considered in Appendix E.

The state of the system Ψ_t evolves according to the Schrödinger equation $i\partial_t \Psi_t = H_t \Psi_t$.

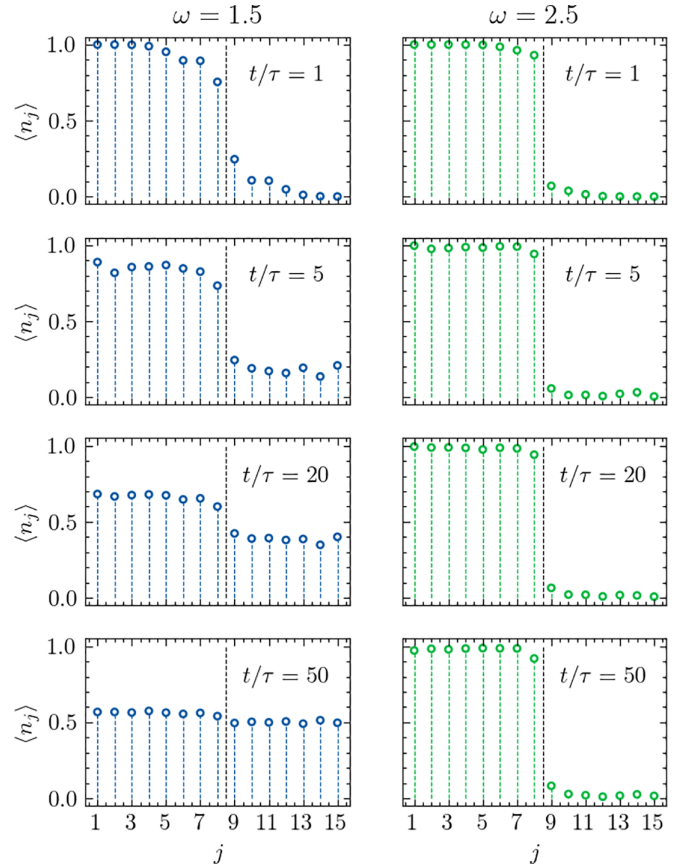


FIG. 2. Evolution of particle density for two driving frequencies. For a smaller frequency (left column), particles tend to distribute uniformly over sites. For a larger frequency (right column), the particle imbalance between the chains is retained. Dashed horizontal line delineates two chains.

We focus on a specific observable, the number of fermions in the right chain:

$$N_R = \sum_{j=L+2}^{2L+1} n_j. \quad (8)$$

Its expectation value $\langle N_R \rangle_t = \langle \Psi_t | N_R | \Psi_t \rangle$ is initially zero and grows as a result of particle flow through the QPC.

At large times, N_R fluctuates around its long-time average $\langle N_R \rangle_\infty \equiv \lim_{t \rightarrow \infty} t^{-1} \int_0^t dt' \langle N_R \rangle_{t'}$. In the absence of degeneracies in the Floquet spectrum, $\langle N_R \rangle_\infty$ can be computed as

$$\langle N_R \rangle_\infty = \sum_{\alpha} |\langle \Phi_{\alpha} | \Psi_0 \rangle|^2 \langle \Phi_{\alpha} | N_R | \Phi_{\alpha} \rangle, \quad (9)$$

where $|\Phi_{\alpha}\rangle$ are the Floquet eigenstates, i.e., eigenstates of the operator of quantum evolution over the driving period $\tau = 2\pi/\omega$.

III. RESULTS

A. Breakdown of thermalization

We use the QuSpin package [17,18] to track real-time dynamics by numerically solving the Schrödinger equation in finite systems of sizes up to $2L + 1 = 15$ (see Fig. 2 for

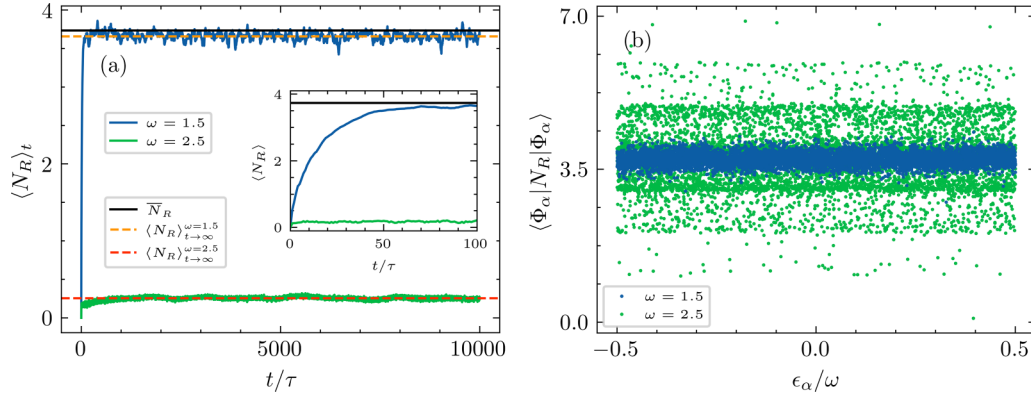


FIG. 3. (a) The number of particles in the right chain, $\langle N_R \rangle_t$, as a function of time for two frequencies. Steady-state values computed according to Eq. (9) are shown by horizontal dashed lines. The horizontal black line shows the value of N_R that corresponds to the uniform distribution of particles over the chains. The interaction strength is $U = 0.5$. The total numbers of sites and fermions are $2L + 1 = 15$ and $N = 8$, respectively. (b) Diagonal matrix elements of the operator N_R in the Floquet eigenbasis.

illustration). We also compute Floquet spectrum and eigenstates for such systems and calculate the long-time average $\langle N_R \rangle_\infty$ according to Eq. (9). The calculations are performed for a range of frequencies.

Let us remind that in the noninteracting case the critical frequency ω_c separating the conducting and insulating behavior of QPC is equal to 2 [6]. In the interacting case, we find a clear conducting behavior of the QPC for frequencies below 1.5, and a clear insulating behavior for frequencies above 2.5. In the present subsection, we focus on these two clear-cut cases. In the next subsection, we discuss a complex, highly system-size-dependent behavior observed for a broad range of intermediate frequencies $1.5 \lesssim \omega \lesssim 2.5$.

Core results are presented in Figs. 3 and 4. For $\omega = 1.5$, the system evolves toward thermalization: Particles flow from the left chain to the right one through the QPC and eventually distribute uniformly across both chains [see Figs. 2 and 3(a)]. The equilibrium value $\langle N_R \rangle_\infty$ coincides, up to small finite-size corrections, with the value

$$\bar{N}_R = NL/(2L + 1) \quad (10)$$

obtained for the uniform distribution of particles over the sites of the system.

In contrast, for $\omega = 2.5$, the QPC effectively turns insulating: The particles stay in the left chain and never come to the right one. In this case, the equilibrium value $\langle N_R \rangle_\infty$ is around zero, as shown in Fig. 3(a) (see also Fig. 2).

In Appendix A, we present the time dependence of the entanglement entropy of one chain. It exhibits behavior qualitatively similar to that of $\langle N_R \rangle_t$: saturates around the thermal equilibrium value at low frequencies and remains low for large frequencies.

The diagonal matrix elements of N_R over the Floquet basis are shown in Fig. 3(b). The patterns for $\omega = 1.5$ and $\omega = 2.5$ are distinctly different. For the smaller frequency, the matrix elements are tightly concentrated around $\bar{N}_R$, consistent with the Floquet ETH [10,11].

The situation is qualitatively different for the larger frequency. As can be seen from Fig. 3(b), in this case matrix elements are scattered over a wide interval, with minimum and maximal instances almost reaching minimal and maximal eigenvalues of N_R , which are 0 and L , respectively. This pattern indicates the breakdown of the Floquet ETH [10,11].

To verify that our results are not artifacts of finite system size but hold in the thermodynamic limit, we perform a finite-size scaling (see Fig. 4). In Fig. 4(a), the scaling

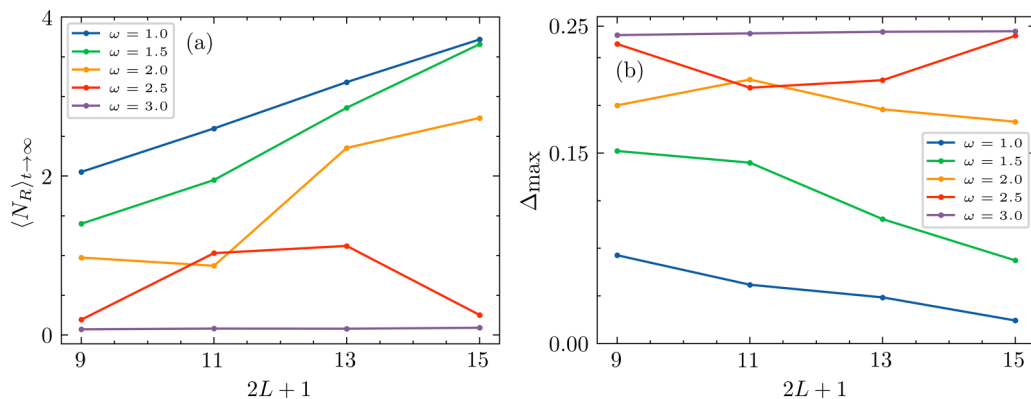


FIG. 4. The scaling with the system size of (a) the steady-state value $\langle N_R \rangle_\infty$ and (b) the maximal deviation of the diagonal matrix element $\langle \Phi_\alpha | N_R | \Phi_\alpha \rangle$ from the uniform value $\bar{N}_R$. The interaction strength is $U = 0.5$. One can see the quantitative difference in the scaling behavior for frequencies below and above the frequency $\omega \simeq 2$.

