

Quantum many-body scarring in 2+1D gauge theories with dynamical matter

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Quantum many-body scarring (QMBS) has emerged as an intriguing paradigm of weak ergodicity breaking in nonintegrable quantum many-body models, particularly lattice gauge theories (LGTs) in $1+1$ spacetime dimensions. However, an open question is whether QMBS has a true high-energy physics origin, and thus exists in higher-dimensional LGTs with dynamical matter. Given that nonergodic dynamics in $d=1$ spatial dimension tend to vanish in $d>1$, it is important to probe this question. Using matrix product state techniques for both finite and infinite systems, we show that QMBS occurs in the 2+1D U(1) quantum link model, as evidenced in persistent coherent oscillations in local observables, a marked slowdown in the growth of the bipartite entanglement entropy, and revivals in the fidelity. Interestingly, we see that QMBS is more robust when the matter degrees of freedom are bosonic rather than fermionic, which we argue is due to the strong nonlocality of the wave function in the latter case. Our results further shed light on the intimate connection between gauge symmetry and QMBS, and highlight the persistence of scarring in higher spatial dimensions. Our findings can be tested in near-term analog and digital quantum simulators, and we demonstrate their accessibility on a recently proposed cold-atom analog quantum simulator.

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I. INTRODUCTION

Quantum many-body scarring (QMBS) is an exciting paradigm of ergodicity breaking in isolated quantum many-body models that are expected to thermalize [1–7]. Despite being ergodic, certain models host special nonthermal *scar* eigenstates that are roughly equally spaced in energy over the whole spectrum [8,9] and exhibit anomalously low bipartite entanglement entropy [10,11]. Upon preparing an initial state with a high overlap with these scar eigenstates, the system avoids thermalization and its dynamics exhibits long-lived coherent oscillations lasting well beyond all relevant timescales, with the time-evolved wave function undergoing persistent periodic revivals [8,12]. QMBS is of great importance in investigations of the eigenstate thermalization hypothesis [13–16], as it facilitates violations of the latter through mechanisms based on spectrum-generating algebras [17–20] and nonthermal-eigenstate embedding [21].

QMBS has also been the subject of various experiments in both analog and digital quantum simulators [1,22–26].

QMBS also has an intimate connection to lattice gauge theories (LGTs) [27–33], which are interacting quantum many-body models hosting gauge symmetries that enforce an intrinsic relation between the local distribution of matter and the allowed corresponding configurations of the gauge fields [34,35]. Indeed, the quantum Ising-like model realized in a Rydberg setup that produced the first instance of QMBS [1] can be effectively described by the PXP model, which Surace *et al.* [27] have shown to map exactly onto the 1+1D spin-1/2 U(1) quantum link model (QLM) [36,37]. The spin-1/2 U(1) QLM is a formulation of the Schwinger model in which the gauge and electric fields are represented by spin-1/2 operators. This truncation has facilitated experimental feasibility in the realization of large-scale quantum simulators of the spin-1/2 U(1) QLM [24,38–40], with proposed algorithms for digital platforms [41], while still capturing salient features of the Schwinger model such as Coleman’s phase transition [42]. Given the ongoing strong drive of realizing quantum simulators of LGTs [43–53], it is important to fully understand the connection between QMBS and LGTs, both from a fundamental point of view and also to facilitate experimental investigations.

In particular, it is interesting to see how QMBS behaves in 2+1D LGTs. Indeed, it is known that nonergodic features

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in 1+1D interacting models tend to vanish in higher spatial dimensions $d > 1$, with many-body localization [54–57] being a prime example. How the *weak* ergodicity-breaking mechanism of QMBS fares for $d = 2$ is a question that has recently received some attention in the context of the PXP model [58,59] and helix states in the XXZ model [4], in addition to other quasi-2+1D systems [25,26]. Nevertheless, scarred models in $d > 1$ are few and far between compared to their counterparts in $d = 1$ [5–7]. This further motivates adding to the collection of such models by investigating LGTs with dynamical matter in higher spatial dimensions.

Importantly, QMBS in LGTs in $d > 1$ has been studied [30,60,61], but only in the case without dynamical matter, where connections to high-energy physics (HEP) phenomena are not as straightforward. In trying to see whether QMBS will persist in the quantum-field-theory limit of gauge theories, such as in quantum electrodynamics, the inclusion of dynamical matter is crucial [31,32]. Indeed, the standard model of particle physics is a framework that includes dynamical matter, and omitting the latter completely removes any insight into how matter interaction and the concomitant gauge-boson scattering affect QMBS. As a testament to this, as we show here, the matter particle statistics has a major effect on QMBS.

In this study, we consider the 2+1D spin-1/2 U(1) QLM with dynamical matter on a square lattice and show that QMBS persists for special far-from-equilibrium quenches. We also showcase how these QMBS regimes can be detected in near-term analog quantum simulators of LGTs in $d = 2$ spatial dimensions.

