**A superconducting quantum simulator based on a photonic-bandgap metamaterial**

Xueyue Zhang1,2†, Eunjong Kim1,2†, Daniel K. Mark3, Soonwon Choi3, Oskar Painter1,2,4∗

Synthesizing many-body quantum systems with various ranges of interactions facilitates the study of quantum chaotic dynamics. Such extended interaction range can be enabled by using nonlocal degrees of freedom such as photonic modes in an otherwise locally connected structure. Here, we present a superconducting quantum simulator in which qubits are connected through an extensible photonic-bandgap metamaterial, thus realizing a one-dimensional Bose-Hubbard model with tunable hopping range and on-site interaction. Using individual site control and readout, we characterize the statistics of measurement outcomes from many-body quench dynamics, which enables in situ Hamiltonian learning. Further, the outcome statistics reveal the effect of increased hopping range, showing the predicted crossover from integrability to ergodicity. Our work enables the study of emergent randomness from chaotic many-body evolution and, more broadly, expands the accessible Hamiltonians for quantum simulation using superconducting circuits.

Realizing a scalable architecture for quantum computation and simulation is a central goal in the field of quantum information science. Although architectures with nearest-neighbor (NN) coupling between quantum particles on a lattice are prevalent, quantum systems with long-range interactions can realize a richer set of computational tasks and physical phenomena (1–4). For instance, in the case of gate-based quantum computation, coupling beyond the NN level enables nonlocal gate operations between qubits, which can reduce the overhead of quantum algorithms and lift the restrictions on code rate and distance of local-interaction-based quantum error-correcting codes (5, 6).

In the case of analog quantum simulation, the inclusion of long-range interactions can alter the behavior of otherwise integrable many-body systems (7, 8), resulting in quantum chaotic dynamics, which is at the root of such topics as quantum thermalization (9) and quantum information scrambling (10). Furthermore, control over the range of lattice connectivity grants access to different physical regimes and the crossover between them, such as in many-body quantum phase transitions (11–13) and the hydrodynamics of nonequilibrium quantum states (14).

For engineered quantum systems consisting of interacting quantum particles on a lattice, it is often challenging to scale to larger lattice sizes while maintaining a high degree of lattice connectivity and single-site control. One common approach, developed for trapped-ion and neutral-atom systems, is to use resonant modes of either vibrational (1) or optical (15) cavities as a quantum bus for mediating interactions between the internal states of atoms across the lattice. Similar schemes have been adopted in superconducting quantum circuits, realizing systems as large as 20 qubits with all-to-all coupling via a common microwave cavity (16). Increasing the number of lattice sites in this case, however, leads to either parasitic coupling arising from dense placement of sites in a fixed-volume cavity or frequency-crowding effects stemming from the increased spectral density of cavity modes when increasing the cavity size (17).

An alternative approach for connecting quantum particles on a lattice is to construct a quantum bus from an intrinsically extensible structure, such as a waveguide. Along this direction, engineered photonic-bandgap waveguides have been proposed as a quantum bus that simultaneously protects quantum particles from radiative damping through the waveguide while allowing for extended-range lattice connectivity (18). The waveguide-bus concept has been investigated in the context of many-body simulation with cold atoms coupled to engineered nanophotonic waveguides (18, 19), and recent experiments have explored qubit-photon bound states in superconducting quantum circuits with microwave photonic-bandgap waveguides (20–24). However, the realization of a scalable many-body quantum simulator, with single-site quantum-particle control and a high level of lattice connectivity, has remained an open challenge.

We demonstrate a scalable many-body quantum simulator consisting of a one-dimensional (1D) lattice of superconducting transmon qubits coupled to a common metamaterial waveguide. This system provides both tunable-range connectivity between qubits and full single-site control and state measurement of individual qubits. The waveguide acts both as a bus for mediating exponentially decaying long-range interactions between qubits and as a Purcell filter enabling multiplexed, rapid readout of the qubit states with high fidelity. This system realizes an extended version of the Bose-Hubbard model with tunable hopping range and on-site interaction. Using our ability to efficiently collect measurement outcomes from many-body quench dynamics—enabled by the fast experimental repetition rate of our system—we perform direct analysis of outcome statistics to learn Hamiltonian parameters in situ and study the effect of hopping range on the evolution of randomness across the system. Specifically, we observe a distribution of outcome bit-string probabilities reflecting the ergodic nature of the Hamiltonian with long-range hopping. This result experimentally confirms the expectation from quantum chaos for interacting many-particle systems, highlighting the connection between ergodic unitary dynamics and its effective statistical description in terms of random matrix theory (25).

