Dissipation-Induced Extended-Localized Transition

Yaru Liu,[1,2] Zeqing Wang,[2] Chao Yang,[3,4,5] Jianwen Jie,[1,*] and Yucheng Wang[3,4,5,†]
[1]Shenzhen Key Laboratory of Ultraintense Laser and Advanced Material Technology, Center for Intense Laser Application Technology, and College of Engineering Physics, Shenzhen Technology University, Shenzhen 518118, China
[2]Department of Physics, Renmin University of China, Beijing 100872, China
[3]Shenzhen Institute for Quantum Science and Engineering, Southern University of Science and Technology, Shenzhen 518055, China
[4]International Quantum Academy, Shenzhen 518048, China
[5]Guangdong Provincial Key Laboratory of Quantum Science and Engineering, Southern University of Science and Technology, Shenzhen 518055, China

(Received 26 October 2023; revised 5 February 2024; accepted 23 April 2024; published 21 May 2024)

A mobility edge (ME), representing the critical energy that distinguishes between extended and localized states, is a key concept in understanding the transition between extended (metallic) and localized (insulating) states in disordered and quasiperiodic systems. Here we explore the impact of dissipation on a quasiperiodic system featuring MEs by calculating steady-state density matrix and analyzing quench dynamics with sudden introduction of dissipation. We demonstrate that dissipation can lead the system into specific states predominantly characterized by either extended or localized states, irrespective of the initial state. Our results establish the use of dissipation as a new avenue for inducing transitions between extended and localized states and for manipulating dynamic behaviors of particles.

DOI: 10.1103/PhysRevLett.132.216301

Introduction.—The investigation of electronic transport properties lies at the heart of condensed-matter physics [1]. Disorder is ubiquitous, and its crucial impact on transport properties was unveiled through Anderson localization (AL) [2,3]. In three-dimensional systems with substantial disorder, a transition can occur from the extended phase to the localized phase. Near this transition point, mobility edges (MEs) may emerge, defining the critical energy that distinguishes extended states from localized ones [3–5]. ME is a vital focus in studying disordered materials and helps in understanding a material's conductivity and electronic properties. The extended-localized (metal-insulator) transition can be induced by altering the position of the Fermi energy across the ME. In addition to random disorder, quasiperiodic potentials can also induce the extended-localized transition (ELT) [6–22], resulting in distinct physical phenomena different from that of disordered potentials. For instance, in one-dimensional (1D) quasiperiodic systems, the ELT and MEs can exist, whereas in disordered systems, these phenomena are expected to occur in dimensions higher than 2 according to scaling theory [23].

With advancements in non-Hermitian physics and the manipulation of both dissipation and quantum coherence in experimental settings, recent years have witnessed a growing interest in studying dissipative open quantum systems. Dissipation can profoundly change the properties of quantum systems, leading to various types of phase transitions [24–34]. The impact of dissipation on the localization and transport properties in disordered [35–44] and quasiperiodic systems [45–51] has also garnered widespread attention. It has been observed that dephasing noise can transform other forms of transport behavior into diffusive [46–48], thus, for subdiffusive and localized systems, dephasing can destroy AL and enhance transport. When considering the presence of MEs, the coupling of the system boundaries to baths may significantly increase environment-assisted quantum transport [49] and energy current rectification [50].

Recently, Yusipov and co-workers applied a dissipative operator [given by Eq. (4) below] to a disordered system and discovered that it can drive AL into a stable state that retains its localized properties without being destroyed [42,43]. In this Letter, we investigate the impact of such dissipation on a 1D quasiperiodic system with MEs and find that the dissipation can drive the system into specific states, which may be extended or localized, regardless of the initial state. This reveals that the dissipation can induce the transition between extended and localized states. Such effects are not achievable with other types of dissipation, such as dephasing, energy decay, particle number decay, and so on. Remarkably, the ELT here does not necessitate the change in either disorder strength or particle density, both of which are believed to be the ways for altering the properties of localization through shifting the relative positions of the Fermi energy and the ME. Thus, the

combination of such dissipation and MEs provides a new approach to induce ELT and manipulate a system's transport properties.

Model.—We consider a dissipative system whose density matrix ρ follows the Lindblad master equation [52,53],

$$\frac{d\rho(t)}{dt} = \mathcal{L}[\rho(t)] = -i[H, \rho(t)] + \mathcal{D}[\rho(t)], \tag{1}$$

where $\mathcal{L}$ is referred to as the Lindbladian, with its dissipative component denoted as

$$\mathcal{D}[\rho(t)] = \Gamma \sum_j \left(O_j \rho O_j^\dagger - 1/2\{O_j^\dagger O_j, \rho\} \right), \tag{2}$$

which contains a set of jump operators O_j, all with the same strength Γ here. Assuming that $\mathcal{L}$ is time independent, we can express $\rho(t) = e^{\mathcal{L}t}\rho(0)$. One can define the steady state as $\rho_{ss} = \lim_{t\to\infty} \rho(t)$, which corresponds to the eigenstate of the Lindblad generator with zero eigenvalue, i.e., $\mathcal{L}[\rho_{ss}] = 0$.

The Hamiltonian we consider in Eq. (1) is denoted as

$$H = J \sum_{j=1} \left(c_j^\dagger c_{j+1} + \text{H.c.} \right) + 2 \sum_{j=1} V_j n_j, \tag{3}$$

where c_j and $n_j = c_j^\dagger c_j$ are, respectively, the annihilation operator and local number operator at site j, and J is the nearest neighbor hopping coefficient. The local potential V_j takes $V \cos[2\pi\omega j + \theta]$ for even sites and vanishes for odd sites, where ω is an irrational number, and V and θ are potential amplitude and phase offset, respectively. The model is referred to as the 1D quasiperiodic mosaic model, which features two MEs at $E_c = \pm J/V$ [13]. These two MEs divide the energy spectrum into three regions, including one extended region ($-J/V < E < J/V$) and two localized regions corresponding to $E > J/V$ and $E < -J/V$, respectively. This model has been recently realized, and the MEs have been detected [54]. Without loss of generality, we set $J = 1$, $V = 1$, $\theta = 0$, and $\omega = (\sqrt{5} - 1)/2$. Unless otherwise stated, we use open boundary conditions in subsequent calculations.

