Entanglement and squeezing in solid-state circuits

Wen Yi Huo\textsuperscript{1} and Gui Lu Long\textsuperscript{1,2}
\textsuperscript{1} Key Laboratory of Atomic and Molecular Nanosciences and Department of Physics, Tsinghua University, Beijing 100084, People’s Republic of China
\textsuperscript{2} Tsinghua National Laboratory For Information Science and Technology, and Institute of Microelectronics, Tsinghua University, Beijing 100084, People’s Republic of China
E-mail: gllong@tsinghua.edu.cn

\textit{New Journal of Physics} \textbf{10} (2008) 013026 (11pp)
Received 21 November 2007
Published 23 January 2008
Online at http://www.njp.org/
doi:10.1088/1367-2630/10/1/013026

Abstract. We investigate the dynamics of a system consisting of a Cooper-pair box and two superconducting transmission line resonators. There exist both linear and nonlinear interactions in such a system. We show that a single-photon entanglement state can be generated in a simple way in the linear interaction regime. In the nonlinear interaction regime, we propose a scheme for generating squeezed states of microwaves using three-wave mixing in solid-state circuits.

Contents

1. Introduction \hfill 2
2. The model \hfill 2
3. Linear interaction and generation of entangled state \hfill 4
4. Nonlinear interaction and generation of squeezed states \hfill 5
5. Discussion and conclusion \hfill 9
Acknowledgments \hfill 10
References \hfill 10
1. Introduction

Superconducting solid-state circuits with Josephson junctions have attracted much attention due to their application in quantum information processing. A series of successful experiments in superconducting charge [1]–[3], flux [4]–[7] and phase [8]–[11] qubits have demonstrated macroscopic quantum coherence and relatively long coherence time. Recently, superconducting qubits coupled to a one-dimensional superconducting transmission line resonator (STLR), LC oscillator or nanomechanical resonator have attracted much attention [12]–[14]. Such systems are of great interest not only as potential candidates for a scalable quantum computer but also in the study of the fundamental quantum mechanics of open systems. Recently, there has been significant progress in realizing quantum optics physics in solid state electrical circuits. This new subject was named ‘circuit quantum electrodynamics (circuit QED)’ [12]–[20].

Entanglement has been considered as one of the most important resources in quantum computation and quantum communication [21]. People are now searching for various possibilities of generating entanglement states in both optical [22, 23] and atomic systems [24]. Squeezed states provide a good example of the interplay between experiment and theory in the development of quantum mechanics. Statistical properties of squeezed states have been widely investigated [25]. The possibility of using squeezed states in quantum communication [26, 27], the study of fundamental quantum physics phenomena, as well as detecting gravitational radiation [28], has been recognized. In quantum optics, nonlinear interaction plays an indispensable role in the generation of squeezed light.

Compared to traditional cavity QED, circuit QED has some advantages, for instance its extremely strong coupling and infinite residence time of the ‘atom’ inside the cavity. Because the ‘atom’ is artificial and can be controlled by externally applied voltage and flux, one can obtain both linear and nonlinear interaction in circuit QED. In this paper, we show that the linear interaction and nonlinear interaction in a circuit QED system can be used to generate entangled states and squeezed states, respectively. Our scheme is based on a system consisting of two STLRs and a superconducting Cooper-pair box (CPB), also a superconducting charge qubit, playing the role of an artificial two-level atom.

In the past few years, two cavities interacting with a two-level atom has been widely investigated [29]–[31]. In [29, 30], the authors discussed the photon distribution concerning the energy flow in two cavities. In [31], the authors discussed the information transfer between a STLR and a nanomechanical resonator considering large detuning between cavities and the atom. Here, we discuss a more practical situation in solid-state circuit experiments and investigate the generation of entanglement and squeezed states. We find that by changing externally biased conditions of the CPB, one can get linear interaction and nonlinear interaction at will. The interaction between the two STLRs can be switched on and off by changing the external bias of the CPB. The squeezed state can be generated in the manner of three-wave mixing in solid-state circuits [32].

