

Observation of emergent $\mathbb{Z}_2$ gauge invariance in a superconducting circuit

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Lattice gauge theories (LGTs) are one of the most fundamental subjects in many-body physics, and has recently attracted considerable research interests in quantum simulations. Here we experimentally investigate the emergent $\mathbb{Z}_2$ gauge invariance in a 1D superconducting circuit with 10 transmon qubits. By precisely adjusting staggered longitudinal and transverse fields to each qubit, we construct an effective Hamiltonian containing an LGT and gauge-broken terms. The corresponding matter sector can exhibit a localization, and there also exists a 3-qubit operator, of which the expectation value can retain nonzero for a long time in low-energy regimes. The above localization can be regarded as the confinement of matter fields, and the 3-body operator is the $\mathbb{Z}_2$ gauge generator. These experimental results demonstrate that, despite the absence of gauge structure in the effective Hamiltonian, $\mathbb{Z}_2$ gauge invariance can still emerge in low-energy regimes. Our work provides a method for both theoretically and experimentally studying the rich physics in quantum many-body systems with emergent gauge invariance.

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Introduction.— Gauge invariance is one of the most fundamental principles of quantum field theories, and lattice gauge theories (LGTs) [1–4] play a significant role in a wide range of modern physics, e.g., helping us understand the confinement of quarks [1]. Recently, as the rapid development of quantum simulations [5,6], studying LGTs in synthetic quantum many-body systems becomes possible and has drawn many interests from both theoretical and experimental physicists [7–13]. On the one hand, many unresolved problems associated with LGTs are potentially solvable via large-scale quantum simulations. On the other hand, quantum simulations provide a new viewpoint to study LGTs, i.e., nonequilibrium dynamics [14–17]. The corresponding experimental studies have been demonstrated in ultracold atoms [18–20].

In condensed-matter physics, gauge invariance is generally not a necessary element and thus absent in the Hamiltonian. However, due to strong quantum fluctuations of quantum many-body systems, gauge invariance may emerge in low-energy regimes leading to various novel quantum phases,

such as quantum spin liquid [21–23] and deconfined quantum critical points [24]. In addition, it is challenging to construct an exact gauge-invariant Hamiltonian in an artificial quantum system [9,18]. Thus one natural question is whether we can construct a quantum many-body system that the gauge structure is absent in the Hamiltonian but emergent in the low-energy physics. Moreover, this emergent gauge invariance can lead to novel dynamics, for instance, the confinement-induced localization of matter degrees of freedom [17,18], which can be detected by state-of-the-art quantum simulators.

In this Letter, we experimentally demonstrate the emergent $\mathbb{Z}_2$ gauge invariance in a superconducting processor with 10 qubits arranged into a chain. Following the scheme in Ref. [25], we apply staggered longitudinal and transverse fields to the system to construct an effective Hamiltonian, which is a mixture of a $\mathbb{Z}_2$ LGT and some gauge-broken terms. To detect the emergent $\mathbb{Z}_2$ gauge invariance, we first study the charge spreading of the matter sector. The experimental result shows that it can exhibit a localization under proper parameters, which is a strong dynamical signature of the confinement induced by emergent gauge invariance. Then, benefiting from the joint readout of three qubits under the arbitrary basis, we can study the time evolution of 3-qubit operators. We find that there is a 3-qubit operator whose expectation value can retain nonzero for a long time. We demonstrate that this operator can be regarded as the $\mathbb{Z}_2$ gauge generator. Therefore this experimental result is direct evidence of the emergent $\mathbb{Z}_2$ gauge invariance. Furthermore, according to relations between charge spreading and initial states, we

