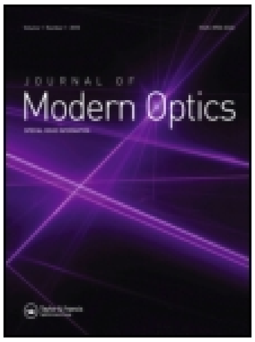

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Preparing entangled states of two atoms through the intensity-dependent interaction with the two-mode squeezed vacuum

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Two four-level atoms resonantly interacting with a two-mode squeezed vacuum field through the intensity-dependent Jaynes–Cummings-like model are considered. It is shown that the atoms can evolve into pure entangled states at certain time. Moreover, the pure entangled states can probabilistically be projected onto extremely entangled states by quantum jump technique. It is shown that the success probability monotonously increases with increasing degree of squeezing without any decoherence processes. However, it is found that if decoherence processes such as atomic spontaneous emission and photon leakage out of the cavity take place, the total success probability first increases but then decreases with increasing degree of squeezing and finally approaches zero as the squeezing becomes stronger and stronger. Therefore, the high degree of squeezing in the initial state of the cavity field is not always helpful for a high probability of creating extremely entangled states when spontaneous emission and cavity decay exist. Given spontaneous emission and photon leakage rates, there is an optimal squeezing degree for which the total success probability becomes maximum.

1. Introduction

Quantum entanglement plays a crucial role in quantum information processing such as quantum teleportation [1], dense coding [2], cryptography [3], computers [4]. Since stable electronic or central-mass motional states of atoms or ions are ideal candidates for quantum bits for various quantum information processes, many schemes for generating entangled states of atoms inside cavities or ions in traps have been proposed [5–14]. Some of them have been implemented in experiments [5, 8]. On the other hand, the generation of entangled photon states is relatively easy realized in down conversion processes with current technology [15–19]. Therefore, it is of great interest to investigate the entanglement transfer between entangled photons and atoms through coherent interactions. Kraus and Cirac [11] showed that two atoms which are trapped in two distant cavities can be entangled when the cavities are coupled to a common broadband squeezed reservoir. Tanas and Ficek [12] considered a system of two two-level atoms located at different spatial positions

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and coupled to the radiation field, whose modes are in a squeezed vacuum state, and showed that the atoms can be in a mixed state with stationary entanglement. Son *et al*. [20] investigated the entanglement transfer from an entangled two-mode continuous variables state to two qubits during the interaction of two two-level atoms with a two-mode squeezed vacuum state. Paternostro *et al*. [21] studied a process of entanglement transfer engineering when two remote qubits respectively interact with modes of local cavities which are driven by a two-mode squeezed vacuum field. Li *et al*. [22] investigated the generation of an entangled state for a pair of two-level atoms trapped in a bad cavity injected with a squeezed vacuum. The enhancement of the dynamical entanglement transfer from cavity field modes in a two-mode squeezed state to an ensembles of two-level atoms was also investigated [23]. In the present paper, we consider that two four-level atoms are trapped in two spatially separated single-mode cavities and the atoms resonantly and intensity-dependently interact with the local cavity modes, respectively. We show that the atoms can evolve into a pure entangled state if the field is initially in the two-mode squeezed vacuum state. We also show that extremely entangled states of the atoms can be created from the pure entangled states by a quantum jump technique with a high probability of success. The influence of decoherent processes such as spontaneous emission and cavity photon leakage on the generation of entangled states is investigated. Without the decoherence effects, the success probability of generating entangled states monotonically increases with increasing squeezing degree of the field. However, if decoherence effects take place, we find that the success probability initially increases and then decreases with increasing squeezing degree of the field. Given rates of spontaneous emission and photon loss, there is an optimum value of squeezing parameter to generate extremely entangled states with the highest success probability.

2. Preparing entangled states without dissipation

As shown in figure 1(*a*), two atoms are trapped in two spatially separated single-mode cavities and resonantly interact with the local cavity modes, respectively. The two local modes of the electromagnetic field are initially driven by a two-mode squeezing light through leaky side-mirrors of the cavities and prepared in the two-mode squeezed vacuum [21]. In the present model, we suppose that each of

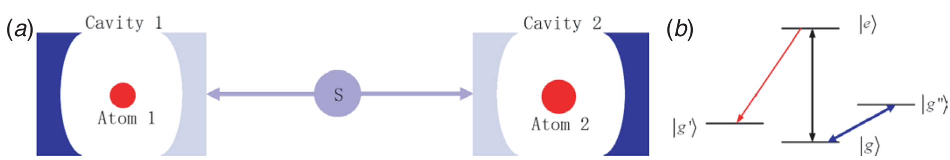

Figure 1. (*a*) Two atoms are trapped in two spatially separated single-mode cavities and each of the atoms resonantly interacts with the local mode of the field. S represents a source of the entangled field. (*b*) Atomic level configuration. (The colour version of this figure is included in the online version of the journal.)

the atoms has four levels. The levels $|g'\rangle$ and $|g''\rangle$ are ancillary and very stable without interacting with the cavity field, but the transition between the ground state $|g\rangle$ and the excited state $|e\rangle$ is coupled to the cavity field. In previous studies [21, 23] the interaction between atoms and local cavity fields is commonly described by the standard Jaynes–Cummings model. In that case, it was found that the atoms and the fields are always entangled and the atoms are statistically in an entangled state [21, 23]. In the present paper, we consider the intensity-dependent interaction between atoms and fields. We show that with this intensity-dependent interaction the atoms can be completely separated from the fields and be in a pure entangled state at a certain time. In contrast to the case of the standard Jaynes–Cummings interaction, the entanglement degree (logarithmic negativity) of the pure state generated by the intensity-dependent Jaynes–Cummings interaction can be monotonically increased by increasing the squeezing degree of the field. Moreover, one can extract extremely entangled states from the pure entangled states by a quantum jump technique with a high probability.

