Dissipative Topological Phase Transition with Strong System-Environment Coupling

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A primary motivation for studying topological matter regards the protection of topological order from its environment. In this work, we study a topological emitter array coupled to an electromagnetic environment. The photon-emitter coupling produces nonlocal interactions between emitters. Using periodic boundary conditions for all ranges of environment-induced interactions, the chiral symmetry inherent to the emitter array is preserved. This chiral symmetry protects the Hamiltonian and induces parity in the Lindblad operator. A topological phase transition occurs at a critical photon-emitter coupling related to the energy spectrum width of the emitter array. Interestingly, the critical point nontrivially changes the dissipation rates of edge states, yielding a dissipative topological phase transition. In the protected topological phase, edge states suffer from environment-induced dissipation for weak photon-emitter coupling. However, strong coupling leads to robust dissipationless edge states with a window at the emitter spacing. Our work shows the potential to manipulate topological quantum matter with electromagnetic environments.

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Introduction.—Vacuum electromagnetic environments can nontrivially change order parameters of matter, producing phase transitions [1,2]. With the advances in cavity quantum electrodynamics (QED) [3–6], vacuum electromagnetic fields are used to manipulate matter [7–10] with strong light-matter interaction. For example, in cavity-interfaced superconductors, a strong coupling with electromagnetic fields changes the superconducting transition temperature [9]. Recently, the vacuum electromagnetic control of matter is receiving growing attention [11–13]. Due to symmetry-protected properties, topological matter is also being studied in the coupling with electromagnetic fields for potential applications [14–16]. The band gap of a kagome metasurface of dipole emitters embedded in a cavity can be tuned by electromagnetic fields [17]. Varying the cavity width can change long-range interactions between emitters and induce topological phase transitions [16].

A prerequisite to make topological protection reliable is to understand dissipative properties of topological systems [18–33]. Energy bands play a pivotal role for topological matter, e.g., in studying topological phases [34–37] and topological criticalities [38–40]. The large gap between energy bands protects topological properties from local disorder [41–49] and thermal noises [50–55]. However, a recent study [56] of time-reversal symmetry protected topological systems with large band gap shows the fragility of topological phases in electromagnetic environments. Via perturbation theory, they find that quantum coherence between edge states in one-dimensional (1D) topological systems is spoiled when system-environment coupling is weak compared to the band gap. This finding shows the challenge of protecting topological quantum matter in electromagnetic environments.

In this work, we study the coupling between a topological emitter array and its electromagnetic environment in the nonperturbative regime, i.e., edge states are coupled to bulk states via the environment. We find that for emitter spacings $d = \lambda_0/4$ and $d = 3\lambda_0/4$, environment-induced interactions have chiral symmetry and produce distinct topological phases. For $d = \lambda_0/4$, environment modifies the topological phase with dissipative edge states. However, the edge states for $d = 3\lambda_0/4$ are protected from dissipation in a parameter space specified by the Lindblad operator. In the thermodynamic limit, a dissipative topological phase transition (DTPT), characterized by a nontrivial change of dissipation of the edge states, occurs at $d = 3\lambda_0/4$ when the single-emitter decay rate induced by the system-environment coupling equals the energy spectrum width of the topological emitter array. These results could be useful for improving topological protection in open quantum systems which have nonlocal dissipations [56].
photons. The emitter-environment coupling is from the imaginary part of the complex permittivity; the Green's function describing the electromagnetic interaction 

does not preserve translational symmetry in the SSH model. The horizontal red dashed line denotes the SSH-type phase transition in the dissipative regime. The vertical green line represents the energy-induced topological phase transition where the decay rate is equal to the spectrum width of the topological system. In particular, the green solid line indicates the dissipative topological phase transition. (d) Photon-mediated interactions for and . The band gap is . (g) The transverse spectrum of the topological emitter array with band gap and spectrum width is shown in Fig. 1(b).