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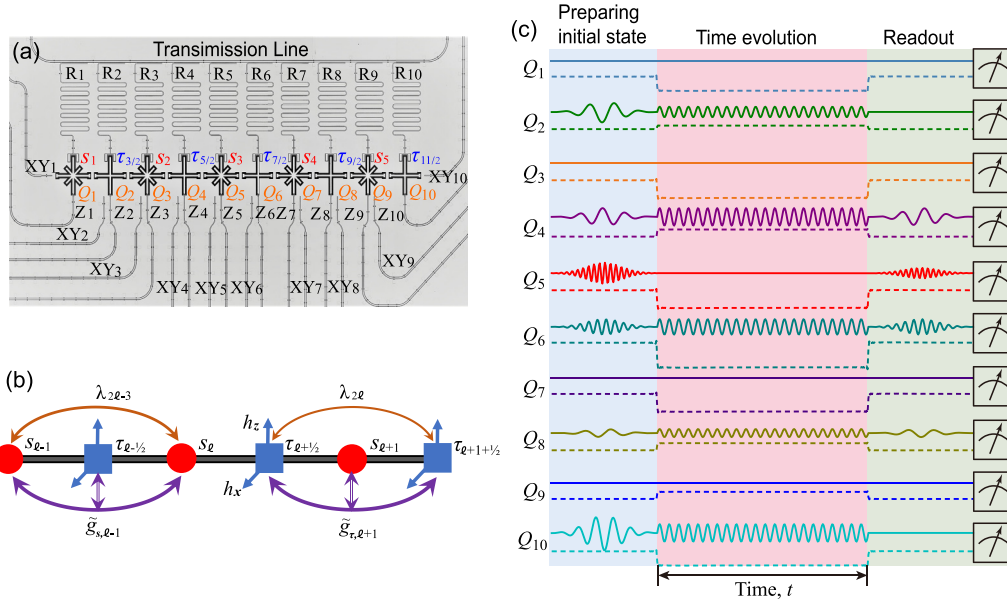


FIG. 1. Set-up and protocol. (a) Optical micrograph of the superconducting circuit. There are 10 transmon qubits arranged into a chain, where the odd and even qubits are used for realizing s and τ spins in the following experiment, respectively. Each qubit couples to a microwave line for the XY driving (XY_j), a flux bias line for the Z pulse (Z_j), and a readout resonator for measurement (R_j). The NN qubits are directly coupled via a capacitor, and the capacitance of each qubit itself contributes to the NNN coupling. The parameters of the device are presented in the Supplemental Material (SM) in detail [44]. (b) Lattice skeleton of the effective Hamiltonian (2). There are two types of spins, i.e., s (red sites) and τ (blue sites), corresponding to matter and gauge fields, respectively. The purple arrows represent the three-body couplings; the orange arrows are direct NNN two-body couplings, and the blue arrows are longitudinal/transverse fields. (c) Pulse sequences of the experiment. The solid and dashed lines represent XY drivings and Z pulses, respectively. First, all qubits are at their idle frequencies for preparing the initial state $|\psi_0\rangle$ via XY driving. Then, each qubit is biased to the working point with Z pulses, and meanwhile, the XY driving is applied to each even qubit. After the system evolves for time t , all qubits are biased back to idle points for readout.

find that the gauge structure mainly emerges in low-energy regimes.

Model and set-up.— Due to the scalability, long coherent time, and high-precision full control, superconducting circuits [26,27] become one of the most competitive candidates for achieving the universal quantum computation [28,29], and they are also a more suitable platform for performing quantum simulations [30–43]. Several theoretical schemes have been proposed to simulate LGTs via specific superconducting circuits [11,12]. However, the corresponding experimental realization is still absent, which may be due to the technological challenges for preparing such superconducting circuits. Here inspired by the experiment in Ref. [25], we find that the realization of LGTs is feasible by using conventional chainlike transmon qubits.

