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Engineering quantum decoherence of charge qubit via a nanomechanical resonator

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Abstract. We propose a theoretical scheme to observe the loss of quantum coherence through the coupling of the superconducting charge qubit system to a nanomechanical resonator (NAMR), which has already been successfully fabricated in experiment and is convenient to manipulate. With a similar form to the usual cavity QED system, this qubit-NAMR composite system with engineered coupling exhibits the collapse and revival phenomenon in a progressive decoherence process. Corresponding to the two components of superposition of the two charge eigenstates, the state of the nanomechanical resonator evolves simultaneously towards two distinct quasi-classical states. Therefore the generalized "which way" detection by the NAMR induces the quantum decoherence of the charge qubit.

PACS. 03.65.-w Quantum mechanics – 74.50.+r Tunneling phenomena; point contacts, weak links, Josephson effects – 03.67.Lx Quantum computation – 85.25.Dq Superconducting quantum interference devices (SQUIDs)

1 Introduction

It is well-known that superposition of quantum states lies at the very heart of modern quantum theory. In an ideal situation, the quantum coherence implied by this superposition results in various fantastic phenomena [1]. However, real systems are never isolated completely from the surrounding environment. Interaction with external systems (such as the environment) leads to entanglement between them. The phase coherence of quantum system is then destroyed [2]. This explains why quantum superposition does not appear in the macroscopic world: this is the transition from the quantum to the classical world [3].

This issue is directly related to the quantum measurement problem where the coupling between the measured system and the measuring apparatus (detector) will cause the reduction of superposition or wave packet collapse [4]. It should be emphasized that the coupling between the measured system and the detector can be controlled to implement a process of quantum measurement. This is quite different from coupling with a real environment, the detailed knowledge of which is usually unavailable. In the past few years, a cavity QED system [5] and trapped ions [6] were utilized to demonstrate how to "engineer" the system-reservoir coupling so that the progressive decoherence can be observed with experimentally accessible technologies. In this paper, we show that the detailed dynamics of quantum decoherence can be illustrated with a solid state system – Josephson junction in the "qubit way"–a two level approximation [7–11].

Recently a variety of qubits based on Josephson junction superconductor systems have been experimentally implemented with a high quality factor [12,13]. Besides these great advances in the implementation of the individual Josephson junction qubits, there have been experiments to demonstrate quantum coherence of two charge qubits system [14] and two flux qubits system [15], and even a conditional gate based on charge qubit was made [16]. In this paper we demonstrate a dynamic decoherence process in the charge qubit-NAMR composite system. As will be shown in the following, this decoherence process can be easily engineered by the bias voltage on the NAMR.

The experimental basis of our NAMR-based scheme relies on the rapid progress in the development of nanomechanical devices. Notably, NAMRs were successfully fabricated with high vibrational frequencies 10 MHz – 1 GHz, tiny mass $10^{-15} - 10^{-16}$ kg and high quality factor $10^3 - 10^5$ [17]. Based on these experimental progresses, some theoretical schemes were proposed to create and detect the superposition of macroscopically distinct quantum states of NAMR by entangling the resonator with

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a Cooper pair box [18–20]. Most recently progress has been made towards a quantum computational architecture, for example the NAMR was integrated with JJ phase qubits and worked as a novel quantum memory [21] like the atomic ensemble [22], and quasi-spin excitons [23].

2 Charge qubit-NAMR composite system

As shown in Figure 1, the JJ charge qubit and the NAMR are connected directly. Here C' is the capacitance of the resonator itself, the value of which depends on the resonator displacement $\hat{x}$, C the capacitance of the junction. C_g the bias capacitance, C_x the distribution capacitance between the resonator and the CPB, and V_g the bias voltage on C_g, V_x the bias voltage on the resonator, E_J the Josephson coupling energy of the CPB. ω_0 is the frequency of the NAMR. The Coulomb energy of the circuit can be written as:

$$H_c = \frac{2e^2}{C_\Sigma(x)} \left(\hat{n} - n_t(\hat{x})\right)^2, \tag{1}$$

where the total capacitance is

$$C_\Sigma(\hat{x}) = C + C_g + C_x(\hat{x}), \tag{2}$$

the total effective gate charge on the CPB is

$$n_t(\hat{x}) = \frac{1}{2e}\left(C_g V_g + C_x(\hat{x}) V_x\right), \tag{3}$$

and

$$C_x(\hat{x}) \equiv \frac{C_x C'}{C_x + C'}. \tag{4}$$

Assuming the distance between the resonator and the CPB is much larger than the amplitude of zero point fluctuation of x, i.e.

