Molecular ensemble-based remote quantum storage for charge qubit via a quasidark state

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We propose a quantum storage scheme independent of the current time-control schemes and study a "quantum data bus" (transmission line resonator) in a hybrid system consisting of a circuit QED system integrated with a cold molecular ensemble. Here, an effective interaction between the charge qubit and molecules is mediated by the off-resonant field in the data bus. Correspondingly, the charge state can be mapped into the collective quasispin state of the molecular ensemble via the standard dark-state-based adiabatic manipulation.

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I. INTRODUCTION

Quantum storage is crucial to quantum information processing [1–3]. Its importance rests with two issues: mapping the qubits with shorter coherence time into a system with relatively longer coherence time [4–6]; and storing the quantum states of the physical system difficultly operated into the states of the system manipulated feasibly [7]. Superconducting qubits, such as the charge qubits, flux qubits, and phase qubits, are hopeful candidates for the basic blocks in the architecture of quantum computer [8] since they exhibit feasibility to be well manipulated and well integrated [9–14]. Currently the decoherence of the superconducting qubits is relatively fast so that no enough logical operations can be performed for a practical quantum computing. Thus people expect that the superconducting qubits can be stored in some systems for which the decoherence time is long enough so that the stored qubits can be operated within enough time.

To overcome the disadvantages of superconducting qubits, a hybrid quantum processor has been suggested interfacing the solid state system and the molecular ensemble [15,16], in which the ensemble of cold polar molecules was used as a long-lived quantum memory storing the information of superconducting quantum qubits by collective spin state via microwave Raman processes in a high-Q stripline cavity [15]. Further more, a holographic quantum computing has been demonstrated using this protocol [16]. The conceptual designs [15,16] used the real photons as the data bus, and the corresponding schemes depended on an exact time control to switch photon-qubit and photon-molecule couplings effectively on and off. Based on the above considerations, to avoid the difficulty in exactly controlling evolution time, we study the data bus role of quantum transmission line. By adiabatically eliminating the variables of the transmission line as quantum data bus, we obtain quasidark state for the effective coupling between the molecular ensemble and the charge qubit. And then adiabatically manipulating the control field, the superconducting charge qubit can be mapped into the collective quasispin state of the molecular ensemble, and a reversing operation in a later time will regenerate the charge qubit state. During the adiabatic passage the transmission line is virtual excited and no real photons in the transmission line are exchanged between the charge qubit and the molecular ensemble. Recently, based on collective coupling of magnetic dipole interaction the cold polar molecular ensemble in the original proposal has been replaced by the spin ensemble [17], for which the coherence time could be longer. However, in our storage scheme, based on the quasidark state of the charge qubit and the molecular ensemble, the storage process has the advantage of being controllable with adiabatically modulating the control field.

The paper is organized as follows. In Sec. II we describe the circuit QED setup hybridized with an embedded molecular ensemble. In Sec. III, we demonstrate how the off-resonant coplanar cavity field induces the effective coupling of the charge qubit to the collective excitation of molecular ensemble and the existence of the quasidark state at parameter resonance. In Sec. IV, we show that using the quasidark state, quantum storage of the charge qubit state can be implemented through an adiabatic manipulation.

II. MODELING THE HYBRID SYSTEM IN TERMS OF COLLECTIVE EXCITATIONS

We consider a hybrid system shown in Fig. 1 that consists of a cold molecular ensemble, a Cooper-pair box (charge qubit) and a stripline cavity formed by a transmission line resonator [16,18]. The molecular ensemble consists of N identical and noninteracting molecules, each with a rotational ground state. The energy level configuration is illustrated in Fig. 1(b) with the electronically excited state $|a\rangle$ and the hyperfine splitting states of the rotational ground state, i.e., the ground state $|b\rangle$ and the metastable lower excited state $|c\rangle$. At the antinodes of quantum electromagnetic field in the stripline cavity, the polar molecules with large electric-dipole moment achieve the strong coupling, similar to that in the circuit QED system [19]. We place the Cooper-pair box and the molecular ensemble in two antinodes at a distance. In some proper conditions [20], the Cooper-pair box can be reduced to a two-level quantum system (charge qubit) with Jaynes-Cummings type coupling to the single mode cavity field. The Hamiltonian of the circuit QED system reads

$$H_1 = \frac{\omega_g}{2}\sigma_z + \omega a^\dagger a + g(a\sigma_+ + \text{H.c.}), \tag{1}$$

where $a(a^\dagger)$ is the quantum cavity field annihilation and (creation) operator, $\sigma_-(\sigma_+)$ is the lowering (raising) operator, and

