Letter

Continuum limit of a qubit-regularized SU(3) lattice gauge theory on a plaquette chain

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We show that a simple qubit-regularized SU(3) lattice gauge theory on a plaquette chain admits a continuum limit with massive gauge-invariant excitations, providing a minimal toy model of strong interactions without quarks. By mapping the plaquette-chain Hamiltonian to the three-state quantum clock model in a magnetic field, we demonstrate that the theory can be tuned to a continuum limit governed at short distances by the $\mathbb{Z}_3$ parafermion conformal field theory, which serves as the UV fixed point. A small relevant magnetic perturbation then drives the system to a massive continuum quantum field theory in the IR, featuring relativistic particles. In this continuum theory, we compute two universal ratios: the mass ratio of the lightest massive particles with opposite charge conjugation, $m^-/m^+ = 1.459(2)$, and $\sqrt{\sigma}/m^+ = 0.2648(2)$, where σ denotes the string tension between a static quark and antiquark.

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Gluons mediate the strong force between quarks, binding them into hadrons such as protons and neutrons. Unlike photons, gluons also interact among themselves and, in the absence of quarks, can form massive bound states known as glueballs. Their dynamics are described by SU(3) Yang-Mills (YM) theory, a non-Abelian gauge theory in the continuum with eight gluons as fundamental degrees of freedom. Despite its conceptual simplicity, determining the glueball spectrum is notoriously difficult. The standard approach is to discretize the theory on a lattice, following Wilson, and compute the energy eigenvalues of the resulting Hamiltonian while taking the thermodynamic and continuum limits [1].

Recently, a new class of lattice gauge theories (LGTs) with finite-dimensional local Hilbert spaces has been proposed using the monomer-dimer tensor-network (MDTN) basis of conventional LGTs [2]. We refer to such constructions as qubit regularizations [3], since they can be encoded in qubits and potentially solved on a quantum computer [4]. In the philosophy of qubit regularization, a gauge theory is constructed from microscopic lattice quantum degrees of freedom that transform under local gauge transformations. The dynamics are then restricted to the gauge-invariant physical Hilbert space. This construction makes it clear that gauge symmetry is not really a symmetry of the theory, but rather a constraint on the local degrees of freedom in a given basis. Indeed, examples abound in highly supersymmetric theories in which the same theory can arise from distinct descriptions using different gauge groups, a property that is probably related to deeper mathematical structures [5]. From this perspective, the gauge group itself is not a fundamental characteristic of the physical theory. Rather, gauge theories are defined using a set of quantum degrees of freedom associated with connections between different points in space. These degrees of freedom transform nontrivially under a gauge group, whereas the physical content resides in the gauge-invariant sector.

It is well known that LGTs are not uniquely determined by their gauge symmetry alone [6]. Qubit regularization highlights a previously underappreciated source of this nonuniqueness, beyond the lattice geometry and the Hamiltonian couplings: the choice of representation of the gauge group carried by the local lattice fields. Since every second-order critical point defines a continuum quantum field theory, qubit-regularized gauge theories suggest the existence of a richer landscape of possible continuum limits.

Because the space of consistent continuum quantum field theories is expected to be smaller than the space of LGTs, it is natural to expect that distinct lattice formulations may flow to the same continuum theory. One exciting possibility is that a sufficiently simple qubit-regularized theory could provide an alternative route to studying $(3+1)$-dimensional YM theory. Equally plausible, however, is the possibility that qubit-regularized LGTs do not generically admit nontrivial continuum limits. Even when

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they do, the resulting continuum theories may belong to universality classes distinct from that of conventional continuum YM theory. Our work represents a small step toward understanding this broader landscape.

The first step in determining whether a qubit-regularized LGT yields a nontrivial continuum quantum field theory (QFT), potentially in the universality class of the conventional YM theory, is to establish the existence of a relativistic quantum critical point. Properties such as asymptotic freedom can arise at such quantum critical points through standard renormalization group (RG) flows or through more exotic RG flows, as discovered recently [7,8]. Assuming such a quantum critical point exists in a pure qubit-regularized LGT without matter fields in $3+1$ dimensions, a relevant perturbation would naturally yield massive relativistic particles in the continuum limit. Identifying these immediately as the glueballs of conventional YM theory would be premature without detailed quantitative comparisons. Nevertheless, such relativistic massive excitations can be defined operationally as "glueball-like" particles: they are massive gauge-invariant particle excitations that emerge near a relativistic continuum limit of a microscopic pure non-Abelian LGT without matter fields. This terminology emphasizes that the excitations are generated purely by gauge fields of the microscopic lattice theory. We emphasize that glueball-like states are not necessarily identifiable with glueballs of conventional YM theory.