II. MODEL

We consider the U(1) QLM on a square lattice described by the Hamiltonian [36,37]

$$\hat{H} = \sum_{\mathbf{r}} \left[-\frac{\kappa}{2} (\hat{\phi}_{\mathbf{r},\mathbf{e}_x}^\dagger \hat{\phi}_{\mathbf{r}+\mathbf{e}_x} + \alpha (-1)^{r_y} \hat{\phi}_{\mathbf{r},\mathbf{e}_y}^\dagger \hat{\phi}_{\mathbf{r}+\mathbf{e}_y} + \text{H.c.}) + m (-1)^{r_x+r_y} \hat{\phi}_{\mathbf{r}}^\dagger \hat{\phi}_{\mathbf{r}} - \alpha^2 J (\hat{U}_{\square_{\mathbf{r}}} + \hat{U}_{\square_{\mathbf{r}}}^\dagger) \right]. \quad (1)$$

The matter degrees of freedom with mass m are represented by the ladder operators $\hat{\phi}_{\mathbf{r}}^{(\dagger)}$ at site $\mathbf{r}$. These can either be fermionic or hard-core bosonic, and we will consider both in this study. The gauge (electric) field is represented by the spin-1/2 operator $\hat{s}_{\mathbf{r},\mathbf{e}_a}^{(z)}$, where the subscript denotes the bond between sites $\mathbf{r}$ and $\mathbf{r} + \mathbf{e}_a$, where $\mathbf{e}_a$ is a unit vector ($a = x, y$). The magnetic interactions of the gauge fields, whose strength is proportional to J , are represented by the plaquette operators $\hat{U}_{\square_{\mathbf{r}}} = \hat{s}_{\mathbf{r},\mathbf{e}_x}^+ \hat{s}_{\mathbf{r}+\mathbf{e}_x,\mathbf{e}_y}^+ \hat{s}_{\mathbf{r}+\mathbf{e}_y,\mathbf{e}_x}^- \hat{s}_{\mathbf{r},\mathbf{e}_y}^-$. The coefficient α is used to tune the ratio of the coupling along the y and x axes, such that the couplings are equally strong at $\alpha = 1$, and at $\alpha = 0$, only the coupling along the x axis remains, and the system behaves like a collection of uncoupled 1+1D QLMs. Thus, by using $0 < \alpha < 1$, we can interpolate between the 2+1D square lattice ($\alpha = 1$) and the 1+1D model ($\alpha = 0$). In 1+1D, this model has been rigorously shown to be nonintegrable [62], and so its 2+1D version is expected to also be nonintegrable.

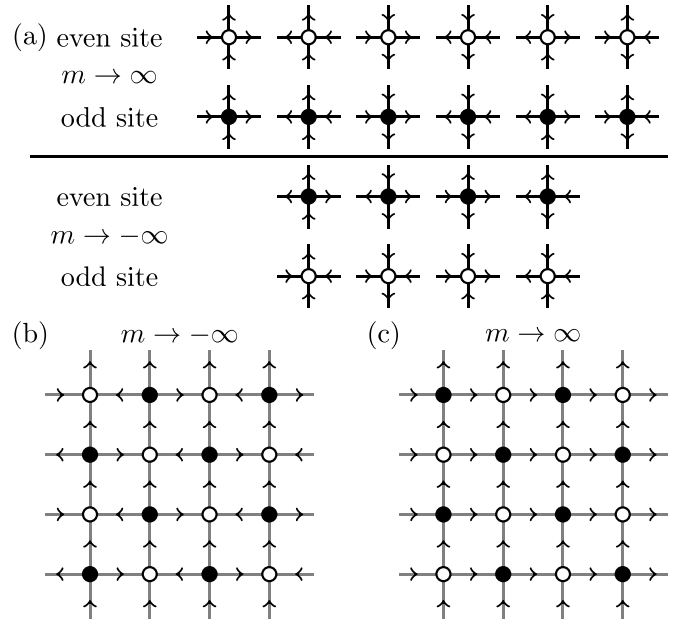


FIG. 1. (a) The gauge-invariant configurations for the links surrounding each matter site in the spin-1/2 U(1) quantum link model (1), according to Gauss’s law (2). The site parity refers to the parity of the sum of the x and y components of the index $r_x + r_y$. Arrows pointing right or up on gauge sites represent eigenstates of $\hat{s}^z$ with eigenvalue $+1/2$, while arrows pointing left or down represent the eigenvalue $-1/2$. (b) A charge-proliferated ground state at $m \rightarrow -\infty$. (c) A vacuum ground state at $m \rightarrow \infty$.

The U(1) gauge symmetry of this model is generated by the operators

$$\hat{G}_{\mathbf{r}} = \hat{\phi}_{\mathbf{r}}^\dagger \hat{\phi}_{\mathbf{r}} - \frac{1 - (-1)^{r_x+r_y}}{2} - \sum_{a=x,y} [\hat{s}_{\mathbf{r},\mathbf{e}_a}^z - \hat{s}_{\mathbf{r}-\mathbf{e}_a,\mathbf{e}_a}^z], \quad (2)$$

which act as a discrete analog of Gauss’s law. We work in the physical sector of states, which are eigenstates of $\hat{G}_{\mathbf{r}}$ for each $\mathbf{r}$ with eigenvalue zero. Given a configuration of the matter sites, this becomes a constraint on the allowed configuration on the gauge sites: The allowed configurations surrounding occupied and unoccupied matter sites are shown in Fig. 1(a).

