Obtaining high-fidelity and robust quantum gates is the key for scalable quantum computation, and one of the promising ways is to implement quantum gates using geometric phases, where the influence of local noises can be greatly reduced. To obtain robust quantum gates, we here propose a scheme for quantum manipulation by combining the geometric phase approach with the dynamical correction technique, where the imperfection control induced X-error can be greatly suppressed. Moreover, to be robust against the decoherence effect and the randomized qubit-frequency shift Z-error, our scheme is also proposed based on the polariton qubit, the eigenstates of the light-matter interaction, which is immune to both errors up to the second order, due to its near symmetric energy spectrum. Finally, our scheme is implemented on the superconducting circuits, which also simplifies previous implementations. Since the main errors can be greatly reduced in our proposal, it provides a promising strategy for scalable solid-state fault-tolerant quantum computation.

Quantum computation$^1$, which uses superposition property of quantum states to perform computation, is believed to be in a position that can solve certain hard computational problems for classical computers. But, during any gate operation, which is the building block of a quantum computer, the noises and decoherence effect will be inevitable, leading to the infidelity of a target quantum gate. Thus, under the noise and decoherence effects, how to build an advisable quantum system for physical implementation of a quantum computer has attracted much attentions. Among the proposed candidates, superconducting quantum circuits system$^5$ has being pursued due to its flexible controllability and easy scalability. Currently, superconducting chip with dozens of transmon qubits$^6$ can be efficiently manipulated to show the quantum advantage$^7$.

Geometric phases, acquired during a cyclic evolution, was firstly discovered by Berry in the adiabatic process$^8$. Then, the Berry phase has been extended to the non-Abelian$^9$ and nonadiabatic$^{10}$ cases. As geometric phases only depend on the global geometric properties of their evolution paths, they are insensitive to certain local noises. Therefore, geometric phases can naturally find important applications in quantum computation$^{11,12}$, where quantum gates are implemented by using geometric phases. As the non-Abelian geometric phase is of the matrix form, it can naturally be used to form universal quantum gates, i.e., the holonomic quantum computation (HQC)$^{13}$. However, its physical implementation is experimentally difficult due to the need of complex interaction between multi-level quantum system$^{14,15}$. Using three level system, simplified nonadiabatic HQC (NHQC) has been proposed$^{16,19}$ and received many renewed theoretical$^{20,15}$ and experimental$^{21,22}$ interests. Recently, further explorations are mainly focus on strengthen the gate robustness$^{23,12}$ and further speedup the gate operation$^{24,25}$ in experimental accessible setups. However, currently, this is a still on-going exploration.

Here, we propose a scheme to implement NHQC, on decoherence-protected polariton qubits, with the dynamical correction technique (DCT)$^{26,27}$. The polariton qubits is formed by the eigenstates of the Jaynes-Cummings (JC) Hamiltonian, which can be implemented in a typical circuit QED setup$^{28}$. Meanwhile, the two eigenstates of the single-excitation subspace are chosen as our logical qubit states, which are decoherence insensitive. Moreover, besides the decoherence effect, the control induced Z-error in superconducting circuits mainly originates from randomized qubit-frequency drifts, which vary in a time scale that is much longer than the gate-time, and thus can be treated as a constant during a quantum gate. Remarkably, due to the near symmetric spectrum of the JC Hamiltonian, our logical qubit states are immune to the static Z-error, up to the second order. Note that, this merit is not shared by previous schemes based on the dressed-state qubits$^{29,30}$ as they have not used the dressed states from the same excitation subspace as logical qubit-states, and the dephasing protection there can only be achieved by further increasing the circuit complexity$^{31}$. From the implementation point of view, this is another merit of our scheme, as it is based on simplified setups and without auxiliaries as in Refs. $^{32,33}$. In addition, in order to be robust against the control induced X-error, from the imperfection control of the amplitude of the driving field and thus can be time-dependent during a quantum gate, we also incorporate the DCT. We show that this correction is better suitable for our polariton qubit than the conventional bare qubits. Note that, decoherence protection can also be obtained in Refs. $^{34,35}$, but the control-error insensitivity is not shared there. Besides, only the implementation of single-qubit gates are considered there. Therefore, our scheme provides a promising way to achieve fault-tolerant and scalable solid-state quantum computation.

