Variable-frequency-controlled coupling in charge qubit circuits: Effects of microwave field on qubit-state readout

Xiao-Ling He,1,2 Yu-xi Liu,3,2 J. Q. You,1,2 and Franco Nori2,3,4
1Department of Physics and Surface Physics Laboratory (National Key Laboratory), Fudan University, Shanghai 200433, China
2Frontier Research System, The Institute of Physical and Chemical Research (RIKEN), Wako-shi 351-0198, Japan
3CREST, Japan Science and Technology Agency (JST), Kawaguchi, Saitama 332-0012, Japan
4Center for Theoretical Physics, Physics Department, Center for the Study of Complex Systems, The University of Michigan, Ann Arbor, MI 48109-1040, USA
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To implement quantum information processing, microwave fields are often used to manipulate superconducting qubits. We study how the coupling between superconducting charge qubits can be controlled by variable-frequency magnetic fields. We also study the effects of the microwave fields on the readout of the charge-qubit states. The measurement of the charge-qubit states can be used to demonstrate the statistical properties of photons.

I. INTRODUCTION

Superconducting quantum circuits are good candidates for implementing quantum information processing.1,2 To construct universal quantum computing, controllable couplings between any pair of qubits are required. Theoretical methods for switchable couplings in charge-qubit circuits have been proposed by changing the amplitude of the bias magnetic flux, e.g., in Refs.2,3,4. However, in experiments, it is much easier to produce fast and precise frequency shifts of the radio-frequency (RF) control signals, as opposed to changing the amplitude of the dc signal. Methods using variable-frequency-controlled couplings in superconducting flux-qubit circuits have been studied5 theoretically, comparing with the coupling approach using the dressed states6,7,8. In this scheme, the two qubits can be coupled to (or decoupled from) each other by modulating the frequencies5 of externally applied variable-frequency magnetic fields to match (or mismatch) the combination of frequencies of the two qubits. The coherent oscillations and conditional gate operations of two superconducting charge qubits with always-on coupling have been demonstrated9 experimentally. Therefore, the next step for charge qubits would be to design superconducting quantum circuits with switchable couplings.

Here, we first generalize our approach5 using the variable-frequency-controlled coupling in flux qubit circuits to the charge qubit circuit proposed in Ref.3. This proposal has the following advantages: (i) the coupling between different charge qubits can be implemented fast by changing the frequency of the externally applied classical field; (ii) these proposed charge qubits always work at their optimal points, and thus the qubits are mostly immune from charge noise10, produced by uncontrollable charge fluctuations; (iii) no additional circuit is needed to realize this controllable coupling.

Besides the controllable coupling, measuring the qubit state is also a very important step in quantum information processing. In superconducting quantum circuits, microwave fields are often used to implement quantum information processing. Here we focus on how microwave fields affect the readouts of the qubit states. In particular, we explore the effect of quantized fields with different statistical properties on measurement results of the qubit states when the charge qubits are placed inside a microcavity, e.g. a three-dimensional cavity11,12 or a superconducting transmission line13.

The paper is organized as follows: in Sec. II, we generalize the variable-frequency-controlled coupling approach in flux-qubit circuits3 to that in charge-qubit circuits4. In Sec. III, we study the effect of the classical and quantized microwave fields on the readout of the qubit states. In Sec. IV, we compare the classical and quantum treatment of the large Josephson junction. Finally, conclusions are presented in Sec. V.