**Metamaterial-based quantum simulator**

The backbone of the many-body quantum simulator in this work is a metamaterial waveguide formed from a chain of lumped-element inductor-capacitor (LC) microwave resonators. The waveguide can be described by a generic model (Fig. 1A) of a 1D cavity array with NN coupling $t$ (26, 27). The corresponding dispersion relation (Fig. 1B) is given by $\omega_k = \omega_k + 2\kappa \cos(kd)$, exhibiting a passband centered around the cavity frequency $\omega_k$ with a bandwidth of $4t$, where $k$ is the wave vector and $d$ is the lattice constant of the array. The bandgap frequencies below $\omega_{-} = \omega_{k} - 2t$ (above $\omega_{+} = \omega_{k} + 2t$) is denoted as the lower (upper) bandgap, abbreviated as LBG (UBG). Inside the bandgaps, the off-resonant coupling between a bare quantum emitter and the waveguide modes gives rise to an emitter-photon bound state (28) whose photonic tail is localized around the emitter. Localization follows a spatial profile $e^{-x^2}$ in the LBG or UBG (27), where $\Delta E$ is the displacement in the number of unit cells from the emitter and $\xi$ is the localization length controlled by the detuning $\Delta$ between the band-edge frequency and the transition frequency of the bound state. The overlap of two bound states results in photon-mediated coupling with a range covering multiple unit cells, i.e., long-range coupling, exhibiting a greater strength and a more extended range $\xi$ at a smaller detuning $|\Delta|$ (Fig. 1C).

The metamaterial waveguide consists of a 42–unit-cell array of capacitively coupled lumped-element microwave resonators (Fig. 1, D and E) and is equipped at both ends with engineered tapering sections, designed to reduce the impedance mismatch to external 50-ohm...
Fig. 1. Metamaterial-based quantum simulator. (A) Schematic showing a 1D array of coupled cavities with nearest-neighbor coupling \( t \). Each cavity is occupied by a quantum emitter (orange ball) with coupling \( g \) to the cavity. (B) Dispersion relation of the coupled cavity array in (A) with a passband between \( \omega_{\pm} \) centered at \( \omega_0 \) (bandwidth of 4f). The LBG (UBG) below (above) the passband is shaded in green (purple). (C) Top (Bottom): Cartoon of two emitter-photon bound states at small (large) detuning \( |\Delta| \), indicated by dark (light) orange arrows in (B), exhibiting an extended (restricted) spatial range and large (small) photonic component in the bound states. (D) Electrical circuit realization of (A) with capacitively coupled LC resonators and transmon qubits corresponding to the cavity array and the quantum emitters, respectively. The coupling capacitors are color coded in accordance with (A). (E) Optical micrograph (false colored) of the fabricated quantum simulator with 42 metamaterial resonators (lattice constant \( d = 292 \mu m \)) colored blue connected to input-output ports (red) via tapering sections (purple). Ten qubits (Q, colored orange), controlled by individual charge drive lines (pink) and flux bias lines (dark blue), and their readout resonators (R, colored green) couple to the 10 inner unit cells of the metamaterial waveguide with a zoomed-in view in the left inset. Detailed view of the coupling region is shown in the right inset. Two auxiliary qubits (yellow) are not used in this experiment. (F) Transmission spectrum through the metamaterial waveguide (red curve) with black arrows indicating the 10 resonances of the readout resonators R.

input-output ports at frequencies lying within the passband of the waveguide (24, 29). Each of the middle 10 metamaterial resonators (unit cells labeled by \( i = 1 \) to 10) couples to a transmon qubit (30), which serves as the quantum emitter. Individual addressing of each qubit is achieved by excitation (XY control) from a flux bias line. Dispersive qubit readout using the passband of the metamaterial waveguide bus, creates a lattice of interacting microwave photons (31). The average decay rate of 10 readout resonators is \( \kappa_{\text{R}}/2\pi = 11.8 \text{ MHz} \), enabling fast, high-fidelity multiplexed readout while maintaining a low level of readout cross-talk. For details of readout methods and characterization, refer to supplementary text IV.