$$d \gg \Delta x = \sqrt{\frac{1}{2m\omega_0}},$$

then

$$C_x(\hat{x}) \cong C_x\left(1 - \frac{x}{d}\right),$$

$$n_t(\hat{x}) \cong n_g + n_x - n_x \frac{\hat{x}}{d}.$$

Here we have defined

$$n_{g,x} = \frac{C_{g,x} V_{g,x}}{2e} \tag{5}$$

Set $n_g + n_x = 0.5$, we can use two-level approximation by just considering $|0\rangle_c$, $|1\rangle_c$ (the state $|0\rangle_c$ ($|1\rangle_c$) represents the state with no (one) cooper pair on the island). The total Hamiltonian reads

$$H = \frac{1}{2}\omega_a \sigma_z - g\left(a + a^\dagger\right)\sigma_x + \omega_0 a^\dagger a, \tag{6}$$

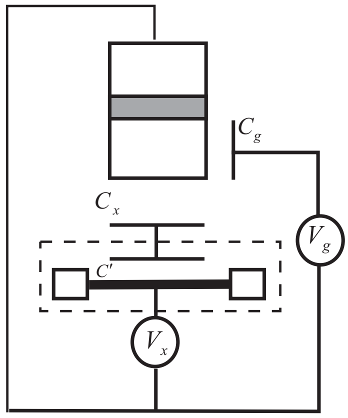

Fig. 1. A charge qubit is connected with a nanomechanical resonator (in the dashed frame) for the controlled coupling of charge qubits. Here C' is the capacitance of the resonator itself. C is the capacitance of the junction. C_g the bias capacitance, C_x the distribution capacitance between the resonator and the CPB, and V_g the bias voltage on C_g, V_x the bias voltage on the resonator.

the coupling coefficient is

$$g = 4E_c n_x \frac{\Delta x}{d} \tag{7}$$

where

$$E_c = \frac{e^2}{C + C_g + C_x} \tag{8}$$

and

$$\omega_a = E_J. \tag{9}$$

Here, the creation operator for the harmonic mode of the NAMR is defined as usual

$$a^\dagger = \sqrt{\frac{2}{m\omega_0}}\left(m\omega_0 x - ip\right). \tag{10}$$

In the model Hamiltonian above we have neglected a constant term. The quasi-spin operators are defined as

$$\sigma_x \equiv \tilde{\sigma}_z = |0\rangle_{cc}\langle 0| - |1\rangle_{cc}\langle 1|,$$

$$\sigma_z \equiv -\tilde{\sigma}_x = |1\rangle_{cc}\langle 0| + |0\rangle_{cc}\langle 1|.$$

Note that the Pauli matrices $\{\tilde{\sigma}_x, \tilde{\sigma}_y, \tilde{\sigma}_z\}$ are defined in charge eigenstate, that is, $|0\rangle_c$ ($|1\rangle_c$) denotes the state with no (one) Cooper pair on the island. But for further convenience we use another set of Pauli matrices as those without tilde in a different representation $\{|0\rangle, |1\rangle\}$ defined as

$$|0\rangle = \frac{1}{\sqrt{2}}\left(|0\rangle_c + |1\rangle_c\right)$$

and

$$|1\rangle = \frac{1}{\sqrt{2}}\left(|0\rangle_c - |1\rangle_c\right).$$

3 The mode squeezing of NAMR

Under normal experiment conditions the energy spacing of the CPB is much larger than that of the NAMR (as to be shown in the subsequent paragraph), thus cannot invoke the rotation-wave-approximation (RWA). The above model is quite similar to a cavity QED model without RWA, which usually describes the interaction of a single mode cavity and an off-resonance two-level atom [24]. In this cavity QED model, when the detuning between the cavity frequency and the $|0\rangle \leftrightarrow |1\rangle$ transition frequency is large enough to avoid any energy transfer between the atom and the cavity, the atoms in different states $|1\rangle$ and $|0\rangle$ will modify the phase of the cavity field in different ways [5,25] and induce the quantum decoherence of atomic states superposition.

