

Disorder-Free Localization and Fragmentation in a Non-Abelian Lattice Gauge Theory

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We investigate how isolated quantum many-body systems dynamically equilibrate under non-Abelian gauge-symmetry constraints. By encoding gauge superselection sectors into static SU(2) background charges, we map out the dynamical phase diagram of a $(1 + 1)$ D SU(2) lattice gauge theory with dynamical matter. We uncover three distinct regimes: (i) an ergodic phase, (ii) a fragmented phase that is nonthermal but delocalized, and (iii) a disorder-free many-body localized regime. In the latter, a superposition of gauge superselection sectors preserves spatial matter inhomogeneities in time, as evidenced by distinct temporal scalings of entropy. We highlight the non-Abelian nature of these phases and argue for potential realizations on qubit processors.

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Introduction—Lattice gauge theories (LGTs) provide a nonperturbative formulation of continuum gauge theories [1–3], where spacetime is discretized into a grid and gauge fields are represented by link variables [4]. Originally developed to investigate quark (de)confinement in quantum chromodynamics [5–7], LGTs have become central to a surge of quantum simulation experiments [8–48] in the last decade. More recently, they have emerged as powerful tools for studying quantum many-body (QMB) dynamics [49]. In particular, due to the presence of gauge symmetry [50,51], which enforces strong local constraints on dynamics, LGTs have been successful in discovering and exploring novel types of nonthermal behavior that circumvent the Eigenstate Thermalization Hypothesis (ETH) [52–60]. In generic interacting isolated quantum systems, if ETH holds, then for initial pure states with a narrow energy distribution, local observables (or equivalently, sufficiently small subsystems) relax at long times to values predicted by the thermal ensemble fixed by the initial state energy [56,61]. A paradigmatic ergodicity-breaking phenomenon in LGTs is QMB scarring [62–75], where special eigenstates, often equally spaced in energy, exhibit anomalously

low entanglement entropy [76–78]. Scars have been experimentally observed in LGTs [18,62,79] and theoretically shown to persist toward the lattice quantum field theory limit [80,81], in one and higher spatial dimensions [82,83], for Abelian [84–93] and non-Abelian gauge groups [75].

Other examples of nonthermal dynamics in LGTs are Hilbert space fragmentation (HSF) [94,95] and disorder-free localization (DFL) [96,97]. In the first case, the Hilbert space, even after resolving all global and local symmetries, shatters into dynamically disconnected Krylov subsectors—subspaces spanned by the repeated action of the Hamiltonian on a given initial state—whose number grows exponentially in system size [69,98]. HSF avoids thermalization and can lead to slow or glassy dynamics confined within strict subspaces from the outset. In the case of DFL, an initial spatial imbalance in the matter fields, which typically melts away in the absence of explicit disorder, survives in time if the system is initialized in a superposition of an extensive number of gauge superselection sectors. DFL has been theoretically demonstrated in Abelian LGTs [99–118] and experimentally observed on a Google Quantum AI quantum processor [46], it but has not been shown in a non-Abelian LGT before now.

The type of underlying symmetries plays a central role in the equilibration of QMB systems [50,51,119,120]. Non-Abelian symmetries, such as SU(2) and SU(3), are ubiquitous in condensed matter and high-energy physics and display noncommuting conserved quantities that cannot be simultaneously resolved, greatly enriching the underlying physics [121–130]. Their role in QMB dynamics has

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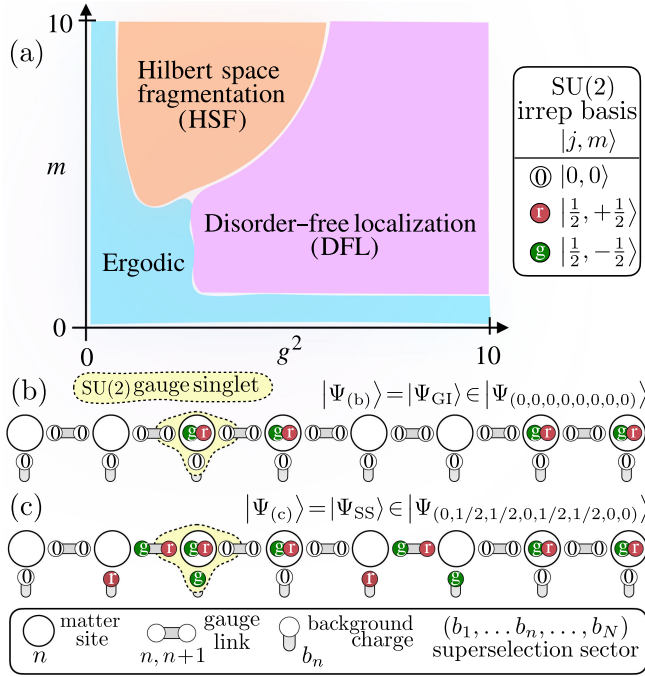


