

Many-Body Non-Hermitian Skin Effect with Exact Steady States in the Dissipative Quantum Link Model

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(Received 14 April 2025; revised 26 October 2025; accepted 21 November 2025; published 26 December 2025)

We introduce a dissipative lattice gauge model that exhibits the many-body version of the non-Hermitian skin effect. The dissipative couplings between dynamical gauge fields on the lattice links and the surrounding environment generate chiral motions of particles residing on lattice sites. Despite the complexity arising from many-body interactions, the local gauge symmetry enables the exact construction of a steady state that displays the many-body non-Hermitian skin effect. Furthermore, our approach can be generalized to realize a new type of many-body non-Hermitian skin effect, dubbed the hierarchical skin effect, where different subsystem degrees of freedom exhibit boundary accumulation of multiple moments at different orders. Our findings can be readily observed by engineering dissipation in state-of-the-art lattice gauge simulators.

DOI: [10.1103/wztw-l8wg](https://doi.org/10.1103/wztw-l8wg)

Introduction—Gauge theory is a cornerstone in modern physics. The lattice formulation of gauge theory provides a versatile framework for investigating fundamental phenomena in quantum field theories and describing novel quantum phases in condensed matter physics [1–3]. While it is difficult to simulate a strongly coupled gauge theory on classical computers, the rapid advancement of quantum technologies allows for experimentally investigating equilibrium and nonequilibrium phases of lattice gauge theory on cutting-edge quantum platforms [4–14]. These novel phenomena includes string dynamics [15–21], quantum many-body scars [22–25], disorder-free localization [26–29], and quantum phase transitions [30–33]. A key ingredient for lattice gauge simulators is to impose gauge constraints for quantum dynamics [34–37]. One of the gauge violation mechanisms stems from the incoherent system-environment couplings of quantum devices [34]. Although the environment is detrimental to gauge-invariant coherent dynamics, the interplay between dissipation and many-body dynamics can lead to unexpected nonequilibrium phenomena that has no analog in closed lattice gauge theory [38,39].

One of the striking phenomena arising in non-Hermitian or open quantum systems is the non-Hermitian skin effect (NHSE), which describes the boundary localization of bulk eigenstates [40–46]. This concept is essential in various aspects of non-Hermitian physics [47–52], such as non-Hermitian topology [53–59] and novel dynamics [60–70]. Recently, there has been a growing interest in exploring the significant role of interactions in shaping NHSE in many-body systems. Considerable efforts have been dedicated to characterizing the NHSE of many-body eigenstates [71–76], engineering many-body NHSE [77–84], and unveiling novel nonequilibrium phases induced by many-body NHSE [85–93].

However, apart from a limited number of integrable systems [94–99], the complexity of many-body interactions generally prevents the exact description of non-Hermitian many-body systems, thereby limiting the analysis of many-body NHSE to numerical methods. Therefore, it is desirable to develop new analytical method to understand the interplay between NHSE and interaction. Moreover, the extensive studies of many-body NHSE are often rooted in the many-body extensions of the Hatano-Nelson model, where an imaginary background gauge field causes the nonreciprocal hoppings of particles [100,101]. Compared to experimental platforms for single-particle NHSE [43–46], it is more challenging to implement imaginary gauge fields in quantum many-body experiments. More fundamentally, a physical explanation of imaginary gauge fields directly from a gauge-theoretic perspective is still lacking.

In this Letter, we address these questions by proposing a dissipative extension of the quantum link model (QLM) in the U(1) lattice gauge theory [4,15,102]. With the

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experimentally accessible dissipation processes acting on the dynamical gauge fields at lattice links, the dissipative QLM effectively induces nonreciprocal hoppings for matter fields at lattice sites, resulting in many-body NHSE under open boundary conditions (OBC). Despite the many-body complexity of this open quantum system, the local gauge symmetry of the quantum-link Hamiltonian enables the analytical construction of the *exact* steady state, which exhibits the many-body NHSE characterized by the accumulation of the dipole moment of matter fields. Notably, not limited to dissipative QLM, our symmetry-based method for constructing exact many-body eigenstates is broadly applicable to a wide class of open quantum many-body systems. Furthermore, the dissipative QLM offers a general method for achieving a *hierarchical skin effect*, characterized by multipole moments for certain subsystem degrees of freedom. From an experimental perspective, our model provides an efficient way to implement nonreciprocal hoppings in state-of-the-art lattice gauge simulators.