The jump operator considered in Eq. (2) is given by [42,43,55–60]

$$O_j = (c_j^\dagger + e^{i\alpha} c_{j+l}^\dagger)(c_j - e^{i\alpha} c_{j+l}), \tag{4}$$

which acts on a pair of sites j and $j + l$. This jump operator does not alter the system's particle number, but it does change the relative phase between the pair of sites with the distance l. For example, this operator synchronizes them from an out-of-phase (in-phase) mode to an in-phase (out-of-phase) mode when the dissipative phase α is set to 0 (π). Such property is important for attaining the intended extended or localized stationary states, as will become evident below.

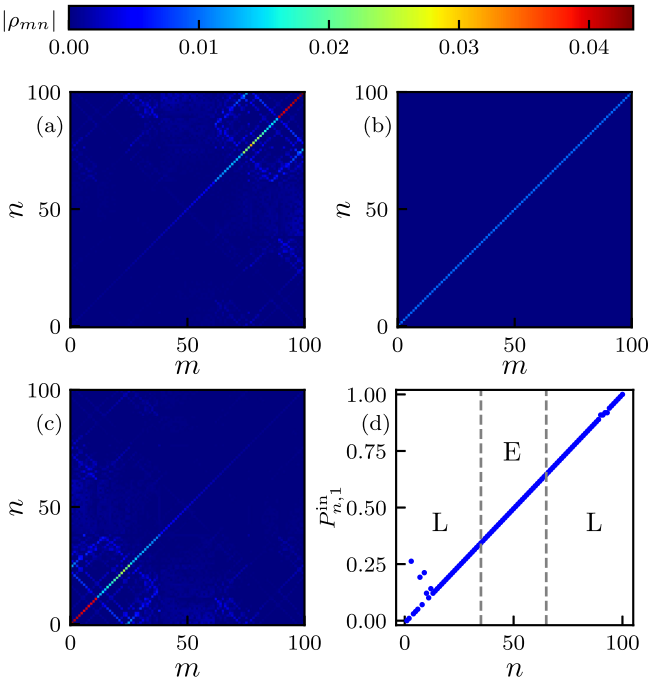

FIG. 1. The absolute values of the density matrix elements for steady states with the dissipative phases (a) $\alpha = 0$, (b) $\alpha = \pi/2$, and (c) $\alpha = \pi$ in the eigenbasis of the Hamiltonian H. (d) The proportion of in-phase site pairs for each eigenstate. Dashed lines mark the MEs, separating eigenstates into localized (L), extended (E), and localized (L) regions as eigenvalues increase. Here we take $l = 1$, $L = 100$, and $\Gamma = 1$.

Dissipation-induced ELT.—We begin by examining the jump operators in Eq. (4) for $l = 1$ and analyzing the properties of the stationary solution ρ_{ss} in the eigenbasis of the Hamiltonian H, that is, $\rho_{mn} = \langle\psi_m|\rho_{ss}|\psi_n\rangle$, where $|\psi_m\rangle$ and $|\psi_n\rangle$ are the eigenstates of the Hamiltonian H. Figures 1(a)–1(c) illustrate the transition of the system's steady state from being predominantly composed of high-energy localized states to predominantly composed of low-energy localized states as the dissipation phase is varied from $\alpha = 0$ to $\alpha = \pi$. The steady state here is independent of the initial state. This implies that, if the initial state is within the extended region located in the middle of the energy spectrum, when the jump operator with $\alpha = 0$ or $\alpha = \pi$ is introduced, the state will ultimately predominantly concentrate on localized regions.

The composition of the steady states can be understood by analyzing the relative phases of neighboring lattice sites. For an arbitrary nth eigenstate $|\psi_n\rangle = \sum_j^L \psi_{n,j} c_j^\dagger |\varnothing\rangle$ with L being the system size, we can calculate the phase difference $\Delta\phi_{i,l}^n$ between the ith and the $(i + l)$th lattice points as $\Delta\phi_{j,l}^n = \arg(\psi_{n,j}) - \arg(\psi_{n,j+l})$. If $\Delta\phi_{j,l}^n = 0$, it implies that they are in phase. Therefore, we can calculate the number of in-phase site pairs $N_{n,l}^{\text{in}}$ with distance l, and

its proportion $P^{\text{in}}_{n,l} = N^{\text{in}}_{n,l}/N_t$, where $N_t = L - l$ represents the total number of site pairs. Figure 1(d) illustrates that eigenstates with higher (lower) energy tends to have larger (smaller) $P^{\text{in}}_{n,1}$ [61], explaining why the steady state predominantly concentrates on high-energy (low-energy) localized eigenstates for $\alpha = 0(\alpha = \pi)$.

We further investigate the effect of the jump operators in Eq. (4) with $l = 2$. We first calculate the proportion of in-phase lattice site pairs $P^{\text{in}}_{n,2}$ and find that it exhibits a V-shaped pattern [Fig. 2(a)]. The localized states on both sides of the energy spectrum tend to exhibit more in-phase site pairs, while the extended states in the middle of the spectrum tend to have more out-of-phase site pairs [61]. Therefore, by appropriately choosing the dissipation phase α, it is possible to control whether the system's steady state is predominantly composed of localized or extended eigenstates, as shown in Figs. 2(b)–2(d). When the dissipative phase is $\alpha = 0$ [Fig. 2(b)], the system is expected to reach a steady state predominantly composed of the states associated with in-phase site pairs, thus primarily concentrating on the localized eigenstates in both higher- and lower-energy regions. Conversely, when $\alpha = \pi$ [Fig. 2(d)], the system is anticipated to attain a steady state mainly composed of the states linked to out-of-phase site pairs, favoring the dominance of extended eigenstates in the midenergy regions. When $\alpha = \pi/2$ [Figs. 2(c) and 1(b)],