Since studying $3+1$ dimensions is difficult, in the present work we focus on constructing an example of a continuum limit in a qubit-regularized SU(3) LGT in a tractable setting. The model we consider is defined on a plaquette chain. We show that it possesses a UV fixed point and use it to define a massive relativistic continuum limit with glueball-like particles as defined above. We compute the ratio of the lowest charge-conjugation-odd and charge-conjugation-even masses and, in the presence of static matter fields, determine a dimensionless ratio involving the string tension.

A similar massive continuum limit was recently identified in a qubit-regularized SU(2) LGT [9]. The present work shows that analogous ideas extend to an SU(3)-motivated construction with the corresponding changes in global symmetry. These results motivate the conjecture that continuum limits of qubit-regularized gauge theories in higher dimensions may emerge as quantum critical points separating confined and deconfined phases.

Although our model is defined on a ladder geometry, its continuum limit should not be identified with pure YM theory in one spatial dimension, which has no local propagating degrees of freedom. By contrast, the plaquette-chain geometry retains a finite transverse structure consisting of two longitudinal chains connected by rungs. The rung links carry information about the transverse gauge degrees of freedom, which can also be interpreted as adjoint scalars. Gauge-invariant operators involving both the rung and longitudinal links can therefore create local excitations that propagate along the longitudinal direction. Such excitations are present in both the conventional Kogut-Susskind LGT and its qubit-regularized counterpart; they are not artifacts of the qubit regularization. It is possible that both theories can be tuned to a continuum limit described by a YM theory coupled to adjoint scalar fields in one spatial dimension.

There is an interesting aspect of these adjoint matter degrees of freedom as additional chains are coupled and the transverse direction becomes more fully resolved. The microscopic transverse gauge degrees of freedom of the lattice theory, which are interpreted as adjoint scalars in the finite-width geometry, become the true transverse gauge degrees of freedom of the higher-dimensional theory. For example, in electromagnetism, these scalar-like degrees of freedom can be identified with the fundamental modes propagating along a thin transmission line, or along waveguides with appropriate boundary conditions, whose quanta are often referred to as one-dimensional photons [10]. In this sense, the adjoint matter fields can be viewed as remnants of the gauge bosons of the higher-dimensional theory. Since these lowest-energy modes remain unaffected by the smallness of the extra dimension, the connection of these adjoint scalar modes with the gauge bosons of the higher-dimensional theory can persist even when all the other modes are pushed to energies near the UV cutoff scale. Thus, in our calculations on the plaquette chain, the rung degrees of freedom retain remnants of the transverse gluons. Such remnants are also familiar in Kaluza-Klein settings. For example, five-dimensional YM theories compactified on S^1 describe a four-dimensional gauge theory coupled to a scalar field in the adjoint representation of the gauge group [11]. Bound states of such remnants in non-Abelian theories are then expected to capture rudimentary glueball-like physics of the higher-dimensional theory. This interpretation has been discussed previously in a different context [12].

The physical Hilbert space of the SU(3) plaquette chain is constructed in the MDTN basis [2], where gauge degrees of freedom reside on links connecting neighboring sites and are labeled by dimer tensor (equivalently, color-flux) states $|\lambda\rangle$ with $\lambda = \mathbf{1}, \mathbf{3}, \bar{\mathbf{3}}$, the irreducible representations of SU(3). A representative physical basis state is shown in Fig. 1.

The ladder geometry of the plaquette chain allows us to distinguish links along the chain and those along the rungs, which we denote by ℓ_c and ℓ_r, respectively. Elementary plaquettes P are defined as the squares formed by four such links. Owing to this geometry, the physical Hilbert space of the theory naturally splits into three topologically distinct sectors. The first sector, $\mathcal{H}_{\mathrm{I}}$, consists of basis states in which the top and bottom chain links of every plaquette P are either $|\mathbf{1},\mathbf{1}\rangle$, $|\mathbf{3},\bar{\mathbf{3}}\rangle$, or $|\bar{\mathbf{3}},\mathbf{3}\rangle$. The second sector, $\mathcal{H}_{\mathrm{II}}$, contains plaquettes in the states $|\bar{\mathbf{3}},\bar{\mathbf{3}}\rangle$, $|\mathbf{1},\mathbf{3}\rangle$, or $|\mathbf{3},\mathbf{1}\rangle$, while the third sector, $\mathcal{H}_{\mathrm{III}}$, contains plaquettes in the states