II. HAMILTONIAN WITH VARIABLE-FREQUENCY CONTROLLED COUPLINGS

We first very briefly review the model Hamiltonian proposed in Ref.4 for two coupled superconducting charge qubits by sharing a large Josephson junction (JJ) (see Fig. 1). The large JJ is classically treated and its charge energy $E_{\text{cl}}$ is neglected. The dc biased magnetic field $\Phi$ is externally applied through the area between the large JJ and the first qubit. Each qubit is also biased by a dc voltage $V_{X_i}$ via the gate capacitance $C_i$ ($i = 1, 2$). The Hamiltonian of the superconducting circuit is

$$H = \sum_{i=1}^{2} \left[ E_i(V_{X_i}) - 2E_{J_i}\cos \left( \frac{n\Phi}{\Phi_0} - \frac{\gamma}{2} \right) \cos \varphi_i \right] - E_{\text{J0}} \cos \gamma$$

(1)

with $E_i(V_{X_i}) = E_{\text{cl}}(n_i - C_iV_{X_i}/2e)^2$. Here $E_{\text{cl}} = 2e^2/(C_i + 2C_{J_i})$ and $E_{J_i}$ are the charge and Josephson energies of the $i$th charge qubit. $E_{\text{J0}}$ is the Josephson energy of the large JJ. The number $n_i$ of excess Cooper pairs in the superconducting island is canonically conjugate to the average phase drop $\varphi_i = (\varphi_{iA} + \varphi_{iB})/2$ of the
FIG. 1: (Color online) Schematic diagram of two charge qubits coupled by a (left) large Josephson junction (JJ) with coupling energy $E_{J0}$ and capacitance $C_{J0}$. For the $i$th charge qubit (where $i = 1, 2$), a superconducting island (denoted by a filled circle) is connected to two identical small JJs (each with coupling energy $E_{Ji}$ and capacitance $C_{Ji}$). Also, this island is biased by the voltage $V_i = V_{X_i} + V_{x_i}(t)$ via a gate capacitance $C_i$, where $V_{X_i}$ is a static (dc) gate voltage and $V_{x_i}(t)$ is a time-dependent (ac) microwave gate voltage. Moreover, a static (dc) magnetic flux $\Phi$ plus a microwave-field-induced magnetic flux $\Phi(t)$ (ac) are applied to the (yellow) region between the large JJ and the first charge qubit.

The phase drop across the large JJ is $\gamma$. Considering that the critical current $I_0 \equiv 2\pi E_{J0}/\Phi_0$ of the large JJ is much larger than the critical currents $I_{ci} \equiv 2\pi E_{Ji}/\Phi_0$ of the charge qubits, the phase $\gamma$ across the large JJ is very small. We can expand the functions of the phase drop $\gamma$ in Eq. (1) into a series and and retain the terms up to second order in the parameters $\eta_i = (I_{ci}/I_0) < 1$. In this case, Eq. (1) can be reduced to

$$H = \sum_{i=1}^{2} \left[ \epsilon_i(V_{X_i})\sigma_z^{(i)} - \bar{E}_{Ji}\sigma_x^{(i)} - \chi_{12}\sigma_x^{(i)}\sigma_x^{(2)} \right]$$

in the spin-$\frac{1}{2}$ representation based on the charge states $|0\rangle \equiv |\uparrow\rangle$ and $|1\rangle \equiv |\downarrow\rangle$ that correspond to zero and one excess Cooper pairs in each Cooper-pair box, where $\epsilon_i(V_{X_i}) = \frac{1}{2}E_{ci}(C_iV_{X_i}/e - 1)$, and

$$\bar{E}_{Ji} = E_{Ji} \cos \left( \frac{\pi \Phi_e}{\Phi_0} \right) \left[ 1 - \frac{3}{8} \sin^2(\frac{\pi \Phi_e}{\Phi_0}) (\eta_i^2 + 3\eta_j^2) \right],$$

with $i, j = 1, 2 (i \neq j)$. The coupling constant $\chi_{12}$ between the two charge qubits is

$$\chi_{12} = L_J I_{c1} I_{c2} \sin^2 \left( \frac{\pi \Phi_e}{\Phi_0} \right),$$

where $L_J = \Phi_0/2\pi I_0$ is the Josephson inductance of the large JJ. It is clear that the coupling between the two qubits is realized via this effective inductance.