This experiment is performed in a superconducting circuit with 10 transmon qubits (Q_1 – Q_{10}) arranged into a chain; see Fig. 1(a). Due to the large and staggered anharmonicity [44], the system can be described by an isotropic 1D XY model with tunable transverse and longitudinal fields [18–20]. In addition, the next-nearest-neighbor (NNN) coupling cannot be neglected in the following experiment. Therefore the Hamiltonian can be written as

$$\hat{H} = \sum_j (g_j \hat{\sigma}_j^+ \hat{\sigma}_{j+1}^- + \lambda_j \hat{\sigma}_j^+ \hat{\sigma}_{j+2}^- + \text{H.c.}) + \sum_{j=1} \left(\frac{V_j}{2} \hat{\sigma}_j^z + h_{x,j} \hat{\sigma}_j^x \right), \quad (1)$$

where $\hat{\sigma}^\pm = (\hat{\sigma}^x \pm i\hat{\sigma}^y)/2$, $\hat{\sigma}^{x,y,z}$ are Pauli matrices, $g_j/2\pi \approx 12$ MHz and $\lambda_j/2\pi \approx 0.7 - 1.1$ MHz are the nearest-neighbor (NN) and NNN coupling strengths [44], respectively, V_j is a longitudinal field tuned by Z pulses, and $h_{x,j}$ is a transverse field controlled by XY drivings.

Here we let the longitudinal field at each odd qubit $V_{2\ell-1}/2\pi = -80$ MHz and the detuning between NN qubits is much larger than the coupling strength, i.e., $V_{2\ell} - V_{2\ell-1} \gg g_j$. In addition, the transverse field is only applied at the even qubits with equal strength h_x . Thus according to Ref. [25], we can obtain an effective Hamiltonian $\hat{H}_{\text{eff}}$, which reads

$$\begin{aligned} \hat{H}_{\text{eff}} &= \hat{H}_1 + \hat{H}_2 + \hat{H}_3, \\ \hat{H}_1 &= \sum_{\ell=1}^4 \tilde{g}_{s,\ell} (\hat{s}_\ell^+ \hat{\tau}_{\ell+\frac{1}{2}}^z \hat{s}_{\ell+1}^- + \text{H.c.}) + \sum_{\ell=1}^5 h_x \hat{\tau}_{\ell+\frac{1}{2}}^x, \\ \hat{H}_2 &= - \sum_{\ell=1}^4 \tilde{g}_{\tau,\ell} (\hat{\tau}_{\ell-\frac{1}{2}}^+ \hat{s}_\ell^z \hat{\tau}_{\ell+\frac{1}{2}}^- + \text{H.c.}), \\ \hat{H}_3 &= \sum_{\ell=1}^4 (\lambda_{2\ell-1} \hat{s}_\ell^+ \hat{s}_{\ell+1}^- + \lambda_{2\ell} \hat{\tau}_{\ell+\frac{1}{2}}^+ \hat{\tau}_{\ell+\frac{3}{2}}^- + \text{H.c.}) + \sum_{\ell=1}^5 h_z \hat{\tau}_{\ell+\frac{1}{2}}^z, \end{aligned} \quad (2)$$

where $\hat{s}_\ell = \hat{\sigma}_{2\ell-1}$ and $\hat{\tau}_{\ell+\frac{1}{2}} = \hat{\sigma}_{2\ell}$ are also Pauli matrices labeling the odd and even qubits, respectively, and $\tilde{g}_{s/\tau,\ell} \approx g_j^2/\Delta \approx -2\pi \times 1.8$ MHz is the effective three-body coupling strength; see Fig. 1(b). Here the effective longitudinal field h_z is the summation of the original longitudinal field of even

qubits ($V_{2\ell}$) and the Lamb shift [44]. In the experiment, $V_{2\ell}$ is adjustable by controlling the detuning between qubit frequencies and XY driving frequencies.