Dissipative quantum link model—We first review the basic notions of U(1) QLM [15], particularly focusing on open boundary conditions. The QLM of (1 + 1)D quantum electrodynamics has a quantum-link Hamiltonian $H = J \sum_{n=1}^{L-1} (\psi_n^\dagger U_{n,n+1} \psi_{n+1} + \text{H.c.})$. The lattice system has L sites and $L - 1$ links. $\psi_n^\dagger$ and ψ_n are the fermionic creation and annihilation operators of matter fields at sites. $U_{n,n+1}$ is the link operator that represents the U(1) gauge degree of freedom, conjugating to the electric field $E_{n,n+1}$ with the relations $[E_{n,n+1}, U_{n,n+1}] = U_{n,n+1}$ and $[E_{n,n+1}, U_{n,n+1}^\dagger] = -U_{n,n+1}^\dagger$. The generators of the local U(1) gauge symmetry in QLM are given by $G_n = \psi_n^\dagger \psi_n - (E_{n,n+1} - E_{n-1,n})$ for $1 < n < L$. In addition, we have $G_1 = \psi_1^\dagger \psi_1 - E_{1,2}$ and $G_L = \psi_L^\dagger \psi_L + E_{L-1,L}$ at the boundary sites. A gauge fixing condition $(G_n + g_n)|\Phi\rangle = 0$ with background charges g_n plays the role of the Gauss law in electrodynamics. In practical quantum simulations, the quantum links are replaced by spin operators with finite Hilbert space dimensions: $U_{n,n+1} \rightarrow s_{n,n+1}^+$, $U_{n,n+1}^\dagger \rightarrow s_{n,n+1}^-$, $E_{n,n+1} \rightarrow s_{n,n+1}^z$. Additionally, we perform the Jordan-Wigner transformation and employ spin-1/2 operators $\tau_n^{x,y,z}$ to describe fermionic operators on lattice sites. The resulting spin version of QLM becomes

$$H = J \sum_{n=1}^{L-1} (\tau_n^+ s_{n,n+1}^+ \tau_{n+1}^- + \text{H.c.}). \quad (1)$$

With $\psi_n^\dagger \psi_n = \tau_n^z + 1/2$, the occupied (unoccupied) states at lattice sites are described by the eigenvalue of τ_n^z being 1/2 (−1/2). The local gauge generator is equivalent to $G_n = \tau_n^z - (s_{n,n+1}^z - s_{n-1,n}^z)$ for $1 < n < L$, with boundary corrections $G_1 = \tau_1^z - s_{1,2}^z$ and $G_L = \tau_L^z + s_{L-1,L}^z$. Additionally, the total particle number $N = \sum_n \psi_n^\dagger \psi_n = \sum_n (\tau_n^z + 1/2)$ of matter fields is also conserved because of the global U(1)

