

Real-Time String Dynamics in a 2 + 1D Non-Abelian Lattice Gauge Theory: String Breaking, Glueball Formation, Baryon Blockade, and Tension Reduction

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Understanding flux string dynamics can provide insight into quark confinement and hadronization. First-principles quantum and numerical simulations have mostly focused on toy-model Abelian lattice gauge theories (LGTs). With the advent of state-of-the-art quantum simulation experiments, it is timely to study string dynamics in non-Abelian LGTs beyond one spatial dimension. Using tensor network methods, we simulate the real-time string dynamics of a 2 + 1D SU(2) Yang-Mills LGT with dynamical matter. In the strong-coupling regime and at resonance, string breaking occurs via a sharp Casimir reduction, along with meson and baryon-antibaryon formation, a distinctively non-Abelian feature. At finite baryon density, we discover a *baryon blockade* mechanism that delays string breaking. Away from resonance, the magnetic term drives purely non-Abelian fluctuations: glueball-like excitations and self-crossed strings that resolve two SU(2) intertwiners with distinct dynamics. For higher-energy strings, we uncover representation-dependent tension-reduction resonances. Our findings serve as a guide for upcoming quantum simulators of non-Abelian LGTs.

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Introduction—In non-Abelian gauge theories such as quantum chromodynamics (QCD), confinement manifests as color-electric flux strings between static charges; at large separations, these strings can break by producing dynamical matter [1–4], thereby flattening the heavy-quark potential and directly connecting to hadronization [5–7]. Such phenomena lie beyond perturbation theory and inspired the development of lattice gauge theories (LGTs) [8,9] in both Lagrangian [10,11] and Hamiltonian formulations [12]. Although Euclidean path integral Monte Carlo (MC) methods for LGTs [13–16] have delivered benchmark results on confinement [17–19], glueball spectra, and k -string tensions [20–22] and static potential [23], exploring finite-density regimes and real-time dynamics has presented significant challenges. There, MC methods suffer from sign problems that obstruct importance sampling and analytic continuation [24–26].

These challenges have motivated alternative Hamiltonian approaches to LGTs that underlie the current major effort to quantum-simulate high-energy physics [27–43]. In recent years, LGT quantum simulators have rapidly progressed from digital [44–46] and analog realizations of Abelian gauge invariance [47–55] to non-Abelian SU(2)/SU(3) hardware implementations [56–68] and scalable variational state preparation [69–72], including real-time observations of string breaking [73–78] and hadron scattering [79–84]. In parallel to this body of experimental work, complementary and benchmarking studies of Hamiltonian-formulated LGTs have been performed using classical variational methods such as tensor networks (TNs) [85–90]. In the last two decades, they have provided fundamental contributions to static properties [91–99] and real-time dynamics of Abelian [100–112] and non-Abelian [113–117] LGTs.

In particular, recent TN simulations have uncovered new string dynamics phenomenology in 2 + 1D Abelian LGTs [118], from the roughening transition [119,120] to string breaking [111,112,121], accompanied by various quantum simulation experiments in 1 + 1D [76–78] and 2 + 1D [73–75]. These results provide a concrete blueprint for nonequilibrium higher-dimensional LGT dynamics and motivate exploring this physics in non-Abelian models, where real-time string dynamics with dynamical matter,

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especially at finite density, has so far only been investigated in one spatial dimension [113,122], and where quantum simulation proposals are emerging [60,123–127].

Our current work advances this frontier. Using tree tensor networks (TTNs), we simulate real-time string dynamics of a $2 + 1\text{D}$ $\text{SU}(2)$ Yang-Mills LGT with dynamical matter at system sizes and timescales beyond current quantum-computing capabilities, while retaining *ab initio* gauge invariance (also with background charges) and controlled access to finite-density sectors. We uncover purely non-Abelian signatures in two spatial dimensions, including baryon-antibaryon production at the breaking point, slowed-down breaking in finite-density regimes, and intertwiner-resolved self-crossed strings away from breaking resonances. In addition, for strings with nonminimal energy, we reveal representation-dependent tension-reduction phenomena. These results connect real-time $\text{SU}(2)$ string dynamics in $2 + 1\text{D}$ to qualitative features relevant to $3 + 1\text{D}$ QCD. They further provide concrete targets for near-term quantum simulators and next-generation TN methods for non-Abelian LGTs beyond one spatial dimension.