Bose-Hubbard model with long-range hopping

The spatially extended bound-state excitations, formed between transmon-qubit excitations and waveguide photons of the metamaterial-waveguide bus, creates a lattice of interacting microwave photons (31). This quantum system is described by an extended version of the 1D Bose-Hubbard model with tunable long-range hopping and on-site interaction. Specifically, each bound state formed from qubit-photons, inheriting the level structure of an anharmonic oscillator from a transmon qubit (30), serves as a bosonic site with local site energy \( \epsilon_i = \omega_{01,i} \) and the on-site interaction \( U_i = \omega_{02,i} - \omega_{01,i} \). Here, \( \omega_{01,i} \) and \( \omega_{02,i} \) are the transition frequencies of the bound state on site Q, from its ground state \( |0\) to the first excited state \( |1\) and that from the first to the second excited state \( |2\), respectively. In addition, the long-range hopping \( J_{ij} \) is enabled by the overlap between a pair of qubit-photon bound states on sites \( Q_i \) and \( Q_j \). The Hamiltonian of this model that captures the basic processes mentioned above can be written as

\[
\hat{H} / \hbar = \sum_{i j} J_{ij} \hat{b}_i^\dagger \hat{b}_j + \sum_i U_i \hat{n}_i (\hat{n}_i - 1) + \sum_i \epsilon_i \hat{n}_i
\]

where \( \hat{b}_j^\dagger (\hat{b}_j) \) is the creation (annihilation) operator and \( \hat{n}_i = \hat{b}_i^\dagger \hat{b}_i \) is the number operator on site \( Q_i \). The parameters of the Hamiltonian realized in this simulator can be learned through experiments enabled by the precise, single-site-level control over qubits.

We measure the on-site interaction \( U_i \) (Fig. 2A) by performing spectroscopy of \( \omega_{02} \) after initializing \( Q_i \) in its first excited state \( |1\). From within either bandgap, \( |U_i| \) decreases as \( \omega_{01} \) approaches the closest band edge owing to dressing from the passband modes of the metamaterial (22)—i.e., the Lamb shift. In the UBG, a wide tuning range of \( |U_i| \) is achievable from the strong hybridization between the \( |1\rangle - |2\rangle \) transition and the band-edge modes at \( (\omega_{01} - \omega_{02})/2\pi < 300 \text{ MHz} \). The magnitude of hopping \( |J_{ij}| \) is measured from vacuum Rabi oscillations between sites \( Q_i \) and \( Q_j \) by
initializing one site with a \( \pi \) pulse and tuning \( \omega_{\text{rot}} \) of both sites on resonance for a duration \( \tau \) with fast flux pulses. For a fixed distance \( |i-j| \), \( |J_{ij}| \) decreases with a decreasing \( |\Delta| \), resulting from larger photonic components of the bound states (Fig. 2B). Compared to the LBG, the UBG exhibits larger \( |J_{ij}| \) at the same \( |\Delta| \), owing to a stronger coupling of the bare qubits to the metamaterial at higher frequencies and the breakdown of the tight-binding cavity array model in the circuit realization (Fig. 1D; also see supplementary text I). At a specific \( \omega_{\text{rot}} \), \( |J_{ij}| \) decreases exponentially as a function of distance \( |i-j| \), resembling the profile of the photonic tail in a qubit-photon bound state (Fig. 2C). From fitting the exponential decay curve we extract the localization length \( \xi \), which ranges from \( \xi = 1.4 \) to \( 4.2 \), with the largest localization length occurring at the smallest achievable band-edge detunings. For even smaller detunings, the eigenstate of the two interacting bound states merges into the passband and becomes radiative to the waveguide.