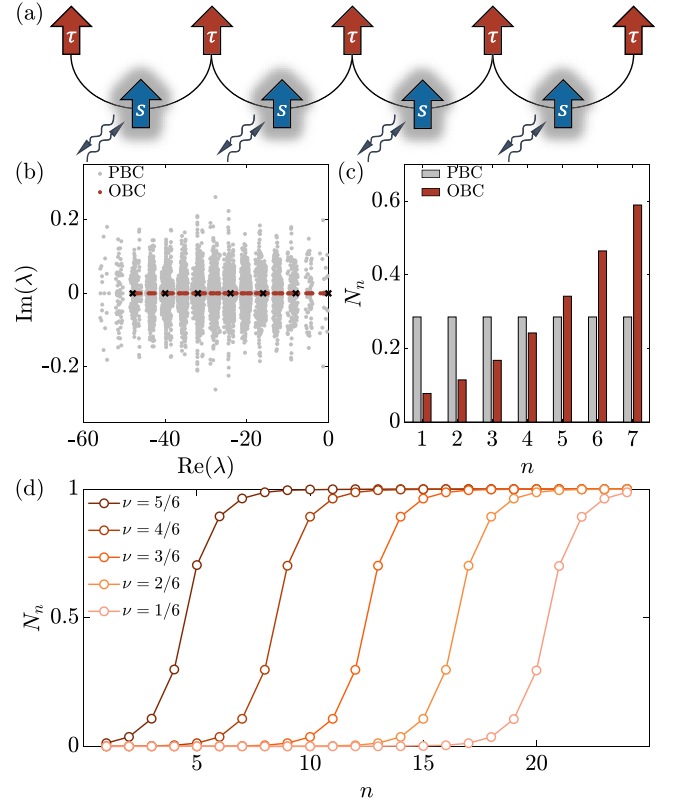


FIG. 1. (a) Dissipative quantum link model with H in Eq. (1). (b) Liouvillian spectrums and (c) steady-state density distributions $N_n = \text{Tr}[\rho_{ss,N}(\tau_n + 1/2)]$ numerically obtained by the exact diagonalization in the sector $\mathcal{G}_n = 0$ and $N = 2$ with parameters $L = 7$, $J = 1$, $\gamma_u = 2.4$, and $\gamma_d = 1.6$. We consider both the periodic (gray color) and open (red color) boundary conditions. Black points in (b) mark the eigenvalues of exact OBC eigenoperators. (d) The asymmetric particle distributions obtained from the exact OBC steady state. $\nu = N/L$ is the filling factor. We take $\beta = \gamma_u/\gamma_d = 3$ and $L = 24$.

symmetry. Without loss of generality, we focus on the simplest $s = 1/2$ case of link spins, with the basis vectors $|\uparrow\rangle$ and $|\downarrow\rangle$ on each link. Our results are also applicable in higher-spin cases.

We are interested in coupling the quantum links with a Markov bath [Fig. 1(a)], which is described by the Lindblad master equation:

$$\frac{d}{dt}\rho = \mathcal{L}[\rho] = -i[H, \rho] + \sum_{\mu} (2L_{\mu}\rho L_{\mu}^\dagger - \{L_{\mu}^\dagger L_{\mu}, \rho\}). \quad (2)$$

The Liouvillian superoperator $\mathcal{L}$ defined in the double Hilbert space is divided into two parts $\mathcal{L} = \mathcal{L}_H + \mathcal{L}_D$, where $\mathcal{L}_H[\rho] = -i[H, \rho]$ describes the coherent dynamics generated by the Hamiltonian H in Eq. (1) and $\mathcal{L}_D[\rho] = \sum_{\mu} (2L_{\mu}\rho L_{\mu}^\dagger - \{L_{\mu}^\dagger L_{\mu}, \rho\})$ denotes the dissipative dynamics associated with jump operators L_{μ} .

In the dissipative QLM, we couple each quantum link to the environment through two quantum jump operators describing biased incoherent spin-flipping processes:

$$L_{n,n+1}^{(u)} = \sqrt{\gamma_u} s_{n,n+1}^+, \quad L_{n,n+1}^{(d)} = \sqrt{\gamma_d} s_{n,n+1}^- \quad (3)$$

These jump operators break the strong U(1) gauge symmetry generated by G_n into a *weak gauge symmetry*, whose generator $\mathcal{G}_n$ acts on the operators O in the Hilbert space as $\mathcal{G}_n[O] = G_n O - O G_n$. Therefore, we get $[\mathcal{G}_n, \mathcal{L}] = 0$. In addition, the global charge N generates a strong symmetry of $\mathcal{L}$, satisfying $[N, H] = [N, L_\mu] = 0$.

These strong and weak symmetries allow for using exact diagonalization to study the many-body spectrums of $\mathcal{L}$. We focus on the weak gauge sector $\mathcal{G}_n[\rho] = 0$, to which the steady states belong. If we add strong dissipators $L_n^{(\text{gauge})} = \sqrt{\Gamma} G_n$ into Eq. (2) as an effective gauge fixing term [34], the relaxation dynamics will quickly converge to the above weak gauge sector. Nevertheless, these dissipators do not affect steady states and are thus not included in our construction. We use exact diagonalization to numerically obtain the spectrums of $\mathcal{L}$ under different boundary conditions [103]. As shown in Fig. 1(b), the OBC spectrum is enclosed by the spectrum under periodic boundary conditions (PBC), which demonstrates the sensitivity of Liouvillian spectrums to boundary conditions [104]. Additionally, the steady state ρ_{ss} , satisfying $\mathcal{L}[\rho_{\text{ss}}] = 0$, shows an asymmetric OBC distribution for the particle number $N_n = \text{Tr}[\rho_{\text{ss}}(\tau_n + 1/2)]$ [Fig. 1(c)]. These numerical results reveal the existence of many-body NHSE.