III. SCARRING DYNAMICS

Here, we consider the dynamics following the global quench of an initial charge-proliferated state $|\psi(t=0)\rangle$, being the gauge-invariant ground state at $m/\kappa \rightarrow -\infty$, using the gauge site configuration shown in Fig. 1(b) (which is chosen such that at $\alpha = 0$, the decoupled chains are in the physical gauge sector of the 1+1D model). We quench to a finite value of the mass m , which is known to lead to scarred dynamics in the 1+1D model [24,63,64], in particular, we quench to $m = 0.84\kappa$, which is the regime considered in the experiment of Ref. [24], while setting $J = 0$. We use numerical time-evolution simulations to obtain these dynamics based on matrix product state (MPS) techniques [65–67]. Specifically, we use the time-dependent variational principle algorithm [68], using single-site updates with adaptive bond dimension expansion [69,70]. Simulations were performed

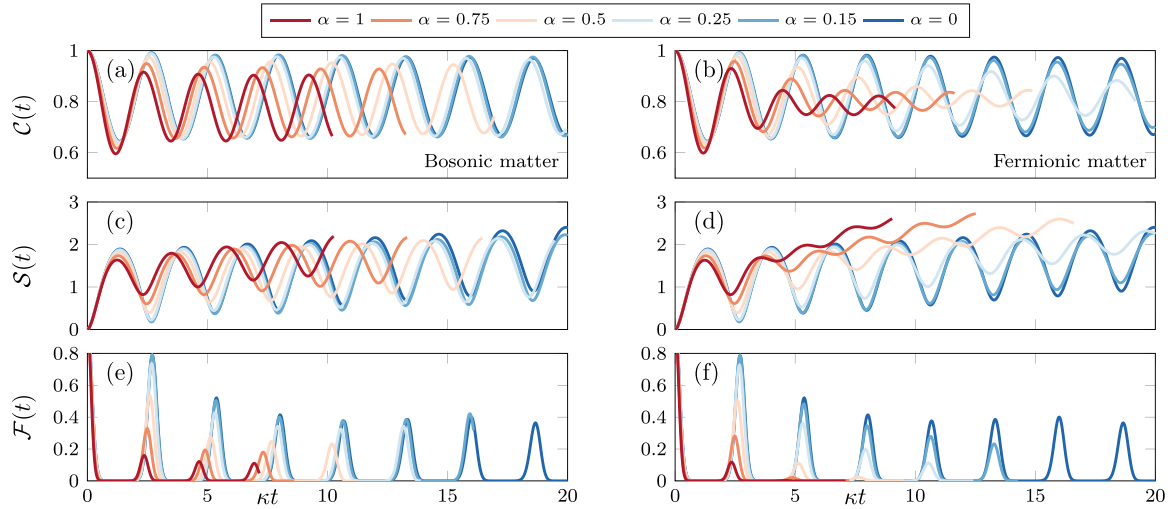


FIG. 2. A numerical time-evolution simulation of the quench of the charge-proliferated state shown in Fig. 1(b) to $m = 0.84\kappa$, using matter with bosonic (left column) and fermionic (right column) statistics on a $L_y = 4$ cylinder, as the coupling ratio α tuned from 1 to 0, effectively tuning between the full two-dimensional (2D) model and the one-dimensional (1D) model. (a), (b) The chiral condensate $\mathcal{C}(t)$ (3) and (c), (d) entanglement entropy across a circumferential slice $\mathcal{S}(t)$, calculated for an infinite-length cylinder. (e), (f) The fidelity $\mathcal{F}(t)$, calculated for a finite cylinder of dimension 4×16 with open boundaries in the x direction.

using a cylindrical lattice geometry, with a circumference of $L_y = 4$ matter sites: We perform calculations using states that are explicitly translation invariant along the x axis, as well as states with a finite length $L_x = 16$ in order to calculate the fidelity with the initial state, using open boundaries at the ends of the cylinder.

The quench simulation results are displayed in Fig. 2, for both bosonic and fermionic matter statistics, as the coupling ratio α is tuned from 1 to 0. As we increase α , the model becomes more “two dimensional,” and more entanglement is generated in the time evolution, and hence a larger MPS bond dimension is required to accurately simulate the dynamics. This therefore limits the largest accessible evolution time as we increase α . In panels (a) and (b), we show the expectation value of the chiral condensate $\mathcal{C}(t) = \langle \psi(t) | \hat{\mathcal{C}} | \psi(t) \rangle$,

$$\hat{\mathcal{C}} = \frac{1}{L_x L_y} \sum_{\mathbf{r}} (-1)^{r_x + r_y} \left(\hat{\phi}_{\mathbf{r}}^\dagger \hat{\phi}_{\mathbf{r}} - \frac{1 - (-1)^{r_x + r_y}}{2} \right), \quad (3)$$

in panels (c) and (d), we show the von Neumann entanglement entropy $\mathcal{S}(t)$ measured using a bipartition formed by a slice along the circumference of the cylinder, and in panels (e) and (f), we show the fidelity with the initial state $\mathcal{F}(t) = |\langle \psi(0) | \psi(t) \rangle|^2$ for the finite system.

For the quenches with bosonic matter, we can see the signatures of QMBS are qualitatively preserved, although less pronounced, as α increases from zero (where the dynamics will correspond to the 1+1D QLM) up to one (where we retrieve the fully fledged 2+1D QLM). The oscillatory behavior of the chiral condensate [Fig. 2(a)] and entanglement entropy [Fig. 2(c)] is clearly visible for all values of α throughout the whole simulation, and the fidelity [Fig. 2(e)] shows revivals of a consistent, though weaker, magnitude. However, in the fermionic matter quenches, the signatures of QMBS clearly weaken as α is increased. The oscillations in the chiral condensate [Fig. 2(b)] quickly decay in magnitude, the growth of the entanglement entropy [Fig. 2(d)] is

significantly more rapid, and the revivals in fidelity [Fig. 2(f)] also swiftly decrease in magnitude.