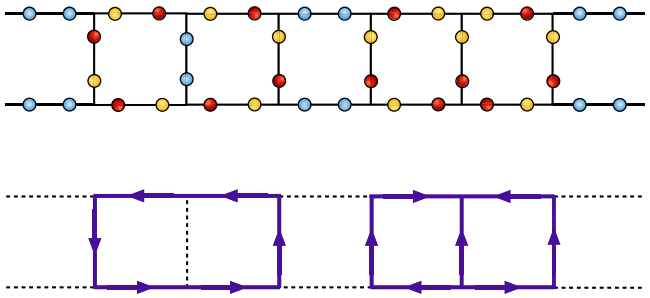


FIG. 1. Physical basis state of the qubit-regularized SU(3) LGT on a plaquette chain in the MDTN basis. Links are represented by dimer tensor states, which may equivalently be viewed as color flux states $|\lambda\rangle$, where $\lambda = \mathbf{1}$ (blue), $\mathbf{3}$ (red), and $\bar{\mathbf{3}}$ (yellow) are irreducible representations of SU(3). The two indices of each tensor are associated with the lattice sites connected by the link and transform in the irreps λ and its conjugate $\bar{\lambda}$. Gauge invariance enforces that the irreps meeting at each site combine to a singlet. For a trivalent site, the allowed combinations are (i) three singlet indices, (ii) a singlet, a triplet, and an antitriplet index, (iii) three triplet indices, or (iv) three antitriplet indices (top). An equivalent representation in terms of color fluxes is shown in the bottom panel, where a dashed line denotes the state $|\mathbf{1}\rangle$, arrows pointing to the right or upward denote $|\mathbf{3}\rangle$, and oppositely oriented arrows denote $|\bar{\mathbf{3}}\rangle$.

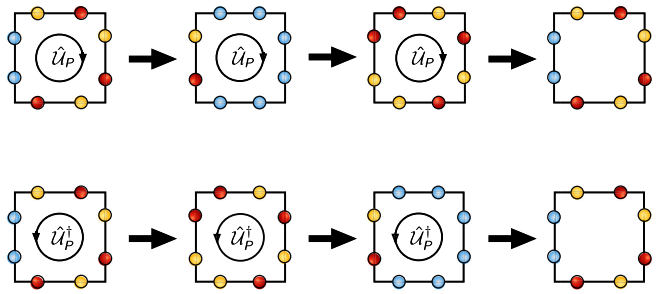


FIG. 2. Action of the plaquette operators $\hat{U}_P$ and $\hat{U}_P^\dagger$ on a representative MDTN state. Traversing the plaquette clockwise or counterclockwise updates the links according to the cyclic rule $\mathbf{1} \to \mathbf{3} \to \bar{\mathbf{3}} \to \mathbf{1}$, applied to the irrep at the first site of each link; the irrep at the opposite end is then fixed.

$|\mathbf{3},\mathbf{3}\rangle$, $|\mathbf{1},\bar{\mathbf{3}}\rangle$, or $|\bar{\mathbf{3}},\mathbf{1}\rangle$. Under charge conjugation $\hat{\mathcal{C}}$, $\mathbf{3}$ and $\bar{\mathbf{3}}$ are exchanged. Consequently, $\mathcal{H}_{\mathbf{1}}$ is invariant under charge conjugation, whereas the sectors $\mathcal{H}_{\mathrm{II}}$ and $\mathcal{H}_{\mathrm{III}}$ are exchanged into one another. Thus, charge conjugation is not a symmetry within either of the latter two sectors separately.

Plaquette chains have been used as a simplified class of quasi-one-dimensional lattice geometries for studying LGTs, particularly in work motivated by quantum computation [13–17]. Specializing to SU(3), the above simple Hilbert space has also been explored previously [18–20]. Most earlier work has focused on models based on the conventional Kogut-Susskind Hamiltonian [6], which typically suffer from sign problems. This makes it difficult to study continuum limits using Monte Carlo methods in higher dimensions. To overcome this difficulty, in Ref. [2] we proposed a class of simple gauge-invariant Hamiltonians for SU(N) LGTs that are free of sign problems. Here, we adapt that construction to the SU(3) plaquette chain, introducing additional couplings to demonstrate the existence of a nontrivial continuum limit. Our plaquette-chain Hamiltonian is given by

$$H_{pc} = \kappa_c \sum_{\ell_c} \hat{\mathcal{E}}_{\ell_c} + \kappa_r \sum_{\ell_r} \hat{\mathcal{E}}_{\ell_r} - g \sum_P (\hat{U}_P + \hat{U}_P^\dagger), \quad (1)$$