The emergence of NHSE can be understood as follows. First, biased spin-flipping processes result in a steady state ρ_{ss} with a preferable direction for link spins. For example, when $J = 0$, $\langle s_{n,n+1}^z \rangle = \text{Tr}[\rho_{\text{ss}} s_{n,n+1}^z] = (\beta - 1)/(2\beta + 2)$ depends on the ratio $\beta = \gamma_u/\gamma_d$. Intuitively, the dissipation-induced polarization of link spins can be statistically interpreted as a unidirectional electric field, rendering polarization-dependent interactions in H to transport particles in a favorable direction [104]. This process effectively becomes nonreciprocal hoppings of matter fields, causing many-body NHSE.

Strong dissipation limit—In the strong dissipation limit $\gamma_{u,d} \gg J$, the emergent nonreciprocal hoppings of matter fields can be seen explicitly using perturbation theory [105]. When $\mathcal{L}_H = 0$, the unperturbed steady states of $\mathcal{L}_D$ are given by $\rho_{\text{ss}}^{(0)} = \rho_\tau \otimes \rho_s$. Here, ρ_τ is an arbitrary density matrix on lattice sites, and the steady-state density matrix on lattice links is

$$\rho_s = \bigotimes_{n=1}^{L-1} \rho_{n,n+1}, \quad \rho_{n,n+1} = \frac{\beta}{1+\beta} |\uparrow\rangle\langle\uparrow| + \frac{1}{1+\beta} |\downarrow\rangle\langle\downarrow|. \quad (4)$$

Here, $\rho_{n,n+1}$ is defined on the $(n, n+1)$ link.

A nonzero $\mathcal{L}_H$ perturbs these degenerate steady states, and the second-order perturbation theory leads to an effective Liouvillian for ρ_τ of matter fields: $\mathcal{L}_{\text{eff}}[\rho_\tau] = \sum_\mu 2L_\mu \rho_\tau L_\mu^\dagger - \{L_\mu^\dagger L_\mu, \rho_\tau\}$. The effective jump operators in $\mathcal{L}_{\text{eff}}$ are given by $L_{n,n+1}^{(r)} = \sqrt{\gamma_r} \tau_n^+ \tau_{n+1}^+$ and $L_{n,n+1}^{(l)} = \sqrt{\gamma_l} \tau_n^+ \tau_{n+1}^-$, with $\gamma_r = \gamma_u J^2 / (\gamma_d + \gamma_u)^2$ and $\gamma_l = \gamma_d J^2 / (\gamma_d + \gamma_u)^2$ satisfying $\beta = \gamma_u/\gamma_d = \gamma_r/\gamma_l$. Surprisingly, the effective Liouvillian $\mathcal{L}_{\text{eff}}$ describes a quantum asymmetric simple exclusion process, whose classical counterpart plays a significant role in statistical physics [106,107]. This effective model is exactly solvable based on the Bethe ansatz, exhibiting operator space fragmentation [108].

Therefore, dissipative QLM offers a practical approach for engineering nonreciprocal hoppings for particles in a quantum many-body system, leading to unidirectional charge transport. Despite its similarity to the imaginary background gauge fields in the Hatano-Nelson model [100,101], the gauge fields in our model are dynamically coupled to the environment in an experimentally accessible way.

Exact steady state—The emergent integrability discussed above only holds when $\gamma_u, \gamma_d \gg J$. It is generally a challenging task to determine the exact many-body Liouvillian spectrum. Nevertheless, the gauge structure of the quantum link Hamiltonian enables analytical construction of the OBC steady states, as well as certain other eigenstates, for arbitrary parameters. The construction directly demonstrates many-body NHSE in the dissipative quantum link model and, more interestingly, provides a general approach to obtaining exact steady states in a wide class of open quantum systems.