These results show that the fate of QMBS in the 2+1D QLM is strongly dependent on the particle statistics of the matter degrees of freedom, as the signs of scarring are preserved more strongly for bosonic matter, but not as much for fermionic matter. In contrast to 1+1D, where particle statistics are functionally irrelevant, in 2+1D we can exchange the position of two particles, which will lead to a change in sign of the wave function for fermionic matter, but not for bosonic matter. The respective strength and weakness of the signatures of scarring for bosonic and fermionic matter is most likely a result of constructive and destructive interference in wave function amplitudes in the coherent revivals of the initial configuration due to this sign change. Intuitively, the inherent nonlocality of the wave function in the case of fermionic matter adds an extra layer of difficulty for the wave function to return to its initial state.

We have investigated quenches to final masses m other than those considered in Fig. 2 (which we show in Appendix D). For larger m , the dynamics become trivial, while for m closer to zero, it becomes numerically infeasible to simulate the dynamics for a meaningful time. Thus, our chosen parameter lies within a sweet spot between these two regimes.

Furthermore, we calculated the results for width-2 cylinders (Appendix E), in which the stark qualitative difference between bosonic and fermionic matter is not visible, which we attribute to a restriction of particle-exchange dynamics. We can then interpret increasing to a width of 4 as causing a respective *strengthening* and *weakening* of the signatures of scarring for bosonic and fermionic matter. We also show results for width-6 cylinders: While the simulations are not able to reach as long times due to the increased growth of entanglement, these do demonstrate qualitatively the same behavior as for width 4. This confirms the applicability of our results on the width-4 cylinder to the full two-dimensional system. Indeed, previous studies have shown that a cylinder

of width 4 already captures the quantum critical point of the quantum Ising square-lattice model in $d = 2$ spatial dimensions [71].

We also find that the plaquette term weakens the QMBS observed in Fig. 2, with more suppressed signs of QMBS the larger J is (Appendix C). This suggests that the origin of the QMBS examined in our work is distinct from previous works focusing on pure gauge theories with only gauge coupling and no matter fields [30,60,61], and that the presence of dynamic matter is vital to the persistence of scarring in our quenches. Indeed, the stark difference in behavior between bosonic and fermionic matter fields further corroborates this.

We have also investigated quenching from the vacuum states of the 2+1D QLM, but we have found that the signs of QMBS are suppressed for both fermionic and bosonic matter quenching to $m = 0$, which is known to exhibit scarring in 1+1D (Appendix A). Our results are inconclusive as to whether there are other regimes where the scarring is more prominent in the 2 + 1 QLM starting from a vacuum initial state.

IV. QUANTUM SIMULATION

In order to demonstrate the relevance of these results to near-term quantum simulators, we show that the dynamics for the bosonic QLM [Figs. 2(a), 2(c), and 2(e)] can be probed using the 2D Bose-Hubbard simulator proposed in Ref. [72] [see Fig. 3(a) for the mapping of the charge-proliferated state], with the Hamiltonian

$$\hat{H}_{\text{BHM}} = \sum_{\mathbf{j}} \left[\tilde{J} \sum_{a=x,y} (\hat{b}_{\mathbf{j}}^\dagger \hat{b}_{\mathbf{j}+\mathbf{e}_a} + \text{H.c.}) + \frac{U_{\mathbf{j}}}{2} \hat{n}_{\mathbf{j}}(\hat{n}_{\mathbf{j}} - 1) + (\boldsymbol{\gamma} \cdot \mathbf{j} - \delta_{\mathbf{j}} - \eta_{\mathbf{j}}) \hat{n}_{\mathbf{j}} \right]. \quad (4)$$

For each matter site in the QLM, we have a corresponding 2×2 unit cell of sites in the Bose-Hubbard model (BHM): The BHM site with index $\mathbf{j} = (j_x, j_y)$ corresponds to a matter site in the QLM if both components are even, to a gauge site if exactly one of them is even, or if both components are odd, the site has no counterpart in the QLM, so it is to be forbidden from contributing to the dynamics. $\hat{b}_{\mathbf{j}}^{(\dagger)}$ and $\hat{n}_{\mathbf{j}} = \hat{b}_{\mathbf{j}}^\dagger \hat{b}_{\mathbf{j}}$ are the bosonic ladder and number operators, respectively, $\tilde{J}$ is the hopping energy, the on-site interaction strength $U_{\mathbf{j}}$ is equal to U on gauge and forbidden sites, but $\tilde{\alpha}U$ on matter sites, for some detuning constant $\tilde{\alpha}$, $\boldsymbol{\gamma} = (\gamma_x, \gamma_y)$ represents a linear tilt in the lattice in both dimensions, and the potential $\delta_{\mathbf{j}}$ ($\eta_{\mathbf{j}}$) is equal to δ (η) only on a gauge (forbidden) site, and zero elsewhere. Following Ref. [72], we use the values of the parameters $\tilde{J} = 30$ Hz, $U = 1300$ Hz, $\tilde{\alpha} = 1.3$, $\gamma_x = 57$ Hz, $\gamma_y = 73$ Hz, $\delta = 649.647$ Hz, and $\eta = 5\delta$. This corresponds to $m \approx -0.84\kappa$ in the QLM {quenches starting from the charge-proliferated initial state [Fig. 1(b)] have the same dynamics for positive and negative m }. The results of the quench of the charge-proliferated initial state in Fig. 3(a) are shown in Figs. 3(b)–3(d), compared with the QLM simulation [Figs. 2(a), 2(c), and 2(e)] using the simulator's effective value of κ , which shows great agreement for the available simulation times. This bodes well for upcoming cold-atom