Emitter modes in the environment are described by , where and are the creation and annihilation operators of photons. The emitter-environment coupling is , where is the dipole moment operator of the th emitter. The electric field operator is 

where is the imaginary part of the complex permittivity; the Green's tensor describes the electromagnetic interaction from to . The dynamics of the topological emitter array is described by the master equation

where the free energy is with (transition frequency of emitters) and the topological emitter array is described by with tunable dimerized interactions . The emitter-environment coupling can be strong compared with the band gap, but much smaller than the energy of emitters, to satisfy the Born-Markov approximation in Eq. (1). The virtual-photon exchange between emitters and environment yields 

where and characterize the strengths of the nonlocal dipole-dipole interactions. In addition to the coherent part of the virtual-photon exchange, the dissipative topological phase transition represented by the green solid vertical line in Fig. 1(c).

Environment-protected chiral symmetry.—As a simple illustration, the electric field operator of the th emitter is . The long-range interaction between the first and the last emitters provides periodic boundary conditions for the NN interaction. Conversely, the long-range interaction between the th and th emitters exhibits translational invariance due to the NN interaction. Therefore, the effective strengths for the NN interaction and the long-range interaction between the th and th emitters are . Moreover, the effective interaction between the th and th emitters (red dashed curves) is . With this protocol, the translational symmetry is preserved for all ranges of interactions induced by the environment at and . For other values of the spacing , the translational symmetry in is broken.
By assuming periodic boundary conditions on $H_{\text{topo}}$, the coherent interaction $H = H_{\text{topo}} + H_{\text{ph}}$ in quasimomentum space is $H/\hbar = \sum_k \Psi_k^\dagger \mathcal{H}(k)\Psi_k$, where $\Psi_k^\dagger = (\sigma_{x,k}^\dagger, \sigma_{y,k}^\dagger)$. Here, $A$ and $B$ denote odd- and even-site emitters, respectively. The 1D symmetry-protected topological system is described by the Su-Schrieffer-Heeger (SSH) model [75]. In the sublattice space, we obtain an effective spin-1/2 Hamiltonian $\mathcal{H}(k) = h_x(k)\tau_x + h_y(k)\tau_y$ with chiral symmetry $\tau_y\mathcal{H}(k)\tau_x = -\mathcal{H}(k)$ [76]. Here, $\sigma_x$, $\tau_y$, $\tau_z$ are Pauli matrices, and

$$h_x(k) = J_1 + J_2 \cos(k) + \frac{g_0}{2} \left[ 1 + \cos\left(\frac{Nk}{2}\right) \right],$$

$$h_y(k) = J_2 \sin(k) + \frac{g_0}{2} \mathcal{F}(k),$$

with $g_0 = \gamma_0/2$ ($-\gamma_0/2$) for $d = \lambda_0/4$ ($3\lambda_0/4$), $\mathcal{F}(k) = \sum_{j=1}^{N/2} 2(-1)^{j-1} \sin(jk) - \sin(Njk/2)$, and energy bands $\epsilon_{\pm}(k) = \pm \sqrt{h_x^2(k) + h_y^2(k)}$. Without the environment, the energy bands are shown in Fig. 2(a). The band gap and spectrum width are

$$\delta\omega = 2|J_1 - J_2|,$$

$$\Delta\omega = 2(J_1 + J_2),$$

respectively. The dimerized interactions $J_{1,2} = J_0(1 \mp \cos\phi)$ yield the band gap $\delta\omega = 4J_0\cos\phi$ and spectrum width $\Delta\omega = 4J_0$. The SSH-type topological phase transition takes place at $k = \pm \pi$ [76] with linear low-energy dispersion. In the electromagnetic environment with emitter spacing $d = 3\lambda_0/4$, the condition

$$\gamma_0 = \Delta\omega,$$

yields a gap closing at $k = 0$ with parabolic dispersion, as shown in Fig. 2(b). The parabolic dispersion [38,39] makes this topological criticality to be different from the one in the SSH model. In the auxiliary space $[h_x(k), h_y(k)]$, the winding number can be defined as $W = (1/2\pi) \int_{BZ} d\theta_k$, with $\theta_k = \arctan[-h_x(k)/h_y(k)]$. For $d = \lambda_0/4$, shown in Fig. 2(c), the system is in a nontopological phase with $W = 0$. However, as $J_0$ is increased, the winding number $W = 1$; i.e., the topological phase is protected when $d = \lambda_0/4$. For $d = 3\lambda_0/4$, in Fig. 2(d), the system has zero winding number for small $J_0/\gamma_0$. However, at a critical point $\gamma_0^c = \Delta\omega$, a topological phase transition takes place. For $\gamma_0 < \gamma_0^c$, the system becomes topological with winding number $W = 1$. Namely, the topological phase is preserved when the spectrum width $\Delta\omega$ is larger than the environment-induced decay $\gamma_0$ of the emitters, as shown in Fig. 1(c).

