Experimental implementation of universal nonadiabatic geometric quantum gates in a superconducting circuit

Y. Xu,1,* Z. Hua,1,* Tao Chen,2,* X. Pan,1 X. Li,1 J. Han,1 W. Cai,1 Y. Ma,1 H. Wang,1 Y.P. Song,1 Zheng-Yuan Xue,2,* and L. Sun1,1

1Center for Quantum Information, Institute for Interdisciplinary Information Sciences, Tsinghua University, Beijing 100084, China
2Guangdong Provincial Key Laboratory of Quantum Engineering and Quantum Materials, GPETR Center for Quantum Precision Measurement, and School of Physics and Telecommunication Engineering, South China Normal University, Guangzhou 510006, China

Using geometric phases to realize noise-resilient quantum computing is an important method to enhance the control fidelity. In this work, we experimentally realize a universal nonadiabatic geometric quantum gate set in a superconducting qubit chain. We characterize the realized single- and two-qubit geometric gates with both quantum process tomography and randomized benchmarking methods. The measured average fidelities for the single-qubit rotation gates and two-qubit controlled-Z gate are 0.9977 and 0.977, respectively. Besides, we also experimentally demonstrate the noise-resilient feature of the realized single-qubit geometric gates by comparing their performance with the conventional dynamic gates with different types of errors in the control field. Thus, our experiment proves a way to achieve high-fidelity geometric quantum gates for robust quantum computation.

In quantum physics, wave functions up to a global phase are equivalent, and thus the important role played by the phase factors had been ignored for a long time. However, the evolution of a quantum state can be traced in some extent by a geometric phase factor. A famous example is the Aharonov-Bohm effect [1], which shows that the phases with a geometric origin can have observable consequences [2]. Different from the dynamical phase, geometric phases [2–4] are gauge invariant and depend only on the global properties of the evolution path. Therefore, besides their fundamental importance, geometric phases have been tested in a variety of settings and have found many interesting applications [5–7].

Recently, there is a renewed interest in applying geometric phases into the field of quantum computation [8–10], which is potentially capable to handle hard problems for classical computers [11]. The reason is that the global properties of the geometric phases can be naturally used to achieve noise-resilient quantum manipulation against certain local noises [12–14], which is essential for practical quantum computation as quantum systems are fragile to both the environment-induced noises and the operational imperfections. By using adiabatic cyclic evolutions, recent experiments have reported the detection of geometric phases [15–23] and the realization of single-qubit operations [24, 25] and a universal set of quantum gates [26, 27] in several physical systems. However, the speed of the quantum gates induced from adiabatic evolution is rather slow, and thus decoherence will introduce considerable errors [28, 29].

To overcome the dilemma between the limited coherence times and the long duration of adiabatic evolution, implementation of quantum gates based on nonadiabatic evolutions has been proposed for both the Abelian [28, 29] and non-Abelian geometric phases [30, 31]. Recently, based on the nonadiabatic non-Abelian geometric phase, single-qubit quantum gates [32–38] and a universal set of quantum gates [39–41] have been experimentally demonstrated in various three-level physical systems. However, the noise-resilience of the geometric phases is not shared by this type of nonadiabatic approach of implementation [42–44]. Robust quantum computation with non-Abelian geometric phase can actually be implemented with two degenerated dark states [45, 46]. However, it is experimentally difficult because of the need of complex control of quantum systems with four energy levels. On the other hand, experimental demonstration of universal quantum computation based on nonadiabatic Abelian geometric phase is also lacking, due to the challenge of exquisite control among quantum systems. In addition, so far there is no direct experimental verification of the noise-resilient feature of geometric quantum gates over the dynamical ones yet.

Here, with a multi-qubit superconducting quantum circuit architecture [47–49], we experimentally demonstrate a robust nonadiabatic geometric quantum computation (GQC) scheme [50]. Superconducting circuit systems possess the distinct merit of large scale integrability and flexibility, and thus are considered as a promising candidate for physical implementation of scalable quantum computation [51, 52]. In the construction of our universal geometric quantum gates, we merely use simple and experimentally accessible microwave controls over the capacitively-coupled superconducting transmon qubits, each of which involves only two states [53]. The leakage of qubit states can be effectively suppressed and the coupling between the two qubits can be parametrically tuned in a large range [54–57]. Meanwhile, our demonstration only utilizes conventional resonant interaction for both single- and two-qubit gates, and thus simplifies the experimental complexity and decreases the error sources. Furthermore, we experimentally demonstrate the noise-resilient feature of the geometric quantum gates over the dynamical ones. Therefore, our experiment proves the way to achieve robust universal GQC on a large-scale qubit lattice.