Now, we study how to apply our variable-frequency-controlled approach [4] to the above charge-qubit circuits [3]. We assume that besides the dc voltages $V_{X_i}$ and the dc magnetic flux $\Phi_e$, an ac microwave voltage $V_{\alpha i}(t) = V_{\alpha i} \cos (\omega_{\alpha i} t)$ with the frequency $\omega_{\alpha i}$ is applied to the superconduction island of the $i$th qubit via its gate capacitance, and an additional variable-frequency (ac) magnetic flux $\Phi(t) = \Phi_e \sin(\omega t)$ is also applied through the area between the large JJ and the first charge qubit (see Fig. 1). To make our proposed charge-qubit more immune from the uncontrollable charge fluctuations, it is also assumed that two charge qubits work at their optimal points, i.e., the applied dc voltages $V_{X_i}$ satisfy the condition $\epsilon_i(V_{X_i}) = 0$. Considering these conditions, the Hamiltonian in Eq. (2) becomes

$$H = \sum_{i=1}^{2} \left[ -\bar{E}_{Ji}\sigma_z^{(i)} + \epsilon_{0i}^{(i)} \cos(\omega_{\alpha i} t)\sigma_x^{(i)} \right] - \chi_{12}\sigma_x^{(i)}\sigma_x^{(2)}$$

and

$$\chi_{12} = L_J I_{c1} I_{c2} \sin \left( \frac{2\pi \Phi_e}{\Phi_0} \right) J_1(\varphi_e),$$

where

$$\xi = J_1(\varphi_e) \left[ 1 - \frac{3(\eta_i^2 + 3\eta_j^2)}{16} \left( 1 - \cos \left( \frac{2\pi \Phi_e}{\Phi_0} \right) J_0(\varphi_e) \right) \right]$$

and

$$g_{12} = L_J I_{c1} I_{c2} \sin \left( \frac{2\pi \Phi_e}{\Phi_0} \right) J_1(\varphi_e),$$

where $\varphi_e = 2\pi \Phi_e/\Phi_0$ and $J_n$ is the $n$th-order Bessel function of the first kind.

In the rotating reference frame at the frequency $\omega_{\alpha i}$ about $\sigma_z^{(i)}$, the Hamiltonian in Eq. (4) is rewritten as

$$H = \sum_{i=1}^{2} \left[ (\hbar \omega_{\alpha i} - \bar{E}_{Ji})\sigma_z^{(i)} + \epsilon_{0i}^{(i)} \sigma_x^{(i)} \right] - \chi_{12}\sigma_x^{(i)}\sigma_x^{(2)}$$

and

$$\xi = J_1(\varphi_e) \left[ 1 - \frac{3(\eta_i^2 + 3\eta_j^2)}{16} \left( 1 - \cos \left( \frac{2\pi \Phi_e}{\Phi_0} \right) J_0(\varphi_e) \right) \right]$$

$$+ \frac{3}{8} \cos^2 \left( \frac{\pi \Phi_e}{\Phi_0} \right) J_0(\varphi_e) J_1(\varphi_e) (\eta_i^2 + 3\eta_j^2).$$

Here, $\varphi_e = 2\pi \Phi_e/\Phi_0$ and $J_n$ is the $n$th-order Bessel function of the first kind.

To eliminate the $\sigma_z^{(i)}$ term in Eq. (5), the frequency $\omega_{\alpha i}$ of the microwave field applied to the gate capacitance is set as $\hbar \omega_{\alpha i} \simeq \bar{E}_{Ji}$. Furthermore, we can tune the flux $\Phi_e$ so that the coupling strength $\chi_{12}$ is less than the coupling strength $g_{12}$. Also, we tune the gate voltage $V_{\alpha i}$ so that the large detuning condition $\epsilon_{0i}^{(i)} - \epsilon_{0i}^{(1)} = \Delta \gg \chi_{12}$ can be satisfied. Under this condition, the always-on interaction $\chi_{12}$ is negligibly small, and the Hamiltonian in Eq. (5) is reduced [14] to