FIG. 1. Model phases and initial-states. (a) Sketch of the explored dynamical phases as a function of mass and coupling strengths. (b),(c) Cartoon picture of two possible charge density wave states (on $N = 8$ site-chain) forming the superposition initial state in Eq. (3): (b) belongs to the SU(2) gauge invariant (GI) sector; (c) belongs to a superselection-sector (SS) where SU(2) Gauss law is modified with SU(2) static background charges on four lattice vertices. Circles, gray ovals, and bottom stubs denote matter sites, gauge links, and static background charges expressed in SU(2) irrep bases.

generated considerable recent interest [131,132], including a trapped-ion experiment [133].

In this Letter, we investigate the role of non-Abelian symmetries in the long-time dynamics of LGTs. We observe for the first time the emergence of DFL in an SU(2) LGT coupled to dynamical matter and find evidence for HSF; see Fig. 1. The uncovered rich (non)thermal dynamics is revealed through several measures, such as matter imbalance, quark population, and entanglement entropy scaling, which reflect the underlying non-Abelian symmetry constraints.

Model—As a prototypical model to observe DFL in non-Abelian LGTs, we consider a $(1+1)$ D hardcore-gluon SU(2) Yang–Mills LGT coupled with flavorless dynamical matter [75,134–136]. This model is described by the following Kogut–Susskind Hamiltonian [137], on an N -site lattice chain with lattice spacing a_0 ($c, \hbar = 1$):

$$\hat{H}_0 = \frac{1}{2a_0} \sum_n \sum_{\alpha,\beta} \left[i\hat{\psi}_{n,\alpha}^\dagger \hat{U}_{n,n+1}^{\alpha\beta} \hat{\psi}_{n+1,\beta} + \text{H.c.} \right] + m_0 \sum_n \sum_{\alpha} (-1)^n \hat{\psi}_{n,\alpha}^\dagger \hat{\psi}_{n,\alpha} + \frac{a_0 g_0^2}{2} \sum_n \hat{E}_{n,n+1}^2. \quad (1)$$

The model in Eq. (1) describes the interaction between flavorless quark matter fields of mass m_0 , living on lattice sites n , and SU(2) gauge fields living on lattice links $(n, n+1)$, coupled with strength g_0 [see Figs. 1(b) and 1(c)]. The quark field is represented as a staggered fermion doublet $\hat{\psi}_{n,\alpha}$, which satisfies the anticommutation relations $\{\hat{\psi}_{n,\alpha}, \hat{\psi}_{n',\beta}^\dagger\} = \delta_{n,n'} \delta_{\alpha,\beta}$ and $\{\hat{\psi}_{n,\alpha}, \hat{\psi}_{n',\beta}\} = 0$ [138], where α, β indices live in the fundamental SU(2) irreducible representation (irrep) ($j = 1/2, m = \pm 1/2$). The single-site matter basis is given by $\{|0\rangle, |r\rangle = \hat{\psi}_r^\dagger |0\rangle, |g\rangle = \hat{\psi}_g^\dagger |0\rangle, |2\rangle = \hat{\psi}_r^\dagger \hat{\psi}_g^\dagger |0\rangle\}$, where $\{r, g\}$ are shorthand notations for $\{m = \pm 1/2\}$, while j is implicit [75,135,136] (see also Fig. 1).