Model—To detect non-Abelian features of string breaking, we consider the $\text{SU}(2)$ Yang-Mills LGT with flavorless dynamical matter [98,115–117] on a $2 + 1\text{D}$ square lattice Λ containing N sites with indices $\mathbf{n} = (n_x, n_y)$ and lattice spacing a_0 . Correspondingly, the lattice link connecting two neighboring sites $\mathbf{n}$ and $\mathbf{n} + \boldsymbol{\mu}_k$ along the direction $k \in \{x, y\}$ is denoted by $(\mathbf{n}, \boldsymbol{\mu}_k)$. This model is described by the following Kogut-Susskind Hamiltonian [12] with *staggered fermions* [128]:

$$\begin{aligned} \hat{H} = & -w \sum_{\alpha, \mathbf{n}, \beta} \left[i \hat{\psi}_{\alpha, \mathbf{n}}^\dagger \hat{U}_{\mathbf{n}, \mathbf{n} + \boldsymbol{\mu}_x}^{\alpha\beta} \hat{\psi}_{\mathbf{n} + \boldsymbol{\mu}_x, \beta} \right. \\ & + (-1)^{n_x + n_y} \hat{\psi}_{\alpha, \mathbf{n}}^\dagger \hat{U}_{\mathbf{n}, \mathbf{n} + \boldsymbol{\mu}_y}^{\alpha\beta} \hat{\psi}_{\mathbf{n} + \boldsymbol{\mu}_y, \beta} + \text{H.c.} \left. \right] \\ & + m \sum_{\alpha, \mathbf{n}} (-1)^{n_x + n_y} \hat{\psi}_{\alpha, \mathbf{n}}^\dagger \hat{\psi}_{\alpha, \mathbf{n}} \\ & + g_E \sum_{\mathbf{n}, k} \hat{E}_{\mathbf{n}, \mathbf{n} + \boldsymbol{\mu}_k}^2 - g_B \sum_{\square} \left(\hat{U}_{\square} + \hat{U}_{\square}^\dagger \right). \end{aligned} \quad (1)$$

The first term describes the staggered hopping interaction, with $w = (1/2a_0)$, between quark matter fields living on neighboring sites and $\text{SU}(2)$ gauge fields on the links in between. The quark field is represented as a staggered fermion $\hat{\psi}_{\alpha, \mathbf{n}}$, which satisfies the canonical anticommutation relations $\{\hat{\psi}_{\alpha, \mathbf{n}}, \hat{\psi}_{\mathbf{n}', \beta}^\dagger\} = \delta_{\mathbf{n}, \mathbf{n}'} \delta_{\alpha, \beta}$ and $\{\hat{\psi}_{\alpha, \mathbf{n}}, \hat{\psi}_{\mathbf{n}', \beta}\} = 0$ [128], where the α, β indices live in the fundamental $\text{SU}(2)$ irreducible representation (irrep). The local matter basis is given by $\{|\Omega\rangle, \hat{\psi}_r^\dagger|\Omega\rangle, \hat{\psi}_g^\dagger|\Omega\rangle, \hat{\psi}_r^\dagger\hat{\psi}_g^\dagger|\Omega\rangle\}$, where $|\Omega\rangle$ represents an empty matter site, while $\{r, g\}$ are shorthand notations for $\{\pm \frac{1}{2}\}$, and j is implicit [98,115–117]. Matter fields have a corresponding staggered mass-energy term

with coupling m . The last two terms of Eq. (1) represent the (chromo)electric and magnetic gauge-field energies coupled with $g_E = (g_0^2/2a_0)$ and $g_B = (1/2g_0^2a_0)$, respectively, where g_0 is the gauge coupling.