Here $\hat{H}_1$ is nothing but a $\mathbb{Z}_2$ lattice gauge field (τ spins) coupled to a matter field (s spins), where the corresponding $\mathbb{Z}_2$ gauge generator reads $\hat{G}_\ell^0 = \hat{\tau}_{\ell-\frac{1}{2}}^x \hat{s}_\ell^z \hat{\tau}_{\ell+\frac{1}{2}}^x$, i.e., $[\hat{G}_\ell^0, \hat{H}_1] = 0$ [45]. However, we can find that additional terms $\hat{H}_2$ and $\hat{H}_3$ both violate this gauge invariance, so the whole Hamiltonian is not a rigorous LGT. However, by numerical simulation via density matrix renormalization group method [46,47] (the details are shown in the SM [44]), we find that a new $\mathbb{Z}_2$ gauge invariance can still emerge in the ground state for a large transverse field h_x and proper longitudinal field h_z . The corresponding new gauge generator is

$$\hat{G}_\ell = \tilde{\tau}_{\ell-\frac{1}{2}}^x \hat{s}_\ell^z \tilde{\tau}_{\ell+\frac{1}{2}}^x, \quad (3)$$

where $\tilde{\tau}^x = \sin \beta \hat{\tau}^z + \cos \beta \hat{\tau}^x$ and $\beta = \arctan(h_z/h_x)$.

Next, we will experimentally investigate quench dynamics of this system to probe the emergent gauge invariance. Here the gauge invariance is almost independent of the filling of s spins [44]. Thus without loss of generality, we consider the system containing only one s spin, i.e., $\sum_{\ell=1}^5 \hat{s}_\ell^+ \hat{s}_\ell^- = 1$. The initial state is chosen as $|\psi_0\rangle = |s\rangle \otimes |\tau\rangle$, where $|s\rangle$ and $|\tau\rangle$ label the states of s and τ spins, respectively. We let $|s\rangle = |00100\rangle$ and $|\tau\rangle = |\Phi_\theta \Phi_\theta \Phi_\theta \Phi_\theta \Phi_\theta\rangle$, where $|\Phi_\theta\rangle = \cos \frac{\theta}{2} |1\rangle + \sin \frac{\theta}{2} |0\rangle$. Here according to the effective Hamiltonian $\hat{H}_{\text{eff}}$, we know that different θ means different energy during the quench dynamics. The experimental procedure can be summarized as follows: First, all qubits are at their idle frequencies, and we use single-qubit rotational gates to prepare the initial state $|\psi_0\rangle$. Then, we bias the frequency of each qubit to the corresponding working point with Z pulses, i.e., let the local potential of each qubit be V_j . Meanwhile, we apply ac drivings with amplitude h_x to even qubits through XY lines to realize transverse fields. Finally, after the system evolves with time t , we bias back all qubits to their idle points and read out the corresponding observable. The pulse sequence of each qubit is shown in Fig. 1(c).

Confinement dynamics.— In 1D $\mathbb{Z}_2$ LGTs, the system is generally in a confined phase. In this case, similar to the confinement of quarks in particle physics, the presence of a gauge field will induce a long-range potential of matter fields, e.g., linear potential [48]. This long-range potential can suppress the transport of matter field and results in a localization [18,45]. Thus we can study the dynamics of s spins to demonstrate whether they can localize, which is a signature for identifying the existence of emergent gauge invariance. Here we measure the density distributions of the photon at odd qubits, i.e., the spin-density distribution of the s sector, defined as

$$P_j(t) := \langle \psi(t) | \hat{\sigma}_j^+ \hat{\sigma}_j^- | \psi(t) \rangle, \quad (4)$$

where $|\psi(t)\rangle = \exp(-i\hat{H}t)|\psi_0\rangle$ is the wave function of the system at time t .