When $\mathcal{L}_H = 0$, the steady-state density matrix ρ_s of $\mathcal{L}_D$ for link spins is a diagonal operator [Eq. (4)]. Following this observation, we introduce a double-space similarity transformation $\mathcal{T} = \mathcal{T}_\tau \otimes \mathcal{T}_s$. While $\mathcal{T}_\tau$ acting on sites will be specified shortly, $\mathcal{T}_s$ acting on links is designed to provide $\mathcal{T}_s[\rho_s] = \mathcal{T}_s \rho_s = I_s$. Here, I_s is the identity operator in link Hilbert space, and ρ_s is given by Eq. (4). Consequently, $\mathcal{T}_s = \bigotimes_{n=1}^{L-1} T_{n,n+1}$ with $T_{n,n+1} = (1+\beta)/\beta |\uparrow\rangle\langle\uparrow| + (1+\beta) |\downarrow\rangle\langle\downarrow| \propto \exp(-\ln \beta s_{n,n+1}^z)$ acting on the $(n, n+1)$ link [109]. It is easy to show that $\mathcal{T}_s \mathcal{L}_D \mathcal{T}_s^{-1} = \mathcal{L}_D^\dagger$ [110] where $\mathcal{L}_D^\dagger[\rho] = \sum_\mu (2L_\mu^\dagger \rho L_\mu - \{L_\mu^\dagger L_\mu, \rho\})$ with L_μ provided by Eq. (3). This result implies that $\mathcal{L}_D$ satisfies the quantum detailed balance condition [111–114]. Notably, the structure of the Liouvillian guarantees that $\mathcal{L}_D^\dagger[I] = 0$ with I being the identity matrix in the whole Hilbert space. A nonzero $\mathcal{L}_H$ also satisfies $\mathcal{L}_H[I] = 0$. Therefore, if there exists a double-space transformation $\mathcal{T}$ that keeps $\mathcal{L}_H$ invariant, i.e., $\mathcal{T} \mathcal{L}_H \mathcal{T}^{-1} = \mathcal{L}_H$, we can conclude that the identity operator I is the steady state of the deformed Liouvillian $\mathcal{T} \mathcal{L} \mathcal{T}^{-1} = \mathcal{L}_H + \mathcal{L}_D^\dagger$. Then the steady state of the original Liouvillian is given by $\rho_{\text{ss}} = \mathcal{T}^{-1}[I]/\text{Tr}(\mathcal{T}^{-1}[I])$.

The fact that the quantum-link Hamiltonian H satisfies the local gauge symmetry $[G_n, H] = 0$ motivates us to construct $\mathcal{T} = \mathcal{T}_\tau \otimes \mathcal{T}_s$ by some generalized gauge transformation. Therefore, the transformation $\mathcal{T} = \mathcal{T}_\tau \otimes \mathcal{T}_s$ compatible with $\mathcal{T}_s$ determined above is provided by $\mathcal{T}[\rho] = T\rho$ where $T = \exp[-\sum_{n=1}^L (\ln \alpha + n \ln \beta) G_n]$. Here α is an arbitrary constant. It is straightforward to show that $\mathcal{T}H\mathcal{T}^{-1} = H$, and that T acting on links is equivalent to the above $\mathcal{T}_s$ up to an overall factor. In the end, we obtain the exact OBC steady state

$$\rho_{ss} = \frac{T^{-1}}{\text{Tr}(T^{-1})} = Z^{-1} \exp \left[\sum_{n=1}^L (\ln \alpha + n \ln \beta) G_n \right] \quad (5)$$

with $Z = \text{Tr}[\exp(\sum_{n=1}^L (\ln \alpha + n \ln \beta) G_n)]$. This is one of the central results of this Letter. We emphasize that the construction in Eq. (5) is only valid under OBC. The lack of translation symmetry makes Eq. (5) incompatible with PBC. The mismatch between two boundary conditions is a signal of many-body NHSE. In the following, we give several remarks regarding this exact many-body steady state.

Recall that the conserved global U(1) charge is $N = \sum_{n=1}^L (\tau_n + 1/2)$. If we define the global dipole charge $D = \sum_{n=1}^L \tau_n [n - (L+1)/2]$ and the total z component of link spins $S^z = \sum_{n=1}^{L-1} s_{n,n+1}^z$, the steady state becomes $\rho_{ss} \propto \exp[\ln \alpha N + \ln \beta (D + S^z)]$. Because N is a conserved quantity, $\ln \alpha$ can be viewed as the chemical potential of a Gibbs ensemble. We can also obtain the N -particle steady state $\rho_{ss,N} \equiv Z_N^{-1} P_N \rho_{ss} P_N$, where P_N is the projection operator into the N -particle Hilbert space and Z_N^{-1} is a normalization factor. Notably, $\rho_{ss,N}$ is independent of the free parameter α . Therefore, the OBC steady states of Eq. (2) have an $(L+1)$ -fold degeneracy.

When $0 < N < L$, the steady state $\rho_{ss,N}$ exhibits many-body NHSE by showing an asymmetric particle density at the lattice sites. As shown in Fig. 1(d), while the parameter β shapes the localization behavior, the filling number N/L determines the locus of the emergent real-space Fermi surface [72]. The link spins, on the other hand, form a uniform configuration with $\langle s_{n,n+1}^z \rangle = (\beta - 1)/(\beta + 2)$, serving as an effective unidirectional electric field that generates the accumulation of dipole charge for particles, i.e., $\text{Tr}(\rho_{ss,N} D) \neq 0$.