Edge state vs dissipative topological phase transition.—Figures 3(a) and 3(b) show the energy spectra of $H$ versus $J_0/\gamma_0$ for (a) $d = \lambda_0/4$ and (b) $d = 3\lambda_0/4$ in a system with an odd number of emitters $N = 21$, where a single edge state appears. In agreement with the topologies in quasimomentum space for these two emitter spacings, a band gap [3(a)] and a band touching [3(b)] are found. In Fig. 3(b), a nontopological edge state is found for the topologically trivial phase. Figures 3(c) and 3(d) show the

FIG. 2. (a) Energy bands of the topological emitter array for $J_1 \neq J_2$ (solid) and $J_1 = J_2$ (dashed). (b) Environment-induced gap closing at emitter spacing $d = 3\lambda_0/4$ and decay rate $\gamma_0 = \Delta\omega$. Topologies from the hybridization between $H_{\text{topo}}$ and $H_{\text{ph}}$ in auxiliary space $[h_x(k), h_y(k)]$ for (c) $d = \lambda_0/4$, and (d) $d = 3\lambda_0/4$. (c) The winding number is zero at $J_0 = 0$ (blue solid), and becomes one for $J_0 > 0$ (green dashed). (d) The winding number is zero for $0 < J_0 \leq \gamma_0/4$ (the blue solid topology denotes $J_0 = \gamma_0/4$), and becomes one for $J_0 > \gamma_0/4$ (green dashed). We consider $\phi = 0.1\pi, N = 6$.

FIG. 3. Energy spectra for (a) $d = \lambda_0/4$, and (b) $d = 3\lambda_0/4$, respectively. Probability distribution $|\psi_n|^2$ of the zero-energy state for (c) $d = \lambda_0/4$ and (d) $d = 3\lambda_0/4$. In (c),(d), red stars, blue triangles, orange squares, and black circles correspond to $J_0/\gamma_0 = 0, 0.25, 1, 5$, and $J_0/\gamma_0 = 0, 0.2, 0.25, 1$, respectively. The inset of (d) shows the IPR of the zero-energy state at different values of $\phi$ for $d = 3\lambda_0/4$. In (a)–(d) we consider $\phi = 0.1\pi, N = 21$. 

250402-3
distributions \( |\psi_0|^2 \) of the edge state. At \( J_0 = 0 \), the edge state is equally distributed at the two edge emitters with wave function \( |\psi_0\rangle = (1/\sqrt{2})(|\sigma_i^+\rangle + |\sigma_i^-\rangle)|G\rangle \), where \( |G\rangle \) is the ground state of the emitter array. In Fig. 3(c), with \( d = \lambda_0/4 \), enlarging \( J_0 \) monotonically increases the component of \( |\psi_0|^2 \) at the left boundary. However, before the critical point, the left-boundary component of the edge state for \( d = 3\lambda_0/4 \) becomes smaller as \( J_0/\gamma_0 \) is increased. At the critical point, the gap of the spectrum closes and the edge state becomes delocalized. By further increasing \( J_0 \), the edge state eventually localizes at the left boundary.

To characterize the changes of the edge state, we study the inverse participation ratio (IPR) [77], \( \text{IPR} = \sum_i |\psi_{0i}|^4/(\sum_i |\psi_{0i}|^2)^2 \), where \( \psi_{0i} \) is the amplitude of the edge state at the \( i \)th emitter. In the inset of Fig. 3(d), we show the IPR versus \( J_0/\gamma_0 \) for \( d = 3\lambda_0/4 \). The IPR of the edge state at \( J_0 = 0 \) is one half due to its equal distribution at two boundaries. A minimum is found at the critical point for different values of \( \phi \), indicating the edge-bulk transition.