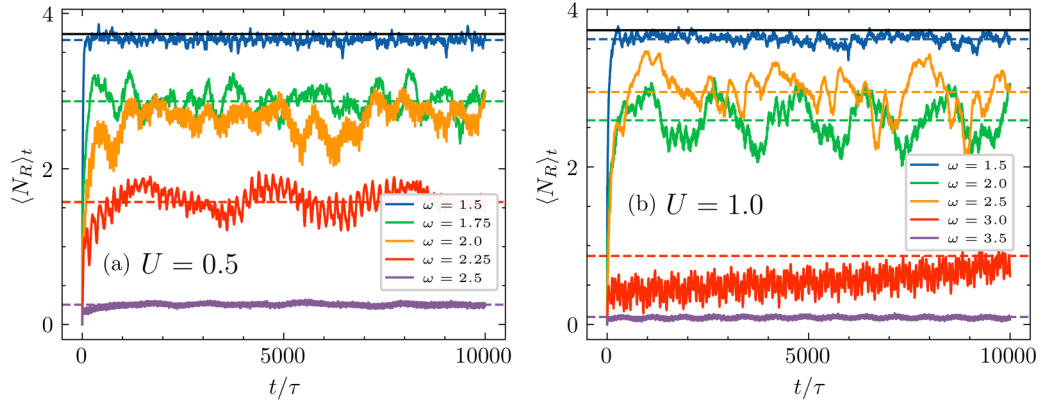


FIG. 5. Illustration of strong finite-size effects at intermediate frequencies. Shown is the number of particles in the right chain, $\langle N_R \rangle_t$, as a function of time for a range of frequencies and two values of interaction strength, (a) $U = 0.5$ and (b) $U = 1.0$.

of the equilibrium value $\langle N_R \rangle_\infty$ is shown. One can see that, depending on the frequency, two types of distinctly different scaling emerge. For smaller frequencies, the equilibrium value scales linearly with the system size, as expected under the assumption of uniform particle distribution in equilibrium. In contrast, for larger frequencies, no growth is observed. The threshold frequency seems to be around $\omega = 2$, which coincides with the critical frequency in the noninteracting case. However, as we discuss in what follows, this observation should be interpreted with care due to complications emerging at intermediate frequencies.

In Fig. 4(b), the Floquet ETH is tested directly. There, we plot the maximal deviation $\Delta_{\max}$ of the diagonal matrix element $\langle \Phi_\alpha | N_R | \Phi_\alpha \rangle$ from the uniform value $\bar{N}_R$ (maximization is performed over all Floquet eigenstates), normalized to the system size:

$$\Delta_{\max} = (2L + 1)^{-1} \max_{\alpha} |\langle \Phi_\alpha | N_R | \Phi_\alpha \rangle - \bar{N}_R|. \quad (11)$$

If the Floquet ETH is valid, $\Delta_{\max}$ should decrease with the system size. This is indeed the case for smaller frequencies, as can be seen from Fig. 4(b). In contrast, for larger frequencies this figure of merit stays approximately constant, indicating the breakdown of the Floquet ETH.

We have verified that the above qualitative picture remains robust under variations of the driving protocol f_t (Appendix B) and tolerates certain on-site time-dependent potentials in QPC (Appendix C). Similarly, the results remain qualitatively unaffected by modifications to the interaction term V_{int} , including next-nearest-neighbor interactions or the suppression of interactions within the QPC (Appendix D). Finally, the physical picture persists for initial states that deviate from full filling of the left chain (Appendix E).

B. Challenge of intermediate frequencies

In the previous subsection, evidence is provided that at sufficiently large frequencies the QPC turns insulating, the uniform particle distribution across the system is never established, and the Floquet ETH is violated. Conversely, at sufficiently small frequencies, the QPC is clearly conducting, fermions distribute uniformly over the system, and no signs of ETH violation are seen. A particularly compelling signature of this frequency-dependent dichotomy is the system-size

scaling of the steady-state particle number in the right chain, $\langle N_R \rangle_\infty$, and the deviation $\Delta_{\max}$ from the ETH prediction, both shown in Fig. 4.

The above dichotomy immediately raises a question about the nature of transition between the two regimes. Is it a sharp dynamical quantum phase transition at some critical frequency ω_c , as in the noninteracting case, or something more complex?

To answer this question, one needs to understand the behavior of the system at intermediate frequencies. This behavior turns out to be quite complex and intricate. In the present subsection, we analyze it and discuss implications that can be derived from this analysis.

Dynamics of $\langle N_R \rangle_t$ at various frequencies is shown in Fig. 5. One can see that across a broad range of intermediate frequencies, the dynamics is highly irregular and exhibits oscillatory features, some with characteristic timescales exceeding the driving period by orders of magnitude. This behavior is highly dependent on the system size, as illustrated in Fig. 19 in Appendix F, and thus is likely to be an artifact of finite size. Note that fluctuations of $\langle N_R \rangle_t$ around $\langle N_R \rangle_\infty$ remain large (i.e., on the order of $\langle N_R \rangle_\infty$ itself) at any time. Similar finite-size effects with high amplitudes and unexpectedly large timescales have also been observed in spin systems [19].

Values of $\langle N_R \rangle_\infty$ computed for intermediate frequencies are also intermediate, lying anywhere between zero and thermal value $\bar{N}_R$ [20].

In the most straightforward scenario, in the thermodynamic limit, there exists a critical frequency ω_c separating two types of dynamics, analogously to the noninteracting case [6]. For frequencies below the critical one, the QPC is conducting and the particle density becomes uniform over the whole system, i.e., $\langle N_R \rangle_\infty = \bar{N}_R$. Conversely, above the critical frequency, the QPC is insulating, the system fails to thermalize and the right chain remains essentially empty, i.e., has zero particle density and $\langle N_R \rangle_\infty = O(1)$ independent of the system size. This scenario is sketched in Fig. 6(a).

In this scenario, intermediate values of $\langle N_R \rangle_\infty$ at intermediate frequencies seen in our finite-size calculations should be interpreted as finite-size fluctuations from either of two limits, $\langle N_R \rangle_\infty = \bar{N}_R$ or $\langle N_R \rangle_\infty \simeq 0$. Given rather moderate

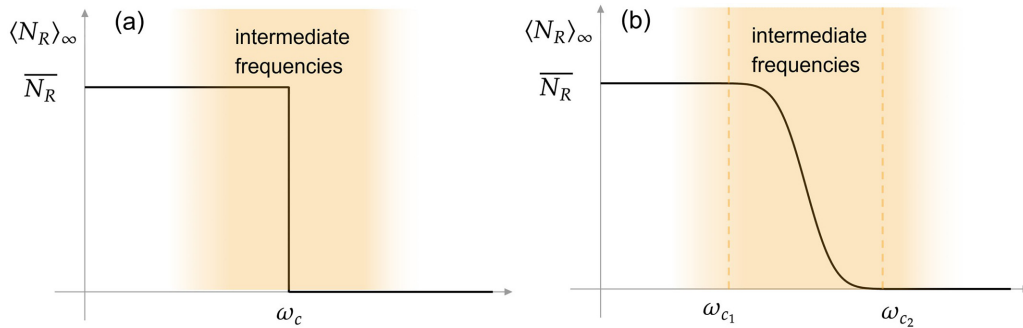


FIG. 6. A sketch of two scenarios for the transition between the completely conducting dynamical phase and the ETH-violating insulating phase. The solid line shows the values of $\langle N_R \rangle_\infty / L$ in the thermodynamic limit. Shaded area marks the range of intermediate frequencies with particularly strong finite-size effects. (a) A sharp dynamical quantum phase transition between a completely uniform particle distribution and the phase with essentially empty right chain. (b) A smooth transition interpolating between the above two extremes within a finite frequency range $\omega_{c1} < \omega < \omega_{c2}$.

system size available within our numerical approach, such interpretation seems reasonable.

If this scenario is realized, it is natural to inquire what the value of the critical frequency is. Unfortunately, due to the broadness of the range of frequencies classified as intermediate, we are unable to determine the value of ω_c with any reasonable precision. In fact, precisely locating the phase transition is a common challenge in studies of finite systems. Such analyses are often capable of revealing two distinct phases with clearly differing properties, but for finite system sizes, these phases are typically separated by a broad crossover region. Only in the thermodynamic limit—inaccessible to exact diagonalization—does this crossover sharpen into a well-defined transition.

Since we cannot precisely identify the critical frequency, we are even less able to determine its dependence on the interaction strength U . For illustration, Fig. 5 compares the dynamics of N_R for $U = 0.5$ and $U = 1$. This comparison reveals a clear trend: The frequencies at which the QPC becomes insulating tend to increase with stronger interparticle repulsion U . This is expected, as stronger fermion repulsion makes it more difficult for the QPC to confine particles and remain insulating. However, for both values of U , the

intermediate-frequency regime remains broad, preventing any reliable quantitative determination of the functional form of $\omega_c(U)$.

Note that our analysis is restricted to moderate interaction strengths, where the system exhibits strong correlations that help suppress finite-size effects. We are unable to meaningfully explore the regimes of very weak or very strong interactions, as proximity to integrable limits further complicates the interpretation of our numerical data.

More complex scenarios can also be consistent with our finite-size numerical results. One possibility is that even in the thermodynamic limit, the steady-state value $\langle N_R \rangle_\infty$ exhibits a smooth crossover from $\bar{N}_R$ to zero as frequency increases between two critical values ω_{c1} and ω_{c2} , as illustrated in Fig. 6(b). In this case, the intermediate values $0 < \langle N_R \rangle_\infty < \bar{N}_R$ obtained in our calculations would represent finite-size approximations of the true thermodynamic behavior.

This scenario suggests that for intermediate frequencies $\omega_{c1} < \omega < \omega_{c2}$, the QPC maintains conductivity until the particle imbalance reaches a certain frequency-dependent value, beyond which it becomes insulating. An example of the QPC whose conductance depends not only on frequency but on particle imbalance is considered in the next section.