In the case of $h_x/2\pi = 2$ MHz, $V_{2\ell}/2\pi = 0$ MHz, and $\theta = -\pi/3$, we can find that s spins delocalize and spread to the whole system very quickly; see Fig. 2(a). However, the situation is different when $h_x/2\pi = 6$ MHz and $V_{2\ell}/2\pi =$

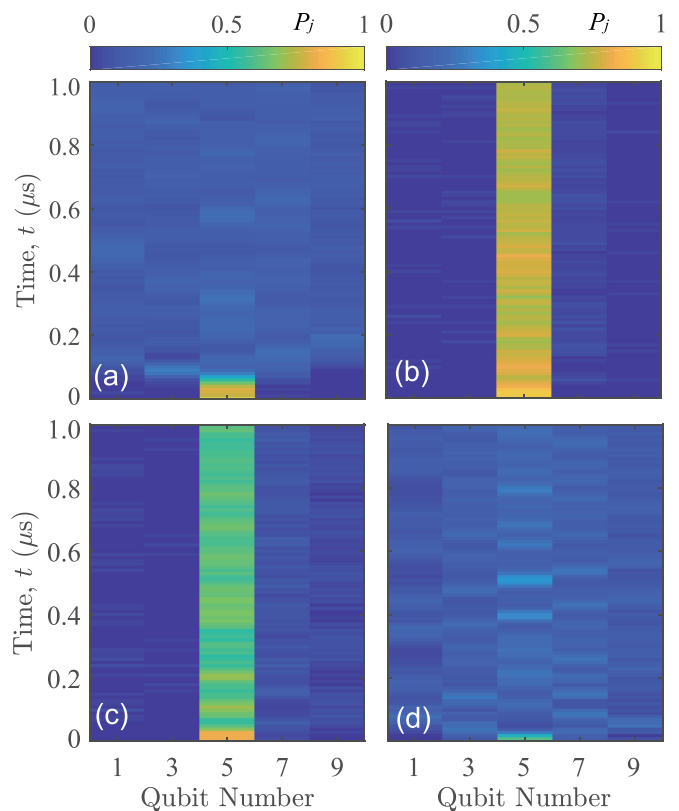


FIG. 2. Spreading of s spins up to 1 μ s for (a) $h_x/2\pi = 2$ MHz, $V_{2\ell}/2\pi = 0$ MHz, $\theta = -\pi/3$; (b) $h_x/2\pi = 6$ MHz, $V_{2\ell}/2\pi = 15$ MHz, $\theta = -\pi/3$; (c) $h_x/2\pi = 6$ MHz, $V_{2\ell}/2\pi = 15$ MHz, $\theta = -\pi/2$; and (d) $h_x/2\pi = 6$ MHz, $V_{2\ell}/2\pi = 15$ MHz, $\theta = \pi$. Here the s spin shows localization in the cases of (b) and (c) indicating the presence of confinement, while it transports freely in (a) and (d) indicating the absence of confinement. Each point is the average of 8000 single-shot readouts.

15 MHz. Figures 2(b) and 2(c) show that the s spin can indeed localize when $\theta = -\pi/3$ or $-\pi/2$, which is an evidence of existing confinement. When we continue to change the initial state, e.g., $\theta = \pi$, the localization of s spins disappears; see Fig. 2(d). These experimental results indicate that the confinement of s spins can indeed emerge for the proper transverse and longitudinal fields and also depend on the system energy.

Next, we further explore the relation between confinement and the system energy, which can be characterized by the localization strength of s spins. Here we define the extended imbalance of s spins [42]

$$\mathcal{I} := \sum_{j=\text{odd}} \eta_j P_j, \quad (5)$$

where $\eta_j = 1/N_1$ ($-1/N_0$) if the initial state of Q_j is $|1\rangle$ ($|0\rangle$), and N_1 (N_0) is the number of $|1\rangle$ ($|0\rangle$) for the initial state of the s sector. In Fig. 3(a), we show the dynamics of extended imbalance for different θ under the condition of $h_x/2\pi = 6$ MHz and $V_{2\ell}/2\pi = 15$ MHz. It shows that $\mathcal{I}$ can stabilize at different values indicating the different localization strength. Now we use the steady value of extended imbalance $\mathcal{I}_\infty$ to quantify the localization strength, which can be calculated as the average of $\mathcal{I}$ during the last 0.8 μ s. Here the larger