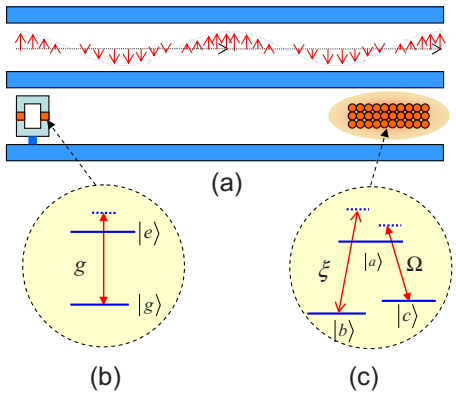

FIG. 1. (Color online) Schematic diagram of a hybrid system where the cold molecular ensemble is coupled to a circuit QED system consisting of a Cooper-pair box and a stripline cavity in (a). The Cooper-pair box [as a two level system (b)] and cold molecular ensemble are at the adjacent antinodes in the stripline cavity. For the molecular ensemble, each molecule (c) with a rotational ground state has the hyperfine splitting ground states $|b\rangle$, $|c\rangle$, and the electronically excited state $|a\rangle$. Here the level $|a\rangle$ is coupled to $|b\rangle$ and $|c\rangle$ by a quantum field in stripline cavity and a classical control field, respectively.

g is the strength of the interaction between the stripling cavity field and the Cooper-pair box. The parameters ω_g and ω denote the frequency of the charge qubit and the quantum cavity field, respectively. Initially we can prepare the polar molecule in the ground state $|b\rangle$, which is coupled to the excited state $|a\rangle$ by the quantum electromagnetic field in the stripline cavity. While the coupling of the metastable excited state $|c\rangle$ to the excited state $|a\rangle$ is realized through the classical control field with the Rabi frequency Ω. Thus the Hamiltonian for the coupled molecular ensemble and stripline cavity can be written as

$$H_2 = \sum_{j=1}^{N} (\omega_a |a\rangle_{jj}\langle a| + \omega_c |c\rangle_{jj}\langle c|) + \omega a^\dagger a + \sum_{j=1}^{N} (\xi a|a\rangle_{jj}\langle b| + \Omega e^{-i\omega_f t}|a\rangle_{jj}\langle c| + \text{H.c.}), \tag{2}$$

where ξ is the interaction strength of the quantum cavity field to the molecule, ω_a and ω_c are the molecule's energy level spacings with respect to the excited state and the metastable excited state respectively, and ω_f is the frequency of the control field.

Here we adopt the quasispin wave approach developed in Ref. [7] and invoke the bosonic operators

$$A = \frac{1}{\sqrt{N}} \sum_{j=1}^{N} |b\rangle_{jj}\langle a|, \tag{3}$$

$$C = \frac{1}{\sqrt{N}} \sum_{j=1}^{N} |b\rangle_{jj}\langle c| \tag{4}$$

to describe the collective excitation of the ensemble. To this end we rewrite the above Hamiltonian $H=H_0+H_I$ in terms of the free part

$$H_0 = \frac{\omega_g}{2}\sigma_z + \omega a^\dagger a + \omega_a A^\dagger A + \omega_c C^\dagger C \tag{5}$$

and the interaction part

$$H_I = g(a\sigma_+ + \text{H.c.}) + \zeta(aA^\dagger + \text{H.c.}) + \Omega(e^{-i\omega_f t}A^\dagger C + \text{H.c.}), \tag{6}$$

where $\zeta=\xi\sqrt{N}$. In the large N limit with the low excitation condition that there are only a few atoms occupying the level $|a\rangle$ or $|c\rangle$ [21], the operator $A(C)$ behaves as a boson which satisfies the bosonic commutation relations

$$[A,A^\dagger]=1, \quad [C,C^\dagger]=1. \tag{7}$$

In the same limit, we also have

$$[A,C]=0, \quad [A,C^\dagger]\to 0, \tag{8}$$

which means that two boson modes described by A and C are independent with each other.