**Many-body Hamiltonian learning**

Beyond the above single- and two-qubit measurements, we perform in situ many-body characterization of Hamiltonian parameters (32, 33), which are otherwise hard to access. For example, the sign of the hopping term \( J_{ij} \) inherits the spatial profile of the photonic component of the bound states. In the case of the bound states in the UBG, the sign of the hopping terms are all uniform (positive), whereas for bound states in the LBG, the hopping terms alternate sign as the distance between lattice sites increases by one (Fig. 2D). This is due to the photonic component of the bound state behaving as a defect mode inside the bandgap, exhibiting a spatial profile resembling the wave vector at the nearest band edge (\( k = 0 \) at the upper band-edge and \( k = \pi/d \) at the lower band edge). Although insignificant in measurements involving only two lattice sites, the sign of the hopping terms does alter the many-body dynamics of the system. Here, we use a many-body fidelity estimator \( F_d \) proposed in (33) to reveal this information. This fidelity estimator, which closely tracks the true many-body fidelity, is obtained for ergodic quench evolution of simple initial states (supplementary text VI).

We follow the sequence described in Fig. 2E to perform the many-body quench evolution. The sequence consists of preparing a set of five randomly chosen sites in their first excited states, followed by flux pulses to align \( \omega_{\text{rot}} \) of all 10 sites for time \( \tau \), and then finally performing site-resolved single-shot measurement on all lattice sites to obtain a 10-bit string \( z = n_1 n_2 \cdots n_{10} \). The many-body fidelity estimator \( F_d \) is calculated by comparing bit-string...
Fig. 3. Two-particle quantum walk with increasing hopping range. (A) Evolution of the population $\langle n_i \rangle$ on sites Q1 to Q10 as a function of normalized evolution time $J_{i,i+1}t$. The system is initialized in $\ket{\psi_0} = 0000100000$ and the evolution occurs at $\omega_{00}/2\pi = 4.50, 4.55, 4.72,$ and $4.80$ GHz with the longest evolution times of 904, 781, 430, and 200 ns from left to right. (B) The second moment $\mu_2^2$ as a function of normalized evolution time $J_{i,i+1}t$. Results calculated from the data in (A) are shown as solid curves with gray scales corresponding to frames in (A) and arrows in Fig. 2B. Result from numerical simulation of the integrable Hamiltonian is shown as the dotted curve, and $\mu_2^2$ for a generic ergodic system is indicated by the red dashed line.

statistics of repeated measurements with numerical simulation of the evolution assuming a set of Hamiltonian parameters in Eq. 1. The maximum $F_2$ is achieved at the parameter values closest to the Hamiltonian realized in the experiment. The fast repetition rate of this experiment enables us to perform a large number of measurements (1.6 $\times$ 10^4 in total), reducing statistical error and increasing sensitivity to small Hamiltonian parameter variations (see supplementary text V for details of qubit control, pulse sequence, and repetition rate).

We compare $F_2$ at $\omega_{00}/2\pi = 4.72$ GHz using three different parameter sets for $J_{i,i+1}$ in Fig. 2F: a first set with amplitudes derived from the two-qubit experiments in Fig. 2B assuming alternating signs (Fig. 2D, left); a second set with the same amplitudes as the first but all positive signs (Fig. 2D, right); and a third set of optimized parameter values that maximize $F_2$. The optimized hopping terms are restricted to be real-valued, with independent $J_{i,i+1}$ for each $i = 1$ to 9 and $J_{i,i}$ for each distance $|i-j| > 1$ (all qubit pairs of the same site having the same $J_{i,j}$). An alternating sign of $J_{i,i+1}$ with distance $i$ is favored, yielding a higher many-body fidelity compared to hopping terms with all positive signs. This is further evidenced by the alternating signs of the resulting optimized parameter set. Although we find small differences between the set of optimized hopping amplitudes and those from the two-qubit experiments with alternating signs (Fig. 2G), $F_2$ of the optimized parameter set is markedly better. The sensitivity of $F_2$ to the hopping terms is highlighted in the inset of Fig. 2F, where the variation of the fidelity versus $J_{i,i+1}$ is shown. For details of the $F_2$ calculation and parameter optimization, refer to supplementary text VI.

