

Suppressing Floquet thermalization by driving transparency in tilted lattices

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(Received 13 December 2024; accepted 5 August 2025; published 2 September 2025)

Suppressing Floquet thermalization is an important goal in the field of Floquet engineering. Here, we propose a driving transparency mechanism to suppress the Floquet thermalization in periodically driven tilted lattices of interacting spinless fermions. By leveraging the interplay between the tilting and interaction, the system can be engineered to prevent direct energy absorption from the driving at particular frequencies. We analytically determine these driving transparency frequencies through the Floquet time-dependent perturbation theory. The system is near-integrable at the driving transparency frequencies. The suppression of thermalization by the driving transparency is demonstrated numerically by the violation of the Floquet eigenstate thermalization hypothesis, the retention of the initial state, and the logarithmic growth of entanglement entropy. Our work provides an intuitive way to suppress Floquet thermalization and offers an opportunity to engineer exotic nonequilibrium Floquet states in tilted lattice systems.

DOI: [10.1103/lhsz-dkmq](https://doi.org/10.1103/lhsz-dkmq)

I. INTRODUCTION

Many-body localization (MBL) in disordered quantum systems has been a subject of extensive study as a robust mechanism for the breakdown of quantum thermalization [1–4]. Recently, it has been demonstrated that thermalization breakdown can also occur in tilted systems without disorder. The tilted lattices of interacting particles exhibit the nonthermal phenomena such as Stark MBL [5–15], Hilbert space fragmentation [16–20], and quantum many-body scars [21], which are induced by the tilting field or the interplay between the tilting field and particle interactions.

When subjected to time-periodic driving, a quantum many-body system generally heats up to a featureless infinite-temperature state. This phenomenon is explained by the Floquet eigenstate thermalization hypothesis (ETH) from a static perspective and that the system indefinitely absorbs energy from the driving from a dynamical perspective [22–24]. Avoiding such a fate is an important goal in Floquet engineering, which is a powerful tool to manipulate and engineer exotic properties of quantum systems through periodic driving [25–29]. To date, several methods have been developed to suppress Floquet thermalization. One approach involves introducing strong disorder to induce the undriven system into the MBL phase, where the system has a finite

bandwidth. When the driving frequency exceeds the bandwidth, energy absorption is prevented and the system remains stable in the MBL phase [30]. Conversely, low-frequency driving invariably leads to Floquet thermalization [30–33]. It shows a monotonous transition between the Floquet thermalization and the MBL phases when the frequency increases [30–33]. For ergodic systems, the energy spectrum is extensive, making Floquet thermalization inevitable. However, the thermalization process can be delayed through the mechanism of Floquet prethermalization [34–39]. If the driving frequency is higher than the effective single-particle excitation energy, the energy absorption requires high-order processes, causing the system to spend a prolonged period in an intermediate nontrivial state described by an effective Hamiltonian [34–39]. Interestingly, a recent study shows that heating can also be suppressed in the low-frequency regime in long-range interacting systems with the energy spectrum split into bands [40]. Energy absorption is suppressed when the driving is off resonance with the gap between the bands [40]. Other methods to hinder Floquet thermalization include using ac electric fields to induce dynamical localization [41,42] and employing multicolor driving protocols to induce interference between the energy absorption processes [43,44]. Notably, the energy spectrum of the tilting systems features regular ladderlike structure for large tilting strengths. Inspired by Ref. [40], it is intriguing to explore whether and how Floquet thermalization can be suppressed in the tilted lattice systems.