We first explain how to construct the single-qubit geometric gate on a superconducting transmon qubit in the \{\ket{0}, \ket{1}\} subspace, where \ket{0} (\ket{1}) denotes the ground (excited) state of the transmon qubit. Conventionally, single-qubit control is realized by applying a microwave drive on resonance with the...
qubit transition $|0\rangle \leftrightarrow |1\rangle$, as described by the Hamiltonian of
\[ H_1 = \frac{1}{2}\Omega(t)e^{i\phi(t)}|0\rangle\langle 1| + H.c., \quad (1) \]
where $\Omega(t)$ and $\phi(t)$ are the time-dependent driving amplitude and phase of the microwave field. To achieve a universal set of arbitrary single-qubit nonadiabatic geometric gates in a single-loop way, we divide the entire evolution time $\tau$ into three intervals: $0 \rightarrow \tau_1$, $\tau_1 \rightarrow \tau_2$, and $\tau_2 \rightarrow \tau$, with the driving amplitude and phase in each component satisfying
\[
\begin{align*}
\int_0^{\tau_1} \Omega(t)dt &= \theta, \quad \phi = \varphi - \frac{\pi}{2}, \quad t \in [0, \tau_1], \\
\int_{\tau_1}^{\tau_2} \Omega(t)dt &= \pi, \quad \phi = \varphi + \gamma + \frac{\pi}{2}, \quad t \in [\tau_1, \tau_2], \\
\int_{\tau_2}^{\tau} \Omega(t)dt &= \pi - \theta, \quad \phi = \varphi - \frac{\pi}{2}, \quad t \in [\tau_2, \tau].
\end{align*}
\]

Consequently, two orthogonal states $|\psi_+\rangle = \cos \frac{\theta}{2}|0\rangle + \sin \frac{\theta}{2}e^{i\varphi}|1\rangle$ and $|\psi_-\rangle = \sin \frac{\theta}{2}e^{-i\varphi}|0\rangle - \cos \frac{\theta}{2}|1\rangle$ undergo a cyclic evolution on the single-qubit Bloch sphere, as shown in Fig. 1(a), resulting in a geometric phase $\gamma$ ($-\gamma$) on the quantum state $|\psi_+\rangle$ ($|\psi_-\rangle$). Therefore, the obtained single-qubit gate of the total geometric evolution is
\[
U_1(\theta, \gamma, \varphi) = \cos \gamma + i\sin \gamma \left( \begin{array}{cc} \cos \theta & \sin \theta e^{-i\varphi} \\ \sin \theta e^{i\varphi} & -\cos \theta \end{array} \right) = \exp \left( i\gamma \vec{n} \cdot \vec{\sigma} \right), \quad (3)
\]
which corresponds to a rotation operation around the axis $\vec{n} = (\sin \theta \cos \varphi, \sin \theta \sin \varphi, \cos \theta)$ by an angle $-2\gamma$. The parameters $\theta$, $\gamma$, $\varphi$ are determined by the applied microwave fields.

Our experiment is performed on a five-Xmon-qubit chain sample \cite{57, 58}, with the simplified circuit schematic shown in Fig. 1(b). Only two adjacent qubits $Q_A$ and $Q_B$ are used in this experiment, with $|0\rangle \leftrightarrow |1\rangle$ transition frequency of $\omega_A/2\pi = 4.602$ GHz and $\omega_B/2\pi = 5.081$ GHz, respectively. Only the lowest two energy levels are considered here due to the large anharmonicity $\alpha_A/2\pi = -202$ MHz and $\alpha_B/2\pi = -190$ MHz for the Xmon qubits $Q_A$ and $Q_B$, respectively. More details about the device parameters can be found in Ref. \cite{59}.