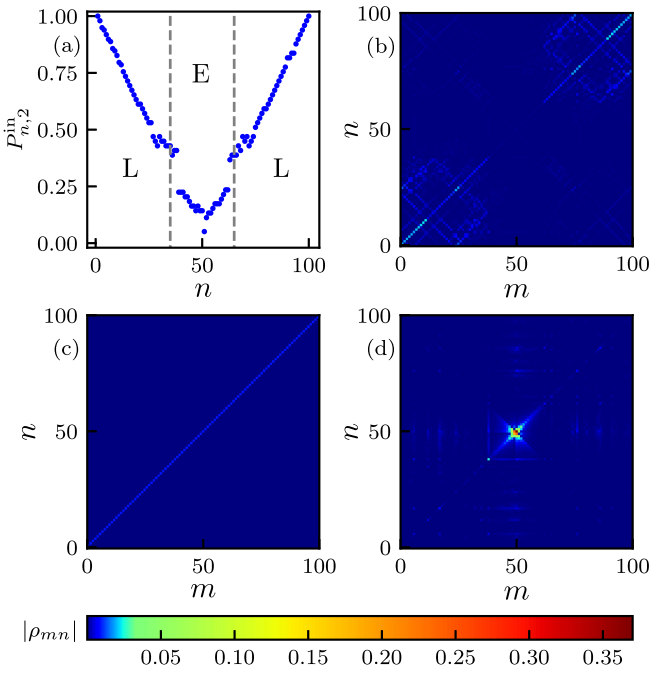

FIG. 2. (a) The proportion of in-phase lattice site pairs for each eigenstate, with the mobility edges marked by the dashed lines. Like in Fig. 1, the dashed lines here represent MEs. The absolute values of the steady-state density matrix elements with the dissipative phases (b) $\alpha = 0$, (c) $\alpha = \pi/2$, and (d) $\alpha = \pi$ in the eigenbasis of Hamiltonian H. Other parameters are $l = 2$, $L = 100$, and $\Gamma = 1$.

the dissipation operator becomes Hermitian, leading to the system reaching the maximally mixed state $(\rho_{ss})_{mn} = \delta_{mn}/L$ as its steady state. According to the diagonal elements of the density matrix ρ_{mn}, for arbitrary α, we can determine the fractions of localized and extended eigenstates in steady states, i.e., $P_l = \sum_i \rho_{ii}$ ($P_e = \sum_j \rho_{jj}$), where $i(j)$ represents the index of the extended (localized) eigenstates $|\psi_i\rangle$ ($|\psi_j\rangle$). When α is tuned from 0 to π, the system's steady state shows the transition from being dominated by localized eigenstates to being dominated by extended eigenstates [Fig. 3(a)], which indicates that dissipation can be used to manipulate the ELT. In Supplemental Material [61], we investigated the variations of P_l and P_e with system size. It is observed that, when $\alpha = 0$, P_e and P_l show minimal changes with size, while for $\alpha = \pi$, as the size increases, P_l (P_e) gradually tends toward 0 (1).

We further examine this transition from a dynamical perspective. We prepare a localized or extended eigenstate as the initial state and introduce dissipation with $l = 2$ at $t = 0$. We then calculate the fidelity, which represents the overlap between the time-evolved state $\rho(t)$ and the initial state ρ_i, denoted as [64,65]

$$F[\rho(t), \rho_i] = \text{Tr}\left[\sqrt{\rho(t)^{1/2}\rho_i\rho(t)^{1/2}}\right]. \tag{5}$$

At $\alpha = 0$ [Fig. 3(b)], the fidelity rapidly approaches zero when the initial state is extended, indicating that the

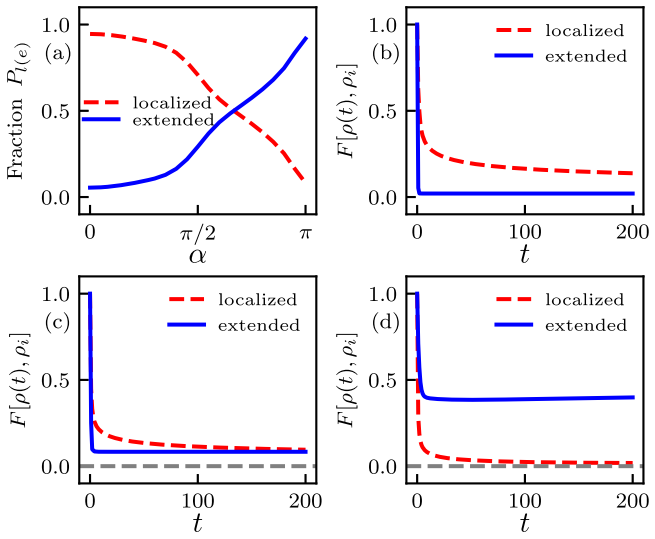

FIG. 3. (a) The fractions of localized eigenstates (P_l) and extended eigenstates (P_e) in steady states as a function of the dissipative phase α. Evolution of the fidelity defined in Eq. (5) after a sudden introduction of dissipation with the strength $\Gamma = 1$ and the phases (b) $\alpha = 0$, (c) $\alpha = \pi/2$, and (d) $\alpha = \pi$. The initial states are set as follows: the system's ground state, which is localized (red dashed line), and the state corresponding to the eigenvalue situated in the center of the energy spectrum of the Hamiltonian H, characterized as extended (blue solid line). Here, $L = 144$ and $l = 2$.

structure of the initial state has been completely modified. Conversely, the fidelity tends to a nonzero value when the initial state is localized, suggesting that certain characteristics of the initial state are preserved. This is because the steady state is primarily composed of localized states, i.e., $P_l \gg P_e$. When $\alpha = \pi/2$ [Fig. 3(c)], the steady state contains extended and localized states, resulting in nonzero fidelity for both. Finally, for $\alpha = \pi$ [Fig. 3(d)], the steady state primarily consists of extended states. Therefore, when the initial state is an extended state, the fidelity is not equal to zero, but it tends to zero when the initial state is localized.

