Generation of two-mode squeezed states for two separated atomic ensembles via coupled cavities

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(Received 11 November 2009; published 25 January 2010)

We propose an efficient scheme for the generation of two-mode squeezed states for two separated atomic ensembles trapped in distant cavities. The scheme is based on selective couplings between the collective atomic modes and two linearly transformed common field modes mediated by an optical fiber or a third cavity. The quanta of the transformed atomic modes are exhausted due to the linear coupling with the transformed field modes, bringing the original atomic modes into the two-mode squeezed states. The experimental implementation of the scheme would be an important step toward quantum communication and networking with continuous variables.

DOI: 10.1103/PhysRevA.81.015804 PACS number(s): 42.50.Dv, 03.65.Ud, 03.67.Mn

Two-mode squeezed states are of crucial importance for quantum communication [1] and nonlocality tests [2] with continuous variable states. Recently, entangled squeezed states of two electromagnetic field modes have been used for the teleportation of quantum states with continuous variables [3]. In the context of cavity QED, schemes have been proposed for the generation of two-mode squeezed vacuum states for two field modes located in a single cavity via interaction with a single driven three-level atom [4–6]. Pielawa *et al.* [7] have proposed a scheme for generating two-mode squeezing in a high-Q microwave resonator with a beam of two-level atoms, which constitute a spin reservoir. Guzman *et al.* [8] have presented a scheme for realizing such continuous variable entanglement for two local cavity modes via interaction with an atomic sample.

On the other hand, single atoms or atomic ensembles with long-lived electronic states are ideal for storing and processing local quantum information. Parkins *et al.* [9] have presented a scheme for the preparation of two-mode squeezed states of the effective modes in a pair of atomic ensembles located in a single cavity through collective interaction with two quantized field modes and laser fields. In order to implement quantum communication or test quantum nonlocality one should prepare entanglement between separate subsystems. Although the scheme of Ref. [9] can also be applied to atomic ensembles in two separated, cascaded two-mode cavities, it is required that the cavity decay rate be much larger than the effective coupling between each atomic mode and the respective cavity modes so that the cavity modes can be adiabatically eliminated from the dynamics. Therefore, the time required to generate the two-mode squeezed state is exteremely long and the fidelity of entanglement may be seriously spoiled by dephasing, which induces an uncontrollable phase noise to Dicke states and destroys the coherence among them.

The coupled atom-cavity-fiber systems are considered as basic building blocks toward scalable quantum networking schemes. Proposals have been suggested for entanglement engineering and quantum communication between single atomic qubits in separate cavities through the exchange of real [10–12] or virtual photons [13] mediated by an optical fiber. In this article, we present an alternative scheme for the generation of two-mode squeezed states for two atomic ensembles trapped in two separate single-mode cavities. We show that, through suitable choice of the detunings and Rabi frequencies of the classical fields driving the atoms, the dynamics of the system is described by the competition between the annihilation operator of one atomic mode and the creation operator of the other atomic mode correlating with the linearly transformed field modes. With the application of a two-mode squeezing operator, the quanta of the transformed atomic modes are absorbed by the field modes, which forces the original atomic modes to evolve to the desired two-mode squeezed vacuum state. Since the two collective atomic modes are coupled to the common field modes, there is no need to eliminate field modes from the dynamics and the rate at which the desired entanglement is produced is much larger than that of the scheme of Ref. [9] when applied to atomic ensembles in two separated cavities.

We consider that two atomic ensembles are trapped in two cavities coupled by a short optical fiber or a third cavity, whose field mode will be referred to as the mediating mode. The atomic number in the ith ensemble is N_i. Suppose that the two cavity modes are resonant with the mediating mode. In the interaction picture, the resonant coupling between the two cavity modes and the mediating mode is given by the interaction Hamiltonian $H_1 = \nu b(a_1^\dagger + a_2^\dagger) + \text{H.c.}$, where b is the annihilation operator for the mediating mode, $a_j^\dagger$ is the creation operator for the jth cavity mode, and ν is the coupling strength.