$$H \approx \sum_{i=1}^{2} \hbar \omega_{\alpha i} \sigma_z^{(i)} + \left[ g_{12}\sigma_x^{(1)}\sigma_x^{(2)} - \sum_{i=1}^{2} g_i \sigma_x^{(i)} \right] \sin(\omega t),$$

with $\hbar \omega_1 = \epsilon_{01}^{(1)} - \chi', \hbar \omega_2 = \epsilon_{01}^{(2)} + \chi', \chi' = \chi_{12}/2\Delta$. 

\[\text{(6)}\]
Let us discuss how the interaction between two qubits can be switched on and off via Eq. (6) by changing the frequency $\omega$ of the variable-frequency magnetic flux $\Phi(t) = \Phi_0 \sin(\omega t)$. Equation (6) shows that the two qubits are selectively rotated around the $z$-axis. When there is no applied ac magnetic flux $\Phi(t)$, the supercurrent $\hat{I}$ can be switched on and off via Eq. (6) by changing the frequency $\omega$ of $\Phi(t)$.

A single-qubit operation can also be implemented via the variable-frequency magnetic flux $\Phi(t)$. For example, if $\omega = \omega_1$, or $\omega = \omega_2$, then one qubit can be selectively rotated around the $x$ axis. When there is no variable-frequency magnetic flux, a rotation around the $z$-axis can be implemented for each qubit. Therefore, any logic gate (see, e.g., Ref. [12]) can be realized by using single-qubit operations and two-qubit operation via the Hamiltonians in Eqs. (7) and (8).

### III. EFFECT OF MICROWAVE FIELDS ON SUPERCURRENTS IN THE MEASUREMENT OF QUBIT STATES

Above, we have shown that the interaction between the two qubits can be switched on and off using a variable-frequency magnetic flux. Two-qubit operations can be implemented, and entangled states between two qubits can also be generated, using Eq. (7) or (8). To implement the readout of two-qubit states, we need to calculate the circulating supercurrent $I$ contributed by the two qubits $|g\rangle$. The operator of the supercurrent $I$ of the two qubits is given by

$$I = \sin\left(\frac{\pi \Phi_e}{\Phi_0}\right) \left( I_{c1} \sigma_x^{(1)} + I_{c2} \sigma_x^{(2)} \right) - \frac{1}{4\Phi_0} \sin\left(\frac{2\pi \Phi_e}{\Phi_0}\right) \left[ I_{c2}^2 + I_{c1}^2 + 2I_{c1} I_{c2} \sigma_x^{(1)} \sigma_x^{(2)} \right].$$

For any given state (e.g., $|\Psi\rangle$) of two qubits, the supercurrent can be obtained by

$$I = \langle \Psi | \hat{I} | \Psi \rangle.$$  

Note that two-qubit operations are always related to the microwave fields. The supercurrent $I$ might be different for different microwave fields with different statistical properties. Below, we study how the different microwave fields affect the supercurrent $I$.

### A. Classical microwave field

We now focus on two-qubit entangled states, created from the ground state $|g_1, g_2\rangle$ via the two-qubit interaction Hamiltonian in Eq. (7). For these created entangled two-qubit states, the contribution of the average values of single-qubit operators $\sigma_i^{(2)}$ ($i = 1, 2$) to the supercurrent is zero, and the supercurrent $I$ is only determined by the two-qubit operator $\sigma_x^{(1)} \sigma_x^{(2)}$ as follows

$$\langle \hat{I} \rangle = -\frac{\eta I_c}{2} \sin\left(\frac{2\pi \Phi_e}{\Phi_0}\right) \left(1 + \langle \sigma_x^{(1)} \sigma_x^{(2)} \rangle\right).$$