2. The model

The system under consideration is shown in figure 1. The first STLR (the green line at the bottom of figure 1) is fabricated coplanar with a CPB. The state of the CPB can be separately controlled by the gate voltage $V_g$ through a gate capacitance $C_g$. The CPB is also coupled to a
Figure 1. Schematic diagram of the combined system of two STLRs and a superconducting CPB, also a superconducting charge qubit. A superconducting transmission line resonator (the green line), shown at the bottom, is coupled to the CPB through the gate capacitance $C_g$, the other STLR (the red line), shown at the top, is coupled to the CPB by a magnetic field induced by its quantized current.

large superconducting reservoir through two identical Josephson junctions with capacitance $C_J$ and Josephson coupling energy $E_J$. This forms a superconducting quantum interference device (SQUID) and is also the basic configuration of a superconducting charge qubit. The SQUID configuration of the charge qubit allows one to apply external flux $\Phi_e$ to control the total effective Josephson coupling energy. For the superconducting charge qubit, the capacitance $C_g$ is much less than $C_J$. In this regime, a convenient basis is formed by the charge states, characterized by the number of Cooper pairs on the CPB. In the neighborhood of $n_g = 1/2$, only two charge states $|N\rangle$ and $|N+1\rangle$ play a role, while all other charge states, having much higher energies, can be ignored. The Hamiltonian of the CPB reads [33]

$$H_q = -\frac{1}{2}E_c(1 - 2n_g)\sigma_z - E_J\cos\left(\frac{\pi \Phi_e}{\Phi_0}\right)\sigma_x,$$

where $E_c = (2e)^2/(2C_\Sigma)$ is the charging energy of one Cooper pair on the CPB, $C_\Sigma = 2C_J + C_g$ is the total capacitance, and $n_g = C_gV_g/(2e)$ is the amount of gate charge induced by the gate voltage.

For a CPB fabricated inside one STLR, there is not only the dc voltage $V_g$ applied on the gate capacitance. The amplitude of the quantized voltage of the first STLR at the antinode $x_1 = L_1/k$ takes its maximum value

$$V_q = \sum_{k} V_0^k (a_k^\dagger + a_k), \quad V_0^k = \sqrt{\frac{\hbar \omega_1^k}{L_1 c_1}}.$$

Here, $\omega_1^k = k\pi/(L_1\sqrt{l_1 c_1})$, with $L_1$, $l_1$ and $c_1$ being the length, the inductance and capacitance per unit length of the first STLR, respectively. At low temperatures, there is only one mode of the first STLR, say $\omega_1^k = \omega_0$, that couples to the CPB. The quantum voltage applied on the gate
applied classical flux, and the amplitude of quantized flux is largest, two STLRs are degenerate, that is the mixing angle. In the linear interaction regime, we focus on the situation where the in equation (1), we chose the representation spanned by $|0\rangle = \sin(\theta/2) |N + 1\rangle + \cos(\theta/2) |N\rangle$ and $|1\rangle = \cos(\theta/2) |N + 1\rangle - \sin(\theta/2) |N\rangle$ which are the eigenstates of the CPB’s Hamiltonian in equation (1) as the work representation. Here, $\theta = \tan^{-1}[2EJ\cos(\pi \Phi_e/\Phi_0)/(E_J(1 - 2n_g))]$ is the mixing angle. In the linear interaction regime, we focus on the situation where the two STLRs are degenerate, that is $\omega_a = \omega_b = \omega$. In the rotating-wave approximation, the Hamiltonian in equation (3) is simplified to

$$H_3 = \hbar \omega (a^+ a + b^+ b) - \frac{1}{2} \hbar \Omega \rho_+ + \hbar g_a (a^+ \rho_- + a \rho_+) + \hbar g_b (b^+ \rho_- - b \rho_+),$$