In this article, we explore the suppression of Floquet thermalization in periodically driven tilted chains of interacting spinless fermions. In the weak driving regime, we derive an effective Floquet Hamiltonian that governs the dynamics at stroboscopic times using the Floquet time-dependent perturbation theory (TDPT). From the effective

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Floquet Hamiltonian, we analytically derive the resonance conditions, which incorporate the interplay between the tilting and interaction. Remarkably, we also identify driving transparency (DT) conditions, under which the system cannot directly absorb energy from the driving. We analyze both the static properties of the Floquet Hamiltonian and the dynamical evolution of the system. We find that at the DT condition, the system is near-integrable and thermalization can be suppressed, with the Floquet states violating the Floquet ETH and the entanglement entropy growing logarithmically with time. This contrasts with the scenario at the resonance conditions, where the Floquet ETH is satisfied and the entanglement entropy grows linearly. In addition, we observe that the resonance and the DT result in a nonmonotonous thermalization behavior as the driving frequency changes. This is in stark contrast to the monotonous Floquet thermalization-MBL phase transition observed in the disordered systems. Our study provides a method to suppress Floquet thermalization and opens up possibilities for realizing exotic nonequilibrium Floquet states in interacting tilted systems.

The rest of the article is organized as follows. In Sec. II, we introduce the model and derive the resonance and DT conditions through the Floquet TDPT. In Sec. III, we analyze the Floquet spectrum properties and the dynamical evolution under the resonance and DT conditions to illustrate the DT mechanism for suppressing Floquet thermalization. In Sec. IV, we summarize our results.

II. MODEL AND DRIVING TRANSPARENCY

We consider a periodically driven one-dimensional tilted lattice of interacting spinless fermions, which is described by

the Hamiltonian

$$\begin{aligned}\hat{H}(t) &= \hat{H}_{\text{on}} + [1 + u(t)]\hat{H}_J, \\ \hat{H}_{\text{on}} &= U \sum_{j=1}^{L-1} \hat{n}_j \hat{n}_{j+1} - g \sum_{j=1}^L j \hat{n}_j, \\ \hat{H}_J &= J \sum_{j=1}^{L-1} (\hat{c}_j^\dagger \hat{c}_{j+1} + \hat{c}_{j+1}^\dagger \hat{c}_j).\end{aligned}\quad (1)$$

Here, $\hat{c}_j^\dagger (\hat{c}_j)$ creates (annihilates) a fermion at site j and $\hat{n}_j = \hat{c}_j^\dagger \hat{c}_j$ is the particle number operator. The parameters J , U , and g are the nearest-neighbor tunneling strength, the nearest-neighbor interaction strength, and the tilting field strength, respectively. $u(t) = -u$ and $+u$ in the first and second halves of each period T , respectively. $\omega = 2\pi/T$ is the driving frequency. L is the total number of lattice sites. We consider $g > 0$ without loss of generality and consider g and U are comparable. We set open boundary condition and half filling with particle number $N = L/2$. We set $\hbar = 1$ and set the energy unit as $J = 1$.

The properties of any time-periodic Hamiltonian $\hat{H}(t)$ at stroboscopic times are determined by its Floquet Hamiltonian

$$\hat{H}_F = \frac{i}{T} \ln \hat{F}, \quad (2)$$

where the Floquet operator $\hat{F} = \mathcal{T} e^{\int_0^T \hat{H}(t) dt}$ is the unitary evolution operator over one period. For our model, it is difficult to exactly obtain $\hat{H}_F$ because $\hat{H}_{\text{on}}$ and $\hat{H}_J$ do not commute with each other. However, when $J, uJ \ll g, U, \omega/2\pi$, we can treat $[1 + u(t)]\hat{H}_J$ as a perturbation to $\hat{H}_{\text{on}}$ and derive an effective Floquet Hamiltonian up to first order through the Floquet TDPT [45–48],

$$\begin{aligned}\hat{H}_F^{(1)} &= \sum_{i=1}^3 \left[J_i \sum_j \hat{\mathcal{P}}_{j-1,j+2}^{(|\Delta_i|)} \hat{c}_j^\dagger \hat{c}_{j+1} + \text{H.c.} \right] + \frac{i}{T} \ln e^{-i\hat{H}_{\text{on}} T}, \\ J_1 &= J\delta_{g,U} - \frac{iJ(1 - \delta_{g,U})}{T\Delta_1} (e^{i\frac{\Delta_1 T}{2}} - 1) [-u + 1 + (u+1)e^{i\frac{\Delta_2 T}{2}}], \\ J_2 &= -\frac{iJ}{T\Delta_2} (e^{i\frac{\Delta_2 T}{2}} - 1) [-u + 1 + (u+1)e^{i\frac{\Delta_2 T}{2}}], \\ J_3 &= -\frac{iJ}{T\Delta_3} (e^{i\frac{\Delta_3 T}{2}} - 1) [-u + 1 + (u+1)e^{i\frac{\Delta_3 T}{2}}],\end{aligned}\quad (3)$$