III. INDUCED EFFECTIVE COUPLING OF CHARGE QUBIT TO COLLECTIVE EXCITATIONS OF MOLECULAR ENSEMBLE

We emphasize again that the current hybrid system-based quantum storage protocol requires an exact time control for couplings of a charge qubit to photons and photons to a molecular ensemble. Namely, one can switch on (off) the couplings exactly in an arbitrarily given instance. Such a time control process includes two steps: (a) couple the photon to the charge qubit first; (b) then switch off the previous coupling and switch on the coupling of the photon to the molecular ensemble again. Due to the "time-energy uncertainty," this kind of time control protocol would inevitably result in experimental errors. It is much appreciated that we can avoid these errors. With these considerations, we detune stripline cavity field slightly from both the charge qubit and the molecular ensemble. Thus the effective interaction between the charge qubit and the molecular ensemble will be induced by eliminating photon degrees of freedom.

Initially, we prepare the microwave field in the vacuum state then the interaction between the charge qubit and the molecular ensemble is induced by the virtual excitation of the transmission line field with no exchange of real photons between the charge qubit (molecular ensemble) and the field. The process can be schematically shown in Fig. 2. Here both the qubit and the molecular ensemble interact simultaneously with the continuum of the cavity field whose spectral structure is illustrated in the top of Fig. 2. We will show as follows that under some conditions related to the large detuning the virtual photon transitions happen between the qubit (molecular ensemble) and the transmission line cavity. The photon-assisted transitions cannot happen in practice, but the effective interaction between qubit and the molecular ensemble is induced, which is similar to the conventional contact interaction. Therefore, our present protocol does not need an exact time control for the coupling of the photon qubits to the charge qubit (to the molecular ensemble).

Next we study in details the virtual photon exchange process by deriving the effective Hamiltonian through a canoni-

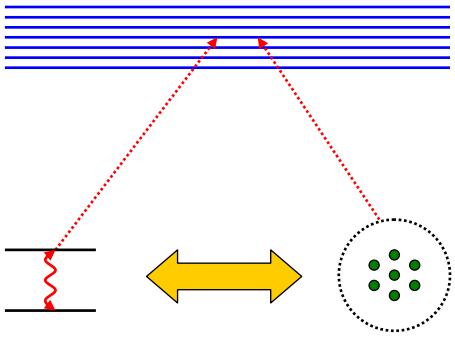

FIG. 2. (Color online) Effective interaction between the charge qubit and the molecular ensemble induced by the virtual photon exchange.

cal transformation. In the rotating frame, in terms of a time-dependent unitary operator $U=\exp(i\omega_f t C^\dagger C)$, we get a time-independent Hamiltonian

$$H^R = U^\dagger H U - iU^\dagger \partial_t U. \tag{9}$$

This time-independent Hamiltonian (9) can be decomposed into tow parts, $H^R=H_0+H_I$, including the "free" part Hamiltonian,

$$H_0 = \frac{\omega_g}{2}\sigma_z + \omega a^\dagger a + \omega_a A^\dagger A + (\omega_c+\omega_f)C^\dagger C + \Omega(A^\dagger C + \text{H.c.}) \tag{10}$$

and the interaction part Hamiltonian,

$$H_I = g(a\sigma_+ + \text{H.c.}) + \zeta(aA^\dagger + \text{H.c.}), \tag{11}$$

which describes the coupling of stripline cavity field to both the charge qubit and the molecular ensemble.

Using Fröhlich's transformation [22–24], we can obtain the effective Hamiltonian

$$H_{eff} = e^{-S}H^R e^S \cong H_0 + \frac{1}{2}[H_I, S], \tag{12}$$

where S is an anti-Hermitian operator satisfying

$$H_I + [H_0, S] = 0. \tag{13}$$

The corresponding anti-Hermitian operator S for Fröhlich transformation adopts the following form:

$$S = (\eta_1 a\sigma_+ + \eta_2 aA^\dagger + \eta_3 aC^\dagger) - \text{H.c.}, \tag{14}$$

where the coefficients η_1, η_2, and η_3 are determined by Eq. (13)

$$\eta_1 = \frac{g}{\Delta_1}, \tag{15}$$

$$\eta_2 = \frac{\zeta\Delta_3}{\Delta_2\Delta_3 - \Omega^2}, \tag{16}$$

$$\eta_3 = \frac{\zeta\Omega}{\Delta_2\Delta_3 - \Omega^2}, \tag{17}$$

with the detunings

$$\Delta_1 = \omega - \omega_g, \Delta_2 = \omega - \omega_a,$$

$$\Delta_3 = \omega - (\omega_c + \omega_f). \tag{18}$$

In order to obtain the effective coupling between the charge qubit and the molecular ensemble, we consider the lightly off resonant case and take appropriate experimental conditions so that the coefficients η_1, η_2, and η_3 are all as small as enough, i.e., $|\eta_1|$, $|\eta_2|$, $|\eta_3| \ll 1$, which are also requisite by the usual second-order perturbation theory. We assume that the microwave cavity field is initially in the vacuum state, i.e., $\langle a^\dagger a\rangle = 0$. Averaged over the vacuum state of the photon field, in the two photon resonance case, the effective Hamiltonian (12) in the interaction picture reads