where $\hat{\mathcal{E}}_\ell$ is diagonal in the link basis $|\lambda\rangle$ associated with link ℓ, with eigenvalue $(1 - \delta_{\lambda,\mathbf{1}})$. It assigns equal energies to the $\mathbf{3}$ and $\bar{\mathbf{3}}$ irreps, which are typically higher than the energy assigned to the $\mathbf{1}$ irrep on each link, implying $\kappa_c, \kappa_r \geq 0$. In contrast, the plaquette operator $\hat{U}_P$ and its Hermitian conjugate are off diagonal in the MDTN basis and modify the irreps on the four links of a plaquette while preserving gauge invariance. Their action is illustrated in Fig. 2. By restricting to $g \geq 0$, we can avoid sign problems and study the model with Monte Carlo methods. Finally, we note that $[H_{pc}, \hat{\mathcal{C}}] = 0$, where $\hat{\mathcal{C}}$ is the charge-conjugation operator introduced earlier. Since the Hamiltonian is local and gauge invariant, it does not mix the three Hilbert spaces $\mathcal{H}_{\mathrm{I}}$, $\mathcal{H}_{\mathrm{II}}$, and $\mathcal{H}_{\mathrm{III}}$. We identify $\mathcal{H}_{\mathrm{I}}$ as the vacuum sector. Introducing heavy static quarks and antiquarks at two sites generates a string of plaquettes connecting them that necessarily lies in one of the other two sectors. Thus, while the Hamiltonian retains the same form as in Eq. (1), the presence or absence of static charges alters the accessible Hilbert space and consequently affects the energy spectrum.

Since the plaquettes admit only three quantum states within each Hilbert space sector, the SU(3) plaquette chain is equivalent, in each sector, to a one-dimensional three-state quantum clock model [20]. The latter is defined in terms of the operators

$$\hat{X} = \begin{pmatrix} 0 & 1 & 0 \\ 0 & 0 & 1 \\ 1 & 0 & 0 \end{pmatrix}, \quad \hat{Z} = \begin{pmatrix} 1 & 0 & 0 \\ 0 & e^{i2\pi/3} & 0 \\ 0 & 0 & e^{-i2\pi/3} \end{pmatrix}, \quad (2)$$

which act on the three states associated with each plaquette P. Restricting the Hamiltonian in Eq. (1) to the sector $\mathcal{H}_{\mathrm{I}}$ yields

$$H^{\mathrm{I}}_{\mathrm{qcm}} = -\sum_P \left\{ J \hat{Z}_P \hat{Z}^\dagger_{P+1} + g \hat{X}_P + \frac{h}{3} \hat{Z}_P + \mathrm{H.c.} \right\}, \quad (3)$$

with $J = \kappa_r/3$ and $h = 2\kappa_c$, up to an overall constant energy shift E_0. In the remaining sectors $\mathcal{H}_{\mathrm{II}}$ and $\mathcal{H}_{\mathrm{III}}$, the Hamiltonians $H^{\mathrm{II}}_{\mathrm{qcm}}$ and $H^{\mathrm{III}}_{\mathrm{qcm}}$ take the same form, except that h is replaced by $h = -\kappa_c$.

The quantum clock model defined in Eq. (3) is critical at $J = g = 1$ and $h = 0$, where its quantum critical behavior belongs to the universality class of the 2D three-state classical Potts model [21]. At this lattice critical point, the IR physics is described by a relativistic $\mathbb{Z}_3$ parafermion conformal field theory (CFT) with central charge $c = 4/5$ [22]. Via the mapping from Eq. (1) to Eq. (3), the same quantum critical point is realized in the qubit-regularized SU(3) plaquette chain at $\kappa_r = 3$, $g = 1$, and $\kappa_c = 0$, where we interpret the parafermionic degrees of freedom as describing the emergent continuum gauge degrees of freedom.[1] Relevant perturbations of the parafermion CFT, in the form of a thermal coupling $\tau > 0$ with $\tau = g - 1$ and a magnetic coupling $h > 0$, generate massive QFTs [23]. Once the critical theory is perturbed to define a massive continuum limit, the same CFT serves as the UV fixed point of the resulting continuum QFT. In this work, we focus on the case $\tau = 0$ and $h > 0$ (equivalently, $\kappa_c > 0$), which arises naturally in the vacuum sector of the plaquette chain.

Since the correlation length associated with the magnetic perturbation diverges as $\xi \sim h^{-15/28}$ for $h \to 0$ [23–25], universal quantities in a finite system of L plaquettes arranged on a circle are functions of the dimensionless variable $\mu = h^{15/28}L$. One particularly interesting universal quantity is the mass ratio m^-/m^+, where m^- and m^+ denote the masses of the lightest stable $\hat{\mathcal{C}}$-odd and $\hat{\mathcal{C}}$-even particles, respectively. This mass ratio can be extracted from the eigenvalues of Eq. (3) as a function of μ from the UV regime where $\mu \to 0$ to the IR regime where $\mu \to \infty$. We will be interested in the IR regime where the lattice model describes the massive QFT.