The large detuning condition

$$\gamma = \frac{g}{|\omega_a - \omega_0|} \ll 1 \tag{11}$$

can be satisfied by taking proper parameters in experiments [13,18,20]. For example, we can take $E_c \sim 160$ μeV, $E_J \sim 50$ μeV, $\omega_0 = 0.5$ μeV, the coupling capacitance $C_x \sim 20$ aF, and $V_x \sim 1$V, so that $n_x \sim 60$. Then we have

$$g \sim 0.1\ \mu\text{eV}, \gamma \sim 2.20 \times 10^{-3} \ll 1.$$

We note that the rotation-wave-approximation condition cannot be satisfied in this system. With the above consideration of rational parameters in experiments, we adiabatically eliminate coherence effects between $|1\rangle$ and $|0\rangle$. Then we obtain an effective Hamiltonian

$$H_{eff} = H_1|1\rangle\langle 1| + H_0|0\rangle\langle 0|$$

which is diagonal with respect to $|0\rangle$ and $|1\rangle$. Here, the effective actions on the NAMR from the two qubit states $|1\rangle$ and $|0\rangle$ are

$$H_k \cong \omega_0 a^\dagger a + (-1)^k \frac{\omega_a}{2} - (-1)^k \frac{g^2}{4\delta}\left(2a^\dagger a + a^2 + a^{\dagger 2} + 1\right) \tag{12}$$

for $k = 0, 1$ respectively and $\delta = \omega_a - \omega_0$. H_{eff} is a typical effective Hamiltonian generating the quantum entanglement of the subsystems. That is, starting from a factorized initial state

$$|\psi(0)\rangle = (c_0|0\rangle + c_1|1\rangle) \otimes |s(0)\rangle,$$

the total system driven by H_{eff} will evolve into an entangled state

$$|\psi(t)\rangle = c_0|0\rangle \otimes |s_0(t)\rangle + c_1|1\rangle \otimes |s_1(t)\rangle \tag{13}$$

where

$$|s_k(t)\rangle = \exp(-iH_k t)|s(0)\rangle \quad (k = 0, 1) \tag{14}$$

and $|s(0)\rangle$ is the initial state of the NAMR. Therefore, the superposition of the charge states $|0\rangle$ and $|1\rangle$ will drive the state of the NAMR to evolve along the two directions $|s_0(t)\rangle$ and $|s_1(t)\rangle$.

It is very interesting to observe that the components Hamiltonian H_1 and H_0 are of Hermitian quadratic form of creation and annihilation operators. Mathematically, they are the same as that used to produce the degenerate parametric amplifier in nonlinear quantum optics with classical pump [26]. Defining a new set of bosonic operators A_k as the linear combinations of a and $a^\dagger$

$$A_k = \mu_k a - \nu_k a^\dagger, \tag{15}$$

where

$$\mu_k = \frac{1}{2}(\sqrt{N_k} + \frac{1}{\sqrt{N_k}})$$
$$\nu_k = \frac{1}{2}(\sqrt{N_k} - \frac{1}{\sqrt{N_k}}) \tag{16}$$

and

$$N_k = \sqrt{\frac{\omega_0\delta}{\omega_0\delta - (-1)^k g^2}}. \tag{17}$$

The Hamiltonian can then be written as

$$H_k = \Omega_k \left(A_k^\dagger A_k + \frac{1}{2}\right) + (-1)^k \left(\frac{\omega_a}{2} - \frac{g^2}{4\delta}\right). \tag{18}$$

Correspondingly the new eigenenergy

$$\Omega_k \cong \omega_0 \left(1 - (-1)^k \frac{g^2}{2\omega_0\delta}\right). \tag{19}$$

Since

$$|\mu_k|^2 - |\nu_k|^2 = 1, \tag{20}$$

we can see the coherent state in terms of A_k is the squeeze state of a. The two operators a and A_k can be transformed into each other though a unitary operator S_k, i.e.