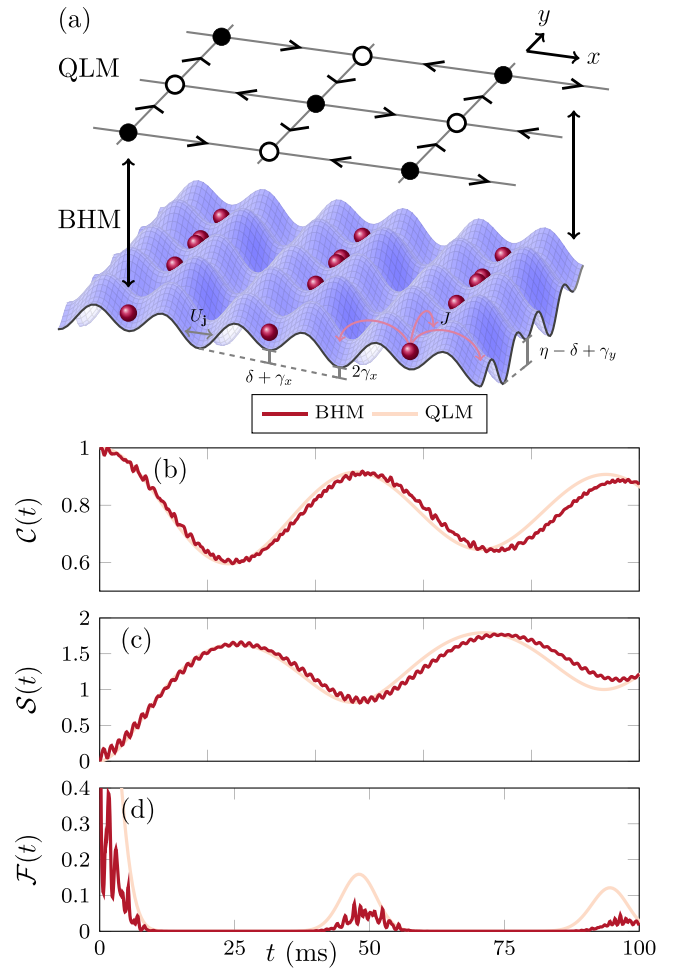


FIG. 3. (a) The mapping between the 2+1D U(1) QLM and the Bose-Hubbard simulator (4) proposed in Ref. [72]. For each matter site in the QLM, there is a corresponding unit cell of 2×2 sites in the BHM, with a shallow site representing the matter site, the deeper sites surrounding it corresponding to gauge sites, and the remaining “forbidden” site, which is detuned by a relatively large potential η . (b)–(d) A numerical time-evolution simulation of the quench of the charge-proliferated state shown in panel (a) to $m = 0.84\kappa$ on a $L_y = 4$ cylinder for Bose-Hubbard simulator (4), compared with the QLM simulation in Fig. 2, showing the chiral condensate $\mathcal{C}(t)$ (3), entanglement entropy $S(t)$, and fidelity with the initial state $\mathcal{F}(t)$.

experiments seeking to realize 2+1D U(1) QLMs, as it brings our findings into the realm of experimental accessibility.

V. DISCUSSION AND OUTLOOK

We have studied QMBS in a spin-1/2 2+1D U(1) QLM with dynamical matter using MPS simulations of quench dynamics. We have found that QMBS is robust when the initial state is the charge-proliferated product state and the quench mass is $m \approx 0.84\kappa$, particularly when the matter degrees of freedom are hard-core bosonic. In the case of fermionic matter, we have found that the signatures of scarring persist only for short times, but then vanish, and the dynamics looks quite ergodic. We have also adopted in our simulations a ratio α between coupling in the x and y