Here, for simplicity, the two qubits are supposed to have identical critical supercurrent $I_{c1} = I_{c2} = I_c$, and then $\eta_1 = \eta_2 = \eta$. The quantum fluctuation of the total supercurrent $I$ is

$$\langle \Delta \hat{I}^2 \rangle = \langle \hat{I}^2 \rangle - \langle \hat{I} \rangle^2,$$

which can be further given by

$$\langle \Delta \hat{I}^2 \rangle = \frac{\eta^2 I_c^2}{4} \sin^2\left(\frac{2\pi \Phi_e}{\Phi_0}\right) \left[1 - \langle \sigma_x^{(1)} \sigma_x^{(2)} \rangle^2 \right] + \frac{2I_c^2}{\Phi_e} \sin\left(\frac{2\pi \Phi_e}{\Phi_0}\right) \left(1 + \langle \sigma_x^{(1)} \sigma_x^{(2)} \rangle \right).$$

Considering that the ratio $\eta$ is small, the first term in Eq. (13) can be neglected in the following calculations. In this case, the supercurrent fluctuation has a similar behavior to the supercurrent $\langle \hat{I} \rangle$ of Eq. (11) and $\langle \Delta \hat{I}^2 \rangle \propto \langle \hat{I} \rangle$, when entangled two-qubit states are created from the ground state $|g_1, g_2\rangle$ through the Hamiltonian in Eq. (7). For convenience, we define a reduced quantity $\kappa$ to describe the supercurrent $\langle \hat{I} \rangle$ and supercurrent fluctuation $\langle \Delta \hat{I}^2 \rangle$ as

$$\kappa(\tau) = 1 + \langle \sigma_x^{(1)} \sigma_x^{(2)} \rangle = \frac{\langle \Delta \hat{I}^2 \rangle}{2I_c^2 \sin^2(\pi \Phi_e/\Phi_0)} = -2\langle \hat{I} \rangle \eta I_c \sin(2\pi \Phi_e/\Phi_0).$$

If two charge qubits are initially in an entangled state $\cos \theta |g_1, g_2\rangle + \sin \theta e^{i\phi} |e_1, e_2\rangle$ and the evolution of the two qubits is governed by the Hamiltonian in Eq. (7), then the reduced supercurrent or supercurrent fluctuation $\kappa(\tau)$ can be given by

$$\kappa_c = 1 + \sin(2\theta) \cos(\tau) \cos \phi + \cos(2\theta) \sin(\tau),$$

which means that $\kappa_c$ is an ac signal. Here, $\tau = t_0$, with the evolution time $t$. If the initial state is $\cos \theta |g_1, g_2\rangle + \sin \theta |e_1, e_2\rangle$, then $\kappa_c = 1 + \sin(\tau + 2\theta)$. When the evolution time $t_0 = -2\theta + 2n\pi - \pi/2$, $\kappa_c = 0$, which gives rise to $\langle \sigma_x^{(1)} \sigma_x^{(2)} \rangle = -1$. Thus, in this case both the total supercurrent and the supercurrent fluctuation become zero.
B. Quantized microwave field

Now let us consider the case when the variable-frequency magnetic flux \( \Phi_c \cos(\omega t) \) is replaced by a quantized magnetic flux, \( \Phi_q a^+ + \Phi_q^* a \), with frequency \( \omega = \omega_1 + \omega_2 \). Following the same way as the above derivation of Eq. (17), we can obtain an interaction Hamiltonian \( H_I \) between the quantized magnetic flux and the two charge qubits

\[
H_I = \xi_{12} a^{+} \sigma_-^{(1)} \sigma_-^{(2)} + \xi_{12} a \sigma_-^{(1)} \sigma_+^{(2)}
\]  

(16)

where

\[
\xi_{12} = -\frac{2\pi \Phi_q \Lambda J \cdot J_e}{\Phi_0} \sin \left( \frac{2\pi \Phi_e}{\Phi_0} \right).
\]

(17)

This model indicates that one photon can flip both qubits simultaneously.