Here, we express gauge link states in 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 . On this basis, the link energy density operator is diagonal and coincides with the quadratic Casimir, $\hat{E}^2|j, m_L, m_R\rangle = j(j+1)|j, m_L, m_R\rangle$ [129]. Correspondingly, the magnetic energy density is expressed in terms of a four-body plaquette operator $\hat{U}_{\square} = \sum_{\alpha, \beta, \gamma, \delta} \hat{U}_{\mathbf{n}, \mathbf{n}_x}^{\alpha\beta} \hat{U}_{\mathbf{n} + \boldsymbol{\mu}_x, \mathbf{n}_y}^{\beta\gamma} \hat{U}_{\mathbf{n} + \boldsymbol{\mu}_y, \mathbf{n}_x}^{\gamma\delta} \hat{U}_{\mathbf{n}, \mathbf{n}_y}^{\delta\alpha}$. In this Letter, we truncate the infinite-dimensional gauge fields [98] to $j \leq j_{\max}$ with $j_{\max} \in \{\frac{1}{2}, 1\}$, where $j_{\max} = \frac{1}{2}$ is the *hardcore-gluon* approximation. This truncation is controlled in the *strong-coupling regime* considered throughout this Letter, $g_E, m \gg g_B, w$: a link carrying irrep j costs an energy $g_E j(j+1)$, so excursions to $j > j_{\max}$ are off resonant and parametrically suppressed under the unitary quench dynamics, as verified in SM [130] (Fig. S5). For each setup, $j_{\max}$ is the smallest truncation containing the initial string and the resonant final manifold: $j_{\max} = \frac{1}{2}$ for Figs. 2 to 4 and $j_{\max} = 1$ for Fig. 5.

Non-Abelian $\text{SU}(2)$ gauge invariance is then locally imposed with the generators of local $\text{SU}(2)$ gauge transformations (Gauss-law operators) $\hat{\mathcal{G}}_{\mathbf{n}} = (\hat{\mathcal{G}}_{\mathbf{n}}^x, \hat{\mathcal{G}}_{\mathbf{n}}^y, \hat{\mathcal{G}}_{\mathbf{n}}^z)$ satisfying $\hat{\mathcal{G}}_{\mathbf{n}}|\Psi_{\{b_{\mathbf{n}}\}}\rangle = b_{\mathbf{n}}|\Psi_{\{b_{\mathbf{n}}\}}\rangle$, $\forall \mathbf{n} \in \Lambda$, where $b_{\mathbf{n}}$ is the background charge at lattice site $\mathbf{n}$. Using a dressed-site approach [98,115–117], we deformionize [131] the original LGT and build a local Hilbert space that retains only those states that satisfy Gauss’s law. In the *hardcore-gluon* approximation (i.e., $j_{\max} = j_{\text{bg}} = \frac{1}{2}$), we have a 3 0-dimensional local basis for each site with zero background charge $|j_{\text{bg}} = 0, m_{\text{bg}} = 0\rangle$, and with a 42-dimensional local basis on each of the two lattice sites where we place the static background charges $\{|j_{\text{bg}} = \frac{1}{2}, m_{\text{bg}} = \pm \frac{1}{2}\rangle\}$. When $j_{\max} = j_{\text{bg}} = 1$, we have a 317-dimensional basis for each site with a static background charge and a 168-dimensional basis for each site with zero background charge; see Supplemental Material (SM) [130].

String initialization—To study the real-time string dynamics of the model, we focus on the strong-coupling regime of Eq. (1), $g_E, m \gg g_B, w$, where the electric term dominates and spontaneous matter creation and plaquette dynamics are suppressed. On top of a *bare vacuum* configuration with empty even and doubly occupied odd sites [116], we set the initial string state $|\psi_0\rangle$ placing two static $\text{SU}(2)$ background charges $j_{\text{bg}} = \frac{1}{2}$ within a patch centered in the lattice bulk and connecting them by a path of active $j = \frac{1}{2}$ links (see Fig. 1 for a sketch).