We first demonstrate the single-qubit nonadiabatic geometric gates on superconducting qubit $Q_B$, with the experimental pulse sequence shown in Fig. 1(c). As a demonstration, here we fix $\theta = \pi/2$, and realize single-qubit $\pi$ and $\pi/2$ rotations around $X$ and $Y$ axes (denoted as $X$, $Y$, $X/2$, and $Y/2$ respectively), which construct a basis set to generate single-qubit Cliffords. The single-qubit geometric gate consists of a $\pi$ rotation sandwiched by two $\pi/2$ rotations with a total width of 80 ns. The envelope of each pulse is a truncated Gaussian pulse with the correction of “derivative removal by adiabatic gate” method in order to suppress the leakage to the undesired energy levels \cite{60}.

We characterize the single-qubit geometric gates by a single-qubit quantum process tomography (QPT) method \cite{59}, with the experimental sequence shown in Fig. 1(c). The experimental process matrices $\chi_{\text{exp}}$ of four specific geometric gates $X$, $X/2$, $Y/2$ and Hadamard $H$ (implemented with a $Y/2$ rotation followed by a $X$ rotation) are shown in Fig. 1(d) with an average process fidelity of 0.9980. The process fidelity is calculated through $F_P = \text{Tr} (\chi_{\text{exp}} \chi_{\text{ideal}})$, where $\chi_{\text{ideal}}$ is the ideal process matrix for the corresponding gate.

Another conventional method, Clifford-based randomized benchmarking (RB) \cite{61–63}, is also used to characterize the geometric gates, with the sequences for both the reference RB and interleaved RB experiments shown in the inset of Fig. 2. The experimental measured sequence fidelity decays as a function of the number of single-qubit Cliffords $m$ for both the reference RB and interleaved RB experiments are shown in Fig. 2. The reference RB experiment gives an average fidelity $F_{\text{avg}} = 0.9977$ for the realized single-qubit nonadiabatic geometric gates in the Clifford group. The measured gate fidelities of the four specific nonadiabatic geometric gates: $X$, $Y$, $X/2$, and $Y/2$, inserted in the random Cliffords in the interleaved RB experiment, are 0.9976, 0.9975, 0.9981, and 0.9975, respectively.

With the realized single-qubit nonadiabatic geometric gates: $X$, $Y$, $X/2$, and $Y/2$, inserted in the random Cliffords in the interleaved RB experiment, are 0.9976, 0.9975, 0.9981, and 0.9975, respectively.
Sequence Fidelity

![Graph of Sequence Fidelity](image)

FIG. 2: RB of single-qubit nonadiabatic geometric gates. Inset is the experimental pulse sequences to perform both the reference RB and interleaved RB experiments. Fit to the reference decay curve gives an average fidelity of 0.9977 for the single-qubit geometric gates in the Clifford group. The difference between the reference and the interleaved decay curves gives the gate fidelity of four specific gates: $X$, $Y$, $X/2$, $Y/2$.

1.0
0.8
0.6
0.4
-0.2 0.0 0.2
Rabi frequency error
Geo $T$
Dyn $T$
1.0
0.9
0.8
0.7
Geo $H$
Dyn $H$
Reference RB
Interleaved RB
Process fidelity
0.4
0.6
0.8
1.0
0.4
0.6
0.8
1.0
m - Number of Cliffords $C_1$
Gate Fidelity
Ref 0.9977
X 0.9976
Y 0.9975
X/2 0.9981
Y/2 0.9975

The geometric gates are realized with two different configuration settings, corresponding to two different geometric evolution trajectories. In configuration A, the geometric gates are realized with the geometric evolution described in Eq. 2 and have distinct advantages over the dynamic gates against additional Rabi frequency error $\epsilon$, as shown in Figs. 3(b-d). In configuration B, the geometric gates are realized by setting the phase $\phi = \varphi + \gamma - \pi/2$ at $[\tau_1, \tau_2]$ interval in Eq. 2, while the unitary of the geometric gate remains the same as that in Eq. 3 when $\theta = \pi/2$. The noise-resilient feature of the geometric gates still persists for different detuning errors, with the experimental results shown in Figs. 3(e-g). All experimental results also agree very well with the numerical simulations. The comparisons clearly illustrate the distinct advantages of the realized nonadiabatic geometric gates. We note that the noise-resilient feature of the geometric gates depends on the types of errors and the cyclic evolution paths of the geometric gates [59]. One can always find a specific evolution path of the control pulse to realize a noise-resilient geometric gate against the dominant error in the system.