In this work, we use the density matrix renormalization group (DMRG) to determine m^-/m^+ as a function of μ. We compute the low-energy spectrum of Eq. (3) at $J = g = 1$ for various values of μ. For each μ, we consider five lattice sizes $L = 32, 48, 64, 80$, and 96, which enables a controlled extrapolation to the thermodynamic limit. In each case, h is tuned to realize the desired value of μ. We denote the ground-state energy at fixed (L, μ) by $\mathrm{E}_0^+(L, \mu)$, and the excited energies by $\mathrm{E}_i^C(L, \mu)$, where $C = \pm$ labels the charge-conjugation eigenvalue and $i = 1, 2, \ldots$ orders the states by increasing energy within each sector. Our DMRG computation was performed using the ITensor software library in Julia [26,27]. Further details of the calculation can be found in the Supplemental Material [28].

We extract the lowest particle masses in the charge-conjugation-even and -odd sectors as

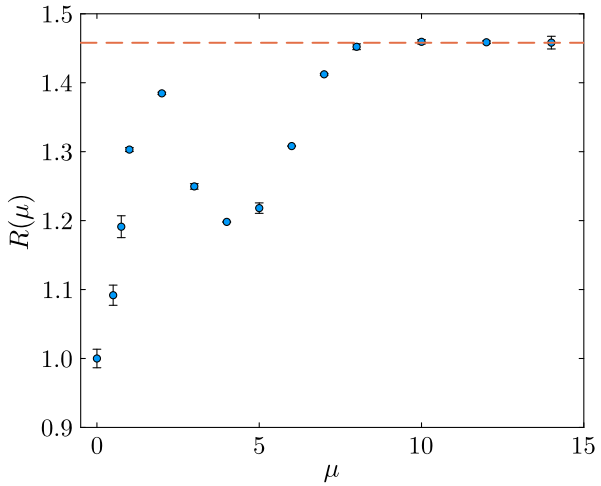


FIG. 3. Plot of the mass ratio $R(\mu)$ in Eq. (5) as a function of the scaling variable $\mu = h^{15/28}L$. The data points are obtained by extrapolating the DMRG data at each fixed value of μ to the joint limit $L \to \infty$ and $h \to 0$. The UV and IR regimes correspond to $\mu \to 0$ and $\mu \to \infty$, respectively. In the UV limit, the degeneracy of the lightest spin primaries related by charge conjugation at the $\mathbb{Z}_3$ parafermion CFT leads to the result that $R(0) = 1$ [23,31]. In the IR limit, no closed-form expression for $R(\infty)$ is known since the theory is nonintegrable.

$$m^C(L, \mu) = \mathrm{E}_1^C(L, \mu) - \mathrm{E}_0^+(L, \mu). \tag{4}$$

To obtain the dimensionless ratio m^-/m^+ in the large-L limit, we assume that for sufficiently large L,

$$\frac{m^-(L, \mu)}{m^+(L, \mu)} = R(\mu) - \frac{A(\mu)}{L}, \tag{5}$$

where $R(\mu)$ and $A(\mu)$ are μ-dependent coefficients. This ansatz provides a good fit at all values of μ [28]. The extracted $R(\mu)$ is shown in Fig. 3. We observe nonmonotonic behavior at intermediate μ, while for $\mu \geq 8$, the ratio begins to saturate toward its IR value. Although we do not have a simple explanation for this nonmonotonicity, similar behavior has been observed in conventional lattice calculations of glueball masses, where it is attributed to the presence of a critical point in the fundamental-adjoint coupling plane at which the scalar glueball mass vanishes [29]. The extracted coefficients in the large μ regime are listed in Table I, yielding the estimate

$$m^-/m^+ \approx 1.459(2). \tag{6}$$

This ratio was previously computed numerically using the truncated conformal space approach in Ref. [23], since the massive QFT emerging in the IR is nonintegrable [30]. Because the earlier results were presented graphically, extracting a precise numerical value is difficult; however, a rough estimate indicates consistency with our result.

[1]Pure Yang-Mills theory in one spatial dimension has no local propagating degrees of freedom. Our model has rungs connecting two chains, and allows transverse gauge degrees of freedom that are expected to be described as adjoint one-dimensional scalars. The parafermions in the UV presumably describe the physics of these bosons.

TABLE I. Fixed-μ scaling-limit estimates of $A(\mu)$ and $R(\mu)$ appearing in Eq. (5). The corresponding dimensionless string-tension ratio, $\sqrt{\sigma}/m^+$, is also shown. The extrapolations use $L = 32$, 48, 64, 80, and 96.