The atomic level configuration is shown in Fig. 1. Each atom has one excited state $|r\rangle$ and two ground states $|e\rangle$ and $|g\rangle$. The transitions $|e\rangle \to |r\rangle$ and $|g\rangle \to |r\rangle$ in the ith cavity ($i = 1$, 2) are coupled to the corresponding cavity mode with the coupling strengths g_{ei} and g_{gi} and detunings Δ_e and Δ_g. Meanwhile, they are driven by two classical fields with the Rabi frequencies Ω_{ei} and Ω_{gi} and detunings $\Delta_g + \delta_e$ and $\Delta_e + \delta_g$. In the interaction picture, the Hamiltonian describing the atom-field interaction is

$$H_2 = \sum_{j=1}^{N_1}\{[g_{e1}a_1 e^{-i\Delta_e t} + \Omega_{e1}e^{i\phi_{e1}}e^{-i(\Delta_g+\delta_e)t}]|r_{1,j}\rangle\langle e_{1,j}| + [g_{g1}a_1 e^{-i\Delta_g t} + \Omega_{g1}e^{i\phi_{g1}}e^{-i(\Delta_e+\delta_g)t}]|r_{1j}\rangle\langle g_{1j}|\}$$

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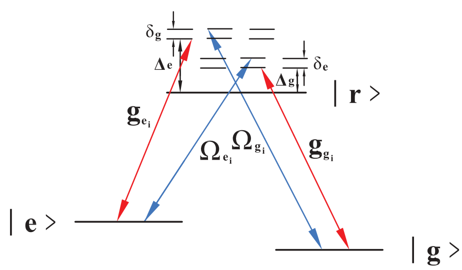

FIG. 1. (Color online) The atomic level configuration and transitions. The transitions $|e\rangle \rightarrow |r\rangle$ and $|g\rangle \rightarrow |r\rangle$ in the ith cavity ($i = 1, 2$) are coupled to the corresponding cavity mode with the coupling strengths g_{ei} and g_{gi} and detunings Δ_e and Δ_g. Meanwhile, they are driven by two classical fields with the Rabi frequencies Ω_{ei} and Ω_{gi} and detunings $\Delta_g + \delta_e$ and $\Delta_e + \delta_g$.

$$+ \sum_{j=1}^{N_2}\{[g_{e2}a_2 e^{-i\Delta_e t} + \Omega_{e2}e^{i\phi_{e2}}e^{-i(\Delta_g+\delta_e)t}]|r_{2j}\rangle\langle e_{2j}| + [g_{g2}a_2 e^{-i\Delta_g t} + \Omega_{g2}e^{i\phi_{g2}}e^{-i(\Delta_e+\delta_g)t}]|r_{2j}\rangle\langle g_{2j}|\} + \text{H.c.}, \quad (1)$$

where ϕ_{ei} and ϕ_{gi} are the phases of the classical fields driving the transitions $|e\rangle \rightarrow |r\rangle$ and $|g\rangle \rightarrow |r\rangle$ for the ith atomic ensemble. Introducing the new bosonic modes $c = \frac{1}{\sqrt{2}}(a_1 - a_2)$, $c_+ = \frac{1}{2}(a_1 + a_2 + \sqrt{2}b)$, and $c_- = \frac{1}{2}(a_1 + a_2 - \sqrt{2}b)$ [12], we can rewrite the Hamiltonian H_1 as $H_1 = \sqrt{2}\nu c_+^\dagger c_+ - \sqrt{2}\nu c_-^\dagger c_-$. In terms of the bosonic modes c, c_+, and c_-, the Hamiltonian H_1 is diagonal. So we can take H_1 as the "free Hamiltonian" mathematically and perform the transformation $e^{iH_1 t}$ to obtain the atom-field interaction Hamiltonian in the "interaction picture" with respect to H_1:

$$\begin{aligned}
H_2' &= e^{iH_{1f}t}H_2 e^{-iH_1 t} \\
&= \sum_{j=1}^{N_1}\left\{\left[\frac{1}{2}g_{e1}(e^{-i\sqrt{2}\nu t}c_+ + e^{i\sqrt{2}\nu t}c_- + \sqrt{2}c)e^{-i\Delta_e t} + \Omega_{e1}e^{i\phi_{e1}}e^{-i(\Delta_g+\delta_e)t}\right]|r_{1,j}\rangle\langle e_{1,j}|\right. \\
&\quad \left. + \left[\frac{1}{2}g_{g1}(e^{-i\sqrt{2}\nu t}c_+ + e^{i\sqrt{2}\nu t}c_- + \sqrt{2}c)e^{-i\Delta_g t} + \Omega_{g1}e^{i\phi_{g1}}e^{-i(\Delta_e+\delta_g)t}\right]|r_{1j}\rangle\langle g_{1j}|\right\} \\
&\quad + \sum_{j=1}^{N_2}\left\{\left[\frac{1}{2}g_{e2}(e^{-i\sqrt{2}\nu t}c_+ + e^{i\sqrt{2}\nu t}c_- - \sqrt{2}c)e^{-i\Delta_e t} + \Omega_{e2}e^{i\phi_{e2}}e^{-i(\Delta_g+\delta_e)t}\right]|r_{2j}\rangle\langle e_{2j}|\right. \\
&\quad \left. + \left[\frac{1}{2}g_{g2}(e^{-i\sqrt{2}\nu t}c_+ + e^{-i\sqrt{2}\nu t}c_- - \sqrt{2}c)e^{-i\Delta_g t} + \Omega_{g2}e^{i\phi_{g2}}e^{-i(\Delta_e+\delta_g)t}\right]|r_{2j}\rangle\langle g_{2j}|\right\} + \text{H.c.} \quad (2)
\end{aligned}$$

Set $\Delta_e, \Delta_g, \Delta_e - \Delta_g \gg g_{ei}, g_{gi}, \Omega_{ei}, \Omega_{gi}, \nu, \delta_e, \delta_g$. Then the upper-level $|r\rangle$ can be adiabatically eliminated and the atoms undergo Raman transitions. Choose the detunings appropriately so that the dominant Raman transitions are induced by the classical field Ω_{ei} and the bosonic mode c or the classical field Ω_{gi} and the bosonic mode c_+, while other Raman transitions are far off-resonant and can be neglected. This leads to the effective Hamiltonian

$$\begin{aligned}
H_e &= [\lambda_1 c e^{i\delta_e t} + \mu_1 c_+^\dagger e^{-i(\delta_g-\sqrt{2}\nu)t}]S_1^+ + \text{H.c.} \\
&\quad + (\alpha_{e1}c^\dagger c + \beta_{e1}c_+^\dagger c_+ + \eta_{e1}c_-^\dagger c_- + \xi_{e1})(S_{z1} + N_1/2) \\
&\quad + (\alpha_{g1}c^\dagger c + \beta_{g1}c_+^\dagger c_+ + \eta_{g1}c_-^\dagger c_- + \xi_{g1})(N_1/2 - S_{z1}) \\
&\quad + \{[-\lambda_2 c e^{i\delta_e t} + \mu_1 c_+^\dagger S_{2j}^+ e^{-i(\delta_g-\sqrt{2}\nu)t}]S_2^- + \text{H.c.}\} \\
&\quad + (\alpha_{e2}c^\dagger c + \beta_{e2}c_+^\dagger c_+ + \eta_{e2}c_-^\dagger c_- + \xi_{e2})(N_2/2 - S_{z1}) \\
&\quad + (\alpha_{g2}c^\dagger c + \beta_{g2}c_+^\dagger c_+ + \eta_{g2}c_-^\dagger c_- + \xi_{g2})(N_2/2 + S_{z1}),
\end{aligned} \quad (3)$$