In order to achieve a universal quantum computation, two-qubit entangling operations are also necessary. In our experiment, a non-trivial two-qubit geometric gate is also realized in a similar way to the single-qubit case by using a parametric modulation drive of one qubit frequency. Considering two adjacent qubits $Q_A$ and $Q_B$ (with anharmonicities $\alpha_A$ and $\alpha_B$) capacitively coupled to each other, the qubit frequency of $Q_A$ is modulated with a sinusoidal form:

$$\omega_A(t) = \omega_A + \varepsilon \sin(\nu t + \Phi),$$

where $\omega_A$ is the mean operating frequency, and $\varepsilon$, $\nu$, and $\Phi$ are the modulation amplitude, frequency, and phase, respectively. Ignoring the higher-order oscillating terms, when the modulation frequency satisfies $\nu = \omega_B - \omega_A + \alpha_B$, the parametric drive will induce a transition operation between the two energy levels $|11\rangle \leftrightarrow |02\rangle$ in the two-qubit subspace with the
The effective Hamiltonian in the interaction picture expressed as

$$H_2 = \frac{1}{2} \tilde{g} e^{i\phi} |11\rangle\langle02| + H.c.,$$

where $\tilde{g} = 2g_{AB} J_1 (\varepsilon/\nu)$ and $\tilde{\phi} = -\Phi + \pi/2$ are the effective coupling strength and phase of the parametric drive. Here, $g_{AB}$ is the static capacitive coupling strength between the two qubits and $J_1 (\varepsilon/\nu)$ is the 1st order Bessel function of the first kind. Similar to the single-qubit geometric gates with the Hamiltonian of Eq. 1, we can realize arbitrary geometric gates $U_1 (\theta, \gamma, \phi)$ in the subspace $\{|11\rangle, |02\rangle\}$ by modulating the effective coupling strength and phase in three time intervals. As a demonstration, we fix $\theta = 0$, resulting in two time intervals of the gate, and realize the geometric phase gate

$$U_2 (\gamma) = \begin{pmatrix} e^{i\gamma} & 0 \\ 0 & e^{-i\gamma} \end{pmatrix}$$

in the subspace $\{|11\rangle, |02\rangle\}$. When only considering the unitary in the two-qubit computational space $\{|00\rangle, |01\rangle, |10\rangle, |11\rangle\}$, the resulting unitary operation corresponds to a controlled-phase gate with an entangled phase $\gamma$:

$$U_2 (\gamma) = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \\ 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & e^{i\gamma} \end{pmatrix}.$$ (6)

The two-qubit geometric controlled-phase gate is performed with two sinusoidal modulation drives applied in series. Each has a square pulse envelope with sine squared rising and falling edges to suppress the adverse impact of sudden phase changes. The modulation frequency $\nu/2\pi = 268.2$ MHz and the modulation amplitude $\varepsilon/2\pi = 150$ MHz lead to an effective coupling strength $\tilde{g}/2\pi \approx 10$ MHz. Thus, the two-qubit controlled-phase gate is implemented within a duration of 112.8 ns. As an example, we here fix $\gamma = \pi$ and realize a controlled-Z (CZ) gate for the two qubits. We first use the two-qubit QPT method to benchmark the performance of the realized CZ gate, with the experimental sequence shown in Fig. 4(a). The experimentally reconstructed process matrix $\chi_{\text{exp}}$ is shown in Fig. 4(b) and indicates a process fidelity of 0.94 for the realized geometric CZ gate.

Besides, a two-qubit Clifford-based RB experiment is also performed to characterize the fidelity of the realized geometric CZ gate. The final measured ground state probability (sequence fidelity) decays as a function of the number of two-qubit Cliffords are displayed in Fig. 4(c) for both the two-qubit reference RB and CZ-interleaved RB experiments. The difference between the fitting results for the two experiments gives the geometric CZ gate fidelity of 0.977. This result is consistent with that from the two-qubit QPT method, when considering the state preparation and measurement error of about 0.03. The infidelity of the gate mainly comes from the qubit-qubit crosstalk, with the detailed error source analysis described in Ref. [59].