where $\Delta_1 = g - U$, $\Delta_2 = g$, $\Delta_3 = g + U$, and the projection operators $\hat{\mathcal{P}}_{j-1,j+2}^{(|\Delta_1|)} = \hat{n}_{j+2}(1 - \hat{n}_{j-1})$, $\hat{\mathcal{P}}_{j-1,j+2}^{(|\Delta_2|)} = 1 - (\hat{n}_{j-1} - \hat{n}_{j+2})^2$, and $\hat{\mathcal{P}}_{j-1,j+2}^{(|\Delta_3|)} = \hat{n}_{j-1}(1 - \hat{n}_{j+2})$. Hamiltonian equation (3) can be understood as a result of the driving-assisted tunneling. The first and last terms in $\hat{H}_F^{(1)}$ depict the effective tunneling processes and the energy detuning [47], respectively. $|\Delta_i|$ are the bare energy barriers for the corresponding tunneling processes.

From Hamiltonian equation (3), we see that there exist resonance conditions for each tunneling process i , under which the energy detuning $\ln e^{-i\Delta_i T} = 0$ and $|J_i|$ is nonzero. The

resonance conditions for the i th tunneling process read [48]

$$|\Delta_i| = 0 \quad \text{or} \quad (2k-1)\omega, \quad k \in \{\text{positive integers}\}. \quad (4)$$

Remarkably, we find that there are DT conditions for each tunneling process, under which $|J_i|$ is zero. Obviously, the DT conditions for the i th tunneling process read

$$|\Delta_i| = 2k\omega, \quad k \in \{\text{positive integers}\}. \quad (5)$$

The resonance and DT can be interpreted from the Fourier expansion of the driving strength

TABLE I. The resonance frequency $\omega_{i,r}$ and DT frequency $\omega_{i,t}$ for the i th tunneling process at $U = 2g$. Here, k are positive integers.

i	1	2	3
$\omega_{i,r}$	$g/(2k-1)$	$g/(2k-1)$	$3g/(2k-1)$
$\omega_{i,t}$	$g/2k$	$g/2k$	$3g/2k$

$Ju(t) = 2iJu \sum_{q \in \text{odd}} (e^{iq\omega}/q\pi)$, where q are odd integers. The driving provides $q\omega$ energy quanta for the system. Thus, resonance between the system and the driving happens when $|\Delta_i|$ is compensated by the provided energy, and the system cannot absorb energy from the driving directly when $|\Delta_i|$ is even integer multiples of ω . The resonance condition $|\Delta_i| = 0$ accounts for the intrinsic resonance between U and g .

In the following study, we show that it is possible to suppress the Floquet thermalization when all the tunneling processes satisfy the DT conditions. For this purpose, we fix $U = 2g$. In this case, there are four characteristic driving frequencies: (1) the full resonance frequencies $\omega_{fr} = g/(2k-1)$ and (2) the full DT frequencies $\omega_{ft} = g/2k$, at which all the tunneling processes satisfy the resonance and DT conditions, respectively; (3) the partial resonance frequencies $\omega_{pr} = 3g, 3g/(6k \pm 1)$ and (4) the partial DT frequencies $\omega_{pt} = 3g/2, 3g/(6k \pm 2)$, at which the third tunneling process satisfies the resonance condition and the DT condition, respectively, and the other processes do not satisfy either of the conditions; see details in Table I. We note that under the full DT conditions, the first-order effective Hamiltonian vanishes and one needs to consider higher-order terms.