**Ergodic many-body dynamics with long-range hopping**

We now use the platform to study the effect of long-range hopping on the many-body dynamics. Specifically, the ergodicity of the 1D Bose-Hubbard model in the hardcore limit ($|U/J| > 1$) depends on the range of hopping, which exhibits integrable behavior with NN hopping, and chaotic behavior with long-range hopping. We study this crossover with various hopping ranges and investigate the resulting dynamics using both conventional one- and two-site correlators, and the statistics of the global bit-strings resulting from qubit-state measurement outcomes across the lattice. This latter technique is particularly useful in identifying universal signatures of ergodicity and the effect of decoherence at long evolution times. The crossover between integrable and ergodic dynamics can be qualitatively visualized by a two-particle quantum walk (34–36) with initial excitations on sites Q5 and Q6 using the sequence shown in Fig. 2E. The measured quantum walk at a few different $\omega_{00}$’s indicated by arrows in Fig. 2B is shown (Fig. 3A) as a function of normalized evolution time $J_{i,i+1}t$, where $J_{i,i+1}$ is the average NN hopping rate (the corresponding numerical simulations are provided in supplementary text VII, showing that the quantum walk patterns are not visibly affected by decoherence). The excitation wave packets smear over the system when $\omega_{00}$ is close to the band-edge frequency. More quantitatively, this trend can be probed by computing the probability $p_z$ of measuring a certain bit-string $z$ in the two-excitation sector at evolution time $t$. For a generic ergodic Hamiltonian, the second moment $\mu_2^2 = \sum z p_z^2$ (25), which reflects the probability fluctuations, converges to $\mu_2^2 = 2/(D+1)$ after initial evolution (33) owing to the chaotic nature of its quantum dynamics ($D = 45$ is the dimension of the two-excitation Hilbert space). No such convergence is expected in an integrable Hamiltonian owing to revivals associated with ballistic propagation of wave packets. As an example, we show in Fig. 3B the results from the spin-1/2 XY model obtained from modifying the Hamiltonian in Eq. 1 by keeping only NN hopping terms in the hardcore limit. When $\omega_{00}$ is closer to the band edge, the measured second moment deviates from the simulated integrable result and converges to $\mu_2^2$ at an earlier normalized evolution time $J_{i,i+1}t$ consistent with the breaking of integrability due to the extended hopping range. With $|U/J| > 36$ for all the measurements illustrated in Fig. 3, finite on-site interactions of the Bose-Hubbard model play a negligible role in the breaking of integrability (supplementary text VIII).

To further probe this ergodic nature of Hamiltonian with long-range hopping, we use the experimental evolution at $\omega_{00}/2\pi = 4.72$ GHz as an example. At a short time ($t = 16$ ns), the excitations remain in their initial sites. This is visualized for a quantum walk with initial excitations on sites Q5 and Q6 in the left panel of Fig. 4C (evolution of population $\langle n_i \rangle$) and in the bottom left panel of Fig. 4D (two-site correlator $\langle n_i n_{i+1} \rangle$). The histogram $P(p_z)$ of experimentally measured bit-string probabilities $p_z$ at this early evolution stage (Fig. 4B, left) shows a distribution with a tail of large $p_z$ values, giving a large $\mu_2$ (Experiment curve in Fig. 4A). This is associated with an insufficient scrambling of the initially localized quantum information. At an intermediate time ($t = 360$ ns), the excitations are more spread out over the entire 1D lattice (Fig. 4D, middle left panel), forming a “speckle” pattern with site-to-site fluctuation that is associated with quantum interference. The quantitative signatures of this speckle pattern manifest in the histogram $P(p_z)$ following the Porter-Thomas (PT) distribution (37) (Fig. 4B, middle) and in the second moment $\mu_2$ settling to the ergodic value $\mu_2^2$. The PT distribution results from the randomness in the distribution of wavefunction magnitudes, which is predicted by Berry's conjecture (38) stating that the single-particle eigenstates of a chaotic system behave like random superpositions of plane waves. Similarly, in the many-body settings, the distribution of wave function magnitudes across basis states also follows the PT distribution. Our observation is the first experimental verification of this many-body version of Berry's conjecture in a Bose-Hubbard system, whose extension in the thermodynamic limit provides the modern theory of quantum thermalization such as eigenstate thermalization hypothesis (39, 40). This draws a connection between quantum many-body chaos and random matrix theory, leading to a deeper understanding of the randomness in many-body dynamics (32). Distinct from the randomness inherent in random circuits (25, 41), the randomness in our case originates from the ergodicity of the time-independent Hamiltonian. In contrast to the experimental results, theoretical calculations at the same evolution time using the integrable Hamiltonian show aggregated excitations on a few sites (Fig. 4E, middle) and the resulting larger value of $\mu_2$ (Integrable theory curve in Fig. 4A). This comparison highlights the effect of long-range hopping in probing the ergodic regime.
In addition, we study the impact of decoherence by juxtaposing the measurement results and the decoherence-free theoretical calculation using the optimal learned Hamiltonian with long-range hopping. Before the evolution time of $\tau = 1 \, \mu s$, the two cases agree in the second moment $\mu_2$ (Experiment and Theory curves in Fig. 4A), the quantum walk population (Fig. 4C, left and middle), and the two-site correlator (Fig. 4D, middle panels), suggesting that these results are not affected by decoherence. After a long evolution time ($\tau = 5.4 \, \mu s$, larger than the averaged Ramsey coherence time $T_{2,1} = 1.16 \, \mu s$), the second moment of the two cases deviates from one another, and the experimental speckle pattern begins to wash out compared to the theoretical modeling (Fig. 4D, top panels). Another probe of the decoherence is the histogram $P(p_z)$ of the measured bit-string probabilities (Fig. 4B, right). Here, the histogram deviates from the PT distribution, narrows substantially, and approaches a uniform distribution corresponding to a completely decohered, maximally mixed state. Additional numerical simulations of $\mu_2$ and $P(p_z)$ for ergodic and integrable systems can be found in supplementary text VIII.