In conclusion, we experimentally demonstrate a nonadiabatic geometric operation scheme, and realize single-qubit geometric gates with an average fidelity of 0.9977. The noise-resilient feature of the realized single-qubit geometric gates is also verified by comparing the performances of both the geometric and dynamic gates with different errors. The distinct advantages of the realized geometric gates indicate they are promising candidates for robust quantum computation. In addition, a nontrivial two-qubit CZ gate is implemented with this geometric method by a parametric modulation of one qubit frequency. The resulted two-qubit CZ gate achieves a fidelity of 0.977 characterized by both QPT and RB methods. Therefore, the demonstrated universal geometric quantum gate set opens the door to implement high-fidelity quantum gates for robust geometric quantum computation.

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*Note added.–* While we were preparing our manuscript, we noticed a similar implementation of nonadiabatic single-qubit geometric gates with a superconducting qubit [64].
These three authors contributed equally to this work.
Electronic address: zyxue83@163.com
Electronic address: luyansun@tsinghua.edu.cn

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Supplementary Material for “Experimental implementation of universal nonadiabatic geometric quantum gates in a superconducting circuit”

Y. Xu,1,* Z. Hua,1,* Tao Chen,2,* X. Pan,1 X. Li,1 J. Han,1 W. Cai,1 Y. Ma,1 H. Wang,1 Y.P. Song,1 Zheng-Yuan Xue,2,∗ and L. Sun1,†

1Center for Quantum Information, Institute for Interdisciplinary Information Sciences, Tsinghua University, Beijing 100084, China
2Guangdong Provincial Key Laboratory of Quantum Engineering and Quantum Materials, GPETR Center for Quantum Precision Measurement, and School of Physics and Telecommunication Engineering, South China Normal University, Guangzhou 510006, China

I. EXPERIMENTAL SETUP

The sample in our experiment is a five-qubit chain device, which consists of five adjacent cross-shaped transmon (Xmon) qubits [1], arranged in a linear array with nearly identical nearest-neighbor coupling strengths. We place the experimental device inside a dilution refrigerator with a base temperature of about 10 mK. The detail of the device has been described in Refs [2, 3].

In the experiment, we have performed the geometric gates with only the first two capacitively coupled qubits QA and QB, whose main parameters are summarized and listed in Table S1. The other three qubits are biased far away from these two operation qubits and thus are nearly completely decoupled.

Details of the measurement circuitry of our experiment is shown in Fig. S1. The XY rotations of the two qubits are controlled by microwave pulses directly generated from a two-channel arbitrary waveform generator AWG5208. The last two channels of AWG5208 are sampling rate. The Z controls of the two qubits are performed through an inversion of the normalized qubit frequency response matrix

\[
M = \begin{pmatrix}
1 & -0.0759 \\
0.0800 & 1
\end{pmatrix}.
\]

Thus by performing the inversion of M to the flux bias controls, each qubit frequency can be independently controlled without changing the other qubit frequency.

With the quantum-limited JPA, both qubits can be readout individually and simultaneously with high fidelity in single-shot measurements. Due to the qubit thermal excitation and relaxation during measurements, there are non-negligible readout infidelities. In order to calibrate the two-qubit readout error, we prepare the system in each computational basis state and simultaneously measure the assignment probability \(\tilde{p} = (p_{00}, p_{01}, p_{10}, p_{11})^T\) of the two qubits. By repeating the experiments for all the two-qubit computational basis states, we obtain the two-qubit simultaneous readout assignment probability matrix \(\mathcal{P}\). Each column represents the two-qubit measurement probabilities after preparing the qubits in the corresponding basis state.

### Table S1: Device parameters of the two operating qubits.

| Parameters                      | QA         | QB         |
|---------------------------------|------------|------------|
| Readout frequency (GHz)         | 6.8386     | 6.8633     |
| Qubit frequency (GHz) (sweet spot) | 4.7813    | 5.1261     |
| Anharmonicity \((\alpha_A/2\pi, \alpha_B/2\pi)\) (MHz) | -202       | -190       |
| \(T_1\) (\(\mu s\)) (sweet spot) | 15.4       | 11.2       |
| \(T_2\) (\(\mu s\)) (sweet spot) | 9.86       | 16.3       |
| \(T_{2E}\) (\(\mu s\)) (sweet spot) | 17.3       | 20.9       |
| \(T_1\) (\(\mu s\)) (operating point) | 20.5       | 26.1       |
| \(T_2\) (\(\mu s\)) (operating point) | 1.73       | 4.86       |

\[\alpha_A = \alpha_B = 58.88 \text{ MHz}, \quad \alpha_A^* = \alpha_B^* = 58.88 \text{ MHz}\]

\[\frac{\pi}{2} \approx 16.68 \text{ MHz} \quad \text{and} \quad \frac{\pi}{2} \approx 0.09 \text{ MHz} \quad \text{and} \quad \frac{\pi}{2} \approx 0.88 \text{ MHz}\]