Firstly, we present the implementation of our polariton qubit and show its decoherence-protection merit. Considering that a transmon qubit is capacitively coupled to a microwave...
The dressed states in the single-excitation subspace \( |\pm\rangle \) are chosen as our logic qubit states, denoted by \(|\pm\rangle\). In order, the noise induced shift of the eigenenergies are \( \delta E_{\pm}\approx \pm (h_x/g)^2 \times g/2 \ll h_x \). That is, for our polariton qubit, the noise effects from the decoherence of the transmon are suppressed to the second order, instead of the leading first order in the conventional case, where the energy shifts will be on the order of \( h_x/g \) for the \( \sigma_x \) noise.

Now, we proceed to present the holonomic manipulation on a polariton logical qubit, which can be achieved when suitable microwave drives applying on the transmon. In the subspace \( S_1 \), \( H_{JC} \) will be a diagonalized matrix with the elements being their eigenvalues, i.e.,

\[
H_0 = \begin{pmatrix}
E_0 & 0 & 0 \\
0 & E_{1-} & 0 \\
0 & 0 & E_{1+}
\end{pmatrix}.
\]

To manipulate the states in \( S_1 \), we consider applying a driven field on the transmon in the form of \( H_d = \sqrt{2f(t)}\sigma_x \). Here, to induce both the transitions between \( |G\rangle \) and \(|\pm\rangle\), the driven field can be chosen as

\[
f(t) = \Omega_1(t)\cos(\omega_1 t - \varphi_1) + \Omega_2(t)\cos(\omega_2 t - \varphi_2),
\]

where \( \Omega_n, \omega_n \) and \( \varphi_n \), with \( n = 1, 2 \), is the amplitude, frequency and phase of the \( n \)th tone of the driven field, respectively. As \( \langle G| \sigma_x^2 |\pm\rangle = (\pm|\pm\rangle|G\rangle = \pm 1/\sqrt{2} \)

\[
H_d = f(t)(|G\rangle \langle + | - |G\rangle \langle - |) + H.c.
\]

In order to investigate the qubit dynamics in a holonomic way, the driven field is set as \( \omega_1 = \omega_- \) and \( \omega_2 = \omega_+ \), with \( \omega_+ = E_{1+} - E_0 \) being the transition frequencies between the corresponding states. Thus, in the interaction picture with respect to \( H_0 \), the interaction Hamiltonian \( H_d \) will be

\[
H_1 = \Omega(t)\frac{e^{i\varphi_2}}{2} \left( \cos \frac{\theta}{2} |+\rangle - \sin \frac{\theta}{2} e^{i\varphi} |-\rangle \right) |G\rangle + H.c.,
\]

where \( \Omega(t) = \sqrt{\Omega_1(t)^2 + \Omega_2(t)^2} \), \( \tan(\theta/2) = \Omega_1(t)/\Omega_2(t) \) with \( \theta \) being a constant, and \( \varphi = \varphi_1 - \varphi_2 \). Here, to get the above interaction, we have set \( \omega_n \gg g \gg \Omega_n \), so that the rotating wave approximation (RWA) can be justified. In this way, the transitions within the artificial atom can be induced, which exhibits the bright and dark states of

\[
|b\rangle = \cos \frac{\theta}{2} |+\rangle - \sin \frac{\theta}{2} e^{i\varphi} |-\rangle,
\]

\[
|d\rangle = \sin \frac{\theta}{2} e^{i\varphi} |+\rangle + \cos \frac{\theta}{2} |-\rangle.
\]

As the dark state \(|d\rangle\) is decoupled, the quantum dynamics under Eq. (7) is captured by the resonant coupling between \(|b\rangle\) and \(|G\rangle\) states. When the cyclic evolution condition \( \int_0^T \Omega(t) dt = 2\pi \) is met, there is no population in state \(|G\rangle\), and thus only the bright state \(|b\rangle\) acquires a geometric phase factor \((\pi - \gamma)\). Consequently, one can obtain a quantum gate in the subspace spanned by \(|\pm\rangle\) depending on the parameters of \( \theta \) and \( \varphi \).
To achieve a fast universal set of single-qubit holonomic gates, single-loop way has been widely applied. Here, on this base, the DCT will also be applied, and then the total cyclic evolution consists of six pulses with different driving phases, i.e.,