Then, to detect string breaking, we measure the overlap $\mathcal{P}_k = |\langle \psi_k | \psi(t) \rangle|^2$ of the time-evolved wave function $|\psi(t)\rangle = e^{-i\hat{H}t}|\psi_0\rangle$ with any flux configuration $|\psi_k\rangle$ of

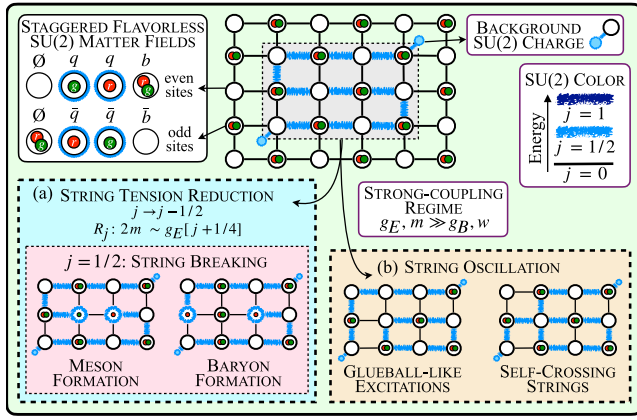


FIG. 1. Real-time string dynamics in a 2 + 1D SU(2) LGT. On a square lattice with staggered SU(2) dynamical matter fields on sites and SU(2) gauge fields on links, a generic string is created by connecting two static background SU(2) charges within a (gray) patch. In the strong-coupling regime, the string undergoes two scenarios: (a) string breaking (and, more generally, tension reduction), where string links downgrade their Casimir $j \rightarrow j - \frac{1}{2}$, allowing for meson and baryon creation; (b) string oscillations, where magnetic effects allow for glueball-like excitations and self-crossing strings. This scheme is for illustration purposes only; throughout this Letter, we consider a 4×4 string patch on an 8×8 lattice.

the same length as the string of the initial state $|\psi_0\rangle$. Correspondingly, we track the expectation value of local observables $\langle \hat{O} \rangle = \langle \psi(t) | \hat{O} | \psi(t) \rangle$ at any time t , properly normalized by the number of lattice sites (N_{patch}) or links ($\mathcal{L}_{\text{patch}}$) in the patch of the string. In particular, we focus on the quadratic Casimir operator and particle densities, such as mesons and (anti)baryons:

$$C = \frac{1}{\mathcal{L}_{\text{patch}}} \sum_{\mathbf{n}, k} \langle \hat{E}_{\mathbf{n}, +\mu_k}^2 \rangle, \quad p_{\text{mes}} = \frac{1}{N_{\text{patch}}} \sum_{\mathbf{n}} \langle \hat{p}_{\mathbf{n}}^{[1]} \rangle, \quad (2a)$$

$$p_{\text{bar}} = \frac{1}{2N_{\text{patch}}} \sum_{\mathbf{n}} \left[1 + (-1)^{n_x + n_y} \left(\langle \hat{p}_{\mathbf{n}}^{[2]} \rangle - \langle \hat{p}_{\mathbf{n}}^{[0]} \rangle \right) \right], \quad (2b)$$

which are defined by means of the local number occupation operators: (zero) $\langle \hat{p}_{\mathbf{n}}^{[0]} \rangle = 1 - \langle \hat{p}_{\mathbf{n}}^{[1]} \rangle - \langle \hat{p}_{\mathbf{n}}^{[2]} \rangle$, (single) $\langle \hat{p}_{\mathbf{n}}^{[1]} \rangle = \sum_{\alpha} \langle \hat{\psi}_{\alpha, \mathbf{n}}^\dagger \hat{\psi}_{\alpha, \mathbf{n}} \rangle - 2\langle \hat{p}_{\mathbf{n}}^{[2]} \rangle$, and (double) $\langle \hat{p}_{\mathbf{n}}^{[2]} \rangle = \langle \hat{\psi}_{\mathbf{n}, r}^\dagger \hat{\psi}_{\mathbf{n}, r} \hat{\psi}_{\mathbf{n}, g}^\dagger \hat{\psi}_{\mathbf{n}, g} \rangle$. Ultimately, we measure the bipartite entanglement entropy $\mathcal{S} = -\text{Tr}[\hat{\rho}_A \log \hat{\rho}_A]$, where $\hat{\rho}_A$ is the reduced density operator of half the system. Throughout this Letter, all results are obtained from TTN simulations of the time evolution under Eq. (1) on an 8×8 lattice and a 4×4 patch, with couplings $g_B = w = 1$.