For gauge link states, we adopt the chromoelectric basis $|j, m_L, m_R\rangle$, where $j \in \mathbb{N}/2$ indicates the spin irreps and $m_R, m_L \in \{-j, \dots, +j\}$ label the states within the spin shell j . In this basis, the link energy density operator is diagonal and coincides with the quadratic Casimir, $\hat{E}^2 |j, m_L, m_R\rangle = j(j+1) |j, m_L, m_R\rangle$ [139]. In this Letter, we implement the hardcore-gluon truncation and restrict $j \in \{0, 1/2\}$, retaining only the states that can be reached from the singlet $|00\rangle$ by applying the parallel transporter $\hat{U}^{\alpha\beta}$ at most once. The resulting truncated basis is $\{|00\rangle, |rr\rangle, |gg\rangle, |gr\rangle, |rg\rangle\}$ [75,135,136].

Non-Abelian SU(2) gauge invariance is locally manifest at each lattice site n through generators of local rotations $\hat{G}_n = (\hat{G}_n^x, \hat{G}_n^y, \hat{G}_n^z)$ that satisfy the following relation:

$$\hat{G}_n |\Psi_{\{b_n^\kappa\}}\rangle = b_n^\kappa |\Psi_{\{b_n^\kappa\}}\rangle, \quad \forall n, \kappa, \quad (2)$$

where b_n^κ is the background charge present on the lattice site n of a QMB state $|\Psi_{\{b_n^\kappa\}}\rangle$ within a specific *gauge superselection sector* κ , i.e., a specific lattice configuration of background charges $\{b_n^\kappa\} = (b_1^\kappa, \dots, b_N^\kappa)$. Similar to the gauge links, each background charge b_n^κ can be expressed in an irrep basis $|j_b, m_b\rangle$, with $j_b \in \mathbb{N}/2$ [see Figs. 2(b) and 2(c)]. Restricting to the first two irreps, i.e., $j_b \in \{0, 1/2\}$ is sufficient to detect DFL. In numerical simulations, we locally enforce the correct background charge in Eq. (2) using a dressed site approach [75,135,136], which yields a deformationized [140] qudit model with a 13-dimensional single-site basis. The first six states with $|b_0\rangle = |j_b = 0, m_b = 0\rangle$ correspond to the sector with no background charges (i.e., the Gauss Law), while the remaining seven states have background charges $\{|b_r\rangle, |b_g\rangle\} = \{|j_b = 1/2, m_b = \pm 1/2\rangle\}$. Its derivation is detailed in Supplemental Material (SM) [134]. For numerical convenience, we rescale the Hamiltonian in Eq. (1) as $(\hat{H}, m, g^2) = 4\sqrt{2}a_0(\hat{H}_0, m_0, 3a_0g_0^2/16)$ [75,136] and use exact diagonalization on a chain of $N = 8$ SU(2) dressed sites with periodic boundary conditions.

Disorder-free localization—To observe DFL, we initialize the dynamics from an equal superposition of all the $\mathcal{N}_{SS}$ gauge superselection sectors (SS) of background charges:

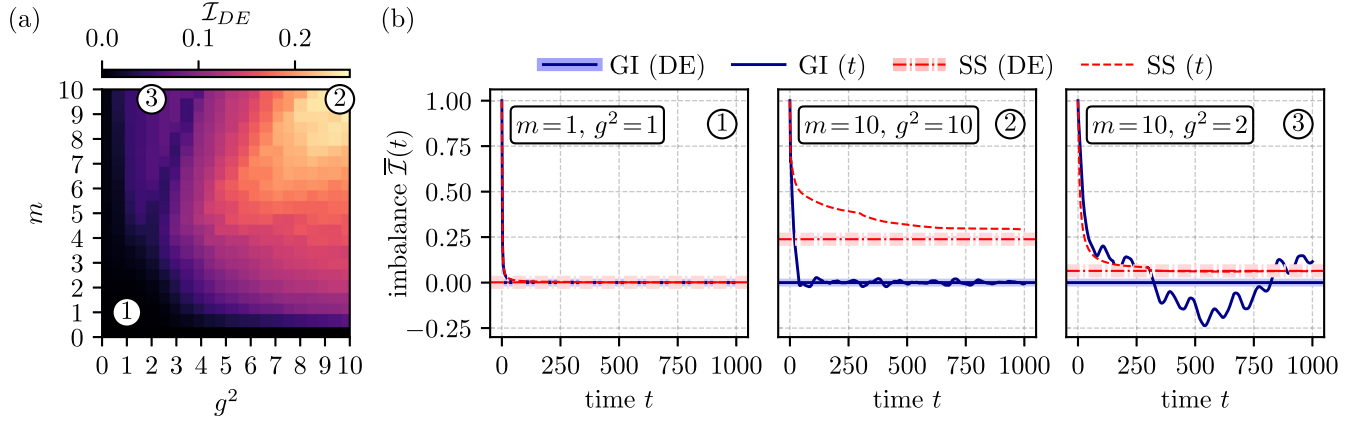


FIG. 2. Matter imbalance. (a) DE prediction of the imbalance in Eq. (5) as a function of mass and coupling strength in a grid with $m \in [0, 10]$ and $g^2 \in [0, 10]$. The three numbered white circles indicate the parameters used in panel (b) to characterize the ergodic (1), DFL (2), and HSF (3) phases in Fig. 1. (b) Time evolution of the (time-averaged) imbalance obtained from the SS and GI initial states in Eqs. (3) (dashed red line) and (4) (solid blue line), respectively, for three different mass-coupling configurations. The averaged imbalance is always compared with the corresponding DE (horizontal lighter line) prediction.

$$|\Psi_{\text{SS}}\rangle = \frac{1}{\sqrt{\mathcal{N}_{\text{SS}}}} \sum_{\kappa=1}^{\mathcal{N}_{\text{SS}}} |\Psi_{\{b_n^\kappa\}}\rangle, \quad b_n^\kappa \in \{0, 1/2\}, \quad \forall n, \kappa. \quad (3)$$

Specifically, each background-charge sector contributes with two simple translation-invariant product states where (i) matter fields are in a two-site charge density wave configuration (staggered chain with two empty sites alternated with two doubly occupied ones) and (ii) gauge fields can be in any of the $j \in \{0, 1/2\}$ SU(2) irreps. In this way, the initial state encodes a finite imbalance in the matter configuration under the Hamiltonian in Eq. (1) and evolves in a superposition of all the possible gauge field configurations. A pictorial representation of some of these contributions is in Figs. 1(b) and 1(c).

For direct comparison, we also consider the same matter configuration of the state in Eq. (3) with no active gauge fields, which belongs to the gauge-invariant (GI) sector without background charges [see Fig. 1(b)]:

$$|\Psi_{\text{GI}}\rangle \in |\Psi_{[0,0,\dots,0]}\rangle. \quad (4)$$

To detect the occurrence of DFL in the system, we introduce the following matter imbalance:

$$\mathcal{I}(t) = \frac{1}{N} \sum_{n=1}^N w_n \rho_n(t), \quad (5)$$

where $0 \leq \rho_n(t) \leq 2$ is the single-site fermion density,

$$\rho_n(t) \equiv \sum_{\alpha} \langle \Psi(t) | \hat{\psi}_{n,\alpha}^\dagger \hat{\psi}_{n,\alpha} | \Psi(t) \rangle, \quad (6)$$

and the coefficients $w_n = \rho_n(0) - 1$ define the weights associated with the considered initial matter configuration. The matter imbalance in Eq. (5) quantifies how much the

system departs from or retains the matter configuration of the initial state. To compute the long-time average of this and any other local observable $\hat{O}(t)$, we use the diagonal ensemble (DE) [56,134], which, for the initial state expressed in the Hamiltonian eigenbasis, $|\Psi(0)\rangle = \sum_s C_s |\Phi_s\rangle$, the long-time average is $O_{\text{DE}} = \sum_s |C_s|^2 \langle \Phi_s | \hat{O} | \Phi_s \rangle$. In Fig. 2(a), we plot the DE imbalance $\mathcal{I}_{\text{DE}}$ for the SS state of Eq. (3) as a function of mass and gauge coupling strength. The plot shows that, for large coupling and moderately large masses, ($g \gg 1, m \gtrsim 1$), the long-time average of the imbalance $\mathcal{I}_{\text{DE}}$ acquires finite values, signaling long-time localization of the initial configuration.