To study the stability of topological features in real space, we here rewrite the Lindblad operator in terms of eigenstates of \( H \),

\[
\mathcal{D}[\rho] = \sum_{m,n} \Gamma_{mn} |\psi_m\rangle\langle\psi_m| \rho - \frac{1}{2} \rho |\psi_m\rangle\langle\psi_m| - \frac{1}{2} |\psi_m\rangle\langle\psi_m| \rho - \frac{1}{2} \rho |\psi_m\rangle\langle\psi_m|,
\]

with \( |\psi_m\rangle = |\psi_m\rangle \langle G \). Here, \( |\psi_m\rangle \) denotes the \( m \)th eigenmode of \( H \). The decay rates are \( \Gamma_{mn} = \sum_i \langle e_i |\psi_m\rangle \langle\psi_n| e_i \rangle \), with \( |e_i\rangle = |\sigma_i^+\rangle |G\rangle \). Specifically, \( \Gamma_{mn} \) denotes the decay rate of the \( m \)th eigenstate to environment; \( \Gamma_{mn} \) is the correlated decay from the \( n \)th state to \( m \)th state. The dissipation of the edge state is governed by \( \Gamma_{m0} \). In Fig. 4(a), we show the scaled decay rate \( \Gamma_{m0}/\gamma_0 \) from edge state to environment versus \( J_0/\gamma_0 \). For \( d = \lambda_0/4 \), \( \Gamma_{m0}/\gamma_0 \) increases with \( J_0/\gamma_0 \), and decreases after reaching the maximum [58]. However, the edge state at \( d = 3\lambda_0/4 \) has a decay rate that decreases in the non-topological phase and that stops decaying at \( J_0 = \gamma_0/4 \). In finite systems, the weak emitter-environment coupling, i.e., small \( \gamma_0/J_0 \), introduces dissipation of the edge state [58], which is responsible for the enhanced photon absorption [78]. However, the edge state for strong coupling is protected against decoherence in the topological phase.

For added clarity, the correlated decays \( \Gamma_{m0} \) (\( m \neq 0 \)) between the edge state and the bulk states are shown in Fig. 4(b). At \( J_0/\gamma_0 = 0.2 \) (blue dots) in the non-topological phase, the edge state not only decays into the environment (\( \Gamma_{00} \neq 0 \)), but also decays into the bulk states of the emitter array. However, at \( J_0/\gamma_0 = 0.3 \) (red squares) in the topological phase, the edge state does not decay to bulk states. At the critical point \( J_0/\gamma_0 = 0.25 \), the dissipations to bulk states are greatly suppressed, except for those of the two bulk states \( m' = \pm(N - 3)/2 \). Near the critical point, the correlated dissipation \( |\Gamma_{m0}| \propto \exp(-\nu_{m'}/N) \), with \( \nu_{m'} > 0 \) in the topological phase. The inset shows \( \ln(|\Gamma_{m0}/\gamma_0|) \) versus \( N \). The values of \( \nu_{m'} \) are 0, 0.005, and 0.0115, for \( J_0/\gamma_0 = 0.25, 0.251, \) and 0.252, respectively. Therefore, in the thermodynamic limit \( N \to +\infty \), the critical point indicates a transition between dissipative and dissipationless edge states, namely, a DTPT, which can be accessed by observing the population dynamics of the emitter array [58]. In Fig. 4(c), the local minima of \( \ln(|\Gamma_{00}/\gamma_0|) \) show the parameter space of edge states protected by the Lindblad operator [79–82] in a small system. This protection is actually attributed to the vanishing overlap between edge states and polarized radiating modes in the Lindblad operator, which exhibits parity property, i.e., dissipations only occur between odd-site (even-site) emitters. Larger systems have broader parameter space for dissipationless edge states [58]. In particular, the condition that the environment-induced decay is half of the spectrum width, i.e., \( J_0/\gamma_0 = 1/2 \), produces a dissipationless edge state

\[
|\psi_0\rangle \left( \frac{J_0}{\gamma_0} \to \frac{1}{2} \right) = \frac{1}{\sqrt{N}} \sum_{n \in \mathbb{N}} (-1)^n \left( \tan \frac{\phi}{4} \right)^{2n} |\psi\rangle_n,
\]

for various localization lengths even near the SSH criticality. Here, \( |\psi\rangle_n = (\sigma_{4n+1} + \sigma_{4n+3})|G\rangle \); namely, the \( (4n + 1) \)th and \( (4n + 3) \)th emitters have the same amplitude.
Hamiltonian in the diagonalized form insensitive to emitter spacing. To confirm this conjecture, $i$ small and energy splitting (in the inset) of the edge states. With appear at the boundaries. Figure 4(d) shows the decay rates arrays with an even number of emitters, two edge states