We now consider that the two qubits are initially in the state \( \cos \theta |g, g\rangle + \sin \theta e^{i\phi} |e, e\rangle \) and the quantum field is initially in a state \( \sum_n D(n) |n\rangle \), here \( D(n) \) will be given below for a given state. From the Hamiltonian (16), the total system evolves to

\[
\Psi(\tau) = \sum_{n=0}^{\infty} \left[ a_n(\tau) |e, e, n\rangle + b_n(\tau) |g, g, n+1\rangle \right] + \cos \theta D(0) |g, g, 0\rangle,
\]

(18)

where

\[
a_n(\tau) = \cos(\tau \sqrt{n+1}) \sin \theta e^{i\phi} D(n) - \sin(\tau \sqrt{n+1}) \cos \theta D(n+1),
\]

\[
b_n(\tau) = \cos(\tau \sqrt{n+1}) \cos \theta D(n+1) + \sin(\tau \sqrt{n+1}) \sin \theta e^{i\phi} D(n),
\]

with the recaled dimensionless time \( \tau = \xi_{12} t \). Using Eq. (14) and Eq. (18), at the time \( \tau \), the reduced supercurrent expectation value or supercurrent fluctuation \( \kappa_q(\tau) \) in the case of the quantized field is

\[
\kappa_q(\tau) = 1+2 \text{Re} \left\{ w_0(\tau) \cos \theta D(0) + \sum_{n=0}^{\infty} [u_n(\tau) v_n(\tau)] \right\},
\]

(19)

where

\[
w_0(\tau) = \cos \tau \sin \theta e^{-i\phi} D^*(0) - \sin \tau \cos \theta D^*(1),
\]

\[
u_n(\tau) = \cos(\tau \sqrt{n+2}) \sin \theta e^{-i\phi} D^*(n+2) - \sin(\tau \sqrt{n+2}) \cos \theta D^*(n+1),
\]

Equation (19) shows that the supercurrent expectation value \( \langle \hat{I} \rangle \) consists of a dc component \(-(\eta J_e/2) \sin(2\pi \Phi_e/\Phi_0)\) and different ac components,

which are modulated by time-dependent factors, e.g. \( \cos(\tau \sqrt{n+1}) \).

We further specify that the quantized field is initially in several different quantum states (16), e.g., (i) the coherent state

\[
|\alpha\rangle = e^{-n/2} \sum_n \frac{\alpha^n}{\sqrt{n!}} |n\rangle,
\]

(20)

with \( \alpha = \sqrt{n} e^{i\phi} \); (ii) the superposition of two coherent states

\[
|\alpha_s\rangle = \frac{1}{N_+} (|\alpha\rangle + |\alpha\rangle) = \sum_n \frac{\alpha^{2n}}{\sqrt{(2n! \cosh b)}} |2n\rangle,
\]

(21)

with \( N_+ = \sqrt{2(1+e^{-2n})} \); and (iii) the squeezed vacuum state

\[
|0, \zeta\rangle = \sum_n \sqrt{\frac{(2n)!}{n! \cosh r}} [-e^{i\beta} \tanh(r/2)]^n |2n\rangle,
\]

(22)

with the squeezing parameter \( \zeta = r e^{i\beta} \) and \( \bar{n} = \sinh^2 r \). Here, \( \bar{n} \) is the average photon number. The photon number distributions \( P(n) \) (e.g., \( P(n) = |D(n)|^2 = |\langle n|\alpha\rangle|^2 \) for a coherent state) of the above three states are shown in Fig. 2. Physically, coherent states display Poissonian distribution and the fluctuations of both quadrature components are equal to the standard quantum fluctuation limit, \( \Delta X_1 = \Delta X_2 = 1/2 \). Squeezed states have sub-Poisson distribution and the fluctuation for one of the quadrature components can be squeezed, e.g., \( \Delta X_1 < 1/2 \).
If the qubits are initially in the ground state $|g_1, g_2\rangle$ and the quantized field is initially in the coherent state $|\alpha\rangle$, the reduced supercurrent expectation value or the supercurrent fluctuation is obtained from Eq. (19):