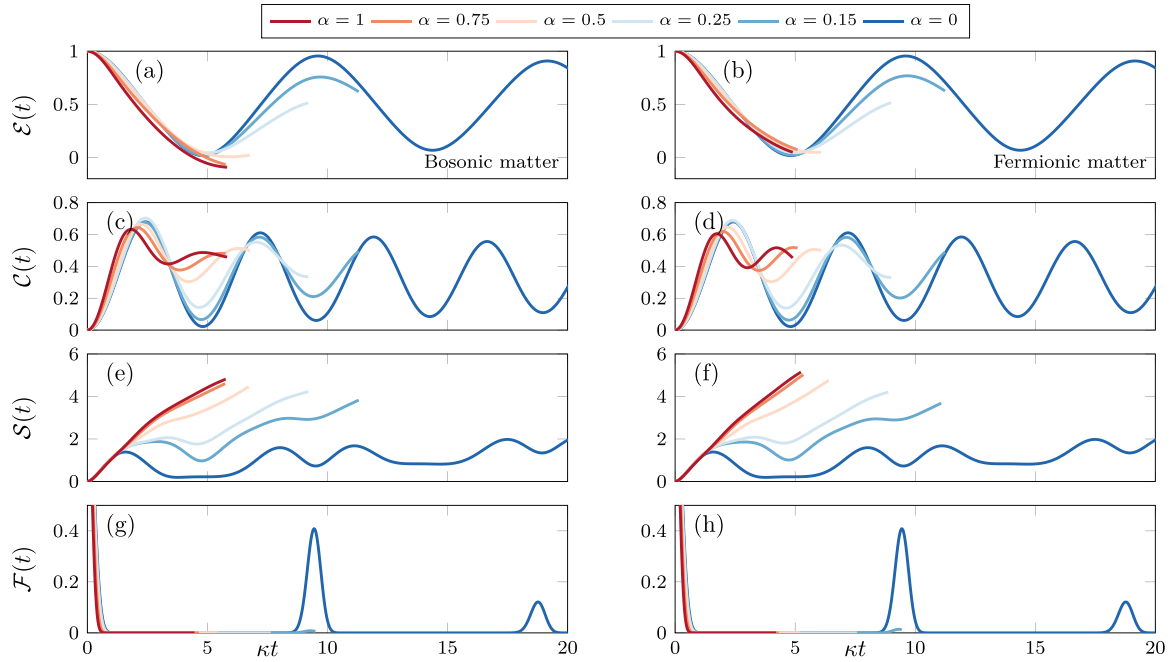


FIG. 4. The time evolution following a quench of the vacuum state shown in Fig. 1(c) to $m = 0$, using matter with bosonic (left column) and fermionic (right column) statistics on a $L_y = 4$ cylinder, as the coupling ratio α tuned from 1 to 0. We plot (a), (b) the flux $\mathcal{E}(t)$ (A1), (c), (d) chiral condensate $\mathcal{C}(t)$ (3), and (e), (f) entanglement entropy $\mathcal{S}(t)$ for an infinite-size system, and the (g), (h) fidelity $\mathcal{F}(t)$ for 4×16 -site system.

directions. By increasing α from 0 to 1, we can interpolate between the dynamics of the 1+1D and the fully fledged 2+1D system, so we can more clearly identify the respective persistence and weakening of scarring in spite of the heightened numerical cost of simulating the quench dynamics of systems in 2+1D. We have also shown that our results do not qualitatively differ over short timescales for larger system sizes (Appendix E), which gives confidence that our results are representative of the two-dimensional thermodynamic limit.

We have also showcased how our findings can be tested on a recently proposed cold-atom quantum simulator of the spin-1/2 U(1) QLM with hard-core bosonic matter degrees of freedom [72], showing great agreement between the simulator dynamics and those of the ideal QLM. However, our results are also amenable for investigation in other proposals that have been put forth [73–76]. Given that going to $d = 2$ spatial dimensions in quantum-simulator realizations of LGTs is the current frontier of the field [48], our work sets the stage for experimentally relevant features that can be probed on such devices once they are available.

While we have focused on dynamical aspects of scarring in terms of consistent revivals and suppressed growth of entropy, it is worth emphasizing that QMBS is primarily a feature of anomalous eigenstates in the spectrum of the Hamiltonian, which results in this behavior when quenching certain initial states. As such, it is vital to also study the spectrum of the QLM in $d = 2$ spatial dimensions. However, since the Hilbert space dimension of such systems can be prohibitively large for exact diagonalization calculations, we leave this to future study. And while we have found a stark qualitative difference in the quench dynamics between bosonic and fermionic matter fields, we have not uncovered

any clear physical mechanism for this difference. We hope that future study into this topic from different approaches may shed more light on the role of matter statistics on the fate of QMBS.

It is important to emphasize that our work does not rule out other QMBS regimes that may still persist in 2+1D but that we have not considered, such as those found in Ref. [77] for pure Abelian gauge theories with arbitrary integer spin. Indeed, our goal in this work was to investigate the fate of QMBS regimes discovered in the 1+1D U(1) QLM and see how well they fare in 2+1D. It was also our intention to find out how the statistics of the matter degrees of freedom affects QMBS, which, as we show, happens in a profound way. One avenue for future work is to study the fate of the QMBS regime we find in 2+1D for higher-level representations ($S > 1/2$) of the electric and gauge fields. In 1+1D, it has been shown that scarring persists and is robust for $S > 1/2$ [31,32]. This is important to assess the fate of QMBS in the limit of 2+1D lattice QED, and a recent proposal may allow experimental observation in a quantum simulator of the spin-1 U(1) QLM [78]. A second venue involves adding a topological θ term [79,80] and studying its interplay with scarring in 2+1D, which has been shown to lead to a plethora of nonergodic behavior in 1+1D [81]. Understanding whether scarring persists even in the 3+1D QLM would also be a very valuable future endeavor.