μ	$A(\mu)$	$R(\mu)$	$\sqrt{\sigma}/m^+$
8.0	1.1(1)	1.452(4)	0.2676(4)
10.0	1.3(1)	1.459(3)	0.2649(2)
12.0	1.57(8)	1.459(2)	0.2648(2)
14.0	1.8(3)	1.458(9)	0.264(1)

To confirm that the massive continuum QFT emerging in the IR is relativistic, we study the dispersion relation of the low-energy spectrum of Eq. (3). Since the system is invariant under discrete translations, its energy eigenstates with energies $\mathrm{E}_i^C(L,\mu)$ can be chosen to have definite lattice momenta $p_n = 2\pi n/L$. Single-particle excitations of distinct particle species below the two-particle threshold can be identified as nondegenerate states at zero momentum. Using this insight, we reorganize the low-energy spectrum $\mathrm{E}_i^C(L,\mu)$ by relabeling the index i as a pair (j,n). We assign the vacuum energy $\mathrm{E}_0^+(L,\mu)$ the label $(j=0,n=0)$. The higher-energy states (in both sectors) correspond to the jth massive particle with lattice momentum p_n. We denote the corresponding energies by $\mathrm{E}_{j,n}^C(L,\mu)$ [28].

We then define the single-particle excitation energies by subtracting the vacuum energy,

$$\epsilon_{j,n}^C(L,\mu) = \mathrm{E}_{j,n}^C(L,\mu) - \mathrm{E}_0^+(L,\mu). \tag{7}$$

The mass of the jth particle is then defined as $m_j^C = \epsilon_{j,0}^C(L,\mu)$. The dispersion relation

$$(\epsilon_{j,n}^C)^2 - (m_j^C)^2 = 2(1-\cos p_n)\zeta^2 + \mathcal{O}(n^4 L^{-4}), \tag{8}$$

where ζ is a constant (independent of n, L, and j) that plays the role of $\hbar c$ converting inverse lattice spacing into energy units, is a hallmark of a relativistic particle. As shown in Fig. 4, our data are well described by this dispersion relation, yielding $\zeta = 2.5907(2)$.

For sufficiently large L and μ, we can also compute the string tension between a static quark-antiquark pair by placing the quark at site $x = 0$ and the antiquark at $x = \mathsf{w}$, both on the lower chain. Although the Hamiltonian remains that of Eq. (1), the Hilbert space is modified: the plaquettes between $x = 0$ and $x = \mathsf{w}$ are forced into sector II, while the remaining plaquettes remain in the vacuum sector I. In the effective quantum clock model of Eq. (3), this corresponds to a modification of the coupling h: on the plaquettes between $x = 0$ and $x = \mathsf{w}$, h is reduced by a factor of 2 and changes sign relative to its value on the remaining plaquettes.

The heavy quark-antiquark potential is defined as

$$V(\mathsf{w}) = \mathrm{E}_0^+(\mathsf{w}) - \mathrm{E}_0^+, \tag{9}$$

where $\mathrm{E}_0^+(\mathsf{w})$ denotes the ground-state energy in the presence of a static quark-antiquark pair separated by w, and E_0^+ is the vacuum ground-state energy. In the confined phase we expect a linear potential,

$$V(\mathsf{w}) = \alpha + \bar{\sigma}\mathsf{w}, \tag{10}$$

where $\bar{\sigma}$ is the string tension (energy per unit length). Figure 5 illustrates how we extract $\bar{\sigma}$ from our data when $L = 96$ and $\mu = 14$.

Defining $\sigma = \bar{\sigma}\zeta$, which has dimensions of energy squared, and using the value of ζ extracted from the dispersion analysis, we compute the dimensionless ratio $\sqrt{\sigma}/m^+$ for each μ and L. Assuming the finite-size corrections are of the form L^{-1}, we extrapolate to the thermodynamic limit [28]. The resulting estimates of $\sqrt{\sigma}/m^+$ are listed in the fourth column of Table I, from which we obtain

$$\sqrt{\sigma}/m^+ = 0.2648(2).$$

Our work demonstrates that simple qubit-regularized SU(3) LGTs can exhibit nontrivial continuum limits with massive relativistic excitations. A similar massive continuum limit was recently identified in a qubit-regularized SU(2) theory [9]. To the best of our knowledge, the continuum limits of the conventional Kogut-Susskind

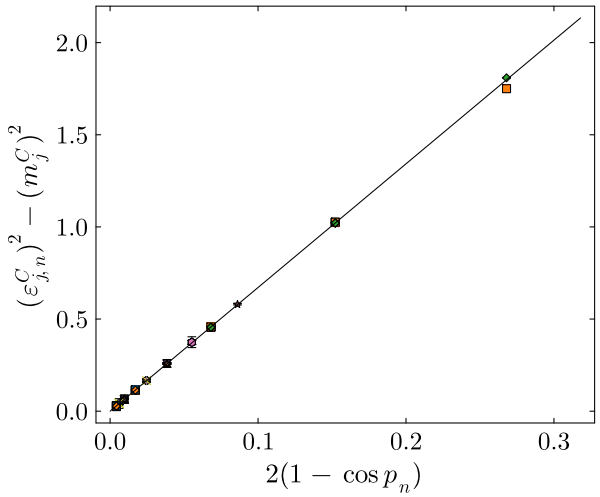


FIG. 4. Plot of the relativistic dispersion relation defined in Eq. (8). The straight line is a linear fit to all 87 data points shown, including lattice sizes $48 \le L \le 96$ and values of μ in the range $10.0 \le \mu \le 14.0$. The data include particles from both $C = \pm$ sectors with quantum numbers $0 \le j \le 2$ and $0 \le n \le 3$. The slope of the fitted line is ζ^2, yielding $\zeta = 2.5907(2)$ with $\chi^2/\mathrm{d.o.f.} = 1.73$. The data underlying this plot, including all points used in the fit, are tabulated in Ref. [28].