\[
\begin{align*}
\int_{T_1}^{T_2} \frac{\Omega(t)}{2} dt &= \frac{\pi}{4} \varphi_1 = \phi + \pi, \varphi_2 = \pi, \quad (9a) \\
\int_{T_2}^{T_1} \frac{\Omega(t)}{2} dt &= \frac{\pi}{2} \varphi_1 = \phi + \varphi_1, \varphi_2 = \varphi_2, \quad (9b) \\
\int_{T_1}^{T_1} \frac{\Omega(t)}{2} dt &= \frac{\pi}{4} \varphi_1 = \phi + \pi, \varphi_2 = \pi, \quad (9c) \\
\int_{T_2}^{T_2} \frac{\Omega(t)}{2} dt &= \frac{\pi}{4} \varphi_1 = \phi + \varphi_2, \varphi_2 = \pi, \quad (9d) \\
\int_{T_3}^{T_3} \frac{\Omega(t)}{2} dt &= \frac{\pi}{4} \varphi_1 = \phi + \varphi_2, \varphi_2 = \pi, \quad (9e) \\
\int_{T_3}^{T_4} \frac{\Omega(t)}{2} dt &= \frac{\pi}{4} \varphi_1 = \phi + \varphi_2, \varphi_2 = \pi, \quad (9f)
\end{align*}
\]

The geometric path of this evolution process in the subspace \{\{b\},\{G\}\} is illustrated in Fig. 1(b). Finally, under the cyclic evolution condition, a certain pure geometric phase can be obtained in the \{b\}, which leads to a holonomic gate in the qubit-state space \{\{-\},\{+\}\} as

\[
U = \begin{pmatrix}
\cos \frac{\gamma}{2} - i \sin \frac{\gamma}{2} \cos \theta & -i \sin \frac{\gamma}{2} \sin \theta e^{i \varphi} \\
-i \sin \frac{\gamma}{2} \sin \theta e^{-i \varphi} & \cos \frac{\gamma}{2} + i \sin \frac{\gamma}{2} \cos \theta
\end{pmatrix}
= \exp \left(-i \frac{\gamma}{2} \cdot \sigma \right)
\]

which describes a rotation operation around the axis \(n = (\sin \theta \cos \varphi, \sin \theta \sin \varphi, \cos \theta)\) by a angle \(\gamma\), thus a universal single-qubit holonomic gate can be obtained. Since the parallel-transport condition \(\langle b(t)| H |d(t) \rangle = 0\) is always satisfied during the evolution and the cyclic condition is also met, \(\gamma\) is of the geometric nature.

The performance of our holonomic gate can be evaluated by considering the influence of dissipation using the the Lindblad master equation

\[
\dot{\rho} = i[\rho, H] + \frac{\kappa}{2} \mathcal{L}(a) + \frac{\Gamma_1}{2} \mathcal{L}(\sigma^-) + \frac{\Gamma_2}{2} \mathcal{L}(\sigma_z), \quad (11)
\]

where \(\rho\) is the density matrix of the system, \(H = H_0 + H_d\), \(\mathcal{L}(A) = 2A \rho A^\dagger - A^\dagger A \rho - \rho A^\dagger A\) is the Lindblad operator, \(\kappa\), \(\Gamma_1\) and \(\Gamma_2\) are decay rate of the cavity, decay and dephasing rate of the qubits, respectively. For example, a NOT and Hadamard gates can be realized by setting \((\gamma, \theta, \varphi) = (1, 1/2, 0)\pi\) and \((\gamma, \theta, \varphi) = (1, 1/4, 0)\pi\) respectively. Considering the limitation of the rotating wave approximation, the coupling strength is set as \(g = \omega_c/20\) with the cavity frequency being \(\omega_c = 2\pi \times 8 \text{ GHz}\), and \(\Omega(t) = \Omega_0 = g/20\). Within current technique, we set \(\kappa = 2\pi \times 0.1 \text{ KHz}\), \(\Gamma_1 = \Gamma_2 = \Gamma = 2\pi \times 4 \text{ KHz}\). Assuming the initial state being \(|+\rangle\), in Fig. 2 we plot the quantum state population and the corresponding fidelity dynamics for holonomic NOT and Hadamard gates, with the state-fidelity being 99.58% and 99.68%, respectively. Furthermore, average over the 2211 possible initial states of \(|\psi\rangle = \cos \theta' |+\rangle + \sin \theta' e^{i \varphi} |\rangle\rangle\) with uniform distributed \(\theta'\) and \(\varphi'\), and the different value are set to be 201 and 11, respectively; as the gate fidelity is sensitive to \(\theta'\) and insensitive to \(\varphi'\). Then, the gate-fidelity for the NOT and Hadamard gates can be obtained as 99.74% and 99.63%, respectively.