where

$$\Omega = \frac{1}{\hbar} \sqrt{E^2_0 (1 - 2n_g)^2 + 4E_0^2 \cos^2 \left(\frac{\pi \Phi_e}{\Phi_0}\right)}.$$
is the transition frequency of the CPB,

\[ g_a = -\frac{e C_g V_0}{\hbar C_z} \sin \theta \quad \text{and} \quad g_b = -\frac{\pi \phi_b E_j \sin(\pi \Phi_e / \Phi_0)}{\hbar \Phi_0} \cos \theta, \]

are the coupling coefficients between the first and second STLR and the CPB, \( \rho_z = |0\rangle \langle 0| - |1\rangle \langle 1| \) and \( \rho_\alpha = \rho_\alpha^\dagger = |1\rangle \langle 0|. \)

When the two degenerate STLRs are resonating with the CPB, the Hamiltonian in equation (4) becomes

\[ H_4 = \hbar \Omega (a^\dagger a + b^\dagger b) - \frac{1}{2} \hbar \Omega \rho_z + \hbar g_a (a^\dagger \rho_- + a \rho_+) + i \hbar g_b (b^\dagger \rho_- - b \rho_+). \]  

(5)

The dynamics of two degenerate modes in an optical fiber resonating with a two-level atom has been widely investigated in classical communication [29, 30], in which the authors concentrated on the energetic flow in the two modes. Our combined system can be cooled to very low temperatures, which indicates that the two STLRs can be prepared in the vacuum states. As a result, it is unnecessary to find out all the eigenstates and eigenenergies in the system, and we can work in the subspace spanned by \( |0\rangle |0\rangle |0\rangle, |0\rangle |0\rangle |1\rangle, |0\rangle |1\rangle |0\rangle \) and \( |1\rangle |0\rangle |0\rangle \), where the order in the states are the first STLR, the second STLR and the CPB.

The ground state of Hamiltonian in equation (5) is \( |0\rangle |0\rangle |0\rangle \) and the corresponding energy is \( E_0 = -\hbar \Omega / 2 \). In this subspace, the dynamics evolution of the system can be evaluated completely by finding out the eigenstates and corresponding eigenenergies. By solving Schrödinger equation, one can find out all the eigenenergies and corresponding eigenstates in the subspace. The eigenenergies are \( E_1 = \hbar \Omega / 2 - \hbar G, E_2 = \hbar \Omega / 2, E_3 = \hbar \Omega / 2 + \hbar G, \) and the corresponding eigenstates are

\[ |\psi_1\rangle = \frac{1}{\sqrt{2}} (|001\rangle - i \sin \alpha |010\rangle - \cos \alpha |100\rangle), \]  

(6a)

\[ |\psi_2\rangle = \cos \alpha |010\rangle + i \sin \alpha |100\rangle, \]  

(6b)

\[ |\psi_3\rangle = \frac{1}{\sqrt{2}} (|001\rangle + i \sin \alpha |010\rangle + \cos \alpha |100\rangle), \]  

(6c)

where \( G = \sqrt{g_a^2 + g_b^2}, \cos \alpha = g_a / G \) and \( \sin \alpha = g_b / G \). If the initial state is \( |\psi(0)\rangle = |001\rangle \), that is two STLRs in vacuum state and CPB in excited state, the state at time \( t \) is

\[ |\psi(t)\rangle = \cos(Gt)|001\rangle + \sin(Gt)(\sin \alpha |010\rangle - i \cos \alpha |100\rangle), \]

here, we neglect a global phase factor \( e^{-i \omega t / 2} \). When \( Gt = (2k + 1)\pi / 2 \), we can get the entangled state of the two STLRs \( |\Psi\rangle = \sin \alpha |01\rangle - i \cos \alpha |10\rangle \). In particular, the Bell state \( |\Psi_-\rangle = (|01\rangle - i |10\rangle) / \sqrt{2} \) can be generated by adjusting the control parameters of \( g_a \) and \( g_b \) so that \( g_a = g_b \).