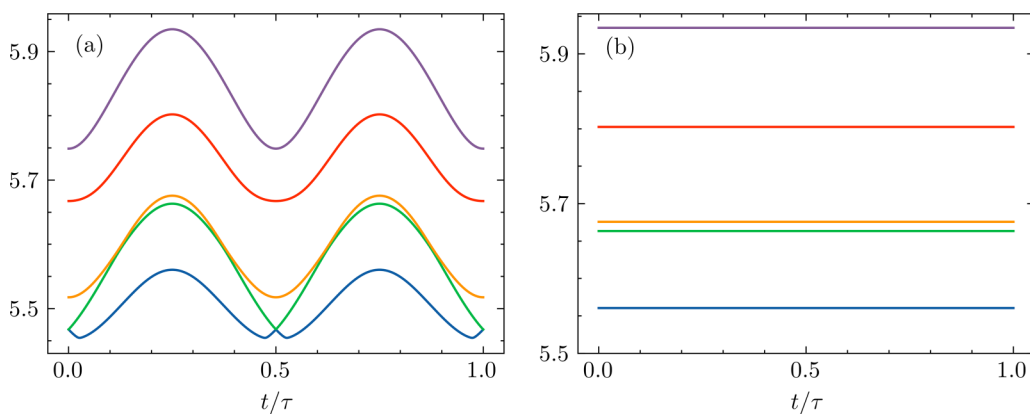


FIG. 7. Instantaneous spectra of the total Hamiltonian for (a) the QPCs (4) and (5) and (b) the isospectral QPC (12) (only lower parts of the spectra are shown). The spectrum in panel (a) shows time dependence, which is a generic behavior. In contrast, the spectrum in panel (b) remains constant throughout the evolution.

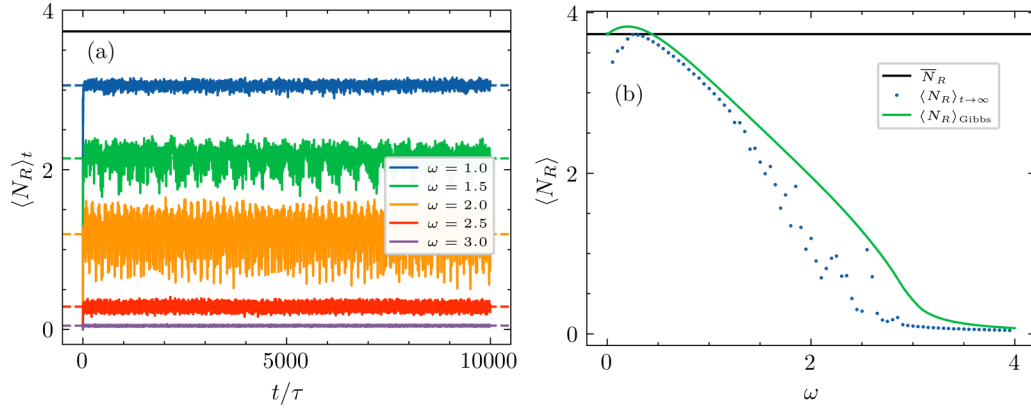


FIG. 8. Results for the isospectral QPC (12) with $2L + 1 = 15$, $N = 8$, and $U = 0.5$. (a) The number of particles in the right chain, $\langle N_R \rangle_t$, as a function of time for a range of frequencies. (b) A comparison between the steady-state value $\langle N_R \rangle_{t \rightarrow \infty}$ and the Gibbs equilibrium value $\langle N_R \rangle_{\text{Gibbs}}$.

To summarize this section, our current numerical data are insufficient to conclusively identify the precise nature of the transition between the conducting dynamical phase and the ETH-violating insulating phase. A complete understanding of this transition remains an important open question for future investigation.

IV. AN ISOSPECTRAL QPC

In this section, we consider a special type of QPC where the Floquet Hamiltonian can be found explicitly. It is given by

$$V_t = -\frac{1}{2}(e^{-i\omega t} c_{L+1}^\dagger c_{L+2} + e^{i\omega t} c_{L+2}^\dagger c_{L+1}). \quad (12)$$

Figure 8(a) shows dynamics of $\langle N_R \rangle_t$ at various frequencies for this QPC.

The QPC (12) is special since the total time-dependent Hamiltonian can be represented as a time-dependent unitary transformation of the time-independent Hamiltonian at fixed time $t = 0$:

$$H_t = e^{i\omega t N_R} H_0 e^{-i\omega t N_R}. \quad (13)$$

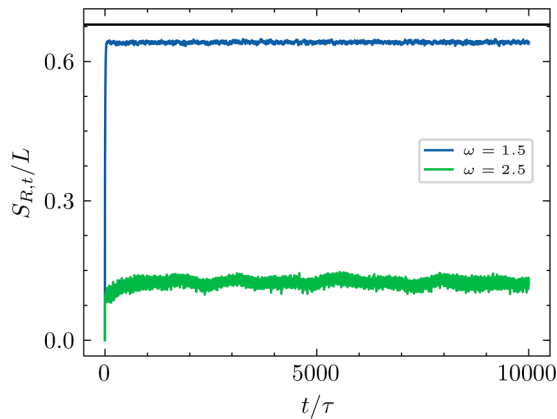


FIG. 9. Entanglement entropy per site of the right chain as a function of time for the QPCs (4) and (5) considered in the main text. System parameters are the same as in Fig. 3. Horizontal solid line shows the thermal value of the entropy per site at infinite temperature.

We refer to such type of time-dependent Hamiltonians as isospectral, since their instantaneous spectra do not vary with time. This is in contrast to the general case, in particular, to the cases considered in the previous section and in Appendixes B–D (see Fig. 7), where the instantaneous spectrum is time-dependent. Isospectral time-dependent Hamiltonians are known to be quite special in various respects (see, e.g., Refs. [21–25]). With some abuse of terminology, we also extend the term “isospectral” to the QPC described by Eq. (12).

As a consequence of Eq. (13), the solution of the Schrödinger equation can be recast in the form $\Psi_t = e^{i\omega t N_R} e^{-it(H_0 + \omega N_R)} \Psi_0$. Hence, the expectation value of our observable of interest, N_R , reads $\langle N_R \rangle_t = \langle \Psi_0 | e^{iH_F t} N_R e^{-iH_F t} | \Psi_0 \rangle$, where

$$H_F = H_0 + \omega N_R \quad (14)$$

is an explicit effective Floquet Hamiltonian.

One observes that the Floquet Hamiltonian (14) introduces an on-site potential bias between the two chains, with a magnitude equal to the driving frequency ω . This leads to an intuitive physical interpretation: The system should thermalize to a state with distinct particle densities in the left and right chains,

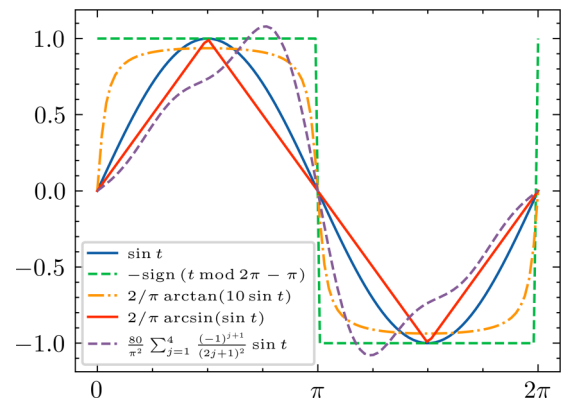


FIG. 10. Plot of the driving protocols f_t as a function of time. Here, the driving frequency is set to 1.