We further demonstrate the robustness of our scheme against the qubit-frequency shift Z-error and driving amplitude X-error, which will introduce additional term in the interaction Hamiltonian in the form of \(\Delta \Omega(t) \sigma_z/2 \) and \(\epsilon \Omega(t) \sigma_x\), respectively, with \(\Delta\) and \(\epsilon\) being the error fractions. As discussed before, the Z-error is assumed to be time-independent and the X-error can be time-dependent, during a holonomic gate on superconducting circuits. We numerically compare our scheme to general NHQC scheme in both error cases. For the Z-error, as shown in Fig. [3a], our scheme is more robust than the single-loop NHQC scheme without encoding. To fight-against the time-dependent X-error, we compare our scheme with DCT and choose \(\Omega(t) = \Omega_0 \sin^2(\pi t/T)\) for all pulses, the enhancement of the gate robustness in this case is shown in Fig. [3b]. As a result, the comparisons show that our scheme is more robust against both errors. We also note that, with the DCT, the gate-time will be prolonged, which will introduce more decoherence induced gate error, and thus there has a crosscut for the two curves in Fig. [3b]. However, with the increasing of the coherent times of the superconducting...
ing quantum circuit\textsuperscript{24}, this effect will be negligible small.

Next, we turn to the implementation of nontrivial two-qubit gates in a similar holonomic way, with two polariton qubits being coupled through their corresponding cavities, which simplified the previous implementation\textsuperscript{57–59} as an auxiliary logical-qubit from the JC coupling is needed to connect two data-logical-qubits there. In general case, the inter-cavity coupling strength \(J(t)\) can be in a time-dependent form\textsuperscript{52}, and the Hamiltonian of this system is

\[
H_{\text{coup}} = H_l + H_r + J(t) (a_l a_r^\dagger + a_r a_l^\dagger) \tag{12}
\]

where the \(H_l\) and \(H_r\) denote the free Hamiltonian in Eq. (4) for the left and right polariton qubits, \(a_{l/r}^\dagger\) and \(a_{l/r}\) denote the creation and annihilation operator for the left/right cavities. Corresponding to the single-qubit notation, our two-qubit gate need to be achieved in the subspace \(S_2 = \{|++, \rangle, |+-\rangle, |-+, \rangle, |--\rangle\}\) with \(|--\rangle = |\rangle_l \otimes |\rangle_r\).

Due to the exchange interaction of photons between two-qubit subspace and double excited subspace \(\{|2-, G\rangle, |2+, G\rangle, |G, 2-\rangle, |G, 2+\rangle\}\) can be driven, where \(|2-, G\rangle = |2-\rangle_l \otimes |G\rangle_r\). Thus, by setting an appropriate time-dependent function of \(J(t)\), only desired transitions can be induced.

As show in Fig. 4a), we implement the gate by using an ancillary state \(|2-, G\rangle\) in the subspace \(S_a = \{|2-, G\rangle, |\rangle\}\).

As \(\langle-\pm | a_1 a_2^\dagger |2-, G\rangle = 1/\sqrt{2}\), the inter-cavity coupling, in the interaction picture with respect to the free Hamiltonian of \(H_l + H_r\), is

\[
H_1 = \frac{1}{\sqrt{2}} \begin{pmatrix}
0 & J(t) e^{-i\phi_1} & J(t) e^{-i\phi_2} \\
J(t) e^{i\phi_1} & 0 & 0 \\
J(t) e^{i\phi_2} & 0 & 0
\end{pmatrix} \tag{13}
\]