Dissipationless subspace of topological edge states.—In arrays with an even number of emitters, two edge states appear at the boundaries. Figure 4(d) shows the decay rates and energy splitting (in the inset) of the edge states. With small $\gamma_0$ (weak coupling), the two localized edge states are coupled by environment-mediated long-range interactions, and the subspace of edge states suffers from decoherence [56, 83]. Surprisingly, when the emitter-environment coupling is strong, i.e., $\gamma_0$ is large, the edge states are decoupled from each other. Therefore, they are both protected from dissipation until the DTPT at $\gamma_0 = \Delta \omega$. Moreover, strong coupling makes the zero splitting between edge states insensitive to emitter spacing around $d = 3\lambda_0/4$ [58].

Dissipationless window.—In Fig. 5(a), we show the dissipation $\Gamma_{(d)}$ of a single edge state versus $d/\lambda_0$. Even though the chiral symmetry is broken for emitter spacings around $3\lambda_0/4$, the edge state can be dissipationless and is insensitive to emitter spacing. To confirm this conjecture, we rewrite the whole system as a non-Hermitian effective Hamiltonian in the diagonalized form $H_{\text{eff}} = \sum_j (\hat{E}_j - i\Gamma_j) |\Psi_j\rangle\langle\Psi_j|$ with the biorthogonal basis $\langle\tilde{\Psi}_j|\Psi_j\rangle = \delta_{j,j'}$. A dissipationless window is found for $d = 3\lambda_0/4$ with strong system-environment coupling, as shown in Fig. 5(b). This window makes the dissipationless edge state robust to disorder in emitter positions [58]. The edge state has finite decay rate around $d = \lambda_0/4$. For emitter spacings $n\lambda_0/2$ ($n = 0, 1, 2, \ldots$) where $H_{\text{ph}}$ is zero, the edge state is more dissipative and shows higher sensitivity to disorder than $d = 3\lambda_0/4$ [58].

Conclusions.—System-environment interplay is fundamental for dissipative topological matter. In this work, we show that a 1D topological emitter array globally coupled to an electromagnetic environment exhibits interesting dissipative properties as the system-environment coupling varies. The energy spectrum width of the emitter array sets a critical value for the system-environment coupling and produces the DTPT. The environment-modified topological edge states are stable and robust due to a dissipationless window in the emitter spacing. Our work paves an avenue for electromagnetic control of topological matter with vacuum fields.

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[1] R. Landig, L. Hruby, N. Dogra, M. Landini, R. Mottl, T. Donner, and T. Esslinger, Quantum phases from competing short- and long-range interactions in an optical lattice, Nature (London) 532, 476 (2016).

[2] Y. Ashida, A. Imamoğlu, J. Faist, D. Jaksch, A. Cavalleri, and E. Demler, Quantum Electrodynamic Control of Matter: Cavity-Enhanced Ferroelectric Phase Transition, Phys. Rev. X 10, 041027 (2020).

[3] S. Haroche and J.-M. Raimond, Exploring the Quantum: Atoms, Cavities, and Photons (Oxford University Press, New York, 2006).

[4] X. Gu, A. F. Kockum, A. Miranowicz, Y.-X. Liu, and F. Nori, Microwave photonics with superconducting quantum circuits, Phys. Rep. 718–719, 1 (2017).

[5] A. F. Kockum, A. Miranowicz, S. De Liberato, S. Savasta, and F. Nori, Ultrastrong coupling between light and matter, Nat. Rev. Phys. 1, 19 (2019).

[6] I. Mivehvar, F. Piazza, T. Donner, and H. Ritsch, Cavity QED with quantum gases: New paradigms in many-body physics, arXiv:2102.04473.

[7] F. Schlawin, A. Cavalleri, and D. Jaksch, Cavity-Mediated Electron-Photon Superconductivity, Phys. Rev. Lett. 122, 133602 (2019).

[8] J. B. Curtis, Z. M. Raines, A. A. Allocca, M. Hafezi, and V. M. Galitski, Cavity Quantum Eliahsberg Enhancement of Superconductivity, Phys. Rev. Lett. 122, 167002 (2019).