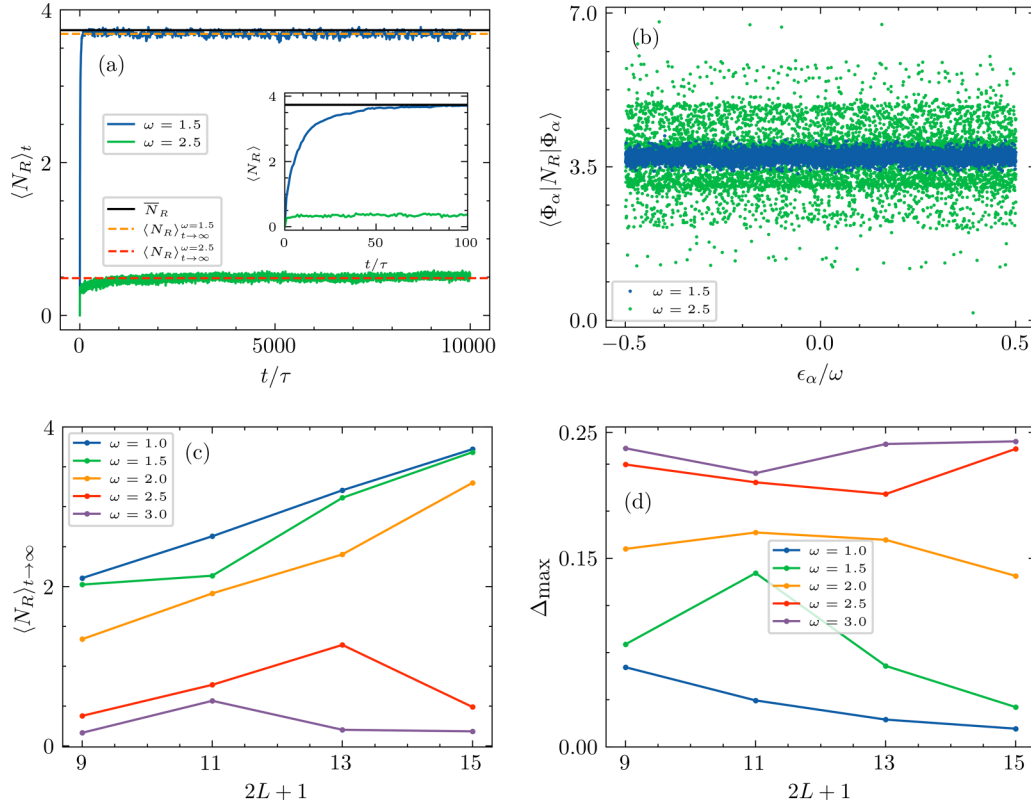


FIG. 11. Results for the driving protocol $f_i = -\text{sign}(t \bmod \frac{2\pi}{\omega} - \frac{\pi}{\omega})$. The interaction strength is $U = 0.5$. (a) The number of particles in the right chain, $\langle N_R \rangle_t$, as a function of time for frequencies below ($\omega = 1.5$) and above ($\omega = 2.5$) the critical frequency $\omega_c \simeq 2$. Horizontal dashed lines correspond to steady-state values computed according to Eq. (9). The horizontal black line shows the value of N_R corresponding to the uniform distribution of particles over the chains. (b) Diagonal matrix elements of the operator N_R in Floquet basis for different Floquet energies ϵ_α . In panels (a) and (b), the system comprises $2L + 1 = 15$ sites and $N = 8$ fermions. (c) The steady-state value $\langle N_R \rangle_\infty$ for the system of various sizes. (d) The maximal deviation of the diagonal matrix element $\langle \Phi_\alpha | N_R | \Phi_\alpha \rangle$ from the uniform value N_R .

where the density difference is governed by both the potential bias ω and the effective temperature.

Specifically, the thermal value of N_R obtained from the Gibbs ensemble reads

$$\langle N_R \rangle_{\text{Gibbs}} = \text{tr} N_R e^{-\beta H_F} / \text{tr} e^{-\beta H_F}. \quad (15)$$

Here, β is the inverse temperature determined by the average Floquet energy of the initial state through the equation $\text{tr} H_F e^{-\beta H_F} / \text{tr} e^{-\beta H_F} = \langle \Psi_0 | H_F | \Psi_0 \rangle$, and the trace is taken over the states with a fixed particle number. In Fig. 8(b), we compare the thermal value $\langle N_R \rangle_{\text{Gibbs}}$ to the steady-state value $\langle N_R \rangle_\infty$ for various frequencies. The agreement between the two is quite good, given that fluctuations around the thermal value scale as $\sim \sqrt{\langle N_R \rangle_{\text{Gibbs}}}$ and thus are quite large for small systems available in our studies.

The isospectral QPC (12) provides an instructive example where the transition between conducting and insulating regimes depends not only on the driving frequency but also on the particle imbalance between the chains, as discussed at the end of the previous section. Indeed, Fig. 8 demonstrates that at any frequency, the QPC initially permits particle flow from the left to the right chain but becomes insulating once the particle imbalance reaches the value determined by the frequency through Eq. (15).

We note that certain nonisospectral Hamiltonians with time-dependent on-site potentials can yield explicitly local Floquet Hamiltonians under more general gauge transformations [21–23,26]. However, the QPC described by Eqs. (4) and (5) in the previous section does not belong to this class, as it involves a time-dependent hopping (non-on-site) term.

V. DISCUSSION AND OUTLOOK

To summarize, we have demonstrated a nonthermal dynamical phase in a periodically driven interacting disorder-free quantum many-body system. The system consists of two fermionic chains connected by a driven QPC. For a sufficiently high driving frequency, the QPC is insulating, and an initial particle imbalance between the chains remains forever, in stark contrast to the prediction of the Floquet ETH. In contrast, for low frequencies, the QPC is conducting, and the fermions eventually spread homogeneously across the chains, in accordance with the Floquet ETH.

The frequency dependence of the effect distinguishes it from other known types of the Floquet thermalization breakdown. In particular, in nongeneric Floquet systems featuring integrability, exact quantum many-body scars, or Hilbert space fragmentation, the violation of the Floquet ETH

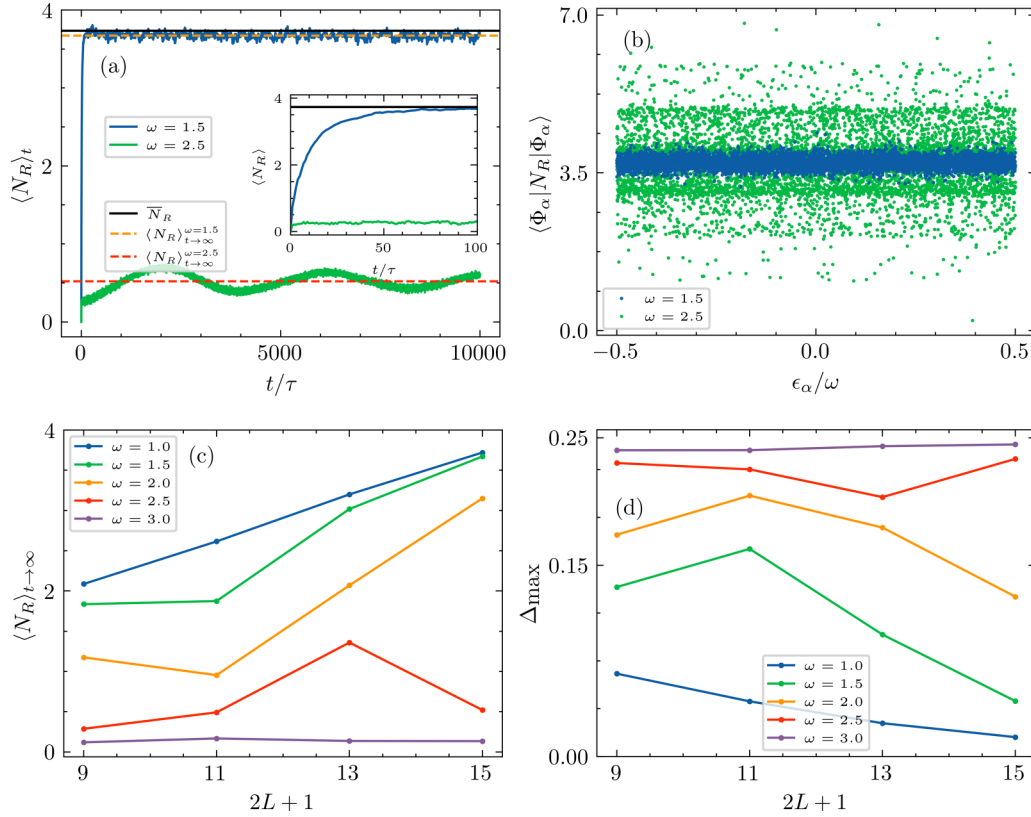


FIG. 12. The same as in Fig. 11 except for the driving protocol $f_t = \frac{2}{\pi} \arctan(\sin 10 \omega t)$.

and the thermalization breakdown typically take place for arbitrary driving frequencies, and athermal eigenstates often are independent of the frequency [27–33]. Alternatively, the breakdown or slowdown of thermalization can occur for a discrete set of frequencies in such systems [34–41]. A completely different frequency dependence of the Floquet ETH violation reported here may indicate a distinct origin of this violation—a question left for further studies.

It should be emphasized that the thermalization breakdown reported here cannot be explained by the Floquet prethermalization [42–46]. The latter phenomenon implies the thermalization slowdown at large frequencies, with the thermalization time depending exponentially on the frequency. If the system under study exhibited “mere” prethermalization, the particle imbalance between the chains would be eventually washed out. We know that this is not the case, since in finite systems under study we have access to arbitrary long times, as well as to the infinite-time averages (9). Furthermore, the critical frequency is not, in fact, large enough for the prethermalization mechanism to be effective, as we demonstrate in Appendix G.

For the same reasons, the effect reported here is different from the dynamical localization [47], coherent destruction of tunneling [48], and dynamical freezing [49,50] phenomena, which are inherently prethermal in nature [51] and occur at discrete frequencies.