### Table S2: Two-qubit simultaneous readout assignment probability matrix \(\mathcal{P}\). Each column represents the two-qubit measurement probabilities after preparing the qubits in the corresponding basis state.

|     | 00     | 01     | 10     | 11     |
|-----|--------|--------|--------|--------|
| 00  | 0.9918 | 0.1058 | 0.1279 | 0.0131 |
| 01  | 0.0031 | 0.8905 | 0.0005 | 0.1137 |
| 10  | 0.0051 | 0.0006 | 0.8686 | 0.0890 |
| 11  | 0.0000 | 0.0032 | 0.0030 | 0.7842 |

*These three authors contributed equally to this work.
*Electronic address: yzyxue83@163.com
*Electronic address: luyansun@tsinghua.edu.cn
we obtain the $4 \times 4$ readout matrix $\mathcal{R}$ as shown in Table S2. Thus the readout errors can be corrected by multiplying the inverse of the readout matrix $\mathcal{R}$ with the measured probability $\tilde{p}$, such that $\tilde{p}^\text{corr} = \mathcal{R}^{-1} \cdot \tilde{p}$ represents the real occupation probabilities of the four computational basis states.

II. QUANTUM PROCESS TOMOGRAPHY

Both single- and two-qubit geometric gates are firstly characterized with quantum process tomography (QPT) method in the main text.

In the single-qubit QPT experiment with the experimental sequence shown in Fig. 1(c) of the main text, we first initialize the qubit with the following four states $\{ |0\rangle, |1\rangle, (|0\rangle + |1\rangle)/\sqrt{2}, (|0\rangle - i|1\rangle)/\sqrt{2} \}$, then apply the nonadiabatic geometric gates, and finally perform the state tomography measurements of the final states. The density matrix of the final two-qubit state is reconstructed from the state tomography measurements with 16 pre-rotations $\{ I, X/2, Y/2, X \}^\otimes 2$, where the basis operators $E_m$ and $E_n$ are chosen from $\{ I, \sigma_x, -i\sigma_y, \sigma_z \}$ with $\sigma_x$, $\sigma_y$, and $\sigma_z$ being Pauli operators.

The two-qubit QPT experiment is similar to that for the single-qubit case but with 16 initial states chosen from $\{ |0\rangle, |1\rangle, (|0\rangle + |1\rangle)/\sqrt{2}, (|0\rangle - i|1\rangle)/\sqrt{2} \}^\otimes 2$ instead. The density matrix of the final two-qubit state is reconstructed from the state tomography measurements with 16 pre-rotations $\{ I, X/2, Y/2, X \}^\otimes 2$.

III. RANDOMIZED BENCHMARKING

We also use another conventional method, Clifford-based randomized benchmarking (RB), to characterize the geometric single-qubit rotation gates and two-qubit CZ gate.

In the single-qubit RB experiment, we perform both the reference RB and interleaved RB experiments with the experimental sequences shown in the inset of Fig. 2 of the main text. In the reference RB experiment, we first apply a random sequence of $m$ quantum gates chosen from the single-qubit Clifford group ($C_1$), then append a recovery gate $C_r$ to invert the whole sequence, and finally measure the ground state probability as the sequence fidelity. The whole experiment is re-

FIG. S1: Details of wiring and circuit components.
peated for \( k = 50 \) different sequences to get the average sequence fidelity. In the interleaved RB experiment, a specific gate \( G \) is interleaved into the \( m \) random Cliffords, and a similar recovery gate is applied to invert the whole sequence. The experimentally measured sequence fidelity curves as a function of the number of Cliffords \( m \) for both reference RB and interleaved RB experiments are fitted to \( F = Ap^m + B \) with different sequence decays \( p = p_{\text{ref}} \) and \( p = p_{\text{gate}} \). The average gate fidelity is given by \( F_{\text{avg}} = 1 - (1 - p_{\text{ref}})(d - 1)/d.1.875, \) with \( d = 2^n \) for \( N \) qubits. The difference between the reference and interleaved RB experiments gives the specific gate fidelity \( F_{\text{gate}} = 1 - (1 - p_{\text{gate}}/p_{\text{ref}})(d - 1)/d. \)

The two-qubit Clifford-based RB experiment is similar, but with the random gates chosen from the two-qubit Clifford group \( (C_2) \) instead. For the interleaved CZ RB experiment, the geometric CZ gate is inserted into the random sequence.