We have revealed that an extended (localized) state can be guided toward a steady state primarily composed of localized (extended) states by applying dissipation. Furthermore, even after removing the dissipation once the system has reached a steady state, its properties continue to persist [66]. This differs from the majority of previous studies on the influence of dissipation on AL, where dissipation disrupts localization, but removing it leads the disordered system back to localization [67]. By introducing a period of dissipation and subsequently removing it, the parameters in the Hamiltonian remain unchanged, but the dynamical properties undergo a profound transformation, as shown in Fig. 4. Consequently, dissipation provides a means to manipulate transitions between localized and extended states.

Experimental realization.—Reference [13] has proposed realizing the quasiperiodic mosaic model based on optical Raman lattices [68–71]. They constructed a spin-1/2 system using two hyperfine states $|\uparrow\rangle = |F_1, m_{F1}\rangle$ and $|\downarrow\rangle = |F_2, m_{F2}\rangle$, and mapped the spin-up (spin-down) lattice sites to the odd (even) sites (see Fig. 5). They applied the Raman coupling via a standing wave field and a plane wave to generate the spin-dependent primary lattice $V_p(x)\sigma_z = V_1 \cos(2k_p x + \phi_p)\sigma_z$ and three standing wave fields together to generate a secondary quasiperiodic lattice $V_s(x) = V_2 \cos(2k_s x + \phi_s)$ only for spin-down atoms. As the depth of the primary lattice significantly exceeds that of the secondary lattice ($V_1 \gg V_2$), we will proceed to discuss the realization of the dissipation operator below, with the secondary lattice's impact considered negligible. We will primarily focus on discussing the realization of the case with $\alpha = 0$ by introducing the auxiliary lattice [56–59]. An arbitrary phase α can be achieved, for instance, through an array of resonators coupled by superconducting qubits [42,60].

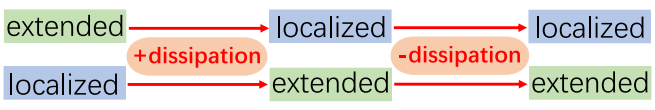

FIG. 4. Schematic illustrating transitions between extended and localized states manipulated by dissipation.

Figure 5(a) shows the realization of the pairwise dissipator with $l = 1$. By coherently coupling two nearly degenerate levels in the system to an auxiliary site in between with antisymmetric Rabi frequencies $\pm\Omega$, which can be obtained by controlling the wavelength of the driving laser to match that of the primary lattice, one can achieve the annihilation part of the dissipative operator. Decay back to the lower sites occurs through spontaneous emission, and this process is isotropic, leading to the form of the creation operator being symmetric [56–59]. In this way, one can realize the jump operator of the form $(c_j^\dagger + c_{j+1}^\dagger)(c_j - c_{j+1})$.

Similarly, the case for $l = 2$ can also be realized, as shown in Fig. 5(b). However, there are several key differences with the case of $l = 1$. (1) The auxiliary lattice must be spin dependent, necessitating the use of two hyperfine states $|\uparrow\rangle = |F_1', m_{F1}'\rangle$ and $|\downarrow\rangle = |F_2', m_{F2}'\rangle$ to construct a spin-1/2 lattice. It is essential to ensure that they satisfy the condition $m_{F1} - m_{F1}' = m_{F2} - m_{F2}'$. (2) It requires a phase difference of π with the primary lattice, namely, that the odd (even) sites correspond to spin-down (spin-up) sites in the auxiliary lattice. (3) By manipulating the polarization of the driving laser, one can attain the specific coupling that spin-up (spin-down) states in the primary lattice only couple with spin-up (spin-down) states in the auxiliary lattice [72,73]. For example, if $m_{F1} - m_{F1}' = 0$, the driving laser needs to be π polarized [74,75]. Moreover, to obtain a π phase shift in the effective Rabi frequency Ω from one spin-up (spin-down) site to the next spin-up (spin-down) site, one needs to set the driving

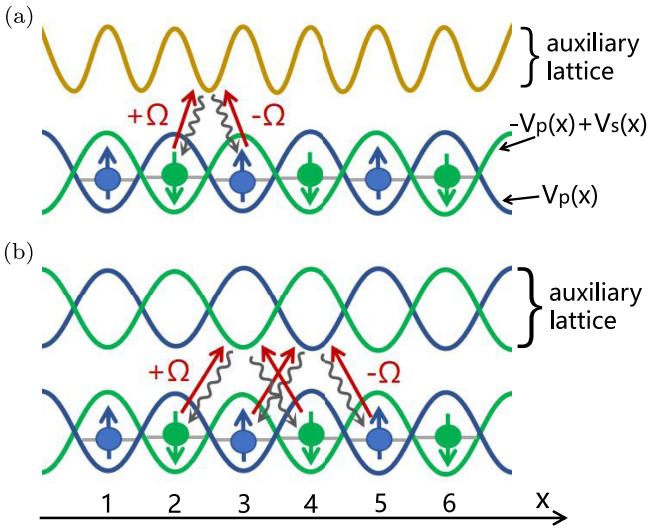

FIG. 5. Schematic realization of the dissipative process of the form $(c_j^\dagger + c_{j+l}^\dagger)(c_j - c_{j+l})$ with (a) $l = 1$ and (b) $l = 2$. The lower (upper) lattice correspond to the physical (auxiliary) lattice. Atoms in the lower sites are coupled to the auxiliary sites in between with opposite Rabi frequency $\pm\Omega$ and then spontaneously decay back to the lower sites.

laser's wavelength to be twice that of the standing wave laser generating the primary lattice.

By comparing the diffusion [76–79] or transport behaviors [15,80–82] of atoms before and after the introduction of dissipation, one can obtain extended and localized information about the initial and final states, thereby detecting the ELT caused by dissipation.