$$\kappa_q(\tau) = 1 - 2 \cos \varphi e^{-\bar{n}} \sum_{n=0}^{\infty} A(n) \sin(\tau \sqrt{n+1}) \cos(\tau \sqrt{n}),$$

(23)

with $A(n) = \bar{n}^{n+1}/(n! \sqrt{n+1})$. Equations (19) and (24) show that the supercurrent expectation value $\langle \hat{I} \rangle$ and the supercurrent fluctuation $(\Delta I)^2$ are very sensitive to the phase $\varphi$ of the coherent state. If $\varphi = \pi/2$, $\kappa_q$ has only a dc component; however, when $\varphi \neq \pi/2$, $\kappa_q$ consists of many different ac components.

If the qubits are initially in the ground state $|g_1, g_2\rangle$, but the quantized fields are initially in the squeezed vacuum states or superposition of coherent states, then from Eq. (19), $\kappa_q(\tau)$ is given by

$$\kappa_q(\tau) = 1 - 2 \sum_{n=0}^{\infty} \text{Re} \left[ B(n) \sin(\tau \sqrt{n+1}) \cos(\tau \sqrt{n}) \right],$$

(24)

with $B(n) = D^*(n+1) D(n)$. Because of $D(2n+1) = 0$, there are $B(n) = 0$ and $\kappa_q(\tau) = 1$. Thus, the oscillatory evolution disappears.

From Eq. (19), we know that the macroscopic supercurrent expectation value $\langle \hat{I} \rangle$ can be described by $\kappa_q$. Figure (3) shows that the supercurrent of the charge qubits are different with the same initial qubit state $(|g, g\rangle + |e, e\rangle)/\sqrt{2}$ but with different initial states of the quantum field. From Eq. (19), in the case of the coherent state $|\alpha\rangle$ with the phase $\varphi = \pi/2$, the total supercurrent $\langle \hat{I} \rangle$ displays a sinusoidal-like evolution, as shown in Fig. (3a). However, when $\varphi = 0$, the total supercurrent is shown in Fig. (3b). If the quantized field is initially in a superposition of coherent states, the total supercurrent $\langle \hat{I} \rangle$, as shown in Fig. (3c), demonstrates the collapse and partial-revival phenomena. In the case of the squeezed vacuum state, the total supercurrent approximately displays an ac current with a quasi-periodic evolution, which is demonstrated by Fig. (3d). All irregular oscillations of the supercurrent expectation or supercurrent fluctuation reflect the coherent interference that comes from the coherent superpositions of the different photon number states. The different initial photon states result in different output of the measurement of the charge-qubit states. Therefore, the measurement of the charge-qubit states can demonstrate the statistical properties of the photons, and charge qubits could be served as photon detectors.

**IV. QUANTIZATION TREATMENT ON LARGE JOSEPHSON JUNCTION**

In the Hamiltonian (11), the Josephson energy term $E_{c0}N^2$ of the large JJ is neglected and the large JJ acts as an effective inductance $L_J$ [4, 17]. We now consider a quantum mechanical treatment for the large JJ. Considering the additional term of charge energy $E_{c0}N^2$, the Hamiltonian of the large JJ can be written as

$$H_0 = E_{c0}N^2 - E_{j0} \cos \gamma,$$

(25)

with the charge energy $E_{c0}$ and the excess Cooper pairs $N$. Because the large JJ works in the phase regime, the spectrum of the large JJ is approximately equivalent to a harmonic oscillator $H_0 = \hbar \omega_p a^\dagger a$, with the plasma frequency

$$\omega_p = \frac{1}{\hbar} \sqrt{8E_f(0)E_c(0)}.$$ 

(26)