Furthermore, ρ_{ss} does not depend on the interaction strength J . Instead, the local gauge symmetry of H plays an indispensable role in constructing the double-space transformation $\mathcal{T}$. This result indicates that ρ_{ss} is robust against the symmetry-preserving disorders in H , such as random potentials $\delta H = \sum_{n=1}^L h_n \tau_n^z + \sum_{n=1}^{L-1} h'_{n,n+1} s_{n,n+1}^z$ or long-range interactions $\delta H' = \sum_{n=1}^{L-2} J'_n (\tau_n^+ s_{n,n+1}^+ s_{n+1,n+2}^+ \tau_{n+2}^- + \text{H.c.})$. This symmetry-protected feature distinguishes the exact steady-state skin mode from the single-particle NHSE in noninteracting systems, where strong disorder typically

destroys skin modes and instead leads to Anderson localization.

The exact construction leading to Eq. (5) offers a general approach to obtaining eigenoperators of a large class of Liouvillian superoperators. With the Liouvillian $\mathcal{L} = \mathcal{L}_H + \mathcal{L}_D$, if we find a double-space similarity transformation $\mathcal{T}_\lambda$ such that

$$\mathcal{T}_\lambda \mathcal{L}_H \mathcal{T}_\lambda^{-1} = \mathcal{L}_H, \mathcal{T}_\lambda \mathcal{L}_D \mathcal{T}_\lambda^{-1} [I] = \lambda I, \quad (6)$$

we can obtain an eigenoperator $\rho = \mathcal{T}_\lambda^{-1} [I]$ of $\mathcal{L}$ with the eigenvalue λ . In the Supplemental Material [104], we show that this approach systematically provides exponentially many eigenoperators of $\mathcal{L}$ in dissipative QLM, with eigenvalues equally spaced on the real axis [black points in Fig. 1(b)]. We remark that the equally spaced eigenvalues are reminiscent of exact many-body scar states in a large class of Hermitian systems [115–117]. We also apply this method to other types of jump operators in dissipative QLM [104]. These demonstrations further indicate the general applicability of our exact method.

Hierarchical skin effect—Figure 1 elucidates that the biased dissipation of link spins causes the many-body skin effect of site spins and increases their dipole moments. Interestingly, this setup provides a general framework to generate a new type of many-body NHSE, called the hierarchical skin effect, where the steady state will accumulate nonzero multipole moments for different subsystems.

We consider a generalized dissipative QLM [Fig. 2(a)], whose Hamiltonian is defined for three species of spin- $\frac{1}{2}$ operators:

$$H' = J_1 \sum_{n=1}^{L-1} (\sigma_n^+ \tau_{n,n+1}^+ \sigma_{n+1}^- + \text{H.c.}) + J_2 \sum_{n=1}^{L-2} (\tau_{n,n+1}^+ s_{n+1}^+ \tau_{n+1,n+2}^- + \text{H.c.}). \quad (7)$$

Here, L is the number of σ spins on the top layer. Likewise, the middle layer has $L-1$ τ spins and the bottom layer consists of $L-2$ s spins. We denote $H_1 \equiv H'|_{J_2=0}$ and $H_2 \equiv H'|_{J_1=0}$, which are quantum-link Hamiltonians between neighboring layers. Like Eq. (1), H_1 has a local U(1) gauge symmetry generated by $G_n^{\sigma\tau} = \sigma_n^z - \tau_{n,n+1}^z + \tau_{n-1,n}^z$ for $1 < n < L$, with boundary corrections $G_1^{\sigma\tau} = \sigma_1^z - \tau_{1,2}^z$ and $G_L^{\sigma\tau} = \sigma_L^z + \tau_{L-1,L}^z$. Similarly, the local U(1) gauge symmetry generators for H_2 are given by $G_{n,n+1}^{\tau s} = \tau_{n,n+1}^z - s_{n+1}^z + s_n^z$ for $1 < n < L-1$, $G_{1,2}^{\tau s} = \tau_{1,2}^z - s_2^z$, and $G_{L-1,L}^{\tau s} = \tau_{L-1,L}^z + s_{L-1}^z$. However, these symmetry generators for H_1 or H_2 are no longer conserved when considering the total Hamiltonian $H' = H_1 + H_2$. Instead, the local symmetry generators in H' involve spin operators in all layers, which are given by $G_n^{H'} = \sigma_n^z - G_{n,n+1}^{\tau s} + G_{n-1,n}^{\tau s}$ for $1 < n < L$, $G_1^{H'} = \sigma_1^z - G_{1,2}^{\tau s}$, and $G_L^{H'} = \sigma_L^z + G_{L-1,L}^{\tau s}$.