4. Nonlinear interaction and generation of squeezed states

Because of its dominance, it is reasonable and also justified that one can only investigate the linear interaction in the system for a wide range of external bias of the CPB. However, the nonlinear interaction should be considered at some external biased points, for instance, \( \Phi_e = m \Phi_0 \), with \( m \) an arbitrary integer. At these biased points, the Josephson coupling energy in equation (2) becomes \( E_j(\Phi_q) = (-1)^m + 1 E_j \cos (\pi \Phi_q / \Phi_0) \), and the second-order term in the expansion should be considered.
An important application of nonlinear interaction is generating squeezed states, which is an outstanding task in quantum mechanics and quantum optics. In the following, we show how we can use the nonlinear interaction in this system to generate squeezed states of the two STLRs. The method used here is three-wave mixing in solid-state circuits. As discussed above, here we also chose the eigenenergy basis of the CPB Hamiltonian to simplify the Hamiltonian in equation (3). Expanding the Josephson coupling energy to the second order in $\Phi_q/\Phi_0$, the Hamiltonian in equation (3) is simplified to

$$H_N = \hbar \omega_a a^+ a + \hbar \omega_b b^+ b - \frac{1}{2} \hbar \Omega \rho_e + \hbar \lambda_a (a^+ \rho_- + a \rho_+) + \hbar \lambda_b (b^+ \rho_- + b \rho_+), \quad (7)$$

where

$$\Omega = \frac{1}{\hbar} \sqrt{E_c^2 (1 - 2n_g)^2 + 4E_j^2}$$

is the transition frequency of the CPB,

$$\lambda_a = -\frac{eC_p V_0}{\hbar C_X} \sin \theta \quad \text{and} \quad \lambda_b = (-1)^{m+1} \frac{E_j \pi^2 \phi_b^2}{2\hbar \Phi_0} \cos \theta$$

are the coupling coefficients between the first and second STLR and the CPB in the nonlinear interaction regime, respectively. The mixing angle is $\theta = \tan^{-1}((-1)^m 2E_j/(E_c(1 - 2n_g)))$.

Assuming the detunings between the CPB and two STLRs satisfy the large detuning limit, that is $|\Delta_a| = |\Omega - \omega_a| \gg \Lambda$ and $|\Delta_b| = |\Omega - 2\omega_b| \gg \Lambda$, here $\Lambda = \sqrt{\lambda_a^2 + \lambda_b^2}$, then the variables of the CPB can be eliminated by performing the Fröhlich transformation [34] on the Hamiltonian in equation (7). Applying a unitary transformation $H_{NS} = e^{-S}H_{NS}e^S$ with $S = \lambda_a (a^+ \rho_- - a \rho_+ + \frac{1}{2}\hbar \lambda_a \lambda_b (1 + \frac{1}{\Delta_b}) b^+ b \rho_+)$, and expanding $H_{NS}$ to the second order in $\lambda/\Delta_i (i = a, b)$, we obtain an effective Hamiltonian

$$H_{NS} = \hbar \left( \omega_a - \frac{\lambda_a^2}{\Delta_a} \right) a^+ a - \frac{1}{2} \hbar \left( \Omega + \frac{2\lambda_b^2}{\Delta_b} + \frac{\lambda_a^2}{\Delta_a} \right) \rho_e + \hbar \left( \omega_b + \frac{2\lambda_b^2}{\Delta_b} - \frac{\lambda_b^2}{\Delta_b} \rho_e - \frac{\lambda_a^2}{\Delta_a} \rho_+ \right) b^+ b
- \frac{1}{2} \hbar \lambda_a \lambda_b \left( \frac{1}{\Delta_a} + \frac{1}{\Delta_b} \right) (b^+ b + a^+ a) \rho_e.$$