String breaking—In the strong-coupling regime, the string energy is well approximated by the number of active links $\mathcal{L}_{\text{active}}$ times their electric cost:

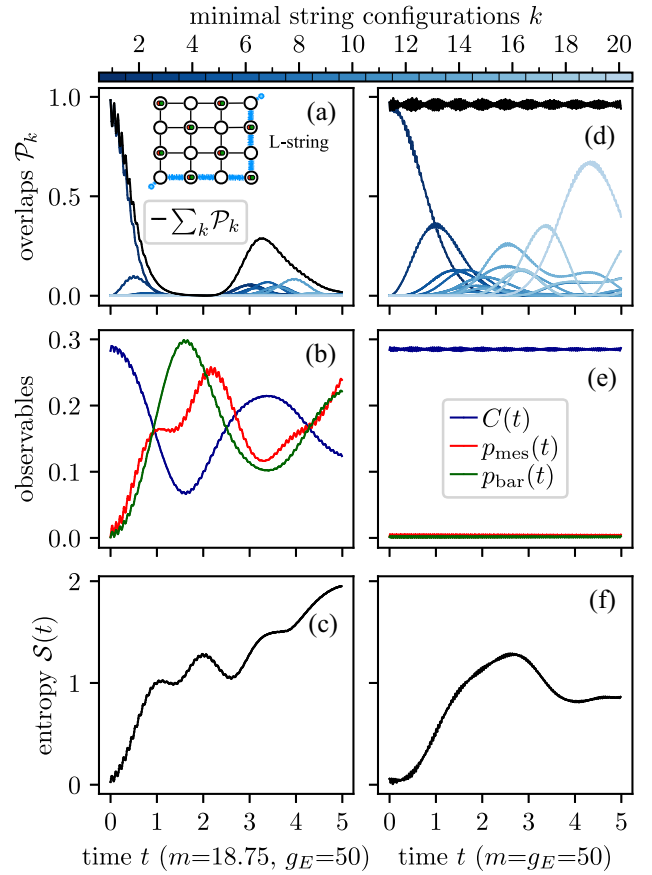


FIG. 2. String breaking vs oscillations. String dynamics starting from a minimal L string. *String breaking* at the resonance $m = \frac{3}{8}g_E$ manifest in the (a) decay of string-manifold overlaps over time, (b) decay of the Casimir and increase in the populations of mesons and baryons, defined in Eq. (2), and (c) a rapid continual growth in the entanglement entropy. Out of resonance (here $m = g_E = 50$, detuned from every R_j resonance), persistent *string oscillations* appear where (d) the probability of being in a string is always close to unity, (e) matter creation is suppressed, and (f) entanglement entropy saturates at later times.

$$E(t=0) = \langle \psi_0 | \hat{H} | \psi_0 \rangle \sim j(j+1)g_E \mathcal{L}_{\text{active}}. \quad (3)$$

Under unitary time evolution, (first-order) string breaking for $j = \frac{1}{2}$ is enabled when the energy $j(j+1)g_E = \frac{3}{4}g_E$ of removing one active link matches the energy $2m$ of creating a particle-hole pair at the sites connecting it. At this resonance point $R_{j=\frac{1}{2}}: m = \frac{3}{8}g_E$, any string configuration of length equal to that of the initial string resonantly hybridizes with shorter strings carrying dynamical matter: the Casimir along the path decreases, the total *string-manifold* weight $\mathcal{P} = \sum_k \mathcal{P}_k$ (sum over degenerate minimal-string configurations) collapses, and matter production proliferates, leading to a rapid and continual growth in the bipartite entanglement entropy [see Figs. 2(a)–2(c) for a minimal initial L string].