The dynamics of the imbalance approaching the long-time average are shown in Fig. 2(b). There, to mitigate the effect of finite-size oscillations, we consider the time-averaged imbalance $\bar{\mathcal{I}}(t) = (1/\Delta t) \int_{t-\Delta t/2}^{t+\Delta t/2} ds \mathcal{I}(s)$, for an intermediate time window $\Delta t = 300$. We focus on three parameter points [circles 1–3 in Fig. 2(a)] representing the ergodic (1), DFL (2), and HSF (3) regions sketched in Fig. 1(a), and compare dynamics from the GI and SS initial states. For the GI initial state, the imbalance $\bar{\mathcal{I}}(t)$ does not remain finite at long times in any of the three regimes; in the HSF region, it relaxes more slowly than in the ergodic case. By contrast, for the SS initial state, the ergodic region shows decay of the imbalance, whereas in the HSF and DFL regions $\bar{\mathcal{I}}(t)$ approaches a nonzero DE value, indicating persistent memory of the initial matter pattern. These results suggest that the model may exhibit DFL in specific parameter regimes, as we show in more detail in the following analysis.

Particle creation—To confirm that the observed localization arises from the superposition of different superselection sectors, we must first ensure that the GI system (without background charges) exhibits ergodic behavior. We define the single-site occupancy operator

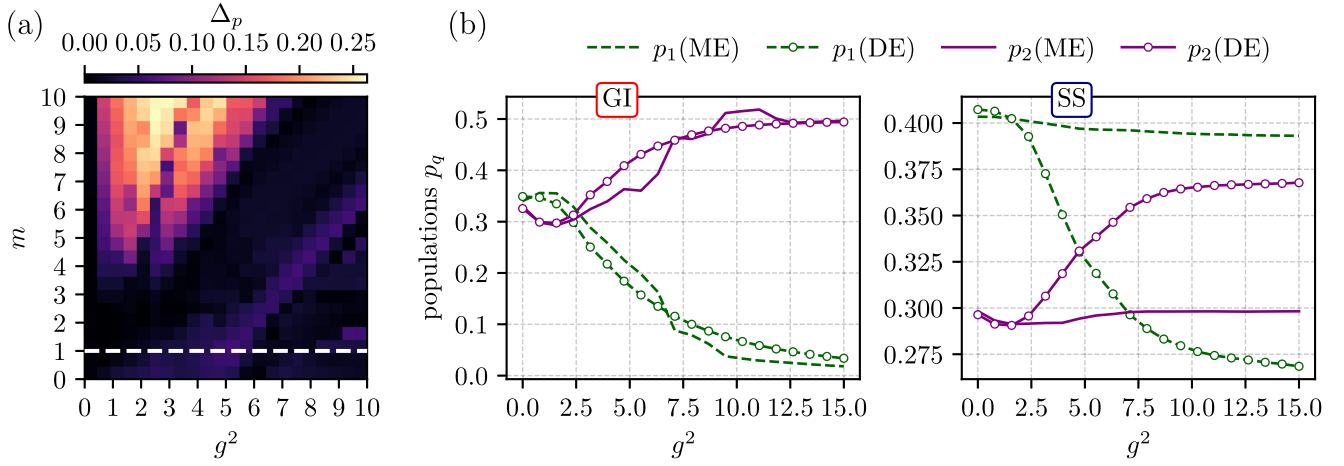


FIG. 3. Long-time dynamics of quark occupancy. (a) Mass-gauge coupling grid of the deviation Δp of the DE and ME predictions, as defined in Eq. (9), of the single p_1 and double p_2 quark populations for the case of the gauge-invariant (GI) initial state. (b) Average quark populations obtained from ME and DE predictions for the GI and SS states at fixed mass $m = 1$ and gauge coupling strength $g^2 \in [0, 15]$ [along the white dashed line in panel (a)].