If the CPB is adiabatically kept in the ground state, the effective Hamiltonian is approximated as

$$H_{NS} = \hbar \omega_a a^+ a + \hbar \omega_b b^+ b - \frac{1}{2} \hbar \lambda_a \lambda_b \left( \frac{1}{\Delta_a} + \frac{1}{\Delta_b} \right) (b^+ b + a^+ a),$$

In the interaction picture, the Hamiltonian $H_{NS}$ reads

$$H_I(t) = -\frac{1}{2} \hbar \lambda_a \lambda_b \left( \frac{1}{\Delta_a} + \frac{1}{\Delta_b} \right) (b^+ e^{i\delta t} + a^+ b^2 e^{-i\delta t}),$$

where $\delta = 2\omega_b - \omega_a$. If $\delta = 0$, that is $\omega_a = 2\omega_b$ and $\Delta_a = \Delta_b = \Delta$, the Hamiltonian $H_I(t)$ becomes

$$H_I = \hbar \kappa (b^+ b + a^+ a). \quad (8)$$
where $\kappa = -\lambda_p/\Delta$ is the nonlinear coupling coefficient between the two STLRs. In the parametric approximation \[25\], the pump field is treated classically and pump depletion is neglected. The Hamiltonian in equation \((8)\) becomes

$$H_1 = \hbar \kappa \beta (b^\dagger b e^{-i\phi} + b^2 e^{i\phi}),$$

where $\beta$ and $\phi$ are the amplitude and phase of the coherent monochromatic microwave field which is used to drive the first STLR. The time evolution operator of the second STLR is

$$U(t) = e^{-i\kappa \beta t (b^\dagger b e^{-i\phi} + b^2 e^{i\phi})}.$$ \(10\)

It is obvious that the time evolution operator in equation \((10)\) is a squeezed operator of the second STLR. For a time duration $\tau$, the squeezed operator reads

$$S(\xi) = e^{-\frac{i\xi}{2} (b^\dagger b e^{-i\phi} + b^2 e^{i\phi})},$$

where $\xi = 2\kappa \beta \tau$ is the squeezed parameter.

For the second STLR initially in the vacuum state $|0\rangle$ and the phase of the pump coherent state $\phi = \pi/2$, the variance in the two quadratures $X_1 = (a + a^\dagger)/2$ and $X_2 = (a - a^\dagger)/2$ can be calculated directly using the transformation $S(\xi) b S(\xi) = b \cosh \xi - b^\dagger \sinh \xi$.

$$\Delta X_1 = \sqrt{\langle X_1^2 \rangle - \langle X_1 \rangle^2} = \frac{e^{-\xi}}{2},$$ \(11a\)

$$\Delta X_2 = \sqrt{\langle X_2^2 \rangle - \langle X_2 \rangle^2} = \frac{e^{\xi}}{2}. $$ \(11b\)

And the second STLR is in the squeezed vacuum state

$$|\xi\rangle = S(\xi)|0\rangle = e^{-(\xi/2)(b^\dagger b - b^2)}|0\rangle.$$ 

We now consider the influence of phase fluctuation on the squeezing. For three-wave mixing in quantum optics, the phase fluctuation is the dominant noise and it affects the variances of two quadratures $X_1$ and $X_2$, and also the squeezing properties. In our proposed system, the phase fluctuation in the pump field is also the dominant noise. In fact, many fluctuations which shift the frequencies $\omega_a$ and $\omega_b$ randomly can be included in this type of noise.