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APPENDIX A: QUENCHES OF THE VACUUM STATE

In Fig. 4, we show quenches of a vacuum initial state [Fig. 1(c)] to $m = 0$. Here, we also show the expectation value of the electric flux $\mathcal{E}(t) = \langle \psi(t) | \hat{\mathcal{E}} | \psi(t) \rangle$,

$$\hat{\mathcal{E}} = \frac{1}{L_x L_y} \sum_{\mathbf{r}} \sum_{a=x,y} \hat{s}_{\mathbf{r},\mathbf{e}_a}^z. \quad (\text{A1})$$

In the 1+1D case ($\alpha = 0$), this quench exhibits QMBS, but for $\alpha > 0$, the signs of QMBS are suppressed for both bosonic and fermionic matter fields, with little qualitative difference between the two different particle statistics in the accessible evolution times.

APPENDIX B: QUENCHES STARTING AT A FINITE MASS

Here, we examine a sudden quench from a finite value of the mass, starting with the gauge-invariant ground state at

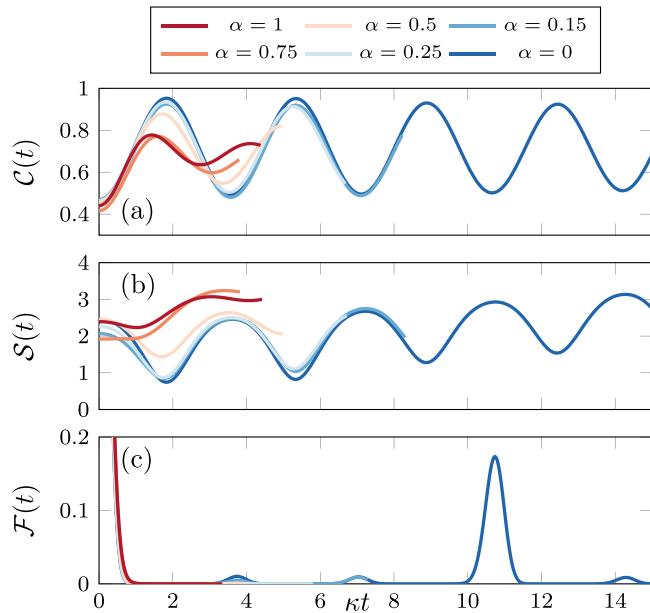


FIG. 5. A quench of the gauge-invariant ground state at $m_i = 0.19\kappa$ to $m_f = -0.4\kappa$, using matter with bosonic statistics on a $L_y = 4$ cylinder, as the coupling ratio α tuned from 1 to 0. We plot (a) the chiral condensate $\mathcal{C}(t)$ (3) and (b) entanglement entropy $\mathcal{S}(t)$ for an infinite-size system, and (c) the fidelity $\mathcal{F}(t)$ for 4×16 -site system.

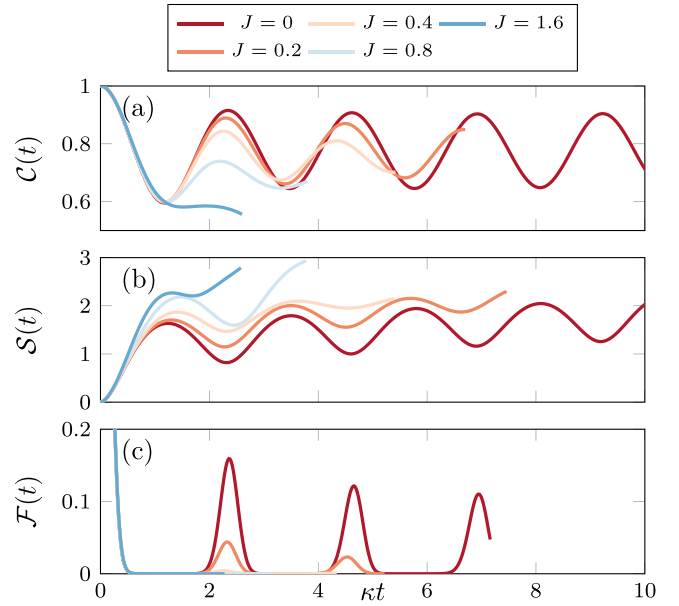


FIG. 6. A quench of the charge-proliferated state shown in Fig. 1(b) to $m = 0.84\kappa$, using matter with bosonic statistics on a $L_y = 4$ cylinder, as the magnetic coupling J is increased, fixing the coupling ratio α to be 1. We plot (a) the chiral condensate $\mathcal{C}(t)$ (3) and (b) entanglement entropy $\mathcal{S}(t)$ for an infinite-size system, and (c) the fidelity $\mathcal{F}(t)$ for 4×16 -site system.

$m_i = 0.19\kappa$, and evolving the state at $m_f = -0.4\kappa$, which was shown to display strong signs of QMBS in the 1+1D model in Ref. [24]. As shown in Fig. 5, however, for $\alpha > 0$ these signs are strongly suppressed. The way the gauge-invariant ground state at $m_i = 0.19\kappa$ changes as α is increased from 0 to 1 would play a significant role here, and requires further investigation.