$$H_{int} = \Delta B^\dagger B + (g_m \sigma_- B^\dagger + \Omega_d B^\dagger D + \text{H.c.}), \tag{19}$$

where the bosonic operators B and D are defined as

$$B = \alpha A + \beta C, \quad D = \alpha C - \beta A, \tag{20}$$

with the coefficients

$$\alpha = \frac{-(\eta_2 g + \eta_1\zeta)}{\sqrt{(\eta_3 g)^2 + (\eta_2 g + \eta_1\zeta)^2}}, \tag{21}$$

$$\beta = \frac{-(\eta_3 g)}{\sqrt{(\eta_3 g)^2 + (\eta_2 g + \eta_1\zeta)^2}}. \tag{22}$$

The effective Hamiltonian (19) shows that the photon degrees of freedom is effectively eliminated and it mediates an effective interaction between the charge qubit and the molecular ensemble. The two photon resonance condition is

$$\omega_B - \omega_g' = \omega_B - \omega_D = \Delta, \tag{23}$$

where ω_B and ω_D are the effective energy level spacings with respect to the states of B-excitation and D-excitation, respectively, and ω_g' is the effective frequency of the charge qubit after the Fröhlich transform. The complicated coefficients in H_{int} are listed as follows:

$$g_m = \frac{1}{2}\sqrt{(\eta_3 g)^2 + (\eta_2 g + \eta_1\zeta)^2}, \tag{24}$$

$$\omega_g' = \omega_g - \eta_1 g, \tag{25}$$

$$\omega_B = (\omega_a - \eta_2\zeta)\alpha^2 + (\omega_c + \omega_f)\beta^2 + 2\left(\Omega - \frac{1}{2}\eta_3\zeta\right)\alpha\beta, \tag{26}$$

$$\omega_D = (\omega_a - \eta_2\zeta)\beta^2 + (\omega_c + \omega_f)\alpha^2 - 2\left(\Omega - \frac{1}{2}\eta_3\zeta\right)\alpha\beta, \tag{27}$$

$$\Omega_d = \alpha\beta[(\omega_c + \omega_f) + \eta_2\zeta - \omega_a] + \left(\Omega - \frac{1}{2}\eta_3\zeta\right)(\alpha^2 - \beta^2). \tag{28}$$

In order to find invariant subspaces of H_{int}, we define the dark state operator

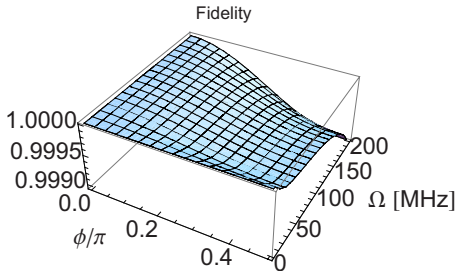

FIG. 3. (Color online) Fidelity during the storage process. For an arbitrary qubit to be stored (ϕ from 0 to $\pi/2$), when the Rabi frequency of the control field adiabatically varies from 0 to 200 MHz, the fidelity is larger than 0.9990.

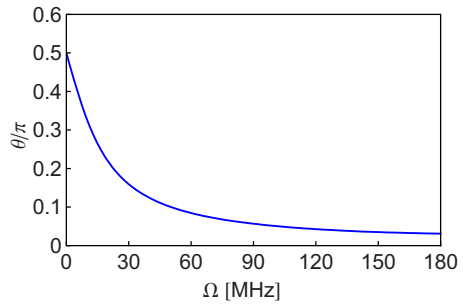

FIG. 4. (Color online) Variance of the mixing angle θ by adiabatic modulation of the control field with the Rabi frequency Ω from 0 to 180 MHz.

$$F = \sigma_- \cos\theta - D \sin\theta, \tag{29}$$

with $\theta(t)$ satisfying $\tan\theta(t) = g_m/\Omega_d(t)$. It mixes the states of the qubit excitation and the collective molecular excitations. We introduce the composite vacuum state $|\mathbf{0}\rangle = |\mathbf{b}\rangle \otimes |g\rangle$, where $|\mathbf{b}\rangle$ is the ground state with all N molecules staying in the same single-particle ground state $|b\rangle$ [7]. Evidently an eigenstate of H_{int} with vanishing eigenvalue is constructed as

$$|\Psi\rangle = F^\dagger |\mathbf{0}\rangle. \tag{30}$$

We notice that $|\Psi\rangle$ is not an ideal dark state, actually it is a superposition state of the A excitation concerning the excited state which will decay during the usual manipulation process. In this sense, we name it quasidark state.