IV. COMPARISON OF GEOMETRIC AND DYNAMIC GATES

We have demonstrated the robustness of the single-qubit geometric gates against two different types of errors (i.e., control amplitude error and qubit frequency shift-induced error) by comparing the geometric gates and the conventional dynamic gates with the same driving strength. The geometric gates are realized with different configuration settings, corresponding to two different geometric evolution trajectories.

In configuration A, the geometric gates are realized with a three-component microwave drive to generate a cyclic geometric evolution. The driving strengths and phases of the three time intervals are described as

\[
\begin{align*}
\int_0^{\tau_1} \Omega(t) dt &= \theta, \quad \phi = \phi - \frac{\pi}{2}, \quad t \in [0, \tau_1], \\
\int_{\tau_1}^{\tau_2} \Omega(t) dt &= \pi, \quad \phi = \phi + \gamma + \frac{\pi}{2}, \quad t \in [\tau_1, \tau_2], \\
\int_{\tau_2}^{\tau_3} \Omega(t) dt &= \pi - \theta, \quad \phi = \phi - \frac{\pi}{2}, \quad t \in [\tau_2, \tau_3].
\end{align*}
\]

We have shown the realized geometric gates in configuration A exhibit distinct advantages over the conventional dynamics gates with additional Rabi frequency error \( \epsilon \), as shown in the main text. However, we note that the geometric gates in configuration A do not always outperform the dynamic gates when introducing additional frequency detuning errors, as shown in Figs. S2(a-c).

In configuration B, the geometric gates are also realized with a three-component microwave drive, but with different phase in the second interval, described as

\[
\begin{align*}
\int_0^{\tau_1} \Omega(t) dt &= \theta, \quad \phi = \phi - \frac{\pi}{2}, \quad t \in [0, \tau_1], \\
\int_{\tau_1}^{\tau_2} \Omega(t) dt &= \pi, \quad \phi = \phi + \gamma - \frac{\pi}{2}, \quad t \in [\tau_1, \tau_2], \\
\int_{\tau_2}^{\tau_3} \Omega(t) dt &= \pi - \theta, \quad \phi = \phi - \frac{\pi}{2}, \quad t \in [\tau_2, \tau_3].
\end{align*}
\]

The unitary gate of this geometric path is the same as that of configuration A when \( \theta = \pi/2 \). We have demonstrated the noise-resilient feature of this gate with different frequency detuning errors in the main text. However, the geometric gates in this configuration do not always perform better than the dynamical gates with different Rabi frequency errors as well, as shown in Figs. S2(d-f).

Therefore, we conclude that the noise-resilient feature of the geometric gates depends on the types of errors and the cyclic evolution paths. Most time one may not find appropriate parameters to realize geometric gates against all kinds of errors, but one can still find a specific evolution path to realize noise-resilient geometric gates against the dominant error in the system.

V. CALIBRATION OF GEOMETRIC CZ

In order to achieve high-fidelity geometric CZ gate, a few experimental parameters should be carefully calibrated. In this section, we will describe in detail the calibration procedure of the realized two-qubit geometric CZ gate.

In our experiment, the two-qubit CZ gate is realized with a parametric modulation drive of one qubit frequency. After initializing the two qubits in \( |11\rangle \) state, we modulate the qubit frequency of \( Q_A \) with a sinusoidal form of

\[
\omega_A(t) = \omega_A + \epsilon \sin(v t + \Phi), \quad \text{(S4)}
\]