$$S_k^\dagger a S_k = A_k. \tag{21}$$

The explicit form of S_k in terms of A_k is

$$S_k \equiv \exp\left[\frac{r_k}{2}A_k^2 - \frac{r_k}{2}A_k^{\dagger 2}\right], \tag{22}$$

where r_k is defined by

$$\mu_k = \cosh r_k,\ \nu_k = \sinh r_k. \tag{23}$$

Suppose the initial state of the NAMR is the coherent state $|\alpha\rangle$, i.e.

$$a\,|\alpha\rangle = \alpha\,|\alpha\rangle\,, \tag{24}$$

According to equation (21),

$$|\alpha\rangle = S_k\,|\alpha\rangle_{A_k}. \tag{25}$$

The component Hamiltonians H_0 and H_1 can create different squeezing of the harmonic mode of the NAMR during the time evolution, i.e.

$$e^{-i\Omega_k A_k^\dagger A_k t}\,|\alpha\rangle = e^{-i\Omega_k A_k^\dagger A_k t} S_k\,|\alpha\rangle_{A_k}$$
$$= e^{-i\Omega_k A_k^\dagger A_k t}\,|\alpha, \mu_k, \nu_k\rangle_{A_k} \tag{26}$$
$$= |\alpha, \mu_k(t), \nu_k(t)\rangle_{A_k}. \tag{27}$$

Here, the state with the subscript A_k is defined in terms of the operator $\hat{A}_k$, that is,

$$\hat{A}_k|\alpha\rangle_{A_k} = \alpha|\alpha\rangle_{A_k},$$
$$\left(\mu_k A_k + \nu_k A_k^\dagger\right)|\alpha,\mu_k,\nu_k\rangle_{A_k} = \alpha|\alpha,\mu_k,\nu_k\rangle_{A_k}. \quad (28)$$

and

$$\mu_k(t) = \mu_k e^{i\Omega_k t}, \nu_k(t) = \nu_k e^{-i\Omega_k t}. \quad (29)$$

Notably, H_0 and H_1 may drive the state of the NAMR from the same coherent state $|\alpha\rangle$ into two different squeezed states $|\alpha,\mu_0(t),\nu_0(t)\rangle_{A_0}$ and $|\alpha,\mu_1(t),\nu_1(t)\rangle_{A_1}$ respectively [27]. With the above considerations, we can evaluate the time evolution of the total system and obtain the wave function of the NAMR at time t

$$|s_k(t)\rangle = e^{i(-1)^k\left(\frac{g^2}{4\delta}-\frac{\omega_a}{2}\right)t - i\frac{\Omega_k}{2}t}|\alpha,\mu_k(t),\nu_k(t)\rangle_{A_k}. \quad (30)$$

The above calculations demonstrate that the off-resonance interaction between the NAMR oscillator mode and the different charge qubit will result in a dynamic squeezing split of the quasi-classical state $|\alpha\rangle$ of NAMR. The two splitting components with different squeezing are represented by different squeezing states.

4 Engineering decoherence of charge qubit

It is known that, dominated by the Hamiltonian H_{eff}, the conditional dynamics process described above tends to an ideal pre-quantum measurement when the overlap $\langle s_1(t)|s_0(t)\rangle$ approaches zero [31]. Physically the pre-measurement implies quantum decoherence of the charge qubit. To see this we consider the reduced density matrix of the charge qubit at time t. Its off-diagonal elements are determined by $c_1^* c_0\langle s_1(t)|s_0(t)\rangle$ and vanish completely as the overlapping $\langle s_1(t)|s_0(t)\rangle$ approaches zero. In this sense, the decoherence factor defined by

$$D(t) = |\langle s_1(t)|s_0(t)\rangle|$$

characterizes the extent of decoherence and the time-dependent behavior of $D(t)$ means a progressive process of decoherence or so called progressive decoherence. The very sharp peaks in $D(t)$ curves may originate from the reversibility of the Schrödinger equation for few body system [4,25].