$$\hat{Q}_n = (-1)^n \sum_{\alpha} \left[\hat{\psi}_{n,\alpha}^\dagger \hat{\psi}_{n,\alpha} - \frac{1 - (-1)^n}{2} \right], \quad (7)$$

which, in the staggered fermion description [75,136,138], counts the number of quarks (on even sites) or antiquarks (on odd sites). Then, we compute the average quark population $p_q(t)$ of each $\hat{Q}_n$ eigenvalue (single-site quark occupancy, $q \in \{0, 1, 2\}$), projecting on the corresponding eigensubspace of $\hat{Q}_n$ and averaging over all sites n :

$$p_q(t) = \frac{1}{N} \sum_n \langle \Psi(t) | \delta_{\hat{Q}_n, q} | \Psi(t) \rangle, \quad (8)$$

where the Kronecker delta $\delta_{\hat{Q}_n, q}$ selects a specific occupancy. In Fig. 3(a), we quantify the deviation between the DE and the microcanonical ensemble (ME) predictions for the quark populations of the GI state Eq. (4). The ME value, which reflects the thermal prediction, is obtained by uniformly averaging $\langle \Phi_s | \delta_{\hat{Q}_n, q} | \Phi_s \rangle$ over eigenstates $|\Phi_s\rangle$ in a narrow energy window around the initial state energy (see SM [134]). In particular, we focus on the average deviation of the single p_1 and double p_2 quark occupancies,

$$\Delta_p \equiv \frac{|p_1^{\text{DE}} - p_1^{\text{ME}}| + |p_2^{\text{DE}} - p_2^{\text{ME}}|}{2}, \quad (9)$$

as a function of the mass and gauge coupling. In the regime of $m \gg 1$, $g^2 \gtrsim 1$, there is a significant region where the ME and DE predictions deviate, indicating the absence of thermalization in the GI sector. In SM [134], we analyze this regime in detail and attribute the observed nonergodic behavior to HSF [69,94], which is known to be prevalent in LGTs [118,141]. Since this region is inherently nonthermal in the GI sector due to the Hamiltonian itself being nonergodic

in this regime, it should be excluded from the preliminary DFL phase diagram shown in Fig. 2(a) since, in our case, DFL is expected to possibly occur only in regimes where the Hamiltonian is ergodic at least in the GI sector. Doing so, we identify the different dynamical phases explored by the initial state Eq. (3) as the ones depicted in Fig. 1(c).

To gain deeper insight into the nature of excitations driving the DFL phase observed in Fig. 2, we compute the average quark occupancy on the lattice, distinguishing between the single p_1 and double p_2 occupancies defined in Eq. (8). In Fig. 3(b), we compare the ME and DE predictions for the GI and SS cases by fixing the mass at a value far from the fragmented phase and varying the gauge coupling. In agreement with our previous analysis, in the GI sector the DE and ME predictions closely follow each other, confirming the thermal behavior of the initial state in Eq. (4). Conversely, when the system is initialized in a superposition of SS in this regime, the ME and DE predictions diverge strongly due to the emergence of the previously discussed DFL phase. As the gauge coupling increases, the dominant contribution of the dynamics comes from the double-occupancy population p_2 , which corresponds to the formation of baryonlike (quark pair) and antibaryonlike (antiquark pair) excitations. This starkly contrasts the ME prediction, which instead anticipates a greater production of mesonlike excitations (adjacent quark-antiquark pairs with an excited shared gauge link) as indicated by the single-occupancy population p_1 . The crucial role of baryonic excitations is a distinctive feature of the non-Abelian SU(2) model, setting this phenomenology apart from its Abelian counterpart [97].

Entanglement entropy—A hallmark feature of DFL is the slow spreading of quantum correlations across the system, reminiscent of that in a many-body localized (MBL) phase [97,142–146]. Indeed, while in ergodic systems governed