The bosonic operators $a$ and $a^\dagger$ are defined by

$$a = \zeta \gamma + i \frac{1}{2\zeta} N, \quad a^\dagger = \zeta \gamma - i \frac{1}{2\zeta} N,$$

(27)

and the phase drop $\gamma$ is expressed as

$$\gamma = \frac{1}{\zeta} (a^\dagger + a),$$

(28)

with $\zeta = (E_f(0)/2E_c(0))^{1/4}$. Due to the large critical supercurrent of the large JJ, one can expand the phase drop $\gamma$...
in Eq. (11) into a series and retain terms to the first order of $\gamma$. Finally, a spin-boson interaction between the two charge qubits and the large JJ is achieved:

$$H = \sum_{i=1}^{2} \left[ \varepsilon_i (V_{e i}) \sigma_z^{(i)} - E_{J i} \cos \left( \frac{\pi \Phi_e}{\Phi_0} \right) \sigma_x^{(i)} \right] + \hbar \omega p a^\dagger a + \sum_{i=1}^{2} \left[ g_{0i} \sigma_x^{(i)} (a^\dagger + a) \right],$$  \hspace{1cm} (29)

where $g_{0i} = - (E_{J i}/2\gamma) \sin (\pi \Phi_e/\Phi_0)$. We assume that the plasma frequency $\omega_p$ of the large JJ is much larger than the splitting of the qubits. Thus the large JJ is always in the ground state when the qubits are operated.

Following the standard technique of adiabatic elimination \[14\], we can eliminate the bosonic mode of the large JJ and obtain an effective interaction Hamiltonian between the two qubits: $\chi_{12} \sigma_z^{(1)} \sigma_x^{(2)}$, with the coupling strength $\chi_{12} = -2g_{01}g_{02}/\hbar \omega_p$. Using the expression of $g_{01}$, $g_{02}$, and $\omega_p$, one can easily confirm that this inter-qubit coupling is the same as that in Eq. (2). Here the large JJ serves as the data bus to virtually mediate the interaction between the two qubits. Therefore, the classical and quantum treatment to the large JJ are equivalent to each other. Generalizing the two-qubit system to the multi-qubit system, the effective inter-qubit coupling term reads $\sum_{i>j} \chi_{ij} \sigma_z^{(i)} \sigma_x^{(j)}$ with $\chi_{ij} = -2g_{0i}g_{0j}/\hbar \omega_p$. Because $g_{0i} = - (E_{J i}/2\gamma) \sin (\pi \Phi_e/\Phi_0)$, the coupling $\chi_{ij}$ is tunable by changing the static magnetic field $\Phi_e$ applied to the loop. We should point out that if the dc magnetic flux $\Phi_e$ is replaced by an ac variable-frequency magnetic flux $\Phi_e(t)$, then the qubit can be selectively coupled to the data-bus by a well-chosen frequency-matching condition between the qubit, data bus, and the variable-frequency magnetic flux.

V. CONCLUSIONS

In summary, we have studied a variable-frequency-control approach in charge-qubit circuits: the switchable coupling between the two charge qubits can be implemented by changing the frequency of the externally applied magnetic flux. Single-qubit operations can also be addressed and operated selectively. The charge qubits are chosen to work at their optimal points, so the effect of the noise, resulted from uncontrollable charge fluctuations, on the charge qubits is much suppressed. Moreover, the effects of the microwave field on the supercurrent of the two qubits are discussed. It is found that the supercurrent of the qubits significantly depends on the states of the microwave field. We also discuss the quantum treatment of the large JJ and find that both quantum and classical treatments are equivalent to each other. If the two-qubit circuit is generalized to many qubits, the interaction $\sigma_z^{(i)} \sigma_x^{(j)}$ can also be achieved.

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