where $S_1^+ = \sum_{j=1}^{N_1}|e_{1j}\rangle\langle g_{1j}|$, $S_2^- = \sum_{j=1}^{N_2}|e_{2j}\rangle\langle g_{2j}|$, $S_{z1} = \frac{1}{2}\sum_{j=1}^{N_1}(|e_{1j}\rangle\langle e_{1j}| - |g_{1j}\rangle\langle g_{1j}|)$, $S_{z2} = \frac{1}{2}\sum_{j=1}^{N_2}(|g_{2j}\rangle\langle g_{2j}| - |e_{2j}\rangle\langle e_{2j}|)$, $\lambda_i = \frac{\Omega_{ei}g_{gi}}{2\sqrt{2}}(\frac{1}{\Delta_g} + \frac{1}{\Delta_g+\delta_e})e^{-i\phi_{ei}}$, $\mu_i = \frac{\Omega_{gi}g_{ei}}{4}(\frac{1}{\Delta_e+\sqrt{2}\nu} + \frac{1}{\Delta_e+\delta_g})e^{i\phi_{gi}}$, $\alpha_{ei} = \frac{g_{ei}^2}{2\Delta_e}$, $\beta_{ei} = \frac{g_{ei}^2}{4(\Delta_e+\sqrt{2}\nu)}$, $\eta_{ei} = \frac{g_{ei}^2}{4(\Delta_e-\sqrt{2}\nu)}$, $\xi_{ei} = \frac{\Omega_{ei}^2}{\Delta_g+\delta_e}$, $\alpha_{gi} = \frac{g_{gi}^2}{2\Delta_g}$, $\beta_{gi} = \frac{g_{gi}^2}{4(\Delta_g+\sqrt{2}\nu)}$, $\eta_{gi} = \frac{g_{gi}^2}{4(\Delta_g-\sqrt{2}\nu)}$, and $\xi_{gi} = \frac{\Omega_{gi}^2}{\Delta_e+\delta_g}$. The two Raman transition channels can also be achieved through dispersive couplings to two different virtual excited states.

Define $d_i = \frac{1}{\sqrt{N_i}}S_i^-$ and $d_i^\dagger = \frac{1}{\sqrt{N_i}}S_i^+$. Then we have $[d_i, d_i^\dagger] = 1 - \frac{2}{N_i}(S_{zi} + N_i/2)$. Suppose that the average number of atoms in the state $|e\rangle$ in the first ensemble and the average number of atoms in the state $|g\rangle$ in the second ensemble are much smaller than N_1 and N_2, respectively. In this case, $S_{zi} \simeq -N_i/2$ and the collective atomic operators d_i and $d_i^\dagger$ can be regarded as the bosonic operators. Then the effective Hamiltonian H_e approximates to

$$\begin{aligned}
H_e &= \sqrt{N_1}[\lambda_1 c e^{i\delta_e t} + \mu_1 c_+^\dagger e^{-i(\delta_g-\sqrt{2}\nu)t}]d_1^\dagger + \text{H.c.} \\
&\quad + N_1(\alpha_{g1}c^\dagger c + \beta_{g1}c_+^\dagger c_+ + \eta_{g1}c_-^\dagger c_- + \xi_{g1}) \\
&\quad + \{\sqrt{N_2}[-\lambda_2 c e^{i\delta_{e2}t} + \mu_2 c_+^\dagger e^{-i(\delta_g-\sqrt{2}\nu)t}]d_2 + \text{H.c.}\} \\
&\quad + N_2(\alpha_{e2}c^\dagger c + \beta_{e2}c_+^\dagger c_+ + \eta_{e2}c_-^\dagger c_- + \xi_{e2}).
\end{aligned} \quad (4)$$

The Stark shifts $(N_1\alpha_{g1} + N_2\alpha_{e2})c^\dagger c$ and $(N_1\beta_{g1} + N_2\beta_{e2})c_+^\dagger c_+$ can be compensated through suitable choice of the detunings δ_e and δ_g. To see this clearly, we perform the further transformation $e^{iH_{e,0}t}$ with $H_{e,0} = \delta_e c^\dagger c + (\delta_g - \sqrt{2}\nu)c_+^\dagger c_+$, and obtain the new engineered Hamiltonian

$$\begin{aligned}
H_{e,i} &= e^{iH_{e,0}t}H_e e^{-iH_{e,0}t} - H_{e,0} \\
&= (N_1\alpha_{g1} + N_2\alpha_{e2} - \delta_e)c^\dagger c + [N_1\beta_{g1} + N_2\beta_{e2} - (\delta_g - \sqrt{2}\nu)]c_+^\dagger c_+ + (N_1\eta_{g1} + N_2\eta_{e2})c_-^\dagger c_- \\
&\quad + \left[c\sqrt{N_1}\lambda_1\left(d_1^\dagger - \frac{\sqrt{N_2}\lambda_2}{\sqrt{N_1}\lambda_1}d_2\right) + \text{H.c.}\right] \\
&\quad + \left[c_+\sqrt{N_2}\mu_2\left(d_2^\dagger + \frac{\sqrt{N_1}\mu_1}{\sqrt{N_2}\mu_2}d_1\right) + \text{H.c.}\right],
\end{aligned} \quad (5)$$

where we have discarded the constant energy terms.