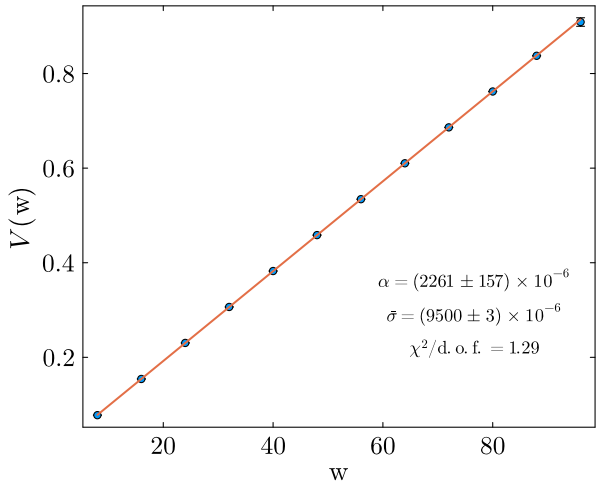


FIG. 5. Static quark potential $V(\mathsf{w})$ versus separation w for $L = 96$ and $\mu = 14.0$. The solid line is a linear fit whose slope gives $\bar{\sigma} = 0.009500(3)$.

LGT on the plaquette-chain geometry have not yet been investigated. An important open question is therefore whether the continuum limit identified in this work can also be realized within the conventional Kogut-Susskind formulation. In the Euclidean spacetime formulation, the Wilson action contains independent couplings multiplying the temporal-chain, temporal-rung, and purely spatial plaquette terms. Exploring the continuum limits obtained by tuning these plaquette couplings would determine whether the conventional and qubit-regularized LGTs can flow to the same continuum theory or instead belong to distinct universality classes.

The continuum limits that we have discovered here and in Ref. [9] can be viewed as arising from confinement–deconfinement transitions. For small nonzero h, the ground state of Eq. (3) lies in the ordered phase with $g/J < 1$, where h selects one of three degenerate vacua, corresponding to confinement. For $g/J > 1$, the ground state is disordered and largely insensitive to h, corresponding instead to deconfinement. It is plausible that analogous quantum critical points associated with confinement–deconfinement transitions exist in higher dimensions. Identifying such points in simple qubit-regularized SU(3) gauge theories and determining whether they give rise to massive continuum QFTs would be an important next step. In three spatial dimensions, such theories would presumably correspond either to conventional YM theory or to a more exotic variant enabled by qubit regularization.

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S. C. conceived of the idea and R. X. S. carried out the investigations of this theory. Both of them created the outline of the manuscript. R. X. S. did all the DMRG calculations. The two together produced the first draft of the paper. T. B. checked the correctness of the physics discussed and provided further interpretation of the model in terms of a pseudo-one-dimensional gauge theory. All three authors refined the manuscript and agree with its final version.

Data availability. There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

[1] A. Athenodorou and M. Teper, The glueball spectrum of SU (3) gauge theory in 3 + 1 dimensions, J. High Energy Phys. 11 (2020) 172.
[2] S. Chandrasekharan, R. X. Siew, and T. Bhattacharya, Monomer-dimer tensor-network basis for qubit-regularized lattice gauge theories, Phys. Rev. D **111**, 114502 (2025).
[3] S. Chandrasekharan, Qubit regularization of quantum field theories, Proc. Sci. LATTICE2024 (**2025**) 001 [arXiv:2502.05716].
[4] C. W. Bauer, Z. Davoudi, N. Klco, and M. J. Savage, Quantum simulation of fundamental particles and forces, Nat. Rev. Phys. **5**, 420 (2023).
[5] A. Kapustin and E. Witten, Electric-magnetic duality and the geometric Langlands program, Commun. Num. Theor. Phys. **1**, 1 (2007).
[6] J. B. Kogut and L. Susskind, Hamiltonian formulation of Wilson's lattice gauge theories, Phys. Rev. D **11**, 395 (1975).