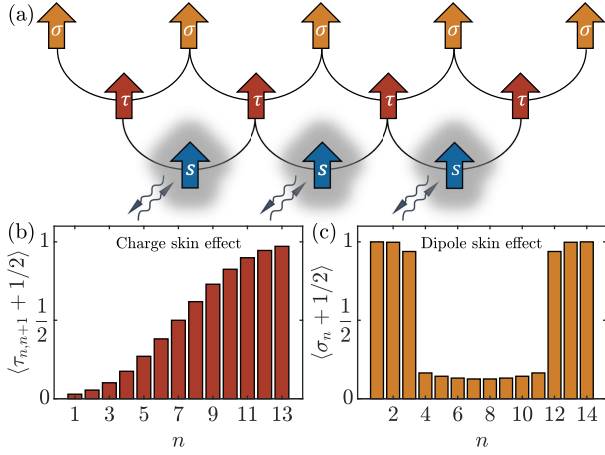


FIG. 2. (a) Generalized dissipative QLM. (b),(c) Steady-state distribution of the middle and top layers. With $L = 14$ and $\beta = 3$, we show the results in the charge sector $(N_{H'}, D_{H'}) = (0, 0)$.

It is easy to show that $[G_n^{H'}, H'] = 0$. Consequently, $G_n^{H'}$ plays a similar role to the local $U(1)$ gauge symmetry generators in the original quantum-link Hamiltonian.

With $\beta = \gamma_u/\gamma_d$, we added biased jump operators $L_n^{(u)} = \sqrt{\gamma_u} s_n^+$ and $L_n^{(d)} = \sqrt{\gamma_d} s_n^-$ to s spins [Fig. 2(a)], where $n = 2, 3, \dots, L-1$. The dynamics of this system is generated by the Lindblad master equation in Eq. (2). For convenience, we express the Liouvillian superoperator as $\mathcal{L} = \mathcal{L}_{H'} + \mathcal{L}_D$, with $\mathcal{L}_{H'}$ given by the Hamiltonian H' and $\mathcal{L}_D$ corresponding to the dissipations of s spins, respectively.

The way to obtain the steady state is to construct a double-space similarity transformation $\mathcal{T}$, defined as $\mathcal{T}[\rho] = (T_\sigma \otimes T_\tau \otimes T_s)\rho$, such that $\mathcal{T}\mathcal{L}\mathcal{T}^{-1} = \mathcal{L}_{H'} + \mathcal{L}_D^\dagger$. When $H_1 = 0$, the top layer is decoupled. The rest of the system is the same as Fig. 1(a). The results above Eq. (5) shows that $T_\tau \otimes T_s = \exp[-\sum_{n=1}^{L-1} (\ln \alpha + n \ln \beta) G_{n,n+1}^{\tau s}]$ with α being a free parameter. A nonzero H_1 constrains the form of T_σ , indicating that $T_\sigma \otimes T_\tau$ should keep H_1 invariant and be compatible with the above $T_\tau \otimes T_s$. As a result, $T_\sigma \otimes T_\tau = \exp\{-\sum_{n=1}^L [\ln \alpha' + \sum_{m=0}^{n-1} (\ln \alpha + m \ln \beta)] G_n^{\sigma \tau}\}$ with α' being another free parameter. Combining these results yields the steady state

$$\rho_{ss} \propto \exp \left[\sum_{n=1}^L \left(\ln \alpha' + n \ln \alpha + \frac{n(n-1)}{2} \ln \beta \right) G_n^{H'} \right]. \quad (8)$$

The free parameters α and α' reveal two strong global symmetries that do not involve s -spin operators and thus commute with $L_n^{(u/d)}$. The first one, $N_{H'} = \sum_{n=1}^L G_n^{H'} = \sum_{n=1}^L \sigma_n^z$, is the total z component of σ spins. The second one, $D_{H'} = \sum_{n=1}^L n G_n^{H'} = \sum_{n=1}^L n \sigma_n^z + \sum_{n=1}^{L-1} \tau_{n,n+1}^z$, consists of the dipole moment of σ spins and the total z component of τ spins. After projection into each symmetry

sector, we obtain the exact symmetry-resolved steady states that exclusively depend on the dissipation parameter β . A typical steady-state distribution is shown in Fig. 2. Intuitively, the polarization of s spins causes the many-body skin effect for τ spins [Fig. 2(b)], generating a nonzero dipole moment in the middle layer. Subsequently, the asymmetric distribution of τ spins effectively serves as a spatially dependent electric field for σ spins. The opposite polarizations of τ spins [Fig. 2(b)] transport the polarizations of σ spins in different directions, making the z component of σ spins accumulate at two boundaries simultaneously [Fig. 2(c)]. As a result, a nonzero quadrupole moment for σ spins reveals the existence of dipole skin effect in the top layer [84].