III. SUPPRESSING FLOQUET THERMALIZATION THROUGH DRIVING TRANSPARENCY

A. Violation of the Floquet ETH

Floquet thermalization is expected to happen due to the Floquet ETH, which states that the Floquet states are fully mixed states and predicts that the Floquet quasienergies satisfy the statistics of circular orthogonal ensemble (COE) [22–24,30,31]. To show that the thermalization can be suppressed by the DT, we study the properties of the Floquet quasienergies and the Floquet states through exact diagonalizing the Floquet Hamiltonian equation (2).

The Floquet quasienergies are obtained by folding the eigenenergies of Hamiltonian equation (2) into one Floquet-Brillouin zone $[-\omega/2, \omega/2)$. Orderly denoting the Floquet quasienergies as $\{\epsilon_\alpha\}$, the ratio of the adjacent quasienergy gaps is defined as

$$r_\alpha = \min \{\delta_\alpha/\delta_{\alpha+1}, \delta_{\alpha+1}/\delta_\alpha\}, \quad (6)$$

with $\delta_\alpha = \epsilon_\alpha - \epsilon_{\alpha-1}$. For systems satisfying the Floquet ETH, such ratio obeys the distribution of COE $P_{\text{COE}}(r)$, which gives the average $\bar{r}_{\text{COE}} \approx 0.527$ [23,30,31]. For generic integrable systems, it obeys the Poisson distribution $P_r(r)$, which gives the average $\bar{r}_p \approx 0.386$ [23,30,31,49]. (Note that nongeneric integrable systems can have energy spacing distribution different from both COE and Poisson [49].) In Fig. 1(a), we plot the averaged adjacent quasienergy gap ratio $\bar{r}$ as a function of g/ω at $g = 6$, $U = 2g$, $u = 0.5$, and $L = 12, 14, 16$, and

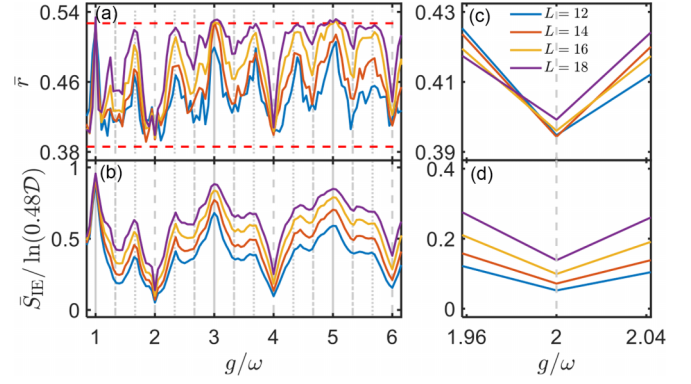


FIG. 1. (a) The averaged quasienergy gap ratio $\bar{r}$ and (b) the averaged and normalized IE $\bar{S}_{IE}/\ln(0.48D)$ vs g/ω for $g = 6$, $U = 2g$, $u = 0.5$, and $L = 12, 14, 16$, and 18 . The solid, dashed, dotted, and dot-dashed gray vertical lines correspond to the (1) full resonance, (2) full DT, (3) partial resonance, and (4) partial DT frequencies, respectively. The two dashed red lines in panel (a) correspond to $\bar{r}_{\text{COE}}$ and $\bar{r}_p$, respectively. (c) and (d) The enlarged view of the curves in panels (a) and (b) around the DT point $g/\omega = 2$, respectively. The energy unit is set as $J = 1$.

18. It shows a rich peak-valley structure with the peak value close to $\bar{r}_{\text{COE}}$ and the valley value approaching $\bar{r}_p$. This differs from the monotonous Floquet thermalization-MBL transition in driven disordered many-body localized systems [30,31,33]. Remarkably, the positions of the peaks (subpeaks) and valleys (subvalleys) well match the full resonance (partial resonance) and full DT (partial DT) frequencies, respectively, which means that the thermalization behavior is well captured by the first-order processes and thermalization is supported and suppressed by the resonance and DT, respectively. Particularly, at the first full DT frequency $\omega = g/2$, $\bar{r}$ is well consistent with $\bar{r}_p$, which means the system is localized. The average $\bar{r}$ as a whole increases and tends to the COE predicted value as ω decreases. This is because the higher-order terms become important as ω decreases [47].