To construct other eigenstates of the interaction Hamiltonian, we introduce the operators as follows:

$$Q_\pm = \sqrt{\frac{\Theta \pm \Delta}{2\Theta}} B \pm \sqrt{\frac{\Theta \mp \Delta}{2\Theta}} (\sigma_- \sin\theta + D \cos\theta), \tag{31}$$

where

$$\Theta = \sqrt{\Delta^2 + 4(g_m^2 + \Omega_d^2)}. \tag{32}$$

Using these operators, we can construct the eigenstates

$$|\Phi_\pm\rangle = Q_\pm^\dagger |\mathbf{0}\rangle, \tag{33}$$

with the corresponding eigenvalues

$$E_\pm = \pm \sqrt{\frac{(\Theta \pm \Delta)(g_m^2 + \Omega_d^2)}{(\Theta \mp \Delta)}}. \tag{34}$$

IV. QUASIDARK STATE-BASED ADIABATIC MANIPULATION FOR QUANTUM STORAGE

In this section, we will show how to store the quantum information of the charge qubit in the molecular ensemble by using the above quasidark state. In this quantum storage scheme, adiabatically modulating the control field, the charge qubit state can be mapped into the molecular ensemble under appropriate experimental conditions.

Now we consider whether the charge qubit state can be mapped into the molecular ensemble effectively. In the storage process, there exist other eigenstates with nonzero eigenvalues besides the quasidark state. To avoid the system to transit to other eigenstates and keep itself within the quasidark state, the manipulation of the control field should satisfy the adiabatic condition $|\langle\Phi_\pm|\partial t|\Psi\rangle/E_\pm| \ll 1$, i.e.,

$$\frac{g_m(\Theta + |\Delta|)}{\sqrt{\Theta(\Theta - |\Delta|)}(g_m + \Omega_d)^3} |\dot{\Omega}| \ll 1, \tag{35}$$

according the quantum adiabatic theorem [25,26]. With the adiabatic decrease in the control field, the mixing angle will increase from 0 to $\pi/2$.

From the expressions Eqs. (21) and (22) of α and β, there is

$$\frac{\alpha}{\beta} = \frac{\Delta_3}{\Omega} + \frac{\Delta_2}{\Omega}\frac{\Delta_3}{\Delta_1} - \frac{\Omega}{\Delta_1}. \tag{36}$$

Here we consider the case of "blue detuning." We assume that the Rabi frequency of the control field is much smaller than the frequency detunings, i.e., $\Omega \ll \Delta_1$, Δ_2, Δ_3, and $\Delta_1 \approx \Delta_3$ which is also required by the two photon resonance condition (23). The values of α and β satisfy $|\alpha| \gg |\beta|$.

Because the composite vacuum state $|g\rangle \otimes |\mathbf{b}\rangle$ is also an eigenstate of the interaction Hamiltonian, if the mixing angle θ changes gradually from 0 to $\pi/2$ by adiabatically decreasing the control field to zero, an arbitrary superposition of charge qubit state

$$|\varphi\rangle = \cos\phi|g\rangle + \sin\phi|e\rangle \tag{37}$$

will undergo the following evolution through the quasidark state during this process,

$$|\varphi\rangle \otimes |\mathbf{b}\rangle \to |\psi\rangle_\theta, \tag{38}$$

with

$$|\psi\rangle_\theta = \sin\phi \cos\theta|e\rangle|\mathbf{b}\rangle + |g\rangle \otimes [\cos\phi|\mathbf{b}\rangle - \sin\phi(\alpha \sin\theta|\mathbf{c}\rangle - \beta \sin\theta|\mathbf{a}\rangle)], \tag{39}$$

where $|\mathbf{a}\rangle = A^\dagger|\mathbf{b}\rangle$ and $|\mathbf{c}\rangle = C^\dagger|\mathbf{b}\rangle$. During the adiabatically modulating process, the values of α and β vary depending on the Rabbi frequency Ω. From relation (36), with the control field decreasing to 0, we can see that $\alpha \to -1$ and $\beta \to 0$. Then we have