More layers with interlayer quantum-link interactions can be included in the open quantum system, such that a hierarchy of many-body NHSE can be established in the exact steady state, characterized by nonzero multipole moments for different layers. This intriguing feature distinguishes the generalized dissipative QLM from the previous result on the many-body NHSE for multipoles in non-Hermitian Hamiltonians [84]. Whereas the latter necessitates a particular m -pole conserving non-Hermitian interaction to produce the $(m+1)$ th multipole moment, our framework for hierarchical skin effect requires merely three-local interactions and on-site dissipations that are more feasible to realize experimentally, simultaneously creating multipole skin effects of any order in different subsystems.

Possible experiments—Our model offers a practical method in engineering nonreciprocal hoppings in a quantum many-body system. One possible platform is the extended Bose-Hubbard simulator in cold-atom systems [10]. By introducing a spatially modulated potential and specifying the model parameters, an effective bosonic Hamiltonian can be engineered in the form $H_{\text{eff}} = J \sum_n (b_{2n-1}^\dagger b_{2n}^\dagger b_{2n+1} + \text{H.c.})$, with the dynamics restricted to the Hilbert subspace with zero or single (double) boson occupations on odd (even) sites. Within this subspace, the model describes a $U(1)$ lattice gauge theory, where the odd (even) sites represent matter (gauge) fields. The detailed derivation and its mapping to the quantum link model [Eq. (1)] can be found in Ref. [10]. Under the same mapping, the spin-flipping jump operators in Eq. (3) can be translated into bosonic gain or loss jump operators [118–120] on even sites, although the dissipative dynamics need to be restricted in the above-mentioned subspace. Another promising platform is the superconducting circuit, where a native three-body quantum-link interaction [Eq. (1)] has recently been realized [14] and there exist known strategies for introducing gain or loss jump operators into superconducting qubits [121].

Discussions—In conclusion, we have constructed gauge-theoretical models with exactly solvable steady states exhibiting NHSE. Our results shed light on novel

nonequilibrium phases tied to gauge symmetries. For instance, the approach to induce the NHSE in dissipative many-body systems and develop the hierarchical skin effect for multipoles can be easily applied to U(1) QLMs in any dimensions with higher spins [23,24,122–125]. An example of 2D NHSE in dissipative QLMs is shown in the Supplemental Material [104]. Another interesting future direction is to extend the dissipative lattice gauge theory to non-Abelian cases [126,127]. It is also intriguing to investigate the dissipative gauge fields using Keldysh field theory [120,128] and topological field theory [129].

Acknowledgments—We thank He-Ran Wang, Fei Song, Shunyu Yao, Zhou-Quan Wan, and Hongzheng Zhao for their helpful discussions. This work is supported by the National Natural Science Foundation of China (Grants No. 12125405), the National Key R&D Program of China (No. 2023YFA1406702), and Quantum Science and Technology-National Science and Technology Major Project (Grant No. 2021ZD0302502). B. L. is supported by the National Science Foundation through Princeton University’s Materials Research Science and Engineering Center DMR-2011750, and the National Science Foundation under Grant No. DMR-2141966. Additional support is provided by the Gordon and Betty Moore Foundation through Grant No. GBMF8685 towards the Princeton theory program.

Data availability—The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

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- [104] See Supplemental Material at <http://link.aps.org/supplemental/10.1103/wztw-l8wg> for the numerical simulation of relaxation dynamics, point-gap topology of many-body Liouvillian, and the construction of exact eigenoperators in different models.
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- [110] Assume that O is an arbitrary operator in the Hilbert space. $T_s \mathcal{L}_D T_s^{-1}$ acting on O represents $T_s \mathcal{L}_D T_s^{-1}[O] = T_s[\mathcal{L}_D[T_s^{-1}[O]]]$. Additionally, the inverse of an invertible superoperator S is defined as $S^{-1}S[O] = SS^{-1}[O] = O$.
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