**Conclusion and outlook**

Our many-body quantum simulator is based on a 1D lattice of transmon qubits connected together using a superconducting metamaterial, which exhibits photonic bandgaps that protect qubit-photon bound states from decoherence and a transmission passband used for high-fidelity multiplexed qubit-state readout. Furthermore, the metamaterial plays the role of a scalable photonic bus to mediate tunable long-range coupling between qubit-photon bound states. This system of interacting bound states realizes a Bose-Hubbard model with long-range hopping. We characterize the system using conventional single- and two-qubit measurements along with a sample-efficient many-body Hamiltonian learning protocol. Lastly, we study the many-body quench dynamics of the system versus the range of the lattice hopping, revealing the ergodic nature of the extended Bose-Hubbard model, as distinct from its NN-coupling counterpart. The major challenge in probing long-time quantum evolution in our experiment is the short Ramsey coherence time $T_{2,1} = 1.16 \, \mu s$, limited by flux-noise–induced dephasing. Incorporating a single refocusing pulse has been shown to increase the coherence time to $T_{2,1} = 5.64 \, \mu s$ at the single-qubit level (supplementary text II). The extended quantum evolution times enabled by further dynamical decoupling, combined with the tunable-range coupling investigated in this work, provide valuable opportunities to explore nonequilibrium dynamics with or without coupling to an environment and quantum phases of matter in the presence of frustration.