Choose the Rabi frequencies, phases, and detunings of the classical fields appropriately so that the following conditions are satisfied:

$$\begin{aligned} |\sqrt{N_2}\lambda_2/\sqrt{N_1}\lambda_1| &= |\sqrt{N_1}\mu_1/\sqrt{N_2}\mu_2| = \tanh(r), \\ \sqrt{N_1}\lambda_1 &= \sqrt{N_2}\mu_2 = \varepsilon, \\ \delta_e &= N_1\alpha_{g1} + N_2\alpha_{e2}, \\ \delta_g &= N_1\beta_{g1} + N_2\beta_{e2} + \sqrt{2}\nu, \\ \phi_{e1} = \phi_{g1} = 0,\ \phi_{g2} &= \theta,\ \phi_{e2} = \pi + \theta. \end{aligned} \tag{6}$$

Then we have

$$\begin{aligned} H_{e,i} &= (N_1\eta_{g1} + N_2\eta_{e2})c_-^{\dagger}c_- + \varepsilon\{c[d_1^{\dagger} + \tanh(r)e^{-i\theta}d_2] \\ &\quad + c_+[d_2^{\dagger} + \tanh(r)e^{-i\theta}d_1] + \text{H.c.}\}. \end{aligned} \tag{7}$$

Performing the unitary transformation $H_{e,i}^{'} = S^{+}(\xi)H_{e,i}S(\xi)$, with $S(\xi)$ being the two-mode squeezing operator $S(\xi) = e^{\xi^* d_1 d_2 - \xi d_1^{\dagger} d_2^{\dagger}}$ and $\xi = re^{i\theta}$, we obtain the linearly coupling Hamiltonian

$$H_{e,i}^{'} = (N_1\eta_{g1} + N_2\eta_{e2})c_-^{+}c_- + \frac{1}{\cosh(r)}\,\varepsilon(cd_1^{\dagger} + c_+d_2^{\dagger} + \text{H.c.}) \tag{8}$$

After an interaction time t, this Hamiltonian leads to the transformations

$$\begin{aligned} d_1 &\to \cos(\Lambda t)d_1 + i\sin(\Lambda t)c, \\ c &\to \cos(\Lambda t)c + i\sin(\Lambda t)d_1, \\ d_2 &\to \cos(\Lambda t)d_2 + i\sin(\Lambda t)c_+, \\ c_+ &\to \cos(\Lambda t)c_+ + i\sin(\Lambda t)d_2, \end{aligned} \tag{9}$$

where $\Lambda = \varepsilon/\cosh(r)$. Suppose that the field modes are initially in the vacuum state (at the optical frequencies, the thermal photons are negligible). With the choice $t = \pi/2\Lambda$, the quantum states of the transformed atomic modes d_1 and d_2 are transferred to field modes c and c_+, respectively; that is, the two transformed atomic modes are driven to the vacuum state. Reversing the unitary transformation $S^{+}(\xi)$, we obtain the two-mode squeezed state $S(\xi)|0,0\rangle$ for the two collective atomic modes.