where \(\omega_{\pm} = E_{2-} - G - E_{\pm}\). When setting \(J(t) = \sqrt{2} J_0 (t) \cos(\omega_{\pm} t - \phi_1) + J_0 (t) \cos(\omega_{-} t - \phi_2)\), choosing \(\omega'_{-} = \omega_{-}\), \(\omega'_{+} = \omega_{+}\), and letting \(\omega_0 \gg \{J_1, J_2\}\) so that the RWA is met, the above Hamiltonian will reduce to

\[
H_2 = \frac{J_0 (t)}{2} e^{i\phi_2} \left( \cos \frac{\theta}{2} |++\rangle + \sin \frac{\theta}{2} e^{i\phi} |--\rangle \right) \langle 2-, G | + \text{H.c.}, \tag{14}
\]

where \(J_c(t) = \sqrt{J_1(t)^2 + J_2(t)^2}\), \(\tan(\theta/2) = J_1(t)/J_2(t)\) with \(\theta\) being a constant, and \(\phi = \phi_1 - \phi_2 + \pi\). Besides, the leakage to other double-excitation subspace may also be occur, due to the oscillating nature of \(J(t)\), e.g., \(\langle++| a_1 a_2^\dagger |2+, G\rangle \neq 0\), which can also be suppressed by the RWA. Thus the harmful transitions will become high frequency oscillation term, which can be removed by RWA. In addition, the manipulation with DCT can still be useful to fight against X-error here, but it will twice the gate-time, which will lead to larger gate-infidelity comparing with the single-qubit gate case. So, in the two-qubit gate cases, we consider the implementation of the single-loop holonomic gates.

Eq. (13) has the same coupling configuration as Eq. (7), and thus holonomic quantum manipulation can be induced in \(S_a\) similar to the single-qubit case. When \(J_c\) is chosen to meet

\[
\int_0^{T_1} J_c(t) dt = \pi, \phi_1 = \phi, \phi_2 = \pi, \tag{15a}
\]

\[
\int_0^{T_2} J_c(t) dt = \pi, \phi_1 = \phi + \alpha - \pi, \phi_2 = \alpha, \tag{15b}
\]

and \(\int_0^T J_c(t) dt = 2\pi\) to form a cyclic evolution, the universal nontrivial two-qubit holonomic gate in \(S_2\) can be obtained as

\[
U_2 = \begin{pmatrix}
1 & 0 & 0 & 0 \\
0 & 1 & 0 & 0 \\
0 & 0 & \cos \frac{\omega_0}{2} - i \sin \frac{\omega_0}{2} \cos \phi & -i \sin \frac{\omega_0}{2} \sin \phi e^{i\phi} \\
0 & 0 & -i \sin \frac{\omega_0}{2} \sin \phi e^{-i\phi} & \cos \frac{\omega_0}{2} + i \sin \frac{\omega_0}{2} \cos \phi
\end{pmatrix}. \tag{16}
\]

Now as a typical example, by setting the parameters of \(\alpha = \pi, \vartheta = \pi/2\) and \(\phi = 0\), once can implement the CNOT gate. The devices are set up as \(\omega_{ql} = 2\pi \times 7.8 \text{ GHz}, \omega_{qr} = 2\pi \times 4.7 \text{ GHz}\) and \(g_{l/r} = \omega_{ql/r}/20\) for the left and right polariton qubit. The inter-cavity coupling strengths are chosen as \(J_1 = J_2 = 2\pi \times 5 \text{ MHz}\) to meet the RWA. The fidelity of this CNOT gate can be 99.37\% with the decoherence of \(\kappa = 2\pi \times 0.1 \text{ kHz}\) and \(\Gamma = 2\pi \times 4 \text{ kHz}\) for both polariton qubits. Furthermore, we plot the CNOT gate fidelity with different decoherence rates of the transmon-qubit, shown in Fig. 4b), which shows that as the decrease of the decoherence rates, holonomic gates with DCT and other optimization strategies can be more and more promising.

In summary, we propose a scheme to implement non-adiabatic holonomic quantum gates on polariton qubits. We have numerically demonstrated that our scheme is more robust compare to the general NHQC case in terms of robust against both Z- and X-errors. Remarkably, this enhancement does not increase the circuit complexity and without auxiliary, and thus simplifies previous explorations. Therefore, it provides a promising strategy toward robust and scalable solid-state quantum computation.

Data Availability Statement

The data that support the findings of this study are available from the corresponding author upon reasonable request.

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