The thermalization suppression is also reflected in the properties of the Floquet states $\{|\psi_\alpha\rangle\}$. The Shannon information entropy (IE) of a Floquet state $|\psi_\alpha\rangle$ is defined as

$$S_{\text{IE}}(|\psi_\alpha\rangle) = - \sum_{m=1}^D |\langle \phi_m | \psi_\alpha \rangle|^2 \ln |\langle \phi_m | \psi_\alpha \rangle|^2, \quad (7)$$

where $\{|\phi_m\rangle\}$ are the eigenstates of the undriven Hamiltonian $\hat{H}_0 = \hat{H}_{\text{on}} + \hat{H}_J$ and $\mathcal{D}$ is the Hilbert space dimension [23,50,51]. It measures the degree of localization of the Floquet state among $\{|\phi_m\rangle\}$. If the system satisfies the Floquet ETH, all the Floquet states distribute uniformly across $|\phi_m\rangle$ and S_{IE} tends to the value $\ln(0.48D)$ [23]. If the system does not react to the driving, the Floquet Hamiltonian is akin to $\hat{H}_0$ and S_{IE} tends to 0. In Fig. 1(b), we plot the averaged and normalized Shannon IE $\bar{S}_{IE}/\ln(0.48D)$ as a function of g/ω at the same parameters in Fig. 1(a). $\bar{S}_{IE}$ shows consistent peak-valley behavior with that of $\bar{r}$. Also, $\bar{S}_{IE}$ tends to $\ln(0.48D)$ and 0 at the full resonance frequencies and the first full DT frequency $\omega = g/2$, respectively.

In Fig. 2, we plot the Floquet states' expectation values of a local observable $\hat{n}_8$ at the full DT frequency $\omega = g/2$

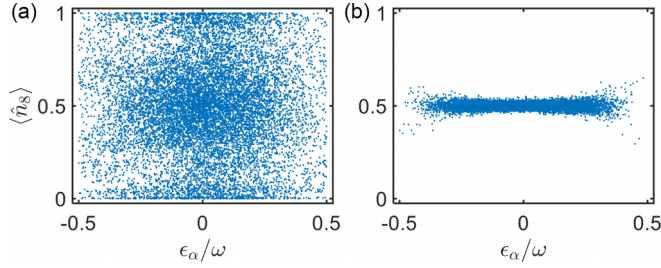


FIG. 2. The expectation value of $\langle \hat{n}_8 \rangle$ vs the scaled Floquet quasienergy ϵ_α/ω for (a) $\omega = g/2$ and (b) $\omega = g$. The other parameters are $g = 6$, $U = 2g$, $u = 0.5$, and $L = 16$. The energy unit is set as $J = 1$.

and the full resonance frequency $\omega = g$. At $\omega = g$, the expectation values $\langle \hat{n}_8 \rangle = \langle \psi_\alpha | \hat{n}_8 | \psi_\alpha \rangle$ are similar for all Floquet states and concentrate on the infinite-temperature ensemble predicted value 0.5, which is consistent with the Floquet ETH [22,23,30]. In contrast, at $\omega = g/2$, the expectation values $\langle \hat{n}_8 \rangle$ fluctuate widely among the Floquet states, which means violation of the Floquet ETH. All the above properties illustrate that Floquet thermalization can be suppressed through the DT and probably mostly suppressed at the first DT frequency $\omega = g/2$.

We now study the finite-size effect on the nonthermal properties at the DT point. For this purpose, we study the Floquet spectrum for system sizes $L = 12, 14, 16$, and 18 . To make a clearer quantitative understanding of the behavior of local observables, we calculate the mean difference of local density between adjacent Floquet states:

$$\bar{\mathcal{M}} = \frac{1}{(\mathcal{D}-1)L} \sum_{\alpha=1}^{\mathcal{D}-1} \sum_{j=1}^L |\langle \psi_{\alpha+1} | \hat{n}_j | \psi_{\alpha+1} \rangle - \langle \psi_\alpha | \hat{n}_j | \psi_\alpha \rangle|. \quad (8)$$