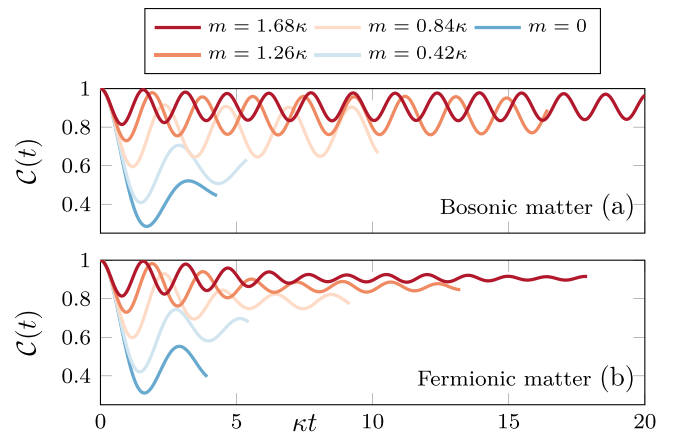


FIG. 7. A quench of the charge-proliferated state shown in Fig. 1(b) to different final masses m , using matter with (a) bosonic and (b) fermionic statistics on a $L_y = 4$ cylinder, fixing the coupling ratio α to be 1, showing the chiral condensate $\mathcal{C}(t)$.

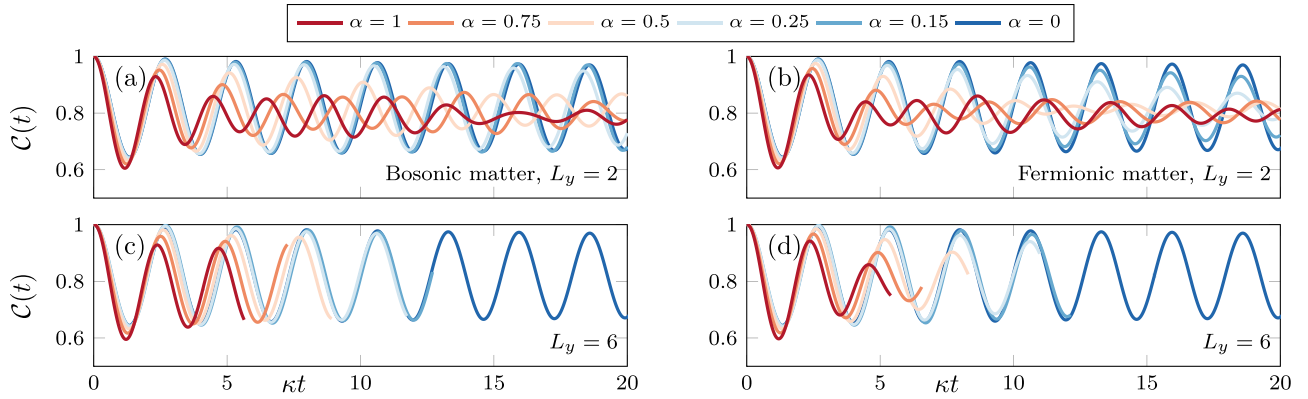


FIG. 8. A quench of the charge-proliferated state shown in Fig. 1(b) to $m = 0.84\kappa$, using matter with bosonic (left column) and fermionic (right column) statistics on a cylinder with $L_y = 2$ (top row) and 6 (bottom row), as we tune the coupling ratio α , showing the chiral condensate $C(t)$.

APPENDIX C: EFFECT OF THE MAGNETIC COUPLING OF THE GAUGE FIELDS

In Fig. 6, we examine the effect of the magnetic coupling of the gauge fields, controlled by the plaquette term with coefficient J in the 2+1D QLM Hamiltonian [Eq. (1)]. We look at the quench from the charge-proliferated state [Fig. 1(b)] to $m = 0.84\kappa$ for bosonic matter, which was shown to display strong signs of QMBS (Fig. 2). However, for $J > 0$, we can see that these signs are quickly suppressed.

APPENDIX D: EFFECT OF CHANGING THE QUENCH MASS

In Fig. 7, we show the effect of changing the final mass m of the quench considered in Fig. 2, fixing the coupling ratio α to be one. As we increase m , the oscillations become smaller as the dynamics become more trivial, since the initial state becomes closer to being an eigenstate (as it is in the limit $m \rightarrow \infty$). Nevertheless, the qualitative distinction between fermionic and bosonic statistics is still clearly visible. As we decrease m toward zero, the oscillations become larger, but the accessible evolution times through MPS numerics are severely limited by the growth of entanglement, so we cannot infer much information about the long-time dynamics in these

cases. Therefore, in a sense, the quench value $m = 0.84\kappa$ lies within a sweet spot where the dynamics are nontrivial, but still numerically tractable for cylinder width $L_y = 4$.

APPENDIX E: RESULTS FOR OTHER CYLINDER WIDTHS

In Figs. 8(a) and 8(b), we plot results of the quench in Fig. 2 performed for cylinder width $L_y = 2$ (as opposed to $L_y = 4$ in the main text). In this case, we do not see as strong a qualitative distinction between the two particle statistics as in Fig. 2, with the strength of oscillations on the width-2 cylinder lying between the fermionic and bosonic width-4 results. We attribute this to the fact that dynamics on a width-2 cylinder is fairly restricted, in the sense that the process of exchanging the position of two particles—which leads to the different physical behavior of bosons and fermions in two, and not one, spatial dimensions—is less likely to occur. We can then interpret the increase to width 4 as leading to a respective *strengthening* and *weakening* of the signatures of scarring for bosonic and fermionic matter.

In Figs. 8(c) and 8(d), on the other hand, we plot the same quench for cylinder width $L_y = 6$. While we cannot observe the dynamics on such timescales as in Fig. 2 for $L_y = 4$, we can indeed confirm that the same qualitative difference between bosonic and fermionic matter fields persists.

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