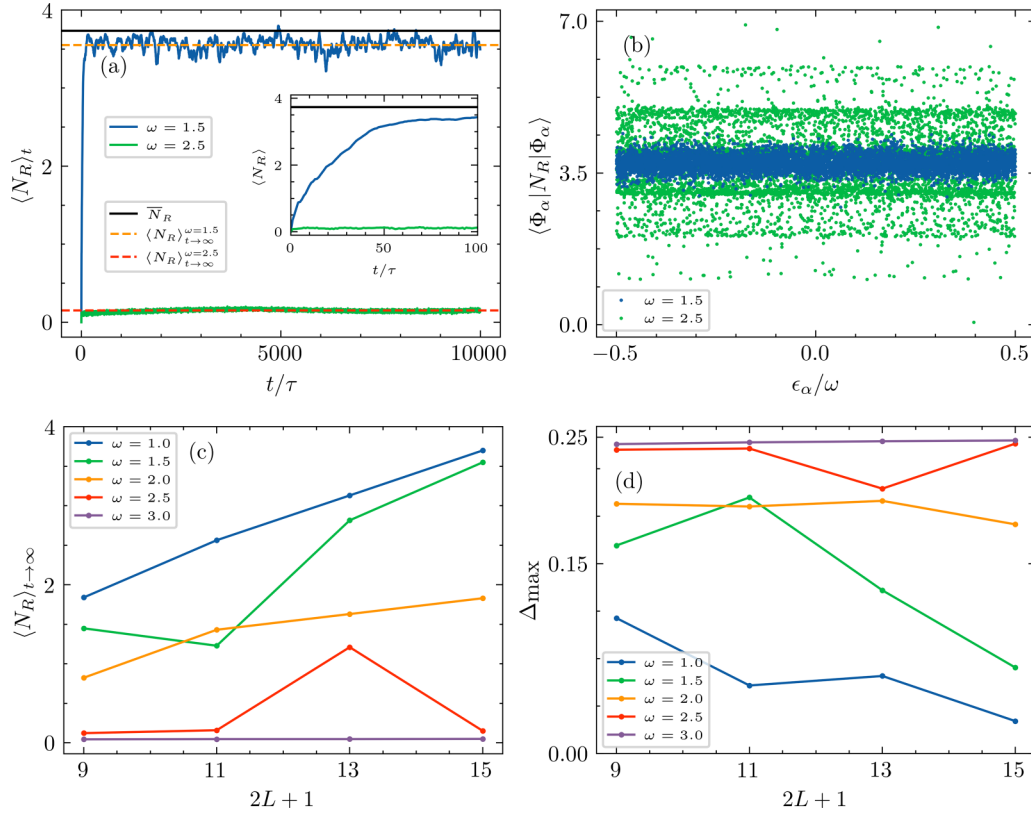
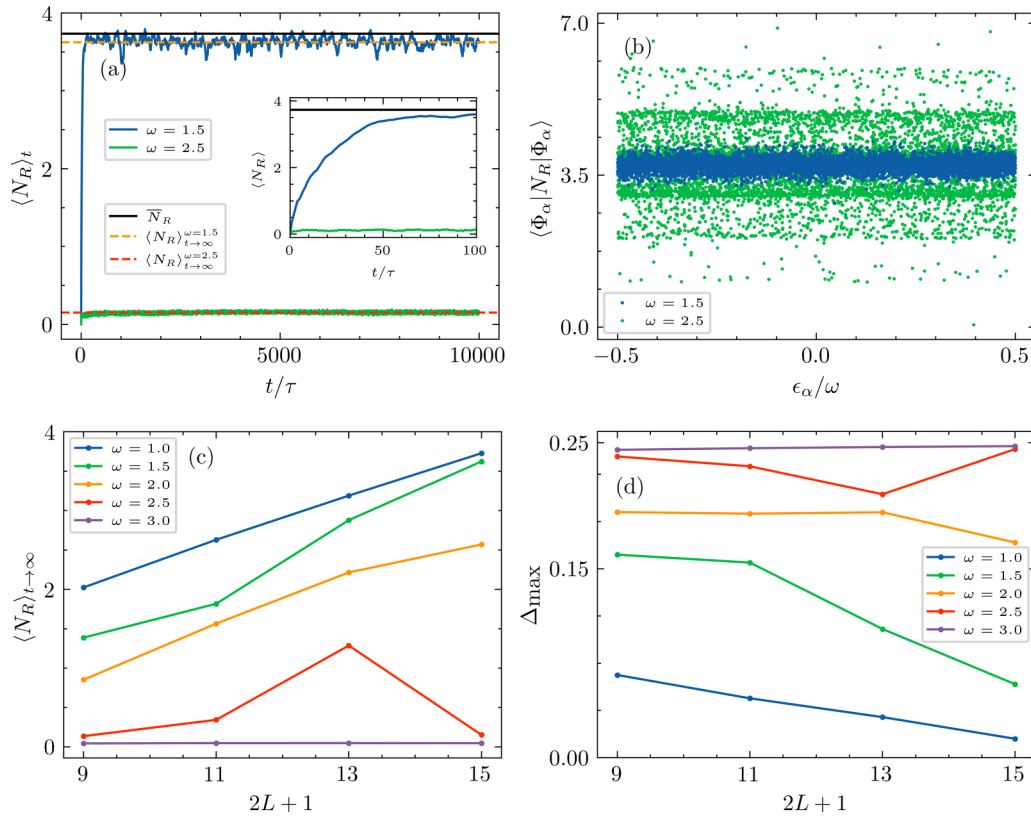
Known reasons for the violation of the Floquet ETH being excluded, the question regarding the nature of our findings becomes pressing. We suggest that addressing this question might require revisiting the foundational reasoning behind

the Floquet ETH. Let us briefly recall this reasoning, which consists of two key parts.

The first part is essentially an application of the conventional ETH for time-independent Hamiltonians extended to the Floquet case. The second part relies on the observation that a many-body Floquet Hamiltonian should be highly nonlocal—a property often justified through the Floquet-Magnus high-frequency expansion. This expansion involves nested commutators between Hamiltonians at different moments of time [52,53], naturally generating complex, nonlocal operators [54]. Such nonlocality is believed to drive thermalization to an effective infinite-temperature state, as evidenced by local observables.

However, the second part of this reasoning contains a potential loophole. Crucially, a Floquet Hamiltonian is not unique, as its eigenenergies are defined only modulo $2\pi/\tau$. The apparent nonlocality arising from the Floquet-Magnus expansion might therefore be spurious; a different choice of Floquet energies could restore locality. Consequently, thermalization to the infinite-temperature state may not be universally justified.

The isospectral model discussed in Sec. IV exemplifies this argument. Its Floquet Hamiltonian (14) is explicitly local, and the system thermalizes in accordance with the ETH for this Hamiltonian—yet not to the infinite-temperature state, contrary to the expectations of the Floquet ETH [10,11]. Notably, had we derived the Floquet Hamiltonian via the Floquet-Magnus expansion, we would have obtained a seemingly nonlocal representation, obscuring the clear physical picture implied by Eq. (14).


 FIG. 13. The same as in Fig. 11 except for the driving protocol $f_t = \frac{2}{\pi} \arcsin(\sin \omega t)$.

 FIG. 14. The same as in Fig. 11 except for the driving protocol $f_t = \frac{80}{\pi^2} \sum_{j=1}^4 \frac{(-1)^{j+1}}{(2j+1)^2} \sin \omega t$.

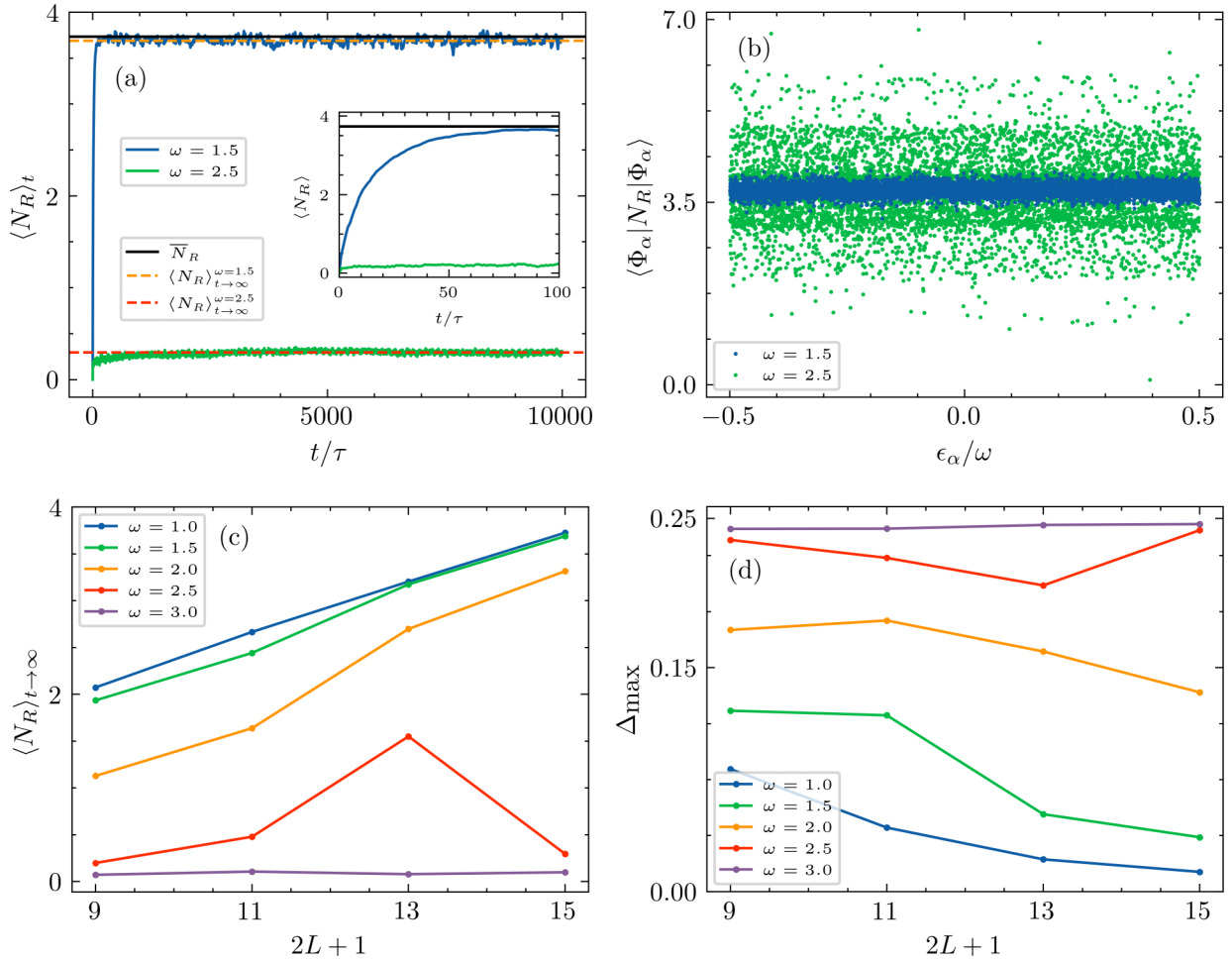


FIG. 15. The same as in Fig. 11 except for the QPC Hamiltonian V_i given by Eq. (C1) with the tunneling amplitude $f_i = \sin \omega t$ and on-site potential $g_i = \sin \omega t$.