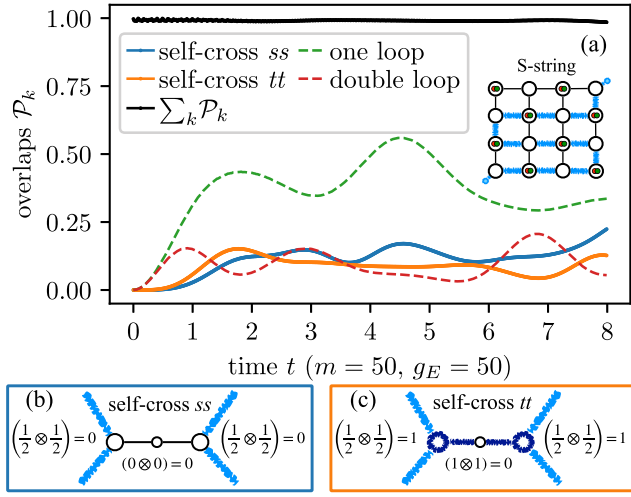


FIG. 3. Glueball-like excitations and self-crossing string formation. String oscillations starting from a maximal S string. Among all the degenerate string configurations $\mathcal{P}_k$ (whose sum stays close to unity) available out of resonance, we highlight in (a) the superpositions of type ss and type tt self-crossing strings and loop configurations (glueball-like excitations) of single and double loops, respectively. The lower panel shows how four $j = \frac{1}{2}$ (cyan) gauge links can fuse into a singlet using the SU(2) intertwiner scheme: the ss type (b), where each link pair intermediately fuses to $j = 0$ (black), and the tt type (c) where each link pair intermediately fuses to $j = 1$ (dark blue).

Unlike the Abelian case [111], matter production at resonance is not limited to mesons: SU(2) fusion rules also enable spontaneous *baryon-antibaryon* formation, visible in Fig. 2(b). Moreover, due to the use of staggered fermions, higher-order string breaking processes (i.e., multiple local breaks) can occur at the same resonance $R_{j=1}$. Videos of this dynamics are available on YouTube [132].

String oscillations: Glueball-like excitations and self-crossing strings—Out of resonance, i.e., for $m \neq \frac{3}{8}g_E$, the initial string cannot break; the plaquette term only reshuffles flux among configurations with the same number of active links. Accordingly, the patch-averaged Casimir remains constant without particle creation. At the same time, the entanglement entropy peaks near maximal superposition of the degenerate string configurations and minimizes when a single configuration dominates; see Figs. 2(d)–2(f).

For initial strings of nonminimal length, like the *snake* S string in Fig. 3, the magnetic term generates fluctuations into shorter strings, pure-gauge loops (which in the strong-coupling limit resemble *glueball-like* excitations [133]), and *self-crossed* strings. While pure-gauge loops and self-crossings also appear in Abelian $\mathbb{Z}_2$ LGTs [111], the SU(2) self-crossing case is richer because a four-valent vertex with four $j = \frac{1}{2}$ links supports two nonequivalent singlet intertwiners [134]. We label them type ss , where the two link pairs fuse as singlets $(\frac{1}{2} \otimes \frac{1}{2}) \rightarrow 0$ and then $(0 \otimes 0) \rightarrow 0$

[see Fig. 3(b)]; and type tt , where the pairs first fuse as triplets $(\frac{1}{2} \otimes \frac{1}{2}) \rightarrow 1$ and then $(1 \otimes 1) \rightarrow 0$ [see Fig. 3(c)]. The two intertwiners share the same external link content (four $j = \frac{1}{2}$ legs) but differ in their internal recoupling— ss bridges through $j = 0$, tt through $j = 1$ —so while they are energy-degenerate, they are coupled to different Clebsch-Gordan coefficients, so they differ in their plaquette and hopping matrix elements and are detuned relative to each other (out of resonance, as shown in Fig. 3(a), type- ss slightly dominates on average). Videos of this dynamics are available on YouTube [132].