Correspondingly, the decoherence factor characterizing quantum decoherence is

$$D(t) \cong \exp\left(-2|\alpha|^2\sin^2\left(\frac{g^2}{2\delta}t\right)\right) \quad (31)$$

where the condition $\frac{g}{\delta} \ll 1$ has been taken into consideration. The decoherence phenomenon with collapse and revival illustrated in Figure 2 is quite typical. It was found theoretically [4,28] in 1993, and the possibility of implementing its observation in cavity QED experiment was also pointed out in reference [25]. In 1997 it was also independently discussed [5] with another cavity QED setup.

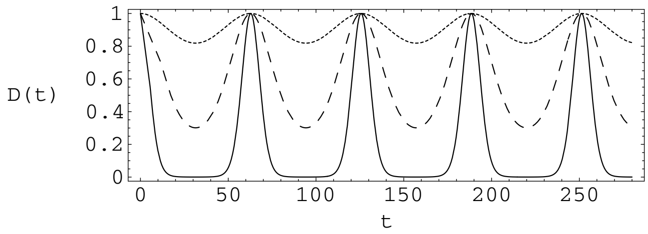

Fig. 2. The time-dependence of decoherence factor with different $|\alpha|^2 = 0.1$ (dot line), $|\alpha|^2 = 0.6$ (dash line), $|\alpha|^2 = 4$ (solid line). The larger $|\alpha|$ means more exact "detection" of this qubit, or the one-mode reservoir is more classical. It leads to a clear vanishing of coherence.

There are two time scales to depict the revival-collapse behaviors of decoherence process. In a very short time departure from $t = 0$,

$$D(t) \approx \exp\left(-\frac{2|\alpha|^2 g^4}{\delta^2}t^2\right), \quad (32)$$

a fast Gaussian decay with large $|\alpha|$ happens in a time scale

$$\tau_1 = \frac{2\delta}{g^2|\alpha|}, \quad (33)$$

which is much shorter than the time scale

$$\tau_2 = \frac{2\pi\delta}{g^2} \quad (34)$$

of the revival of the coherence. With the above parameters implemented in the experiments, we estimate $\tau_1 \simeq 1.38 \times 10^{-6}$ s and $\tau_2 \simeq 9.55 \times 10^{-6}$ s for $|\alpha| = 30$. The time-dependence of decoherence factor with different $|\alpha|^2 = 0.1$ (dot line), $|\alpha|^2 = 0.6$ (dash line), $|\alpha|^2 = 4$ (solid line). The larger $|\alpha|$ means the more exact "detection" about this qubit or the one-mode reservoir is more classical. It leads to an evident vanishing of coherence.

However, the decoherence factor $D(t)$ can not be directly observed. Its implied decoherence phenomenon with collapse and revival can only be reflected in an indirect way, which is similar to the cavity QED case in which the level populations were measured. A way to observe the engineered quantum decoherence phenomenon of this system is to detect the current through the probe junction connected to the Cooper pair box electrode [29]. When the charge qubit is in the high level state $|1\rangle_c$, there are two tunnelling electrons passing the probe junction. In fact, under a proper bias condition, the state decays into $|0\rangle_c$ via two single-electron tunnelling through the probe junction. In comparison with the case of atomic cavity QED, the advantage using charge qubit to test one-bit reservoir induced decoherence is due to the macroscopical nature of the superconductive system and the well-controlled nature of the coupling to one-bit engineered reservoir.

As usual, it is difficult to observe the two electrons via a single trial, but one can see an average effect of this tunnelling process. Another role of the probe junction is the preparation of the initial state by relaxing the Cooper

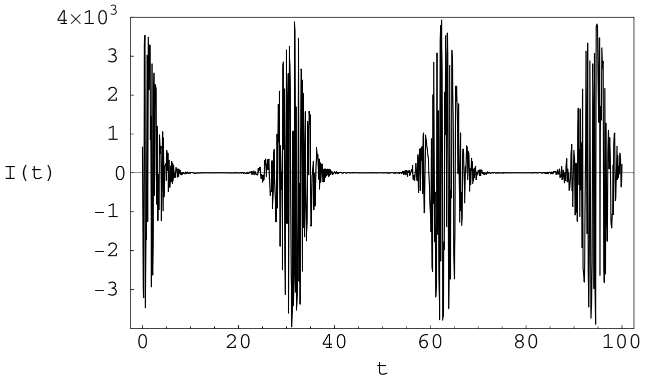

Fig. 3. The oscillation of the charge current with $|\alpha|^2 = 4$.