[9] A. Thomas, E. Devaux, K. Nagarajan, T. Chervy, M. Seidel, D. Hagenmüller, S. Schütz, J. Schachenmayer, C. Genet, G. Pupillo, and T. W. Ebbesen, Exploring superconductivity under strong coupling with the vacuum electromagnetic field, arXiv:1911.01459.
[48] Y. Wang, Y.-H. Lu, F. Mei, J. Gao, Z.-M. Li, H. Tang, S.-L. Zhu, S. Jia, and X.-M. Jin, Direct Observation of Topology from Single-Photon Dynamics, Phys. Rev. Lett. 122, 193903 (2019).

[49] W. Nie, Z. H. Peng, F. Nori, and Y.-X. Liu, Topologically Protected Quantum Coherence in a Superatom, Phys. Rev. Lett. 124, 023603 (2020).

[50] M. Gong, G. Chen, S. Jia, and C. Zhang, Searching for Majorana Fermions in 2D Spin-Orbit Coupled Fermi Superfluids at Finite Temperature, Phys. Rev. Lett. 109, 105302 (2012).

[51] O. Viyuela, A. Rivas, and M. A. Martin-Delgado, Uhlmann Phase as a Topological Measure for One-Dimensional Fermion Systems, Phys. Rev. Lett. 112, 130401 (2014).

[52] Z. Huang and D. P. Arovas, Topological Indices for Open and Thermal Systems Via Uhlmann’s Phase, Phys. Rev. Lett. 113, 076407 (2014).

[53] B. Monserrat and D. Vanderbilt, Temperature Effects in the Band Structure of Topological Insulators, Phys. Rev. Lett. 117, 226801 (2016).

[54] C.-E. Bardyn, L. Wawer, A. Altland, M. Fleischhauer, and S. Diehl, Probing the Topology of Density Matrices, Phys. Rev. X 8, 011035 (2018).

[55] R. Unanyan, M. Kiefer-Emmanouilidis, and M. Fleischhauer, Finite-Temperature Topological Invariant for Interacting Systems, Phys. Rev. Lett. 125, 215701 (2020).

[56] M. McGinley and N. R. Cooper, Fragility of time-reversal symmetry protected topological phases, Nat. Phys. 16, 1181 (2020).

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

[58] See the Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevLett.127.250402 for additional details about the derivation of the master equation, environment-induced topological phase transition, topological protection of quantum coherence, and experimental scheme for detecting the dissipative topological phase transition.

[59] Y. Chen, C. Neill, P. Roushan, N. Leung, M. Fang, R. Barends, J. Kelly, B. Campbell, Z. Chen, B. Chiaro, A. Dunsworth, E. Jeffrey, A. Megrant, J. Y. Mutus, P. J. J. O’Malley, C. M. Quintana, D. Sank, A. Vainsencher, J. Wenner, T. C. White et al., Qubit Architecture with High Coherence and Fast Tunable Coupling, Phys. Rev. Lett. 113, 220502 (2014).

[60] L. Knöll, S. Scheel, and D.-G. Welsch, QED in dispersing and absorbing media, arXiv:quant-ph/0006121.

[61] H. T. Dung, L. Knöll, and D.-G. Welsch, Resonant dipole-dipole interaction in the presence of dispersing and absorbing surroundings, Phys. Rev. A 66, 063810 (2002).

[62] D. Dzsotjan, A. S. Sørensen, and M. Fleischhauer, Quantum emitters coupled to surface plasmons of a nanowire: A Green’s function approach, Phys. Rev. B 82, 075427 (2010).

[63] A. Gonzalez-Tudela, D. Martin-Cano, E. Moreno, L. Martin-Moreno, C. Tejedor, and F. J. García-Vidal, Entanglement of Two Qubits Mediated by One-Dimensional Plasmonic Waveguides, Phys. Rev. Lett. 106, 020501 (2011).

[64] D. Martín-Cano, A. González-Tudela, L. Martín-Moreno, F. J. García-Vidal, C. Tejedor, and E. Moreno, Dissipation-driven generation of two-qubit entanglement mediated by plasmonic waveguides, Phys. Rev. B 84, 235306 (2011).