[7] T. Bhattacharya, A. J. Buser, S. Chandrasekharan, R. Gupta, and H. Singh, Qubit regularization of asymptotic freedom, Phys. Rev. Lett. **126**, 172001 (2021).
[8] S. Maiti, D. Banerjee, S. Chandrasekharan, and M. K. Marinkovic, Asymptotic freedom at the Berezinskii-Kosterlitz-Thouless transition without fine-tuning using a qubit regularization, Phys. Rev. Lett. **132**, 041601 (2024).
[9] R. X. Siew, S. Chandrasekharan, and T. Bhattacharya, Asymptotic freedom and massive glueballs in a qubit-regularized SU(2) gauge theory, Phys. Rev. D **113**, 114522 (2026).
[10] E. Shahmoon, I. Mazets, and G. Kurizki, Giant vacuum forces via transmission lines, Proc. Natl. Acad. Sci. U.S.A. **111**, 10485 (2014).
[11] L. Del Debbio, A. Hart, and E. Rinaldi, Light scalars in strongly-coupled extra-dimensional theories, J. High Energy Phys. 07 (2012) 178.
[12] F. Antonuccio and S. Dalley, Glueballs from $(1+1)$-dimensional gauge theories with transverse degrees of freedom, Nucl. Phys. **B461**, 275 (1996).
[13] J. Nyhegn, C.-M. Chung, and M. Burrello, $\mathbb{Z}_N$ lattice gauge theory in a ladder geometry, Phys. Rev. Res. **3**, 013133 (2021).
[14] S. Pradhan, A. Maroncelli, and E. Ercolessi, Discrete Abelian lattice gauge theories on a ladder and their dualities with quantum clock models, Phys. Rev. B **109**, 064410 (2024).
[15] X. Yao, SU(2) gauge theory in $2+1$ dimensions on a plaquette chain obeys the Eigenstate thermalization hypothesis, Phys. Rev. D **108**, L031504 (2023).
[16] L. Ebner, A. Schäfer, C. Seidl, B. Müller, and X. Yao, Entanglement entropy of $(2+1)$-dimensional SU(2) lattice gauge theory on plaquette chains, Phys. Rev. D **110**, 014505 (2024).
[17] P. P. Popov, E. Ballini, A. Bottarelli, M. Burrello, P. Silvi, M. M. Wauters, and P. Hauke, When the center matters: Color screening and gluelumps in dihedral lattice gauge theories, arXiv:2512.00152.
[18] A. Ciavarella, N. Klco, and M. J. Savage, Trailhead for quantum simulation of SU(3) Yang-Mills lattice gauge theory in the local multiplet basis, Phys. Rev. D **103**, 094501 (2021).
[19] S. V. Kadam, A. Naskar, I. Raychowdhury, and J. R. Stryker, Loop-string-hadron approach to SU(3) lattice Yang-Mills theory: Hilbert space of a trivalent vertex, Phys. Rev. D **111**, 074516 (2025).
[20] V. Chen, B. Müller, and X. Yao, Minimally truncated SU(3) lattice gauge theory and string tension, Phys. Rev. D **114**, 074502 (2026).
[21] S. Elitzur, R. B. Pearson, and J. Shigemitsu, Phase structure of discrete Abelian spin and gauge systems, Phys. Rev. D **19**, 3698 (1979).
[22] A. B. Zamolodchikov and V. A. Fateev, Nonlocal (parafermion) currents in two-dimensional conformal quantum field theory and self-dual critical points in Z_N-symmetric statistical systems, Sov. Phys. JETP **62**, 215 (1985).
[23] L. Lepori, G. Z. Tóth, and G. Delfino, Particle spectrum of the 3-state Potts field theory: A numerical study, J. Stat. Mech. (2009) P11007.
[24] J. Cardy, *Scaling and Renormalization in Statistical Physics*, Vol. 5 (Cambridge University Press, Cambridge, England, 1996).
[25] R. S. K. Mong, D. J. Clarke, J. Alicea, N. H. Lindner, and P. Fendley, Parafermionic conformal field theory on the lattice, J. Phys. A **47**, 452001 (2014).
[26] M. Fishman, S. R. White, and E. M. Stoudenmire, The ITensor software library for tensor network calculations, SciPost Phys. Codebases, 4 (2022).
[27] M. Fishman, S. R. White, and E. M. Stoudenmire, Codebase release 0.3 for ITensor, SciPost Phys. Codebases, 4-r0.3 (2022).
[28] See Supplemental Material at http://link.aps.org/supplemental/10.1103/xtsm-rw6d for the fitted data and plots.
[29] U. M. Heller, SU(3) lattice gauge theory in the fundamental adjoint plane and scaling along the Wilson axis, Phys. Lett. B **362**, 123 (1995).
[30] V. Fateev, Integrable deformations in Z_N-symmetrical models of the conformal quantum field theory, Int. J. Mod. Phys. A **06**, 2109 (1991).
[31] J. L. Cardy, Operator content of two-dimensional conformally invariant theories, Nucl. Phys. **B270**, 186 (1986).