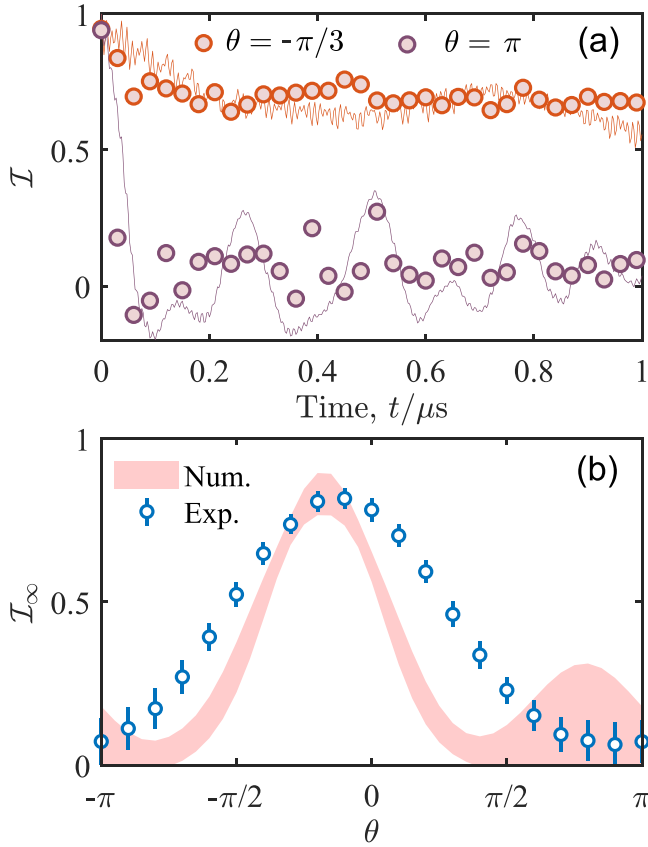


FIG. 3. Extended imbalance $\mathcal{I}$ versus initial states. (a) Time evolution of $\mathcal{I}$ with $h_x/2\pi = 6$ MHz, $V_{2\ell}/2\pi = 15$ MHz, and different θ , i.e., cases in Figs. 2(b) and 2(c). The dots are experimental data, while the solid curves are the corresponding numerical results. (b) The relation between $\mathcal{I}_\infty$ and θ . Here $\mathcal{I}_\infty$ is calculated as the average of $\mathcal{I}$ from 0.2 to 1 μs , and the standard deviation is the corresponding variance. Using Gaussian fitting, we obtain the peak of experimental result is at $\theta = \theta_m \approx -0.35$. The numerical results are obtained via Hamiltonian in Eq. (1), where decoherence and dephasing are both neglected.

$\mathcal{I}_\infty$ means the stronger localization strength. In Fig. 3(b), the relation between $\mathcal{I}_\infty$ and θ is presented. We can find that when $\theta = \theta_m \approx -0.35$, the corresponding $\mathcal{I}_\infty$ is the largest, indicating the strongest localization strength in this case. In the following discussion, we will verify that this is because the emergent gauge invariance can only exist in a low-energy regime, and the initial state is close to the ground state when $\theta = \theta_m$.

Gauge invariance.— Now we start to directly study the gauge invariance of the system. Here when deriving the effective Hamiltonian $\hat{H}_{\text{eff}}$, the correction of longitudinal field h_z from high-order terms can hardly be directly confirmed accurately. Thus the β or $\mathbb{Z}_2$ gauge generator $\hat{G}_\ell$ is in fact unknown. However, we can define an ansatz of $\mathbb{Z}_2$ gauge generator as

$$\hat{G}_\ell(\alpha) := \hat{T}_{\ell-\frac{1}{2}}(\alpha) \hat{s}_\ell^z \hat{T}_{\ell+\frac{1}{2}}(\alpha), \quad (6)$$