pair box to the ground state for a second trial. The current is proportional to the charging rate of the occupation probability

$$P_1(t) = Tr(\rho|1\rangle_{cc}\langle 1|) \tag{35}$$

of Cooper pair in $|1\rangle_c$. The corresponding current

$$I(t) = \frac{\partial}{\partial t}(-2eP_1(t)) \tag{36}$$

is explicitly expressed for the case with $c_0 = c_1 = \frac{1}{\sqrt{2}}$ as

$$I(t) \approx eD(t)\left\{\frac{g^2|\alpha|^2}{\delta}\left[\sin\frac{g^2}{\delta}t\cos\left(\omega_a t + |\alpha|^2\sin\frac{g^2}{\delta}t\right)\right.\right.$$
$$\times(1-2\cos 2\omega_0 t) + \cos\frac{g^2}{\delta}t\sin\left(\omega_a t + |\alpha|^2\sin\frac{g^2}{\delta}t\right)$$
$$\left.\left.\times(1-2\sin 2\omega_0 t)\right] + \omega_a\sin\left(\omega_a t + |\alpha|^2\sin\frac{g^2}{\delta}t\right)\right\}. \tag{37}$$

The oscillation of the charge current with $|\alpha|^2 = 4$. The oscillation of the current is demonstrated in Figure 3. It can be seen that the current oscillates sinusoidally, and the coupling to external reservoir adds the periodical amplitude modulation as the direct manifestation of decoherence. Experimentally, one can use the ratio of envelope width and the fixed period to measure the extent of decoherence quantitatively.

5 Conclusion with remarks

To observe this engineered quantum decoherence process, the system should be screened from decoherence and dissipation caused by the environment. The quality factor actually plays an important role in the experimental realization of the proposal. Due to the realistic quantum decoherence resulting from various mechanisms, such as background charge fluctuation, voltage-current damping and so on, the quality factor of the Cooper pair box in the present experiment is not high enough for demonstration of the subtle behaviors in our engineered decoherence. For the coupled system, the quality factor might be even worse. However, with the rapid progress of experimental technologies on Josephson junction qubits, the coherence of this kind of charge qubit might be enhanced greatly in the near future.

Another difficulty in the experimental implemention of the above theoretical setup is to initially prepare the large junction in a coherent state. The scheme to cool the NAMR to ground state has been proposed [20]. If the external sources can add the linear force proportional to x or p, they may force the NAMR oscillation mode to evolve into a coherent state from the vacuum state. For example, we apply a magnetic field along the direction of the beam of the NAMR with charge distributed on it. When the beam oscillates harmonically, the Lorentz force will act on the resonator. The force is just proportional to p. The magnitude of the coherent state is proportional to the external magnetic field. In practice, the initial state may be easily prepared in a thermal equilibrium at finite temperature, but this state is described by a diagonal density matrix in the coherent-state representation ("Q-representation"). Thus, the collapse and revival of coherence described above can still be observed and the increase of the temperature can enhance it. For the cavity QED case we have shown this enhancement effect by straightforward calculations [25]. The same calculations can be done here for the charge qubit.

Finally we give two remarks. (1) We note that the relevant quantum measurement problem of Josephson junction qubit has been considered theoretically by Averin [30]. He extends the concept of quantum non-demolition (QND) measurement to coherent Rabi oscillation of JJ qubit. The advantage of such QND measurement is that the detector induced decoherence can be avoided during the observation of the oscillation spectrum. This suggested that a scheme combining flux and charge qubit may be used in our setup to detect the engineered quantum decoherence without "additional quantum decoherence". (2) As normally understood [31], the quantum decoherence reflects the complementarity since the nanomechanical resonator mode plays the role of carrying away information about the phase of the charge qubit. The phase uncertainty appears when enough information of qubit is determined by the nanomechanical resonator in a classical state. The more precise the information about the phase of the qubit we obtain, the stronger the influence the resonator mode will exert on the qubit. The revival of coherence results substantially from the fact that the reservoir is only of a single mode. Its origin lies in the reversibility of the time evolution for systems of few degrees of freedom, governed by the Schrödinger equation.

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