Our findings thus suggest that a revision of the widely accepted understanding of thermalization in generic quantum Floquet systems can be required [55]. A complete theoretical explanation of these results within the framework of quantum thermalization remains an important open question for future research.

Another intriguing question for further studies concerns the scope of the reported effect. While we have demonstrated its robustness in one-dimensional geometry, our numerical tests do not extend to higher dimensions. It is not uncommon when exotic out-of-equilibrium phenomena in one dimension have no counterparts in higher dimensions [57–62]. It is therefore reasonable to ask whether the reported type of the Floquet thermalization breakdown extends beyond one-dimensional geometry. We note that this question, while challenging for numerical exploration, can perhaps be addressed with the state-of-the-art quantum simulators with planar arrangement of qubits [63–67].

As a word of caution, we note that, in the case of noninteracting fermions, dissipation (caused, e.g., by phonon-fermion interactions) destroys the phase transition and makes the QPC conducting at any frequency [68]. Likely it will have the same effect in the interacting case. This should be accounted for while conceiving experimental tests of the effect.

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DATA AVAILABILITY

The numerical data supporting the findings of this study are available from the corresponding author upon reasonable request.

APPENDIX A: ENTANGLEMENT ENTROPY

Here, we show the time dependence of the entanglement entropy $S_{R,t} = -\text{tr}(\rho_{R,t} \ln \rho_{R,t})$, where $\rho_{R,t}$ is the density matrix of the right chain. Qualitatively, it is similar to that of $\langle N_R \rangle_t$: It saturates about a thermal value for low frequencies but stays small for large frequencies (see Fig. 9).

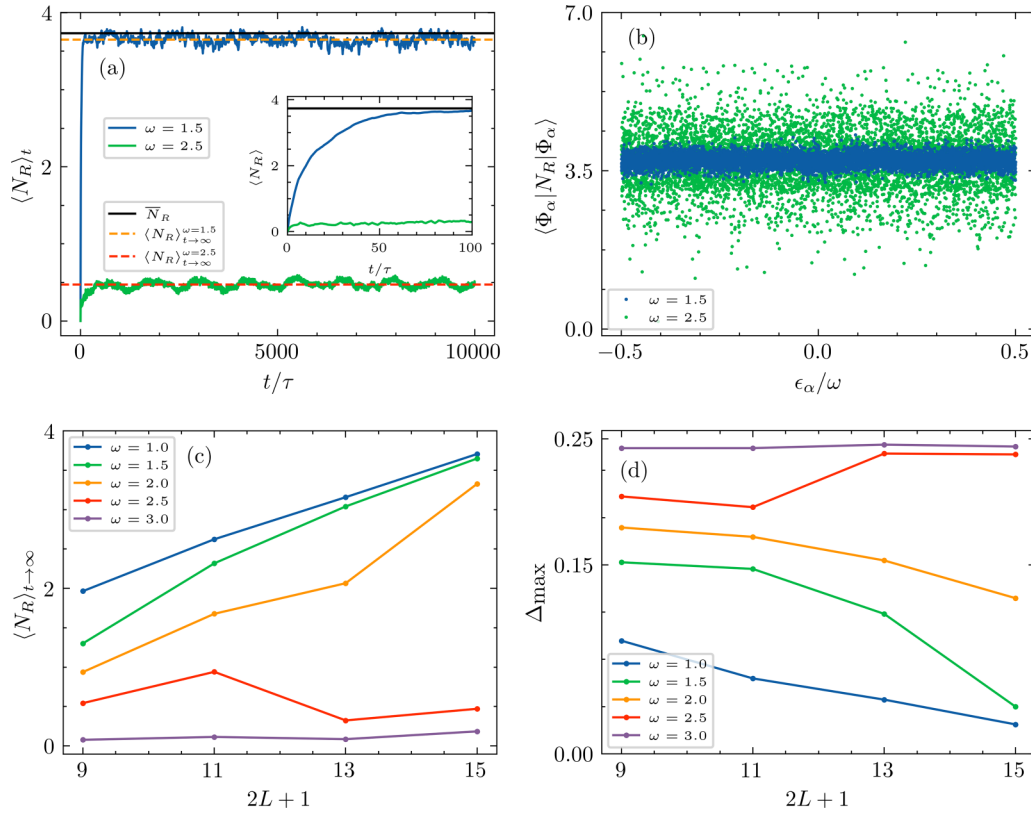


FIG. 16. Results for interactions with the next-to-nearest-neighbor terms included. The same as in Fig. 11 except for the driving protocol $f_t = \sin \omega t$ and the interaction term V_{int} given by Eq. (D1) with $U = U' = 0.5$.

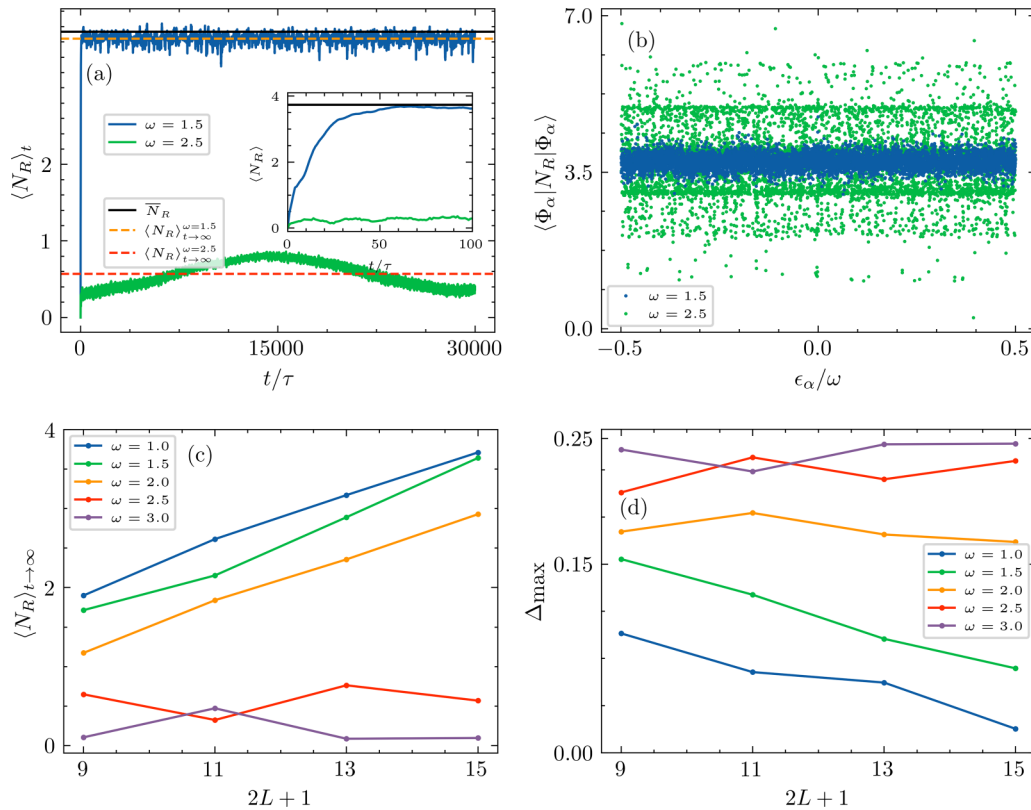


FIG. 17. Results for the case of no interactions within the QPC. The same as in Fig. 11 except for the driving protocol $f_t = \sin \omega t$ and the interaction term V_{int} given by Eq. (D2) with $U = 0.5$.

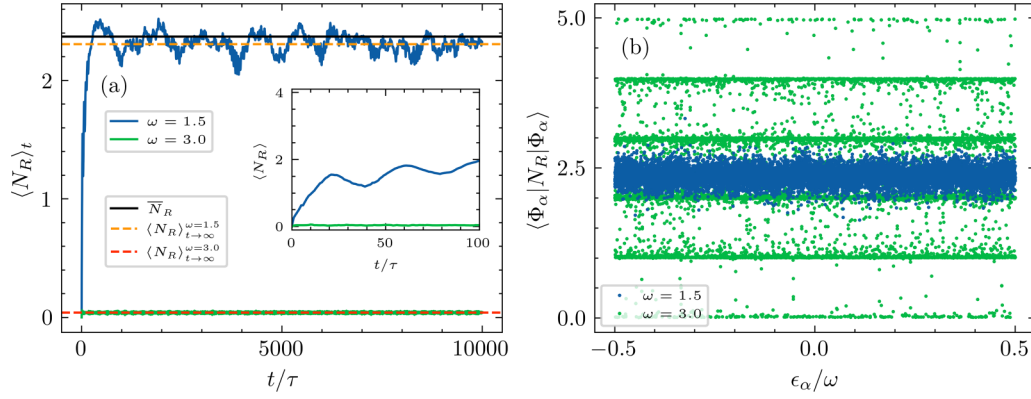


FIG. 18. Results for the initial state away from the full filling. The same as in Figs. 11(a) and 11(b) except for the driving protocol $f_t = \sin \omega t$, the system size $2L + 1 = 19$, and the number of fermions $N = 5$. The frequencies are given by $\omega = 1.5$ and $\omega = 3.0$.

APPENDIX B: VARIOUS DRIVING PROTOCOLS

Here, we present results for various driving protocols f_t in Eq. (5). The protocols under study are shown in Fig. 10. The results are shown in Figs. 11–14. One can see that they are qualitatively similar to the results reported in the main text.