We now briefly discuss the experimental feasibility of the scheme. For the case that the two cavities are coupled via a third cavity, parameters $g_{e,i} = g_{g,i} = g = 2.5\times10^9\ \text{s}^{-1}$, $\nu = 4\times10^7\ \text{s}^{-1}$, $\Gamma = 1.6\times10^7\ \text{s}^{-1}$, and $\kappa = 0.4\times10^5\ \text{s}^{-1}$ are experimentally available [14,15], where Γ and κ are the atomic spontaneous emission rate and cavity decay rate, respectively. Set $\Delta_g = 1.0\times10^{12}\ \text{s}^{-1}$, $\Delta_e = \Delta_g + \Delta^{'} = 1.3\times10^{12}\ \text{s}^{-1}$, $\Omega_{e1} = 0.1g$, $\Omega_{e2} = 0.076g$, $\Omega_{g1} = 0.14g$, and $\Omega_{g2} = 0.184g$. Then the probability for the Raman transition $|e\rangle \longleftrightarrow |g\rangle$ induced by the transformed field mode c_- and the classical fields is on the order of $(g\Omega_e)^2/(\Delta_g\nu)^2 \sim 3.91\times10^{-6}$ and thus the effective Hamiltonian H_e of Eq. (2) is a good approximation. The coupling strength Λ is about $1.77\times10^7\ \text{s}^{-1}$ and the two-mode squeezed state with $r = 1$ can be obtained after an interaction time $t = \pi/2\Lambda \simeq 0.89\times10^{-7}$ s. The decoherence rate of the atomic sample due to atomic spontaneous emission, which is given by the effective single-atom spontaneous emission rate [16], is about $\Gamma_e = \Gamma\Omega_{e1}^2/\Delta_e^2 + \Gamma\Omega_{e2}^2/\Delta_e^2 \simeq 3.94\ \text{s}^{-1}$. This leads to an infidelity $\Gamma_e t \simeq 3.5\times10^{-7}$. The infidelity induced by the field decay is on the order of $\kappa t \simeq 3.56\times10^{-3}$. According to the Hamiltonian (8), the two-mode squeezed state can also be produced in the steady state. In this case the field decay plays a positive role in the prepration of entangled states. As $\Lambda > \kappa$, the time needed to reach the steady state is about $2/\kappa \simeq 5\times10^{-5}$ s [9]. Due to the uncertainty in the atomic numbers and the fluctuations in the Rabi frequencies the conditions of Eq. (6) cannot be strictly satisfied. In this case the obtained state is not a pure entangled state. However, the entanglement is only slightly affected by moderate deviations of these conditions. For example, even if $\sqrt{N_1}\lambda_1/\sqrt{N_2}\mu_2$ is deviated from unity by 10%, the reduction in the variance of the quadrature $\frac{1}{2\sqrt{2}}(d_1 + d_1^{\dagger} + d_2 + d_2^{\dagger})$ is only degraded by about 1–2 dB [9].

It should be noted that when the scheme of Ref. [9] is applied to preparation of entanglement between two atomic ensembles trapped in separated cavities, the two collective atomic modes are coupled to different field modes. It is required that the coupling strength Λ between each atomic mode and the corresponding field modes be much smaller than the cavity decay rate κ so that the field modes can be eliminated from the dynamics. In this case, the time needed to obtain the steady state is about κ/Λ^2, which is extremely long. In this case, the entanglement may be seriously destroyed by dephasing, one of the main decoherence resources in the atomic system. In distinct contrast, in the present scheme the collective atomic modes are coupled to common effective field modes so that it is unnecessary to eliminate the field modes from the dynamics to obtain the entanglement between two remote atomic ensembles and the needed time is greatly shortened for the same value of Λ.

In conclusion, we have proposed a scheme for the production of two-mode squeezed states for two atomic ensembles trapped in separated cavities connected by an optical fiber or a third cavity. The scheme is based on selective couplings between the atomic ensembles and common field modes. With suitable choice of the parameters of the classical fields and interaction time the quanta of the transformed atomic modes are transferred to the field modes, driving the original atomic modes to the desired two-mode squeezed state. Due to the collective enhancement of the atom-field coupling, the required interaction time decreases as the atomic numbers increase, which is important in view of decoherence. The scheme may provides the basic blocks for quantum communication and nonlocality test with continuous variables.

This work was supported by the National Natural Science Foundation of China under Grant No. 10974028, the Doctoral Foundation of the Ministry of Education of China under Grant No. 20093514110009, the Natural Science Foundation of Fujian Province under Grant No. 2009J06002, and funds from the State Key Laboratory Breeding Base of Photocatalysis, Fuzhou University.

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