**REFERENCES AND NOTES**

1. K. Mølmer, A. Sørensen, Phys. Rev. Lett. 82, 1835–1838 (1999).
2. V. D. Vaity et al., Phys. Rev. X 8, 011002 (2018).
3. A. Perival et al., Nature 600, 630–635 (2021).
4. D. Bluemink et al., Nature 604, 451–456 (2022).
5. S. Bravyi, D. Poulin, B. Terhal, Phys. Rev. Lett. 104, 050503 (2010).
6. N. Doffes, M. E. Beverland, M. A. Tremblay, arXiv:2109.14599 (quant-ph) (2021).
7. P. Jurcic et al., Nature 511, 202–205 (2014).
8. P. Richerme et al., Nature 511, 198–201 (2014).
9. R. Nandkishore, D. A. Huse, Ann. Rev. Condens. Matter Phys. 6, 15–38 (2015).
10. M. K. Joshi et al., Phys. Rev. Lett. 124, 240505 (2020).
11. R. Islam et al., Science 340, 583–587 (2013).
12. R. Landig et al., Nature 532, 476–479 (2016).
13. S. Ebadi et al., Nature 595, 227–232 (2021).
14. M. K. Joshi et al., Science 376, 720–724 (2022).
15. S.-B. Zheng, G.-C. Guo, Phys. Rev. Lett. 85, 2392–2395 (2000).
16. Q. Guo et al., Nat. Phys. 17, 234–239 (2021).
17. Y. P. Zhang et al., Nat. Phys. 15, 741–744 (2019).
18. J. S. Douglas et al., Nat. Photonics 9, 326–331 (2015).
19. D. E. Chang, J. S. Douglas, A. Gonzalez-Tudela, C.-L. Hung, J. H. Kimble, Rev. Mod. Phys. 90, 031002 (2018).
20. Y. Liu, A. A. Houck, Nat. Phys. 13, 48–52 (2017).
21. M. Motesansiz et al., Nat. Commun. 9, 3705 (2018).
22. N. M. Sundaresan, R. Lundgren, G. Zhu, A. V. Gorshkov, A. A. Houck, Phys. Rev. X 9, 011021 (2019).
23. M. Scigluzo et al., Phys. Rev. X 12, 031036 (2022).
24. E. Kim et al., Phys. Rev. X 11, 01105 (2021).
25. S. Boixo et al., Nat. Phys. 14, 595–600 (2018).
26. M. J. Hartmann, F. G. S. L. Brandão, M. B. Plenio, Nat. Phys. 2, 849–855 (2006).
27. G. Calaf, F. Ciccarello, D. Chang, P. Rabl, Phys. Rev. A 93, 033833 (2016).
28. S. John, J. Wang, Phys. Rev. Lett. 64, 2418–2421 (1990).
29. V. S. Ferrante et al., Phys. Rev. X 11, 041043 (2021).
30. J. Koch et al., Phys. Rev. A 76, 042319 (2007).
31. A. A. Houck, H. E. Türeci, J. Koch, Nat. Phys. 8, 292–299 (2012).
32. J. Choi et al., arXiv:2103.03535 (quant-ph) (2021).
ACKNOWLEDGMENTS
The authors thank A. Gorshkov, A. González-Tudela, D. Chang, O. Motrunich, R. Ma, F. Brandão, G. Refael, S. Meesala, V. Ferreira, G. Kim, A. Butler, and Z. Zheng for helpful discussions. We appreciate MIT Lincoln Laboratories for the provision of traveling-wave parametric amplifiers used for both spectroscopic and time-domain measurements in this work, and the AWS Center for Quantum Computing for the Eccosorb filters installed in the cryogenic setup for infrared filtering. We also thank the Quantum Machines team for technical support and discussions on the Quantum Orchestration Platform. Funding: This work was supported by the AFOSR Quantum Photonic Matter MURI (grant FA9550-16-1-0323), the DOE-BES Quantum Information Science Program (grant DE-SC0020152), the Institute for Quantum Information and Matter, an NSF Physics Frontiers Center (grant PHY-1125565) with support of the Gordon and Betty Moore Foundation, the Kavli Nanoscience Institute at Caltech, and the AWS Center for Quantum Computing. D.K.M. acknowledges support from the NSF QLCI program (2016245) and the DOE Quantum Systems Accelerator Center (contract no. 7568717).

Author contributions: X.Z. and E.K. planned the experiment, performed the device design and fabrication, and performed the measurements. X.Z., E.K., D.K.M., and S.C. analyzed the data. O.P. supervised the project. All authors contributed to the writing of the manuscript. Competing interests: The authors declare no competing interests. Data and materials availability: Experimental data shown in the main text and supplementary materials, as well as the simulation code, are available in Zenodo (42). License information: Copyright © 2023 the authors, some rights reserved; exclusive licensee American Association for the Advancement of Science. No claim to original US government works. https://www.sciencemag.org/about/science-licenses-journal-article-reuse

SUPPLEMENTARY MATERIALS
science.org/doi/10.1126/science.ade7651
Materials and Methods
Supplementary Text
Figs. S1 to S14
Tables S1 and S2
References (43–88)
Submitted 6 September 2022; accepted 16 December 2022
10.1126/science.ade7651