[65] G. Angelatos and S. Hughes, Entanglement dynamics and Mollow nonuplets between two coupled quantum dots in a nanowire photonic-crystal system, Phys. Rev. A 91, 051803 (2015).

[66] A. Asenjo-Garcia, J. D. Hood, D. E. Chang, and H. J. Kimble, Atom-light interactions in quasi-one-dimensional nanostructures: A Green’s-function perspective, Phys. Rev. A 95, 033818 (2017).

[67] S. A. H. Gangaraj, G. W. Hanson, and M. Antezza, Robust entanglement with three-dimensional nonreciprocal photonic topological insulators, Phys. Rev. A 95, 063807 (2017).

[68] A. Asenjo-Garcia, M. Moreno-Cardoner, A. Albrecht, H. J. Kimble, and D. E. Chang, Exponential Improvement in Photon Storage Fidelities Using Subradiance and Selective Radiance in Atomic Arrays, Phys. Rev. X 7, 031024 (2017).

[69] P. Doyeux, S. A. H. Gangaraj, G. W. Hanson, and M. Antezza, Giant Interatomic Energy-Transport Amplification with Nonreciprocal Photonic Topological Insulators, Phys. Rev. Lett. 119, 173901 (2017).

[70] T. Shi, D. E. Chang, and J. I. Cirac, Multiphoton-scattering theory and generalized master equations, Phys. Rev. A 92, 053834 (2015).

[71] G. Calajó, F. Ciccarello, D. Chang, and P. Rabl, Atom-field dressed states in slow-light waveguide QED, Phys. Rev. A 93, 033833 (2016).

[72] M. Mirhosseini, E. Kim, X. Zhang, A. Sipahigil, P. B. Dieterle, A. J. Keller, A. Asenjo-Garcia, D. E. Chang, and O. Painter, Cavity quantum electrodynamics with atom-like mirrors, Nature (London) 569, 692 (2019).

[73] P. Y. Wen, K.-T. Lin, A. F. Kockum, B. Suri, H. Ian, J. C. Chen, S. Y. Mao, C. C. Chiu, P. Delsing, F. Nori, G.-D. Lin, and I.-C. Hoi, Large Collective Lamb Shift of Two Distant Superconducting Artificial Atoms, Phys. Rev. Lett. 123, 233602 (2019).

[74] M. Zanner, T. Orell, C. M. Schneider, R. Albert, S. Oleschko, M. L. Juan, M. Silveri, and G. Kirchmair, Coherent control of a symmetry-engineered multi-qubit dark state in waveguide quantum electronics, arXiv:2106.05623.

[75] W. P. Su, J. R. Schrieffer, and A. J. Heeger, Solitons in Polyacetylene, Phys. Rev. Lett. 42, 1698 (1979).

[76] J. K. Asbóth, L. Oroszlány, and A. Pályi, A Short Course on Topological Insulators (Springer, New York, 2016).

[77] N. C. Murphy, R. Wortis, and W. A. Atkinson, Generalized inverse participation ratio as a possible measure of localization for interacting systems, Phys. Rev. B 83, 184206 (2011).

[78] W. Nie, T. Shi, F. Nori, and Y.-X. Liu, Topology-Enhanced Nonreciprocal Scattering and Photon Absorption in a Waveguide, Phys. Rev. Applied 15, 044041 (2021).
[79] V. V. Albert, B. Bradlyn, M. Fraas, and L. Jiang, Geometry and Response of Lindbladians, Phys. Rev. X 6, 041031 (2016).
[80] F. Minganti, A. Biella, N. Bartolo, and C. Ciuti, Spectral theory of Liouvillians for dissipative phase transitions, Phys. Rev. A 98, 042118 (2018).
[81] F. Minganti, A. Miranowicz, R. W. Chhajlany, and F. Nori, Quantum exceptional points of non-Hermitian Hamiltonians and Liouvillians: The effects of quantum jumps, Phys. Rev. A 100, 062131 (2019).
[82] S. Lieu, M. McGinley, and N. R. Cooper, Tenfold Way for Quadratic Lindbladians, Phys. Rev. Lett. 124, 040401 (2020).
[83] T.-S. Deng, L. Pan, Y. Chen, and H. Zhai, Stability of Time-Reversal Symmetry Protected Topological Phases, Phys. Rev. Lett. 127, 086801 (2021).