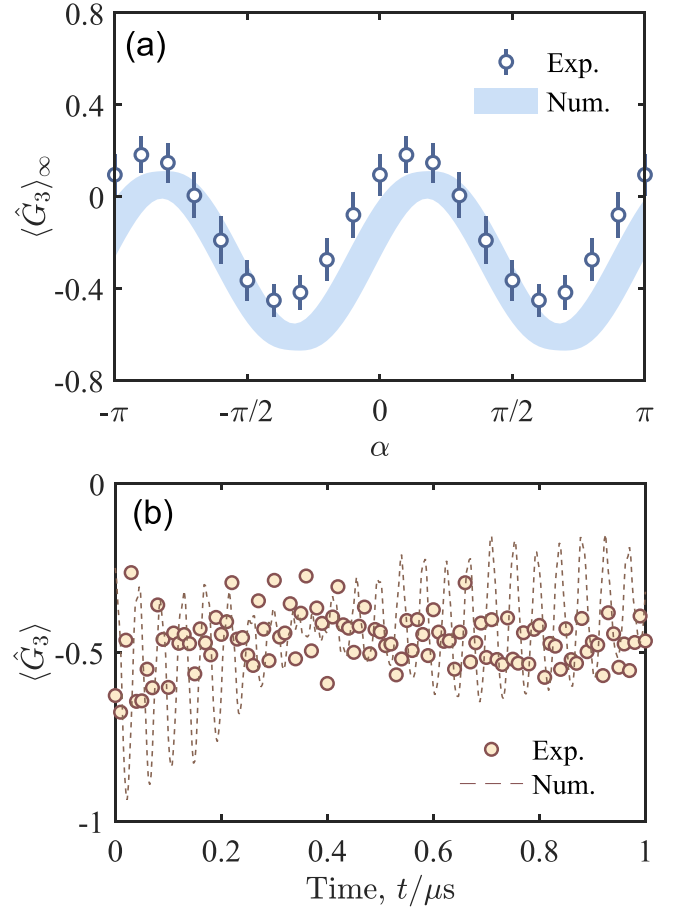


FIG. 4. Gauge invariance. (a) The relation between $\langle \hat{G}_3 \rangle_\infty$ and α with the parameters identical to the cases in Figs. 2(b). Here $\langle \hat{G}_3 \rangle_\infty$, calculated as the average of $\mathcal{I}$ from 0.2 to 1 μs with the standard deviation being the error bar, represents the steady expectation value of 3-qubit operator $\hat{G}_3(\alpha)$, defined in Eq. (6). The curve is almost a sine-shape curve with the valley at $\alpha \approx -1.19$. (b) Time evolution of 3-qubit operator $\langle \hat{G}_3(\alpha = -1.19) \rangle$ up to 1 μs . It shows that $\langle \hat{G}_3 \rangle$ can retain a nonzero value for a long time, where the oscillation is from the high-order term [44].

where $\hat{T}_{\ell-\frac{1}{2}}(\alpha) = \cos(\alpha) \hat{\tau}_{\ell-\frac{1}{2}}^x + \sin(\alpha) \hat{\tau}_{\ell-\frac{1}{2}}^z$, so $\hat{T}_{\ell+\frac{1}{2}}(\beta) = \hat{\tau}_{\ell-\frac{1}{2}}^x$ and $\hat{G}_\ell(\beta) = \hat{G}_\ell$ is the emergent $\mathbb{Z}_2$ gauge generator defined in Eq. (3).