At the DT point $g/\omega = 2$, $\bar{r}$ and $\bar{S}_{\text{IE}}$ increase with increasing L (for $\bar{r}$, this increasing relationship holds for $L > 12$); see Figs. 1(c) and 1(d), respectively. This behavior of $\bar{r}$ is different from that in the stable nonthermal phase, where $\bar{r}$ decreases with increasing L and converges to the Poisson statistic predicted value ~ 0.38 [52]. $\bar{\mathcal{M}}$ decreases exponentially with L , which is much slower than that at the resonance condition $g/\omega = 1$ and the full ETH predicted scaling $\sim 1/\sqrt{\mathcal{D}}$ [53,54]; see Fig. 3(b). These results mean that at the DT point, the nonthermal behavior is not robust for increasing system size

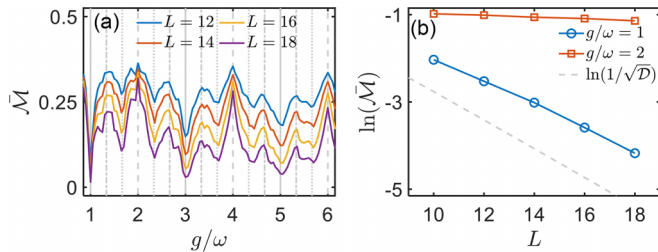


FIG. 3. (a) $\bar{\mathcal{M}}$ vs g/ω for $L = 12, 14, 16$, and 18 . (b) $\ln(\bar{\mathcal{M}})$ vs L at $g/\omega = 1$ and 2 . The gray dashed line represents $\ln(1/\sqrt{\mathcal{D}})$ vs L . The other parameters are $g = 6$, $U = 2g$, and $u = 0.5$. The energy unit is set as $J = 1$.

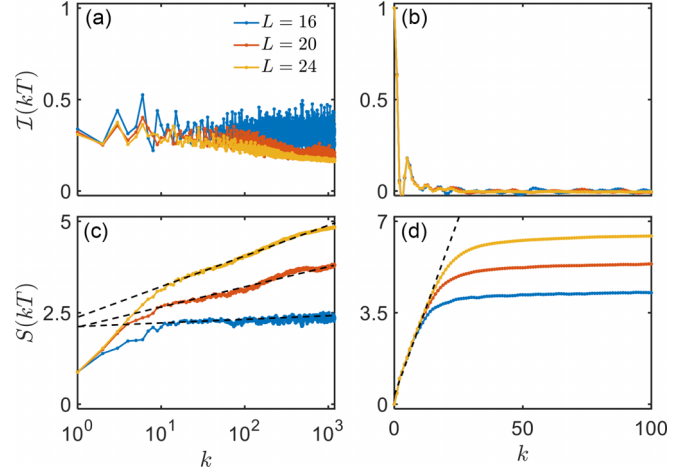


FIG. 4. Dynamical evolution of (a) and (b) the imbalance $\mathcal{I}(kT)$ and (c) and (d) the half-chain EE $S(kT)$ for $\omega = g/2$ (the left column) and $\omega = g$ (the right column), when quenching from the charge density wave state $|10 \dots 10\rangle$. The black dashed lines in panels (c) and (d) exhibit the logarithmic and linear growth of the EE, respectively. The other parameters are $g = 6$, $U = 2g$, and $u = 0.5$. The energy unit is set as $J = 1$.

and the system is not real integrable. Around the DT point, not all the curves of $\bar{r}$ for different L intersect [see Fig. 1(c)]; the curves of $\bar{S}_{\text{IE}}$ and $\bar{\mathcal{M}}$ move up and down as L increases, respectively [see Figs. 1(d) and 3(a), respectively], which means that there does not exist a nonthermal parameter range.