for different time durations, and finally measure the \( |11\rangle \) state probability of the two qubits. After repeating the above experiment with different modulation frequency \( v \), we can obtain a chevron oscillation pattern of the two qubits \( |11\rangle \) state probability as a function of both modulation frequency and time duration, from which the effective coupling strength \( g \) and the resonant transition frequency between \( |11\rangle \) and \( |02\rangle \) can be extracted.
In order to realize the geometric path evolution in \(|11\rangle, |02\rangle\) subspace, we apply two sinusoidal modulation drives in series and each one has a duration of \(\pi/\tilde{g}\), corresponding to a \(\pi\) rotation in \(|11\rangle, |02\rangle\) subspace. Therefore, in the two-qubit computational subspace \(|00\rangle, |01\rangle, |10\rangle, |11\rangle\), the resulted unitary gate can be expressed as \([5]\)

\[
U = \begin{pmatrix}
1 & 0 & 0 & 0 \\
0 & e^{i\phi_{01}} & 0 & 0 \\
0 & 0 & e^{i\phi_{10}} & 0 \\
0 & 0 & 0 & e^{i\phi_{11}}
\end{pmatrix},
\]

(S5)

where \(\phi_{01}\) and \(\phi_{10}\) are the single-qubit phases, and \(\gamma = \phi_{11} - \phi_{01} - \phi_{10}\) is the two-qubit entangled phase.

We characterize these single- and two-qubit phases by firstly initializing these two qubits in a product state \(|\psi_0\rangle = (|00\rangle + |01\rangle + |10\rangle + |11\rangle)/2\), then performing two sinusoidal modulation drives in series with a relative phase \(\Delta\Phi\) between them, and finally implementing the two-qubit state tomography to extract all these phases. The measurement results shown in Fig. S3(a) indicate that the two-qubit entangled phase \(\gamma\) is linearly related to the relative phase \(\Delta\Phi\), while the single-qubit phases \(\phi_{01}\) and \(\phi_{10}\) nearly remain the same. By choosing appropriate \(\Delta\Phi\), we thus realize a two-qubit CZ gate with the entangled phase \(\gamma = \pi\). The single-qubit phases can be compensated by rotating the reference frame in software.

We further verify the realized two-qubit CZ gate by performing different numbers of CZ gates, with the experimental result shown in Fig. S3(b). The measured single- and two-qubit phases are linearly related to the number of CZ gates. The small and finite offsets of the linear fits of these phases mainly come from the non-negligible crosstalk during the state preparation and measurement.

\[\Omega_{ZZ} = -\frac{2g^2_{AB}(Q_A + Q_B)}{(\omega_A - \omega_B - \alpha_A)(\omega_A - \omega_B + \alpha_B)} = 1.17 \text{ MHz},\]

(S6)

where \(\omega_A\) and \(\omega_B\) are the qubit frequencies of \(Q_A\) and \(Q_B\) at the operating spot, respectively. Thus this crosstalk interaction will induce an additional \(\sim 20^\circ\) phase shift of the \(|11\rangle\) state for an average single-qubit gate time of 40 ns, which vastly suppresses the two-qubit gate fidelity.

To further demonstrate the impact of ZZ crosstalk error in our system, we perform the simultaneous RB experiments for both qubits \(Q_A\) and \(Q_B\), in which we perform RB experiment on each qubit individually and operate both qubits simultaneously \([7]\). The experimental results are shown in Fig. S4. The crosstalk effect on \(Q_A\) when operating \(Q_B\) can be determined by firstly performing RB of \(Q_A\) alone [red markers in Fig. S4(a)], then benchmarking both qubits simultaneously and tracing out the contribution of \(Q_B\) [blue markers in Fig. S4(a)]. The difference of these two decay curves corresponds to the crosstalk-induced error on \(Q_A\) by controlling \(Q_B\). Similarly, we also find the crosstalk effect on \(Q_B\) when operating \(Q_A\), with the experimental result shown in Fig. S4(b). The extra infidelity when both qubits are operated simultaneously mainly causes the fidelity of the two-qubit reference RB to fall below expectation in the main text. Therefore, a tunable coupler that can switch off the coupling between adjacent qubits is desirable.

**VI. CROSSTALK-INDUCED ERROR**

The two-qubit CZ gate fidelity extracted from the two-qubit reference RB experiment in the main text is lower than expected, which mainly results from the two-qubit non-negligible ZZ crosstalk-induced error. With the two-qubit parameters in Table S1, we can estimate the phase shift rate of the \(|11\rangle\) state caused by the ZZ crosstalk with \([6]\)
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