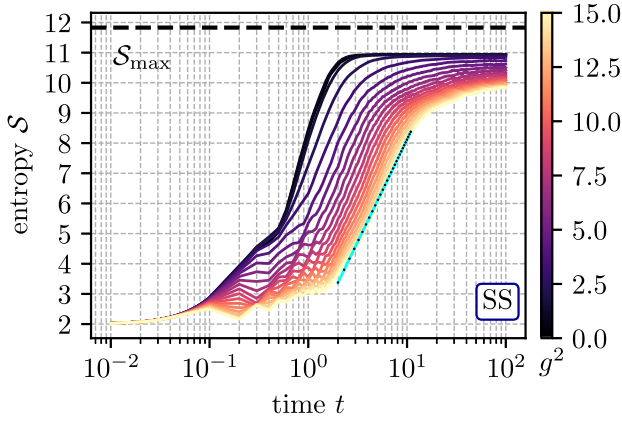


FIG. 4. Entanglement scaling. Half-chain bipartite entanglement entropy in base two as a function of time (in logarithmic scale) for fixed mass $m = 1$ and a range of gauge couplings $g^2 \in [0, 15]$ in the superposition of superselection sectors (SS) case. The cyan dashed line is the linear fit of the $g^2 = 15$ curve and corresponds to $S = 2.94(2) \cdot \log(t) + 1.33(3)$. The black dashed line represents the maximal bipartite entanglement entropy value $S_{\max}$ corresponding to an equal distribution over all accessible system configurations. Notice that the SS initial state given in Eq. (3) under periodic boundary conditions is not a product state, as in the Abelian case [97], rather, is an exact matrix product state with bond dimension $\chi = 4$ [134]; thus the entanglement entropy at the initial time starts from $S = \log \chi$.

by the ETH, quantum correlations spread linearly with time, in MBL phases, entanglement entropy grows logarithmically with time. This phenomenon can be quantified by measuring the bipartite von Neumann entanglement entropy, defined as $S = -\text{Tr}[\hat{\rho}_A \log \hat{\rho}_A]$, where A and B denote the two halves of the chain described by the density operator $\hat{\rho}_{AB}$, and $\hat{\rho}_A = \text{Tr}_B[\hat{\rho}_{AB}]$ is the reduced density operator of subsystem A . In Fig. 4, we show the entanglement entropy scaling for the SS initial state across a range of gauge couplings $g^2 \in [0, 15]$ at a fixed mass $m = 1$. Aside from late-time saturation effects due to the finite system size, the SS case exhibits a clear linear entanglement growth in $\log(t)$ across almost all the coupling values, compatible with an MBL phase and similar to what was observed in the Abelian scenario [97]. In detail, stronger gauge couplings limit the spread of entanglement, leading to lower saturation values. Indeed, as g^2 increases, gauge-link excitations are strongly penalized, and hopping processes are thereby suppressed. This constrained dynamics affects the spread of quantum correlations in each superselection sector, including the GI case, where the entanglement entropy scales linearly in time (see SM [134]), which is compatible with ergodicity.

Summary and outlook—We have investigated the dynamical phases of a $(1 + 1)\text{D}$ truncated $\text{SU}(2)$ LGT with dynamical matter. To explore the dynamics across different gauge superselection sectors, we developed a formalism based on static $\text{SU}(2)$ background charges, enabling the

preparation of superpositions across distinct sectors. Within this framework, we identified three different regimes: in addition to a thermal ergodic phase, we found a nonthermal yet delocalized fragmented regime and a DFL phase that emerges when superpositions of gauge superselection sectors are considered. In the latter, an initial imbalance in the matter distribution persists over time, and the entanglement entropy grows logarithmically, reminiscent of MBL. The dynamics in these regimes are strongly shaped by non-Abelian excitations such as baryons. In the DFL regime, characterized by large gauge coupling and mass, the gauge-field and background-charge truncations are well controlled. We expect the qualitative features of the dynamics to remain robust under larger truncation cutoffs as supported by recent results on QMB scars in the same model [75].

Several immediate directions follow from our Letter; these include the stability of the DFL phase in $(2 + 1)\text{D}$ and under different non-Abelian gauge groups. Another interesting avenue is to analyze the robustness of DFL to coherent error terms that explicitly break the non-Abelian gauge symmetry, which are commonly encountered in quantum simulators. The effect of decoherence on this phase is also a critical direction worth investigating.

Finally, we point out that a qudit formulation with a 13-dimensional local (single-site) Hilbert space faithfully captures the DFL phenomenology described in this Letter. This makes the model particularly suitable for digital quantum simulation on qudit-based quantum processors, as recently proposed for the same model in the absence of background charges [136]. Platforms supporting such large local dimensions are already experimentally accessible, as demonstrated by recent implementations using 13-level trapped-ion qudits based on $^{137}\text{Ba}^+$ [147], as well as other qudit architectures employing metastable states of ion isotopes [148] and circular Rydberg atoms [149,150].

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Data availability—The data that support the findings of this article are openly available [152].

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