B. Nonthermal dynamics

We now study the dynamical evolution to further illustrate the thermalization suppression through the DT. We evolve the system using the exact diagonalization method when $L < 20$ and the time-evolving block decimation (TEBD) method [55,56] when $L \geq 20$. In the TEBD algorithm, we use a second-order Trotter decomposition with time step $\Delta t = 2\pi/J \times 10^{-3}$ to discretize the time evolution operator. The bond dimension at each step is adjusted such that the discarded singular value weight is smaller than 10^{-7} at each singular value decomposition. We consider quenching from the charge density wave state $|\text{CDW}\rangle = |10 \dots 10\rangle$ with odd lattice sites occupied and even lattice sites empty, and investigate the evolution of the density imbalance

$$\mathcal{I}(kT) = \frac{1}{N} [\langle \psi(kT) | \hat{N}_o | \psi(kT) \rangle - \langle \psi(kT) | \hat{N}_e | \psi(kT) \rangle] \quad (9)$$

and the bipartite von Neumann entanglement entropy

$$S(kT) = -\text{Tr}[\hat{\rho}_A(kT) \ln \hat{\rho}_A(kT)]. \quad (10)$$

Here, $|\psi(kT)\rangle = e^{-i\hat{H}_F kT} |\text{CDW}\rangle$ is the evolving state at time kT , $\hat{N}_{o(e)} = \sum_{j \in \text{odd(even)}} \hat{n}_j$ is the total particle number operator of the odd (even) sites, and $\hat{\rho}_A(kT) = \text{Tr}_{A^c} [|\psi(kT)\rangle \langle \psi(kT)|]$ is the reduced density matrix of subsystem A , with A being the left half of the chain and $\text{Tr}_{A^c}(\dots)$ denoting the partial trace over the complementary sites of A .

In Figs. 4(a) and 4(b), we plot the evolution of $\mathcal{I}(kT)$ at the full DT frequency $\omega = g/2$ and the full resonance frequency $\omega = g$, respectively. In the full resonance condition,

the imbalance quickly decays to zero for all system sizes, which is a hallmark of thermalization dynamics. In the full DT condition, the imbalance remains finite for a long range of time, which means absence of thermalization. For relatively small system size, the imbalance saturates to a finite value. For larger system sizes, the results show that the imbalance does not converge for the numerical accessible time. This means that the suppression of thermalization at the full DT condition may be a prethermal phenomenon for large system sizes, which is consistent with that the DT point does not correspond to a real integrable behavior.

In Figs. 4(c) and 4(d), we plot the evolution of $S(kT)$ at $\omega = g/2$ and $\omega = g$, respectively. In the full resonance condition, the EE initially grows linearly and then rapidly saturates due to finite system sizes. In the full DT condition, the EE grows logarithmically, which is qualitatively similar to that in many-body localized systems [31,57–59] and gives further support of the suppression of Floquet thermalization.

IV. SUMMARY

In summary, we have discovered a driving transparency mechanism to suppress Floquet thermalization in periodically driven tilted lattices of interacting spinless fermions. We analytically derive the resonance conditions and the driving transparency conditions for specific tunneling processes through the Floquet time-dependent perturbation theory. The system can achieve full transparency to the driving due to the interplay between the tilting and the interaction. Under the full

driving transparency conditions, the system cannot directly absorb energy from the driving, thereby significantly suppressing Floquet thermalization. This suppression of Floquet thermalization is evidenced by the properties of the Floquet spectrum and the dynamical evolution, which exhibit violation of the Floquet ETH, retention of the initial state, and logarithmic growth of entanglement entropy. Our findings offer an approach to suppress Floquet thermalization and may lead to the emergence of exotic Floquet nonequilibrium states in tilted lattice systems, such as Floquet time crystal and Floquet topological insulator.

ACKNOWLEDGMENTS

This study was supported by the National Natural Science Foundation of China (Grants No. 12305048 and No. 12025509), Shenzhen Fundamental Research Project (Grant No. JCYJ20230808105009018), the National Key Research and Development Program of China (Grant No. 2022YFA1404104), and Guangdong Provincial Quantum Science Strategic Initiative (GDZX2305006). Y.K. was supported by the National Natural Science Foundation of China (Grant No. 12275365) and the Natural Science Foundation of Guangdong (Grant No. 2023A1515012099).

DATA AVAILABILITY

The data supporting this study's findings are available within the article.

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