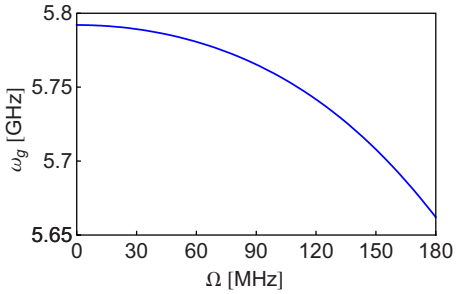

FIG. 5. (Color online) Modulation of the frequency of the charge qubit with the change in the control field.

$$|\psi\rangle_{\pi/2} = |g\rangle \otimes (\cos\phi|\mathbf{b}\rangle + \sin\phi|\mathbf{c}\rangle). \tag{40}$$

This means that we can transfer the information of the charge qubit state into the molecular ensemble without any real photon exchange. Different from the usual dark state storage, in our storage process the unwanted noise term (collective excitation state $|\mathbf{a}\rangle$) emerges. By the above theoretical analysis that $|\alpha| \gg |\beta|$, the effect of the noise caused by the emergence of the collective A excitation is very small. To further illustrate this point, we consider the fidelity (see, e.g., Fig. 3) by some experimental data

$$F = |{}_\theta\langle\psi'|\psi\rangle_\theta|^2, \tag{41}$$

with

$$|\psi'\rangle_\theta = \sin\phi\cos\theta|e\rangle|\mathbf{b}\rangle + |g\rangle \otimes (\cos\phi|\mathbf{b}\rangle + \sin\phi\sin\theta|\mathbf{c}\rangle), \tag{42}$$

which corresponds to the evolution of the initial state when we use the dark state to replace the quasidark state. From the Refs. [16,19,27], we choose the experimental data: $g/2\pi = 110$ MHz, $\zeta/2\pi = 50$ MHz, $\omega/2\pi = 6.044$ GHz, $(\omega_c + \omega_f)/2\pi = 5.744$ GHz, and $\omega_a/2\pi = 5.544$ GHz. Just as shown in Fig. 3, the fidelity can reach 0.9990.

Next we will show how the coefficients change with respect of the control field in details by applying these experimental data. The magnitudes of effective coupling g_m vary from 14.8 to 16.4 MHz with the Rabi frequency Ω from 0 to 180 MHz. The mixing angle θ changes from $\pi/2$ to 0 by adiabatically modulating the control field with the Rabi frequency Ω varying from 0 to 180 MHz, (see, e.g., Fig. 4). The frequency of charge qubit needs to be modulated (see, e.g., Fig. 5) with the variance of the control field to satisfy the resonance condition (23).

The decoherence time of the CPB is about 1 μs [10,20] and about 4 μs [28] when replacing the traditional Cooper pair box with a transmon design [29]. And the loss rates of the stripline cavity can be as low as $2\pi \times 5$ kHz [27]. Thus the implementation time of the state transfer is much less than the decoherence time of the charge qubit and the cavity. All these results indicate that the charge qubit can be mapped into the molecular ensemble effectively without any real photon exchange.

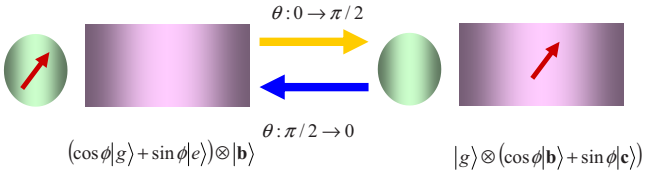

FIG. 6. (Color online) Mapping the charge qubit into the molecular ensemble by virtual photon transition.

This storage process can be schematically shown in Fig. 6. Here, rotating the mixing angle θ from 0 to $\pi/2$ by changing the strength of the control field adiabatically, the information of the charge qubit state is stored in the molecular ensemble. And the reverse manipulation will make the molecular ensemble release the information to the charge qubit again.

V. CONCLUSIONS

In summary, we have explored and studied a quantum storage scheme for a charge qubit, which is based on the virtual photon exchange between the charge qubit and the collective excitations of polarized molecules. Our investigation concerns a hybrid system consisting of a superconducting circuit QED system and an ensemble of polarized molecules. Through the adiabatic elimination using Fröhlich's transformation, we explicitly derive an effective coupling between the charge qubit and the molecular ensemble. It shows that the quasidark state can be dynamically produced as an entanglement of the charge qubit plus the molecular ensemble, which means that the qubit state can directly be mapped into the collective excitation state of the molecular ensemble without middle time control. The corresponding numerical results show the feasibility of this experimental scheme.

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