Conclusion and discussion.—We have investigated the influence of dissipation on the 1D quasiperiodic mosaic model, which possesses exact MEs, and proposed its experimental realization. By calculating the distribution of the steady-state density matrix and the characteristics of quench dynamics, we revealed that dissipation can drive the system into specific states primarily composed of either extended or localized states, regardless of the initial states. Hence, dissipation can be utilized to induce transitions between extended and localized states, thereby enabling the manipulation of particle transport behaviors. In addition to its applications in condensed-matter physics, our results also have potential applications in quantum simulation, specifically in controlling the dynamical behavior of particles and preparing desired states. Given a specific example, when simulating systems with MEs using cold atoms, preparing atoms near the ground state is challenging when they are localized, but introducing dissipation can aid in achieving this goal.

The manipulation of the ELT is expected to be common in various systems featuring MEs [61]. When considering the impact of this dissipation on the anomalous MEs separating critical states from localized or extended ones, it is found that dissipation can induce critical-localized or critical-extended transitions (see Supplemental Material [61]). Our results pose several interesting issues. How do such dissipation affect a many-body system with MEs? Can the dissipation be employed to manipulate the transitions between thermalized states and many-body localized states? Can the similar ELT exist in a three-dimensional dissipative disordered system with MEs? Further, the dissipative operator in Eq. (4) utilizes the phase distribution characteristics of different states in the energy spectrum to select specific states as steady states. This provides a new perspective, suggesting that we can also explore other distinguishing features to construct experimentally feasible dissipative operators for the purpose of selecting different states.

We thank Long Zhang, Xin-Chi Zhou, Shi Yu, and Bing Yang for valuable discussions. This work was supported by National Key R&D Program of China under Grant No. 2022YFA1405800, the National Natural Science Foundation of China (Grants No. 12104210 and No. 12104205), and the Natural Science Foundation of Top Talent of SZTU(GDRC202202). Y. L. acknowledges support from the Fundamental Research Funds for the Central Universities, and the Research Funds of Renmin University of China (22XNH100).

*Corresponding author: Jianwen.Jie1990@gmail.com
†Corresponding author: wangyc3@sustech.edu.cn

[1] A. Lagendijk, B. Tiggelen, and D. S. Wiersma, Fifty years of Anderson localization, Phys. Today **62**, No. 8, 24 (2009).
[2] P. W. Anderson, Absence of diffusion in certain random lattices, Phys. Rev. **109**, 1492 (1958).
[3] P. A. Lee and T. V. Ramakrishnan, Disordered electronic systems, Rev. Mod. Phys. **57**, 287 (1985).
[4] F. Evers and A. D. Mirlin, Anderson transitions, Rev. Mod. Phys. **80**, 1355 (2008).
[5] B. Kramer and A. MacKinnon, Localization: Theory and experiment, Rep. Prog. Phys. **56**, 1469 (1993).
[6] G. Roati, C. D'Errico, L. Fallani, M. Fattori, C. Fort, M. Zaccanti, G. Modugno, M. Modugno, and M. Inguscio, Anderson localization of a non-interacting Bose-Einstein condensate, Nature (London) **453**, 895 (2008).
[7] C. M. Soukoulis and E. N. Economou, Localization in one-dimensional lattices in the presence of incommensurate potentials, Phys. Rev. Lett. **48**, 1043 (1982).
[8] S. Das Sarma, S. He, and X. C. Xie, Mobility edge in a model one-dimensional potential, Phys. Rev. Lett. **61**, 2144 (1988).
[9] J. Biddle, B. Wang, D. J. Priour Jr, and S. Das Sarma, Localization in one-dimensional incommensurate lattices beyond the Aubry-André model, Phys. Rev. A **80**, 021603(R) (2009); J. Biddle and S. Das Sarma, Predicted mobility edges in one-dimensional incommensurate optical lattices: An exactly solvable model of Anderson localization, Phys. Rev. Lett. **104**, 070601 (2010); D. Vu and S. Das Sarma, Generic mobility edges in several classes of duality-breaking one-dimensional quasiperiodic potentials, Phys. Rev. B **107**, 224206 (2023).
[10] X. Li, X. Li, and S. Das Sarma, Mobility edges in one dimensional bichromatic incommensurate potentials, Phys. Rev. B **96**, 085119 (2017).
[11] H. Yao, H. Khoudli, L. Bresque, and L. Sanchez-Palencia, Critical behavior and fractality in shallow one-dimensional quasiperiodic potentials, Phys. Rev. Lett. **123**, 070405 (2019).
[12] S. Ganeshan, J. H. Pixley, and S. Das Sarma, Nearest neighbor tight binding models with an exact mobility edge in one dimension, Phys. Rev. Lett. **114**, 146601 (2015).
[13] Y. Wang, X. Xia, L. Zhang, H. Yao, S. Chen, J. You, Q. Zhou, and X.-J. Liu, One dimensional quasiperiodic mosaic lattice with exact mobility edges, Phys. Rev. Lett. **125**, 196604 (2020).
[14] X.-C. Zhou, Y. Wang, T.-F. J. Poon, Q. Zhou, and X.-J. Liu, Exact new mobility edges between critical and localized states, Phys. Rev. Lett. **131**, 176401 (2023).
[15] Y. Wang, L. Zhang, W. Sun, T.-F. J. Poon, and X.-J. Liu, Quantum phase with coexisting localized, extended, and critical zones, Phys. Rev. B **106**, L140203 (2022).
[16] T. Liu, X. Xia, S. Longhi, and L. Sanchez-Palencia, Anomalous mobility edges in one-dimensional quasiperiodic models, SciPost Phys. **12**, 027 (2022).
[17] S. Longhi, Topological phase transition in non-Hermitian quasicrystals, Phys. Rev. Lett. **122**, 237601 (2019); S. Longhi, Metal-insulator phase transition in a non-Hermitian Aubry-André-Harper model, Phys. Rev. B **100**, 125157

(2019); S. Longhi, Non-Hermitian control of localization in mosaic photonic lattices, Appl. Phys. Lett. **123**, 161102 (2023).