Now we consider the effect of phase fluctuation associated with a finite linewidth $D$ on the squeezing properties of the second STLR. For the linewidth $D$, the variances in the two quadratures are given by \[35\]

$$\Delta X_1 = \frac{1}{2} \left[ \frac{2 \Omega_p D e^{-2Dt}}{2 \Omega_p^2 - D^2/8} - \frac{(4 \Omega_p + 5D)e^{(2\Omega_p - 3D/2)t}}{8 \Omega_p + 2D} + \frac{(4 \Omega_p - 5D)e^{-(2\Omega_p + 3D/2)t}}{8 \Omega_p - 2D} \right] + \frac{D \sinh(2\Omega_p t)}{\sqrt{16 \Omega_p^2 + D^2}} e^{D t/2} \right]^{1/2},$$ \(12a\)

$$\Delta X_2 = \frac{1}{2} \left[ \frac{2 \Omega_p D e^{-2Dt}}{2 \Omega_p^2 - D^2/8} + \frac{(4 \Omega_p + 5D)e^{(2\Omega_p - 3D/2)t}}{8 \Omega_p + 2D} - \frac{(4 \Omega_p - 5D)e^{-(2\Omega_p + 3D/2)t}}{8 \Omega_p - 2D} \right] + \frac{D \sinh(2\Omega_p t)}{\sqrt{16 \Omega_p^2 + D^2}} e^{-D t/2} \right]^{1/2},$$ \(12b\)

New Journal of Physics 10 (2008) 013026 (http://www.njp.org/)
Figure 2. \( \Delta X_1 \) versus \( \xi \) for different ratios (a) \( D/\Omega_p = 0 \), (b) \( D/\Omega_p = 0.001 \), (c) \( D/\Omega_p = 0.01 \) and (d) \( D/\Omega_p = 0.1 \).

in the limit \( D \ll \Omega_p \), where \( \Omega_p = 2\kappa\beta = \xi/t \) is the effective Rabi frequency. In figure 2, we plot \( \Delta X_1 \) versus \( \xi \) for various ratios of \( D/\Omega_p \). The variance in the amplitude of \( X_1 \) increases due to the existence of phase fluctuation and there is a minimum in \( X_1 \) which decreases with increasing \( D/\Omega_p \). Equations (12a) and (12b) are then simplified to

\[
\Delta X_1 = \frac{1}{2} \sqrt{e^{-2\Omega_p t} + \left( \frac{D}{2} t \right) e^{2\Omega_p t}}, \tag{13a}
\]

\[
\Delta X_2 = \frac{1}{2} e^{\Omega_p t} \sqrt{1 - 2Dt}. \tag{13b}
\]

in the limit \( D \ll t^{-1} \ll \Omega_p \). It is evident that the phase fluctuation affects the squeezed properties severely with increasing time \( t \) due to the existence of the term \( \exp(\Omega_p t) \) in equation (13a). Therefore to observe the squeezing in a noisy environment, it is important to choose the right \( \xi \).

Noise in the CPB, such as background charge fluctuation [36] and \( 1/f \) noise [37], leads to a short decoherence time of the CPB. As a result, the short decoherence time of the CPB also affects the squeezing efficiency. The noise in the CPB and the fluctuation in the amplitude of the pump field together lead to fluctuation in the effective Rabi frequency \( \Omega_p \). According to the study [35], the influence of the fluctuation in \( \Omega_p \) on the squeezing efficiency can be written explicitly,

\[
\Delta X_1 = \frac{1}{2} \exp \left[ 2D' t + 2(e^{-D' t} - 1) - \Omega_p t \right], \tag{14a}
\]

\[
\Delta X_2 = \frac{1}{2} \exp \left[ 2D' t + 2(e^{-D' t} - 1) + \Omega_p t \right], \tag{14b}
\]

where \( D' \) is the amplitude linewidth induced by the fluctuation in the effective Rabi frequency \( \Omega_p \). In figure 3, we plot the variance \( \Delta X_1 \) versus squeezing parameter \( \xi \) for different ratios of \( D'/\Omega_p \). We chose two large ratios of \( D'/\Omega_p \) to show the influence of the fluctuations in \( \Omega_p \) on the squeezing efficiency. As a matter of fact, our numerical estimation (not shown here) indicates that the influence of decoherence of the CPB on the squeezing efficiency can be neglected even if the decoherence time is chosen as 50 ns.
Figure 3. $\Delta X_1$ versus $\xi$ for different ratios (a) $D'/\Omega_p = 0$, (b) $D'/\Omega_p = 0.1$, (c) $D'/\Omega_p = 0.5$.