Baryon blockade mechanism—Unlike MC simulations, where finite (baryon) density regimes are challenging because of the sign problem [26], here the baryon number is a quantum number associated with the global particle-hole U(1) symmetry of Eq. (1), and can be easily fixed in exact diagonalization (ED) [135] and TN simulations [90,136]. It is then of interest to investigate the role of finite baryon density in the string-breaking dynamics.

A good measure for the string breaking is the break time t_{break} corresponding to the minimum of the average Casimir operator. Indeed, as previously shown in Figs. 2(a)–2(c), the Casimir minimum occurs in parallel to the minimum of the string-manifold weight $\mathcal{P}$ and the peak of meson and baryon formation (see also Ref. [130]).

As shown in Fig. 4 for an initial S string, adding SU(2) baryons at the resonance $m = \frac{3}{8}g_E$ progressively delays t_{break} . We attribute this to a *baryon blockade* effect: finite density reduces the availability of final states for meson creation and prunes plaquette-mediated resonant paths, suppressing the string breaking mechanism.

String tension reduction—The *hardcore-gluon* truncation at $j_{\text{max}} = \frac{1}{2}$ and the minimal background charge $j_{\text{bg}} = \frac{1}{2}$ already captures pure non-Abelian string dynamics in the strong-coupling regime. Enabling larger gauge-field

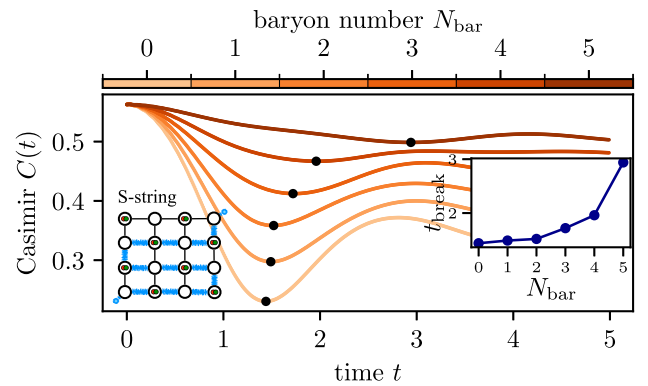


FIG. 4. Baryon blockade. Effect of the baryon number sector N_{bar} on the string-breaking mechanism in terms of the first observed minimum (t_{break}) in the time evolution of the Casimir operator. The inset panel shows the scaling of t_{break} as a function of the baryon number sector N_{bar} . We explore the sectors $N_{\text{bar}} = 0, \dots, 5$ supported by the 4×4 patch.

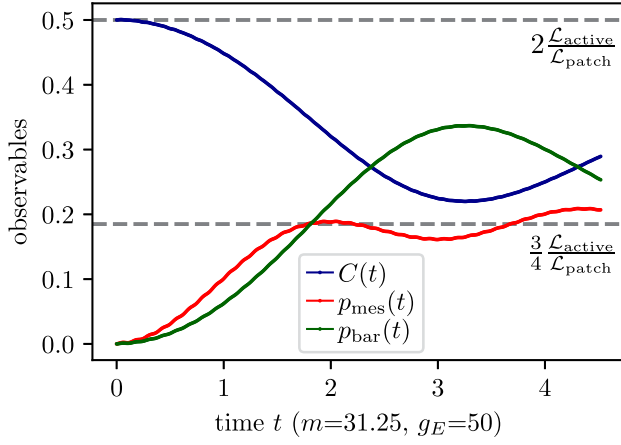


FIG. 5. String tension reduction mechanism. At the resonance $R_{j=1}$: $m = \frac{5}{8}g_E$, the $j = 1$ initial L string reduces its tension from $C = 2(\mathcal{L}_{\text{active}}/\mathcal{L}_{\text{patch}})$ (upper dashed line) toward $C \gtrsim \frac{3}{4}(\mathcal{L}_{\text{active}}/\mathcal{L}_{\text{patch}})$ (lower dashed line) through meson and baryon creation while not being broken.