[18] H. Jiang, L.-J. Lang, C. Yang, S.-L. Zhu, and S. Chen, Interplay of non-Hermitian skin effects and Anderson localization in non-reciprocal quasiperiodic lattices, Phys. Rev. B **100**, 054301 (2019); Y. Liu, Y. Wang, X.-J. Liu, Q. Zhou, and S. Chen, Exact mobility edges, PT-symmetry breaking and skin effect in one-dimensional non-Hermitian quasicrystals, Phys. Rev. B **103**, 014203 (2021).

[19] M. Gonçalves, B. Amorim, E. V. Castro, and P. Ribeiro, Hidden dualities in 1D quasiperiodic lattice models, SciPost Phys. **13**, 046 (2022); M. Gonçalves, B. Amorim, E. V. Castro, and P. Ribeiro, Renormalization-group theory of 1D quasiperiodic lattice models with commensurate approximants, Phys. Rev. B **108**, L100201 (2023); M. Gonçalves, B. Amorim, E. V. Castro, and P. Ribeiro, Critical phase dualities in 1D exactly-solvable quasiperiodic models, Phys. Rev. Lett. **131**, 186303 (2023).

[20] H. P. Lüschen, S. Scherg, T. Kohlert, M. Schreiber, P. Bordia, X. Li, S. D. Sarma, and I. Bloch, Single-particle mobility edge in a one-dimensional quasiperiodic optical lattice, Phys. Rev. Lett. **120**, 160404 (2018); T. Kohlert, S. Scherg, X. Li, H. P. Lüschen, S. D. Sarma, I. Bloch, and M. Aidelsburger, Observation of many-body localization in a one-dimensional system with single-particle mobility edge, Phys. Rev. Lett. **122**, 170403 (2019).

[21] F. A. An, E. J. Meier, and B. Gadway, Engineering a flux-dependent mobility edge in disordered zigzag chains, Phys. Rev. X **8**, 031045 (2018); F. A. An, K. Padavić, E. J. Meier, S. Hegde, S. Ganeshan, J. H. Pixley, S. Vishveshwara, and B. Gadway, Observation of tunable mobility edges in generalized Aubry-André lattices, Phys. Rev. Lett. **126**, 040603 (2021).

[22] Y. Wang, J.-H. Zhang, Y. Li, J. Wu, W. Liu, F. Mei, Y. Hu, L. Xiao, J. Ma, C. Chin, and S. Jia, Observation of interaction-induced mobility edge in an atomic Aubry-André wire, Phys. Rev. Lett. **129**, 103401 (2022).

[23] E. Abrahams, P. W. Anderson, D. C. Licciardello, and T. V. Ramakrishnan, Scaling theory of localization: Absence of quantum diffusion in two dimensions, Phys. Rev. Lett. **42**, 673 (1979); B. Hetényi, S. Parlak, and M. Yahyavi, Scaling and renormalization in the modern theory of polarization: Application to disordered systems, Phys. Rev. B **104**, 214207 (2021).

[24] T. Prosen and I. Pižorn, Quantum phase transition in a far-from-equilibrium steady state of an *XY* spin chain, Phys. Rev. Lett. **101**, 105701 (2008).

[25] H. T. Mebrahtu, I. V. Borzenets, H. Zheng, Y. V. Bomze, A. I. Smirnov, S. Florens, H. U. Baranger, and G. Finkelstein, Observation of Majorana quantum critical behaviour in a resonant level coupled to a dissipative environment, Nat. Phys. **9**, 732 (2013); H. T. Mebrahtu, I. V. Borzenets, D. E. Liu, H. Zheng, Y. V. Bomze, A. I. Smirnov, H. U. Baranger, and G. Finkelstein, Quantum phase transition in a resonant level coupled to interacting leads, Nature (London) **488**, 61 (2012).

[26] M. V. Medvedyeva, M. T. Čubrović, and S. Kehrein, Dissipation-induced first-order decoherence phase transition in a noninteracting fermionic system, Phys. Rev. B **91**, 205416 (2015).

[27] K. Shastri and Francesco Monticone, Dissipation-induced topological transitions in continuous Weyl materials, Phys. Rev. Res. **2**, 033065 (2020).

[28] M. Soriente, T. L. Heugel, K. Arimitsu, R. Chitra, and O. Zilberberg, Distinctive class of dissipation-induced phase transitions and their universal characteristics, Phys. Rev. Res. **3**, 023100 (2021).

[29] K. Yamamoto, M. Nakagawa, N. Tsuji, M. Ueda, and N. Kawakami, Collective excitations and nonequilibrium phase transition in dissipative fermionic superfluids, Phys. Rev. Lett. **127**, 055301 (2021).

[30] W. Nie, M. Antezza, Y.-X. Liu, and F. Nori, Dissipative topological phase transition with strong system-environment coupling, Phys. Rev. Lett. **127**, 250402 (2021).

[31] K. Kawabata, T. Numasawa, and S. Ryu, Entanglement phase transition induced by the non-Hermitian skin effect, Phys. Rev. X **13**, 021007 (2023).

[32] Ezequiel I. Rodriguez Chiacchio, A. Nunnenkamp, and M. Brunelli, Nonreciprocal Dicke model, Phys. Rev. Lett. **131**, 113602 (2023).

[33] L.-N. Wu, J. Nettersheim, J. Feß, A. Schnell, S. Burgardt, S. Hiebel, D. Adam, A. Eckardt, and A. Widera, Dynamical phase transition in an open quantum system, arXiv:2208.05164.

[34] R. Kuzmin, N. Mehta, N. Grabon, R. A. Mencia, A. Burshtein, M. Goldstein, and V. E. Manucharyan, Observation of the Schmid-Bulgadaev dissipative quantum phase transition, arXiv:2304.05806.

[35] S. A. Gurvitz, Delocalization in the Anderson model due to a local measurement, Phys. Rev. Lett. **85**, 812 (2000).