5. Discussion and conclusion

It is notable that the two STLRs should be degenerate in the linear interaction regime, but one needs $\omega_a = 2\omega_b$ to generate a squeezed state in the nonlinear interaction regime. To meet this requirement in the same system, one can couple different modes of the two STLRs. For example, in the linear interaction regime, one can couple the second mode of the first STLR to the first mode of the second STLR, and in the nonlinear interaction regime, one can couple the fourth mode of the first STLR to the first mode of the second STLR.

In order to estimate the strength of the coupling between the two STLRs and the CPB, we select the following experimental parameters [12, 15, 17, 19]: $E_J/\hbar = 2\pi \times 9$ GHz, $\omega_b = 2\pi \times 10$ GHz, $C_g/C_\Sigma = 0.1$, $\alpha = S/r = 20 \mu$m, the ratio $\phi_b/\Phi_0$ is about $10^{-4}$. In the linear interaction regime, the coupling strengths $g_a$ and $g_b$ are about $2\pi \times 40$ MHz and $2\pi \times 1.5$ MHz, respectively. In the nonlinear interaction regime, the coupling strengths $\lambda_a$ and $\lambda_b$ are about $2\pi \times 40$ MHz and $2\pi \times 0.25$ kHz, here the transition frequency of the CPB is chosen as $2\pi \times 22.5$ GHz. The nonlinear coupling constant $\kappa$ is about $2\pi \times 4$ Hz, and the effective Rabi frequency $\Omega_p$ can be as large as $2\pi \times 40$ MHz with current microwave techniques.

In conclusion, we have investigated the dynamics of a system consisting of a CPB coupled to two STLRs. There are both linear and nonlinear interactions existing in such a system as a result of the nonlinear function of the coupling between the CPB and the STLR. In the linear interaction regime, we focused on the generation of the single-photon entangled state between the two STLRs.

In the nonlinear interaction regime, we discussed mainly the generation of squeezed states. Different from the previous proposals for generating squeezed states [38]–[42], our scheme does not need any operations on the CPB and does not use dissipation and measurement to generate the needed nonlinear interaction. Just keeping the CPB in the ground state, the squeezed states can be generated by driving the first STLR with a coherent microwave field and biasing the CPB at some special flux points.
Acknowledgments

This work is supported by the National Fundamental Research Program grant no. 2006CB921106, China National Natural Science Foundation grant nos 10325521 and 60635040.