truncations (and static charges) enriches the physics by opening *representation-changing* resonances where the string *reduces* its tension without fully breaking. Consider an initial string with $j = 1$, i.e., $j(j+1) = 2$, between two static charges with $j_{\text{bg}} = 1$, whose energy still follows Eq. (3). Here $j_{\text{bg}} = 1$ denotes an external triplet probe source fixing the flux representation at the endpoints of the simulated segment, rather than a static version of the fundamental staggered fermion. Then, a single particle-hole pair creation on neighboring sites along the path can reduce a link from $j = 1$ to $j = \frac{1}{2}$, lowering its chromoelectric energy by $\Delta E(1 \rightarrow \frac{1}{2}) = (2 - \frac{3}{4})g_E = \frac{5}{4}g_E$ at the cost of $2m$. Crucially, such a process only occurs at the *tension-reduction resonance* $R_{j=1}$: $m = \frac{5}{8}g_E$ and is not compatible with the subsequent breaking ($\frac{1}{2} \rightarrow 0$), whose energy cost $\Delta E(\frac{1}{2} \rightarrow 0) = \frac{3}{4}g_E$ does not match $2m = \frac{5}{4}g_E$. Hence, the string does not fully break; it relaxes at lower tension ($j = \frac{1}{2}$), as further shortening requires off-resonant (plaquette-assisted) or higher-order processes.

As shown in Fig. 5 for an initial L string at $R_{j=1}$, tension reduction manifests as a reduction of the patch-averaged Casimir [$C: 2(\mathcal{L}_{\text{active}}/\mathcal{L}_{\text{patch}}) \rightarrow \frac{3}{4}(\mathcal{L}_{\text{active}}/\mathcal{L}_{\text{patch}})$] accompanied by meson and baryon-antibaryon production. In general, for arbitrarily large- j strings, (first-order) string-tension reduction ($j \rightarrow j - \frac{1}{2}$) occurs for a family of resonances R_j : $2m = g_E(j + \frac{1}{4})$, which coincides with string breaking only when $j = \frac{1}{2}$.

Summary and outlook—We have performed the first TN real-time simulations of string dynamics in a $2+1\text{D}$ $\text{SU}(2)$ Yang-Mills LGT with dynamical matter. In a strong-coupling regime, we have identified a resonance where string states anticross with broken-string manifolds,

forming mesons and baryons. We uncovered a many-body, non-Abelian effect—baryon-blockade string breaking—where finite baryon density delays string breaking at fixed resonance. Out of resonance, the magnetic term generates pure non-Abelian glueball-like excitations and self-crossed strings that resolve two $\text{SU}(2)$ intertwiners at a four-valent vertex, with type ss favored by electric-energy detuning and plaquette flips. Finally, at higher gauge field truncations, we observe representation-dependent *tension reduction* ($j \rightarrow j - \frac{1}{2}$) resonances.

Together, these results provide *ab initio*, real-time, many-body signatures of confinement dynamics in two spatial dimensions, beyond the reach of MC methods. It is important to note that the evolution times and system sizes we access in our TN simulations go significantly beyond the state-of-the-art quantum simulation experiments on string dynamics in $2+1\text{D}$ LGTs [73,74,81]. Furthermore, several features of the phenomenology uncovered in this work are distinctively non-Abelian, which not only goes beyond the accessible physics of the aforementioned experiments, but also beyond previous TN works since we probe far-from-equilibrium non-Abelian string dynamics in $2+1\text{D}$. Consequently, our results push the state of the art in TN simulations of higher-dimensional non-Abelian LGTs and delineate concrete targets for near-term quantum simulators of these phenomenologically rich models.

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Views and opinions expressed are, however, those of the author(s) only and do not necessarily reflect those of the European Union or the European Research Council Executive Agency. Neither the European Union nor the granting authority can be held responsible for them.

Data availability—The data and the scripts that reproduce all the figures of this Letter and of SM [130] are openly available on Zenodo [139].

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