[36] B. Nowak, J. J. Kinnunen, M. J. Holland, and P. Schlagheck, Delocalization of ultracold atoms in a disordered potential due to light scattering, Phys. Rev. A **86**, 043610 (2012).

[37] K. Rayanov, G. Radons, and S. Flach, Decohering localized waves, Phys. Rev. E **88**, 012901 (2013).

[38] A. G. Yamilov, R. Sarma, B. Redding, B. Payne, H. Noh, and H. Cao, Position-dependent diffusion of light in disordered waveguides, Phys. Rev. Lett. **112**, 023904 (2014).

[39] M. Balasubrahmaniyam, S. Mondal, and S. Mujumdar, Necklace-state-mediated anomalous enhancement of transport in Anderson-localized non-Hermitian hybrid systems, Phys. Rev. Lett. **124**, 123901 (2020).

[40] S. Weidemann, M. Kremer, S. Longhi, and A. Szameit, Coexistence of dynamical delocalization and spectral localization through stochastic dissipation, Nat. Photonics **15**, 576 (2021).

[41] D. A. Huse, R. Nandkishore, F. Pietracaprina, V. Ros, and A. Scardicchio, Localized systems coupled to small baths: From Anderson to Zeno, Phys. Rev. B **92**, 014203 (2015).

[42] I. Yusipov, T. Laptyeva, S. Denisov, and M. Ivanchenko, Localization in open quantum systems, Phys. Rev. Lett. **118**, 070402 (2017).

[43] O. S. Vershinina, I. I. Yusipov, S. Denisov, M. V. Ivanchenko, and T. V. Laptyeva, Control of a single-particle localization in open quantum systems, Europhys. Lett. **119**, 56001 (2017); I. I. Yusipov, T. V. Laptyeva, and M. V. Ivanchenko, Quantum jumps on Anderson attractors, Phys.

Rev. B **97**, 020301 (2018); I. Vakulchyk, I. Yusipov, M. Ivanchenko, S. Flach, and S. Denisov, Signatures of many-body localization in steady states of open quantum systems, Phys. Rev. B **98**, 020202(R) (2018).

[44] S. Longhi, Anderson localization in dissipative lattices, Ann. Phys. (Amsterdam) **535**, 2200658 (2023).

[45] A. Purkayastha, A. Dhar, and M. Kulkarni, Nonequilibrium phase diagram of a one-dimensional quasiperiodic system with a single-particle mobility edge, Phys. Rev. B **96**, 180204(R) (2017).

[46] M. Saha, B. P. Venkatesh, and B. K. Agarwalla, Quantum transport in quasiperiodic lattice systems in the presence of Büttiker probes, Phys. Rev. B **105**, 224204 (2022).

[47] A. M. Lacerda, J. Goold, and G. T. Landi, Dephasing enhanced transport in boundary-driven quasiperiodic chains, Phys. Rev. B **104**, 174203 (2021).

[48] C. Chiaracane, A. Purkayastha, M. T. Mitchison, and J. Goold, Dephasing-enhanced performance in quasiperiodic thermal machines, Phys. Rev. B **105**, 134203 (2022).

[49] D. Dwiputra and F. P. Zen, Environment-assisted quantum transport and mobility edges, Phys. Rev. A **104**, 022205 (2021).

[50] V. Balachandran, S. R. Clark, J. Goold, and D. Poletti, Energy current rectification and mobility edges, Phys. Rev. Lett. **123**, 020603 (2019).

[51] C. Chiaracane, M. T. Mitchison, A. Purkayastha, G. Haack, and J. Goold, Quasiperiodic quantum heat engines with a mobility edge, Phys. Rev. Res. **2**, 013093 (2020).

[52] G. Lindblad, On the generators of quantum dynamical semigroups, Commun. Math. Phys. **48**, 119 (1976).

[53] H.-P. Breuer and F. Petruccione, *The Theory of Open Quantum Systems* (Oxford University Press, Oxford, 2002).

[54] J. Gao, I. M. Khaymovich, X.-W. Wang, Z.-S. Xu, A. Iovan, G. Krishna, A. V. Balatsky, V. Zwiller, and A. W. Elshaari, Experimental probe of multi-mobility edges in quasiperiodic mosaic lattices, arXiv:2306.10829.

[55] We note that in the jump operator, Eq. (4), there is no negative sign in front of both α. It is easy to prove that when $\alpha = \pi/2$, the jump operator is Hermitian.

[56] S. Diehl, A. Micheli, A. Kantian, B. Kraus, H. P. Büchler, and P. Zoller, Quantum states and phases in driven open quantum systems with cold atoms, Nat. Phys. **4**, 878 (2008).

[57] B. Kraus, H. P. Büchler, S. Diehl, A. Kantian, A. Micheli, and P. Zoller, Preparation of entangled states by quantum Markov processes, Phys. Rev. A **78**, 042307 (2008).

[58] S. Diehl, A. Tomadin, A. Micheli, R. Fazio, and P. Zoller, Dynamical phase transitions and instabilities in open atomic many-body systems, Phys. Rev. Lett. **105**, 015702 (2010); S. Diehl, E. Rico, M. A. Baranov, and P. Zoller, Topology by dissipation in atomic quantum wires, Nat. Phys. **7**, 971 (2011).

[59] C.-E. Bardyn, M. A. Baranov, C. V. Kraus, E. Rico, A. Imamoğlu, P. Zoller, and S. Diehl, Topology by dissipation, New J. Phys. **15**, 085001 (2013).

[60] D. Marcos, A. Tomadin, S. Diehl, and P. Rabl, Photon condensation in circuit quantum electrodynamics by engineered dissipation, New J. Phys. **14**, 055005 (2012).

[61] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevLett.132.216301 for details on (I) a simple example to understand the distribution of phase differences; (II) P_l and P_e change with the system size; (III) the impact of dissipation on another 1D system with exact MEs; (IV) the impact of dissipation on the systems with anomalous MEs; (V) the impact of dissipation on the Dyson model, which include Refs. [12,14–16,62,63].