APPENDIX C: A QPC WITH ON-SITE POTENTIAL

It was found in Ref. [6] that, in the noninteracting case, an alternating on-site potential within the QPC typically destroys the effect and ensures the current flow at any frequency. However, a notable exception was found for the QPC term V_t given by

$$V_t = -\frac{1}{2}(c_{L+1}^\dagger \quad c_{L+2}) \begin{pmatrix} g_t & f_t \\ f_t & -g_t \end{pmatrix} \begin{pmatrix} c_{L+1}^\dagger \\ c_{L+2} \end{pmatrix}, \quad (\text{C1})$$

with the tunneling amplitude $f_t = \sin(\omega t)$ and on-site potential $g_t = \sin(\omega t)$. We have found the same qualitative picture in the interacting case. In particular, the current halts completely for the above V_t , as shown in Fig. 15.

APPENDIX D: VARIOUS INTERACTION TERMS

We have verified that the physical picture remains the same for various types of interaction terms. As an illustration, we show in Fig. 16 the results for the next-to-nearest-neighbor

interactions,

$$V_{\text{int}} = U \sum_{j=1}^{2L} n_j n_{j+1} + U' \sum_{j=1}^{2L-1} n_j n_{j+2}, \quad (\text{D1})$$

where U and U' are two interaction constants. In addition, we consider the case where the interparticle interactions are switched off within the QPC:

$$V_{\text{int}} = U \sum_{j=1}^L n_j n_{j+1} + U \sum_{j=L+2}^{2L} n_j n_{j+1}. \quad (\text{D2})$$

The results are presented in Fig. 17.

APPENDIX E: INITIAL STATES AWAY FROM FULL FILLING

We have examined initial states away from the full filling of the left chain. Here, we present results for the ground state of the Hamiltonian

$$H_L + U \sum_{j=1}^L n_j n_{j+1}. \quad (\text{E1})$$

The physical picture does not change qualitatively (see Fig. 18).

We note that reliable finite-size scaling cannot be established in this case, as the maximum achievable particle

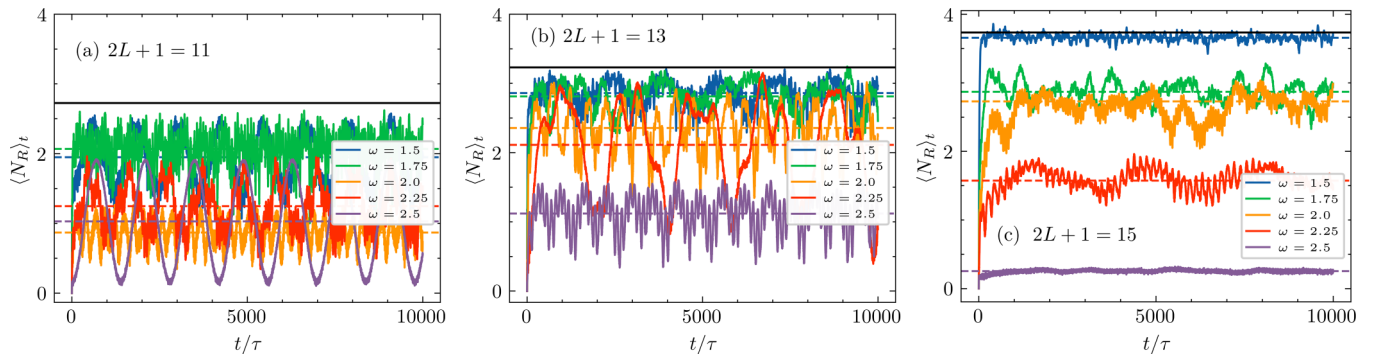


FIG. 19. Dependence of the dynamics on the system size: (a) $2L + 1 = 11$ and $N = 6$, (b) $2L + 1 = 13$ and $N = 7$, and (c) $2L + 1 = 15$ and $N = 8$. The rest of the parameters are the same as in Fig. 3.

number at approximate half-filling ($N = 5$) is already small. For even smaller system sizes, the system would effectively become few-body, rendering it an inadequate proxy for many-body behavior.

APPENDIX F: FINITE-SIZE EFFECTS

Figure 19 illustrates finite-size effects, which turn to be particularly strong at intermediate frequencies.

APPENDIX G: ESTIMATE OF THE PRETHERMALIZATION TIMESCALE AND ENERGY SCALE

A conservative estimate Λ of the prethermalization energy scale is obtained as the sum of operator norms of Hamiltonian

terms with a support covering a fixed site [46]. For our system (1)–(6), this implies $\Lambda = 1 + 2U$.

Prethermalization is expected when $\omega \gg \Lambda$. In contrast, in our calculations $\omega \sim \Lambda$. Therefore, prethermalization cannot explain the observed halting of the current.

One can make this argument somewhat more quantitative. The ratio of frequency-dependent thermalization timescales $\tau_{\text{Gibbs}}(\omega)$ for two different frequencies is estimated as $\tau_{\text{Gibbs}}(\omega_2)/\tau_{\text{Gibbs}}(\omega_1) \sim \exp((\omega_2 - \omega_1)/\Lambda)$. For $\omega_1 = 1.5$, $\omega_2 = 2.5$, and $U = 0.5$, this estimate reads $\tau_{\text{Gibbs}}(\omega_2)/\tau_{\text{Gibbs}}(\omega_1) \sim \sqrt{e} \simeq 1.6$. Thus, the prethermalization timescales for these two frequencies should be of the same order of magnitude. However, we clearly observe thermalization at the frequency ω_1 , with $\tau_{\text{Gibbs}}(\omega_1) \sim 10$, while no thermalization is observed at the frequency ω_2 up to times $\sim 1000 \tau_{\text{Gibbs}}(\omega_1)$. This argument confirms that the observed lack of thermalization cannot be explained by “mere” prethermalization.

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- [1] Y. V. Nazarov and Y. M. Blanter, *Quantum Transport: Introduction to Nanoscience* (Cambridge University Press, Cambridge, UK, 2009).
 - [2] G. Platero and R. Aguado, Photon-assisted transport in semiconductor nanostructures, *Phys. Rep.* **395**, 1 (2004).
 - [3] S. Kohler, J. Lehmann, and P. Hänggi, Driven quantum transport on the nanoscale, *Phys. Rep.* **406**, 379 (2005).
 - [4] D. A. Ryndyk, *Theory of Quantum Transport at Nanoscale An Introduction*, Springer Series in Solid-State Sciences, Vol. 185 (Springer, Cham, 2016).
 - [5] G. T. Landi, D. Poletti, and G. Schaller, Nonequilibrium boundary-driven quantum systems: Models, methods, and properties, *Rev. Mod. Phys.* **94**, 045006 (2022).
 - [6] O. Gamayun, A. Slobodeniuk, J.-S. Caux, and O. Lychkovskiy, Nonequilibrium phase transition in transport through a driven quantum point contact, *Phys. Rev. B* **103**, L041405 (2021).
 - [7] See also Ref. [8] for a very similar phase transition in a classical system of anharmonic oscillators, and Ref. [9] for a somewhat similar phase transition driven by the frequency of the periodic external voltage bias.
 - [8] M. Romero-Bastida, Effect of external potential on the energy transport in harmonically driven segmented Frenkel-Kontorova lattices, *Phys. Rev. E* **110**, 054115 (2024).
 - [9] M. Žnidarič, B. Žunkovič, and T. Prosen, Transport properties of a boundary-driven one-dimensional gas of spinless fermions, *Phys. Rev. E* **84**, 051115 (2011).
 - [10] A. Lazarides, A. Das, and R. Moessner, Equilibrium states of generic quantum systems subject to periodic driving, *Phys. Rev. E* **90**, 012110 (2014).
 - [11] L. D’Alessio and M. Rigol, Long-time behavior of isolated periodically driven interacting lattice systems, *Phys. Rev. X* **4**, 041048 (2014).
 - [12] K. Seetharam, P. Titum, M. Kolodrubetz, and G. Refael, Absence of thermalization in finite isolated interacting Floquet systems, *Phys. Rev. B* **97**, 014311 (2018).
 - [13] B. Ye, F. Machado, C. D. White, R. S. K. Mong, and N. Y. Yao, Emergent hydrodynamics in nonequilibrium quantum systems, *Phys. Rev. Lett.* **125**, 030601 (2020).
 - [14] A. Pizzi, D. Malz, G. De Tomasi, J. Knolle, and A. Nunnenkamp, Time crystallinity and finite-size effects in clean Floquet systems, *Phys. Rev. B* **102**, 214207 (2020).
 - [15] T. N. Ikeda and A. Polkovnikov, Fermi’s golden rule for heating in strongly driven Floquet systems, *Phys. Rev. B* **104**, 134308 (2021).
 - [16] A. Morningstar, D. A. Huse, and V. Khemani, Universality classes of thermalization for mesoscopic Floquet systems, *Phys. Rev. B* **108**, 174303 (2023).
 - [17] P. Weinberg and M. Bukov, QuSpin: A Python package for dynamics and exact diagonalisation of quantum many body systems. Part I: Spin chains, *SciPost Phys.* **2**, 003 (2017).
 - [18] P. Weinberg and M. Bukov, QuSpin: A Python package for dynamics and exact diagonalisation of quantum many body systems. Part II: Bosons, fermions and higher spins, *SciPost Phys.* **7**, 020 (2019).
 - [19] I. Ermakov, O. Lychkovskiy, and B. V. Fine, Periodic classical trajectories and quantum scars in many-spin systems, *arXiv:2409.00258*.
 - [20] Observe, however, that $\langle N_R \rangle_\infty$ does not necessarily monotonically decrease with the frequency, as one could expect. For example, $\langle N_R \rangle_\infty|_{\omega=2} < \langle N_R \rangle_\infty|_{\omega=2.5}$ at $U = 1$, as one can see from Fig. 5(b).
 - [21] F. Yan, L. Yang, and B. Li, Formal exact solution for the Heisenberg spin system in a time-dependent magnetic field and Aharonov-Anandan phase, *Phys. Lett. A* **251**, 289 (1999).
 - [22] F. Yan, L. Yang, and B. Li, Invariant Hermitian operator and geometric phase for the Heisenberg spin system in a time-dependent magnetic field, *Phys. Lett. A* **259**, 207 (1999).
 - [23] O. V. Gamayun and O. V. Lychkovskiy, A map between time-dependent and time-independent quantum many-body Hamiltonians, *Proc. Steklov Inst. Math.* **313**, 41 (2021).
 - [24] N. Il’in, A. Aristova, and O. Lychkovskiy, Adiabatic theorem for closed quantum systems initialized at finite temperature, *Phys. Rev. A* **104**, L030202 (2021).
 - [25] X. Wang, X. He, and J. Wu, Quantum geometry and many-body Landau-Zener tunneling in time-dependent quantum systems with instantaneous quantum integrability, *arXiv:2503.18396*.