If the specific eigenstate has emergent $\mathbb{Z}_2$ gauge invariance and the initial state has a large overlap with this state, then the expectation value of $\hat{G}_\ell$ during the quench dynamics will be nearly time independent. Furthermore, the steady expectation value of $\hat{G}_\ell(\alpha)$ during the dynamics approaches the minimum/maximum when $\alpha = \beta$ [44]. Therefore based on these considerations, we can determine β and thus fix $\hat{G}_\ell$ in the experiment. To measure the expectation value of $\hat{G}_\ell(\alpha)$, we need the joint readout of adjacent three qubits under the specific basis, which is accessible in superconducting circuits. We should measure the expectation values of $\hat{\tau}_{\ell-\frac{1}{2}}^z \hat{s}_\ell^z \hat{\tau}_{\ell+\frac{1}{2}}^z$, $\hat{\tau}_{\ell-\frac{1}{2}}^z \hat{s}_\ell^z \hat{\tau}_{\ell+\frac{1}{2}}^x$, $\hat{\tau}_{\ell-\frac{1}{2}}^x \hat{s}_\ell^z \hat{\tau}_{\ell+\frac{1}{2}}^z$, and $\hat{\tau}_{\ell-\frac{1}{2}}^x \hat{s}_\ell^z \hat{\tau}_{\ell+\frac{1}{2}}^x$, respectively. Then, we combine these four values linearly according to Eq. (6).

In Fig. 4(a), we show the relation between $\langle \hat{G}_3 \rangle_\infty$ and α when $h_x/2\pi = 6$ MHz, $V_{2\ell}/2\pi = 15$ MHz, and $\theta = -\pi/3$. Here $\langle \hat{G}_3 \rangle_\infty$ is the average of $\langle \hat{G}_3 \rangle$ during the last $0.8 \mu\text{s}$ representing the steady value of $\langle \hat{G}_3 \rangle$. We find that $\langle \hat{G}_3 \rangle_\infty$ is α dependent, and it can approach the minimum when $\alpha = \beta \approx -1.19$. According to Fig. 4(b), we can find that $\langle \hat{G}_3(-1.19) \rangle$ is stabilized at a nonzero value (the oscillation is from the high-order term and absent for effective Hamiltonian $\hat{H}_{\text{eff}}$ [44]). This result provides evidence that some specific eigenstates of $\hat{H}_{\text{eff}}$ are almost the eigenstates of $\hat{G}_\ell$. That is, these eigenstates of the effective Hamiltonian can emerge as a $\mathbb{Z}_2$ gauge invariance.

According to $\hat{H}_{\text{eff}}$, when the external field $\sqrt{h_x^2 + h_z^2}$ is much larger than the coupling strength, the τ sector will become the leading contribution of the energy, and the initial state is much closer to the ground state when $\theta = -\pi/2 - \beta$. From Figs. 3(b) and 4(a), we can find that θ_m almost equals $-\pi/2 - \beta$, and the localization of the s sector mainly exists at the vicinity of θ_m , i.e., in low-energy regimes. Therefore we conjecture that the $\mathbb{Z}_2$ gauge invariance mainly emerges in the low-energy states of $\hat{H}_{\text{eff}}$. In the SM [44], by obtaining wave functions of all eigenstates of $\hat{H}_{\text{eff}}$ by the exact diagonalization method, we can find that the gauge invariance indeed mainly emerges in low-energy regimes.

Summary.— In conclusion, we have experimentally investigated the emergent $\mathbb{Z}_2$ gauge invariance in a 10-qubit

superconducting processor. Our experimental results demonstrate that $\mathbb{Z}_2$ gauge invariance can indeed emerge in a low-energy regime, even though the $\mathbb{Z}_2$ gauge structure is absent in the effective Hamiltonian. Moreover, this emergent gauge invariance can lead to exotic dynamical behaviors, for instance, confinement-induced localization, which has been observed in this experiment. Our results can scale up to larger quantum systems and enable the further study of emergent LGTs in superconducting circuits. For instance, the dynamics of string breaking [16], the thermalization of the effective Hamiltonian $\hat{H}_{\text{eff}}$, and whether existing disorder-free many-body localization [49,50] are interesting issues. In addition, how to realize a truly gauge-invariant Hamiltonian on superconducting circuits is another relevant question.

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