[62] X. Deng, S. Ray, S. Sinha, G. V. Shlyapnikov, and L. Santos, One-dimensional quasicrystals with power-law hopping, Phys. Rev. Lett. **123**, 025301 (2019).

[63] F. J. Dyson, The dynamics of a disordered linear chain, Phys. Rev. **92**, 1331 (1953).

[64] M. Nielsen and I. Chuang, *Quantum Computation and Quantum Information* (Cambridge University Press, Cambridge, England, 2000).

[65] P. Zanardi, H. T. Quan, Xiaoguang Wang, and C. P. Sun, Mixed-state fidelity and quantum criticality at finite temperature, Phys. Rev. A **75**, 032109 (2007).

[66] After the system reaches a steady state and dissipation is removed (i.e., Γ becomes 0), it is easy to obtain that the diagonal elements of ρ_{mn} do not change with time. Therefore, the fraction of localized eigenstates (P_l) or extended eigenstates (P_e) remains constant over time. The off-diagonal elements evolve with time according to $\rho_{mn}(t) = e^{i(E_m - E_n)t}\rho_{mn}$, where E_m and E_n are the eigenvalues of the Hamiltonian H. When α is 0 or π, ρ_{mn} is very small when E_m and E_n are in the extended and localized regions, respectively. Furthermore, if we observe the system's dynamical behavior for a long period of time, the effects generated by off-diagonal elements will vanish. In our Letter, we focus on P_l and P_e, thus concentrating only on the properties of the diagonal elements.

[67] The disruption of localization discussed here diverges from the principles employed in previous studies. For example, many previous studies considered dephasing effects, and this disruption occurs by undermining the interference phenomenon crucial for AL. When dissipation is removed, the system's interference effects are restored, leading to the reoccurrence of AL. In our Letter, we disrupt localization by transforming particles from localized to extended states. After removing dissipation, particles stay in extended states, preventing the reappearance of localization.

[68] X.-J. Liu, Z.-X. Liu, and M. Cheng, Manipulating topological edge spins in a one-dimensional optical lattice, Phys. Rev. Lett. **110**, 076401 (2013); X.-J. Liu, K. T. Law, and T. K. Ng, Realization of 2D spin-orbit interaction and exotic topological orders in cold atoms, Phys. Rev. Lett. **112**, 086401 (2014).

[69] B. Song, L. Zhang, C. He, T. F. J. Poon, E. Hajiyev, S. Zhang, X.-J. Liu, and G.-B. Jo, Observation of symmetry-protected topological band with ultracold fermions, Sci. Adv. **4**, eaao4748 (2018); B. Song, C. He, S. Niu, L. Zhang, Z. Ren, X.-J. Liu, and G.-B. Jo, Observation of nodal-line semimetal with ultracold fermions in an optical lattice, Nat. Phys. **15**, 911 (2019).

[70] L. Zhang and X.-J Liu, Spin orbit coupling and topological phases for ultracold atoms, arXiv:1806.05628.

[71] Z. Wu, L. Zhang, W. Sun, X.-T. Xu, B.-Z. Wang, S.-C. Ji, Y. Deng, S. Chen, X.-J. Liu, and J.-W. Pan, Realization of two-dimensional spin-orbit coupling for Bose-Einstein condensates, Science **354**, 83 (2016); Z.-Y. Wang, X.-C. Cheng, B.-Z. Wang, J.-Y. Zhang, Y.-H. Lu, C.-R. Yi,

S. Niu, Y. Deng, X.-J. Liu, S. Chen, and J.-W. Pan, Realization of ideal Weyl semimetal band in ultracold quantum gas with 3D spin-orbit coupling, Science **372**, 271 (2021).

[72] E. Arimondo, M. Inguscio, and P. Violino, Experimental determinations of the hyperfine structure in the alkali atoms, Rev. Mod. Phys. **49**, 31 (1977).

[73] D. Budker, W. Gawlik, D. F. Kimball, S. M. Rochester, V. V. Yashchuk, and A. Weis, Resonant nonlinear magneto-optical effects in atoms, Rev. Mod. Phys. **74**, 1153 (2002).

[74] A. Jenkins, J. W. Lis, A. Senoo, W. F. McGrew, and A. M. Kaufman, Ytterbium nuclear-spin qubits in an optical tweezer array, Phys. Rev. X **12**, 021027 (2022).

[75] N. Chen, L. Li, W. Huie, M. Zhao, I. Vetter, C. H. Greene, and J. P. Covey, Analyzing the Rydberg-based optical-metastable-ground architecture for ^{171}Yb nuclear spins, Phys. Rev. A **105**, 052438 (2022).

[76] R. Ketzmerick, K. Kruse, S. Kraut, and T. Geisel, What determines the spreading of a wave packet?, Phys. Rev. Lett. **79**, 1959 (1997).

[77] M. Larcher, F. Dalfovo, and M. Modugno, Effects of interaction on the diffusion of atomic matter waves in one-dimensional quasiperiodic potentials, Phys. Rev. A **80**, 053606 (2009).

[78] Y. Wang, L. Zhang, S. Niu, D. Yu, and X.-J. Liu, Realization and detection of non-ergodic critical phases in optical Raman lattice, Phys. Rev. Lett. **125**, 073204 (2020).

[79] T. Shimasaki, M. Prichard, H. E. Kondakci, J. E. Pagett, Y. Bai, P. Dotti, A. Cao, T.-C. Lu, T. Grover, and D. M. Weld, Anomalous localization and multifractality in a kicked quasicrystal, arXiv:2203.09442.

[80] M. Saha, S. K. Maiti, and A. Purkayastha, Anomalous transport through algebraically localized states in one-dimension, Phys. Rev. B **100**, 174201 (2019).

[81] J.-P. Brantut, J. Meineke, D. Stadler, S. Krinner, and T. Esslinger, Conduction of ultracold fermions through a mesoscopic channel, Science **337**, 1069 (2012).

[82] C.-C. Chien, S. Peotta, and M. D. Ventra, Quantum transport in ultracold atoms, Nat. Phys. **11**, 998 (2015).