References

[1] Nakamura Y, Pashkin Y A and Tsai J S 1999 Nature 398 786
[2] Pashkin Y A, Yamamoto T, Astafiev O, Nakamura Y, Averin D V and Tsai J S 2003 Nature 421 823
[3] Yamamoto T, Pashkin Y A, Astafiev O, Nakamura Y and Tsai J S 2003 Nature 425 941
[4] Mooij J E, Orlando T P, Levitov L, Tian L, van der Wal C H and Lloyd S 1999 Science 285 1036
[5] Izmalkov A, Grajcar M, Il'ichev E, Wagner Th, Meyer H-G, Smirnov A Yu, Amin M H S, Maassen vanden Brink A and Zagoskin A M 2004 Phys. Rev. Lett. 93 037003
[6] Majer J B, Paauw F G, ter Haar A C J, Harmans C J P M and Mooij J E 2005 Phys. Rev. Lett. 94 090501
[7] Plourde B L T, Robertson T L, Reichardt P A, Hime T, Linzen S, Wu C E and Clarke J 2005 Phys. Rev. B 72 060506
[8] Yu Y, Han S, Chu X, Chu S-I and Wang Z 2002 Science 296 889
[9] Xu H et al 2005 Phys. Rev. Lett. 94 027003
[10] Berkley A J, Xu H, Ramos R C, Gubrud M A, Strauch F W, Johnson R R, Anderson J R, Dragt A J, Lobb C J and Wellstood F C 2003 Science 300 1548
[11] McDermott R, Simmonds R W, Steffen M, Cooper K B, Cicak K, Osborn K d, Oh S, Pappas D P and Martinis J M 2005 Science 307 1299
[12] Blais A, Huang R-S, Wallraff A, Girvin S M and Schoelkopf R J 2004 Phys. Rev. A 69 062320
[13] Chiorescu I, Bertet P, Semba K, Nakamura Y, Harmans C J P M and Mooij J E 2004 Nature 431 159
[14] Cleland A N and Geller M R 2004 Phys. Rev. Lett. 93 070501
[15] Wallraff A, Schuster D I, Blais A, Frunzio L, Huang R-S, Majer J, Kumar S, Girvin S M and Schoelkopf R J 2004 Nature 431 162
[16] Wallraff A, Schuster D I, Blais A, Frunzio L, Majer J, Devoret M H, Girvin S M and Schoelkopf R J 2005 Phys. Rev. Lett. 95 060501
[17] Schuster D I, Wallraff A, Blais A, Frunzio L, Huang R-S, Majer J, Girvin S M and Schoelkopf R J 2005 Phys. Rev. Lett. 94 123602
[18] Blais A, Gambetta J, Wallraff A, Schuster D I, Girvin S M, Devoret M H and Schoelkopf R J 2007 Phys. Rev. A 75 032329
[19] Schuster D I et al 2007 Nature 445 515
[20] Huo W Y and Long G L 2007 Int. J. Quantum Inf. 5 829 (Preprint quant-ph/0702104)
[21] Nielsen M A and Chuang I L 2000 Quantum Computation and Quantum Information (Cambridge: Cambridge University Press)
[22] Kwiat P G, Mattle K, Weinfurter H, Zeilinger A, Sergienko A V and Shih Y 1995 Phys. Rev. Lett. 75 4337
[23] Lamas-Linares A, Howell J C and Bouwmeester D 2001 Nature 412 887
[24] Rowe M A, Kielpinski D, Meyer V, Sackett C A, Itano W M, Monroe C and Wineland D J 2001 Nature 409 791
[25] Scully M O and Zubairy M S 1997 Quantum Optics (Cambridge: Cambridge University Press)
[26] Cochrane P T, Ralph T C and Milburn G J 2002 Phys. Rev. A 65 062306
[27] Marino A M and Stroud C R Jr 2006 Phys. Rev. A 74 022315
[28] Bocko M F and Onofrio R 1996 Rev. Mod. Phys. 68 755
[29] Benivegna G and Messina A 1994 J. Mod. Opt. 41 907
[30] Xie Y B 1994 J. Mod. Opt. 42 2239
[31] Sun C P, Wei L F, Liu Y X and Nori F 2006 Phys. Rev. A 73 022318
[32] Huo W Y and Long G L 2007 Preprint 0704.0960
[33] Makhlin Y, Schön G and Shnirman A 2001 Rev. Mod. Phys. 73 357
[34] Fröhlich H 1950 Phys. Rev. 79 845
[35] Wódkiewicz K and Zubairy M S 1983 Phys. Rev. A 27 2003
[36] Astafiev O, Pashkin Y A, Nakamura Y, Yamamoto T and Tsai J S 2004 Phys. Rev. Lett. 93 267007
[37] Paladino E, Faoro L, Falci G and Fazio R 2002 Phys. Rev. Lett. 88 228304
[38] Moon K and Girvin S M 2005 Phys. Rev. Lett. 95 140504
[39] Zhou X X and Mizel A 2006 Phys. Rev. Lett. 97 267201
[40] Rabl P, Shnirman A and Zoller P 2004 Phys. Rev. B 70 205304
[41] Tian L and Simmonds R W 2004 Preprint cond-mat/0606787
[42] Ruskov R, Schwab K and Korotkov A N 2005 Phys. Rev. B 71 235407