- [26] A. Colcelli, G. Mussardo, G. Sierra, and A. Trombettoni, Integrable Floquet Hamiltonian for a periodically tilted 1D gas, *Phys. Rev. Lett.* **123**, 130401 (2019).
- [27] V. Gritsev and A. Polkovnikov, Integrable Floquet dynamics, *SciPost Phys.* **2**, 021 (2017).
- [28] K. Mizuta, K. Takasan, and N. Kawakami, Exact Floquet quantum many-body scars under Rydberg blockade, *Phys. Rev. Res.* **2**, 033284 (2020).
- [29] H. Yarloo, A. Emami Kopaei, and A. Langari, Homogeneous Floquet time crystal from weak ergodicity breaking, *Phys. Rev. B* **102**, 224309 (2020).
- [30] V. Khemani, M. Hermele, and R. Nandkishore, Localization from Hilbert space shattering: From theory to physical realizations, *Phys. Rev. B* **101**, 174204 (2020).
- [31] B. Mukherjee, S. Nandy, A. Sen, D. Sen, and K. Sengupta, Collapse and revival of quantum many-body scars via Floquet engineering, *Phys. Rev. B* **101**, 245107 (2020).
- [32] S. Moudgalya, B. A. Bernevig, and N. Regnault, Quantum many-body scars and Hilbert space fragmentation: A review of exact results, *Rep. Prog. Phys.* **85**, 086501 (2022).
- [33] A. Teretenkov and O. Lychkovskiy, Duality between open systems and closed bilayer systems, and thermofield double states as quantum many-body scars, *Phys. Rev. B* **110**, L241105 (2024).
- [34] A. Das, Exotic freezing of response in a quantum many-body system, *Phys. Rev. B* **82**, 172402 (2010).
- [35] S. Bhattacharyya, A. Das, and S. Dasgupta, Transverse Ising chain under periodic instantaneous quenches: Dynamical many-body freezing and emergence of slow solitary oscillations, *Phys. Rev. B* **86**, 054410 (2012).
- [36] A. Haldar, R. Moessner, and A. Das, Onset of Floquet thermalization, *Phys. Rev. B* **97**, 245122 (2018).
- [37] A. Haldar, D. Sen, R. Moessner, and A. Das, Dynamical freezing and scar points in strongly driven Floquet matter: Resonance vs emergent conservation laws, *Phys. Rev. X* **11**, 021008 (2021).
- [38] C. Fleckenstein and M. Bukov, Prethermalization and thermalization in periodically driven many-body systems away from the high-frequency limit, *Phys. Rev. B* **103**, L140302 (2021).
- [39] S. Sugiura, T. Kuwahara, and K. Saito, Many-body scar state intrinsic to periodically driven system, *Phys. Rev. Res.* **3**, L012010 (2021).
- [40] S. Aditya and D. Sen, Dynamical localization and slow thermalization in a class of disorder-free periodically driven one-dimensional interacting systems, *SciPost Phys. Core* **6**, 083 (2023).
- [41] S. Ghosh, I. Paul, and K. Sengupta, Prethermal fragmentation in a periodically driven fermionic chain, *Phys. Rev. Lett.* **130**, 120401 (2023).
- [42] D. A. Abanin, W. De Roeck, and F. Huveneers, Exponentially slow heating in periodically driven many-body systems, *Phys. Rev. Lett.* **115**, 256803 (2015).
- [43] T. Mori, T. Kuwahara, and K. Saito, Rigorous bound on energy absorption and generic relaxation in periodically driven quantum systems, *Phys. Rev. Lett.* **116**, 120401 (2016).
- [44] A. Rubio-Abadal, M. Ippoliti, S. Hollerith, D. Wei, J. Rui, S. L. Sondhi, V. Khemani, C. Gross, and I. Bloch, Floquet prethermalization in a Bose-Hubbard system, *Phys. Rev. X* **10**, 021044 (2020).
- [45] W. Beatriz, O. Janes, A. Akkiraju, A. Pillai, A. Oddo, P. Reshetikhin, E. Druga, M. McAllister, M. Elo, B. Gilbert, D. Suter, and A. Ajoy, Floquet prethermalization with lifetime exceeding 90 s in a bulk hyperpolarized solid, *Phys. Rev. Lett.* **127**, 170603 (2021).
- [46] W. W. Ho, T. Mori, D. A. Abanin, and E. G. Dalla Torre, Quantum and classical Floquet prethermalization, *Ann. Phys.* **454**, 169297 (2023).
- [47] D. H. Dunlap and V. M. Kenkre, Dynamic localization of a charged particle moving under the influence of an electric field, *Phys. Rev. B* **34**, 3625 (1986).
- [48] F. Grossmann, T. Dittrich, P. Jung, and P. Hänggi, Coherent destruction of tunneling, *Phys. Rev. Lett.* **67**, 516 (1991).
- [49] H. Guo, R. Mukherjee, and D. Chowdhury, Dynamical freezing in exactly solvable models of driven chaotic quantum dots, *Phys. Rev. Lett.* **134**, 226501 (2025).
- [50] R. Mukherjee, H. Guo, K. Lewellen, and D. Chowdhury, Arresting quantum chaos dynamically in transmon arrays, *Newton* **1**, 100011 (2025).
- [51] D. J. Luitz, Y. Bar Lev, and A. Lazarides, Absence of dynamical localization in interacting driven systems, *SciPost Phys.* **3**, 029 (2017).
- [52] T. Kuwahara, T. Mori, and K. Saito, Floquet–Magnus theory and generic transient dynamics in periodically driven many-body quantum systems, *Ann. Phys.* **367**, 96 (2016).
- [53] M. Bukov, L. D’Alessio, and A. Polkovnikov, Universal high-frequency behavior of periodically driven systems: From dynamical stabilization to Floquet engineering, *Adv. Phys.* **64**, 139 (2015).
- [54] D. E. Parker, X. Cao, A. Avdoshkin, T. Scaffidi, and E. Altman, A universal operator growth hypothesis, *Phys. Rev. X* **9**, 041017 (2019).
- [55] We also note recent experimental evidence supporting the breakdown of the Floquet ETH [56].
- [56] Y. Guo, S. Dhar, A. Yang, Z. Chen, H. Yao, M. Horvath, L. Ying, M. Landini, and H.-C. Nägerl, Observation of many-body dynamical localization, *Science* **389**, 716 (2025).
- [57] M. B. Zvonarev, V. V. Cheianov, and T. Giamarchi, Spin dynamics in a one-dimensional ferromagnetic Bose gas, *Phys. Rev. Lett.* **99**, 240404 (2007).
- [58] M. Knap, C. J. M. Mathy, M. Ganahl, M. B. Zvonarev, and E. Demler, Quantum flutter: Signatures and robustness, *Phys. Rev. Lett.* **112**, 015302 (2014).
- [59] E. Burovski, V. Cheianov, O. Gamayun, and O. Lychkovskiy, Momentum relaxation of a mobile impurity in a one-dimensional quantum gas, *Phys. Rev. A* **89**, 041601 (2014).
- [60] O. Lychkovskiy, Perpetual motion of a mobile impurity in a one-dimensional quantum gas, *Phys. Rev. A* **89**, 033619 (2014).
- [61] O. Lychkovskiy, Perpetual motion and driven dynamics of a mobile impurity in a quantum fluid, *Phys. Rev. A* **91**, 040101(R) (2015).
- [62] B. Bertini, F. Heidrich-Meisner, C. Karrasch, T. Prosen, R. Steinigeweg, and M. Žnidarič, Finite-temperature transport in one-dimensional quantum lattice models, *Rev. Mod. Phys.* **93**, 025003 (2021).
- [63] F. Arute, K. Arya, R. Babbush, D. Bacon, J. C. Bardin, R. Barends, R. Biswas, S. Boixo, F. G. S. L. Brandao, D. A. Buell

- et al.*, Quantum supremacy using a programmable superconducting processor, *Nature (London)* **574**, 505 (2019).
- [64] M. Gong *et al.*, Quantum walks on a programmable two-dimensional 62-qubit superconducting processor, *Science* **372**, 948 (2021).
- [65] S. Ebadi, T. T. Wang, H. Levine, A. Keesling, G. Semeghini, A. Omran, D. Bluvstein, R. Samajdar, H. Pichler, W. W. Ho *et al.*, Quantum phases of matter on a 256-atom programmable quantum simulator, *Nature (London)* **595**, 227 (2021).
- [66] P. Scholl, M. Schuler, H. Williams, A. Eberharter, D. Barredo, K. Schymik, V. Lienhard, L.-P. Henry, T. Lang, T. Lahaye *et al.*, Quantum simulation of 2D Ising models using Rydberg atoms, *Nature (London)* **595**, 233 (2021).
- [67] H. Bernien *et al.*, Probing many-body dynamics on a 51-atom quantum simulator, *Nature* **551**, 579 (2017).
- [68] I. Ermakov and O. Lychkovskiy, Effect of dephasing on the current through a periodically driven quantum point contact, *JETP Lett.* **119**, 40 (2024).