In the interaction picture, the Hamiltonian of the atom-field coupled system is ($\hbar = 1$) taken in the form

$$H = \sum_{j=1,2} g_j \left[a_j \sqrt{a_j^\dagger a_j} \sigma_{j+} + \sqrt{a_j^\dagger a_j} a_j^\dagger \sigma_{j-} \right], \tag{1}$$

where $\sigma_{j+} = |e\rangle_j \langle g|$, $\sigma_{j-} = (\sigma_{j+})^\dagger$ and $a_j(a_j^\dagger)$ is the annihilation(creation) operator for photons in mode j and g_j is the coupling constant between atom j and field j. This modified Jaynes–Cummings interaction is often employed in the literature [24–29, 31]. It is noticed that three operators $J_-^j \equiv a_j \sqrt{a_j^\dagger a_j}, J_+^j \equiv \sqrt{a_j^\dagger a_j} a_j^\dagger$, $J_0^j \equiv (1/2) + a_j^\dagger a_j$ in (1) constitute the Holstein–Primakoff representation of the $su(1,1)$ Lie algebra [26], which satisfy the commutation relations

$$[J_0^j, J_\pm^j] = \pm J_\pm^j, \tag{2}$$

$$[J_-^j, J_+^j] = 2J_0^j. \tag{3}$$

On the other hand, as is well known, three operators $L_-^j \equiv a_j^2/2, L_+^j \equiv a_j^{\dagger 2}/2, L_0^j \equiv (1/2 + a_j^\dagger a_j)/2$ also form a representation of the $su(1,1)$ Lie algebra. Thus, in the representation space of the $su(1,1)$ Lie algebra, the Hamiltonian (1) is completely equivalent to the two-photon interaction and represents a nonlinear interaction. The form of the Hamiltonian (1) could be realized in the laser-assisted interaction between the electric transitions and the central-mass vibrational modes of trapped ions with special mode functions of the driving electromagnetic field [29].

Suppose that initially the atoms are in state $|\phi\rangle_{12}$ and the field is in the two-mode squeezed vacuum state

$$|\psi\rangle_{12} = \sqrt{1-\lambda^2} \sum_{m=0}^{\infty} \lambda^m |m\rangle_1 \otimes |m\rangle_2, \tag{4}$$

where $\lambda = \tanh r$ and r is the squeezing parameter. At time t, the atom–field coupled system is in the state $U(t)|\phi\rangle_{12} \otimes |\psi\rangle_{12}$ with the evolution operator

$U(t) = \exp[-itH]$. For the initial states $|\phi\rangle_{12} = |e\rangle_1|e\rangle_2$, $|g\rangle_1|g\rangle_2$, and $|e\rangle_1|g\rangle_2$, it is easily shown that at $t_1 = \pi/(2g)$ the atoms and the field are separable, and evolve into the states

$$U(t_1)|e\rangle_1|e\rangle_2 \otimes |\psi\rangle_{12} = -\sqrt{1-\lambda^2}(|gg\rangle_{12} - \lambda|ee\rangle_{12})\sum_{m=0}^{\infty}\lambda^{2m}|2m+1\rangle_1|2m+1\rangle_2, \tag{5}$$

$$U(t_1)|g\rangle_1|g\rangle_2 \otimes |\psi\rangle_{12} = \sqrt{1-\lambda^2}(|gg\rangle_{12} - \lambda|ee\rangle_{12})\sum_{m=0}^{\infty}\lambda^{2m}|2m\rangle_1|2m\rangle_2, \tag{6}$$

$$U(t_1)|g\rangle_1|e\rangle_2 \otimes |\psi\rangle_{12} = -i\sqrt{1-\lambda^2}(|gg\rangle_{12} + \lambda|ee\rangle_{12})\sum_{m=0}^{\infty}\lambda^{2m}|2m\rangle_1|2m+1\rangle_2, \tag{7}$$

where $g_1 = g_2 = g$ is assumed. Thus, if the cavity field is switched off at $t_1 = \pi/2g$, the atoms are in the pure states

$$|\phi\rangle_{12} = \frac{1}{\sqrt{1+\lambda^2}}(|gg\rangle_{12} \pm \lambda|ee\rangle_{12}). \tag{8}$$

Once the atoms are in the states (8), in order to avoid the effects of atoms spontaneously decaying from the excited state $|e\rangle$, two π laser pulses are simultaneously applied to the atoms, respectively, as shown in figure 1(*b*) and the atoms are transferred from the state $|e\rangle$ to the stable state $|g'\rangle$ without affecting the state $|g\rangle$. In this way, the atoms are prepared in the stable states

$$|\phi\rangle_{12} = \frac{1}{\sqrt{1+\lambda^2}}(|gg\rangle_{12} \pm \lambda|g'g'\rangle_{12}). \tag{9}$$

The states (9) are entangled states as long as the field is initially in the squeezed state. The negative eigenvalue of the partial transposed density matrix corresponding to the states (9) is

$$\mathcal{N} = -\frac{\lambda}{1+\lambda^2}. \tag{10}$$

As a measure of entanglement, the logarithmic negativity is then given by [30]

$$E_{\mathcal{N}} = \log_2\left(\frac{2\lambda}{1+\lambda^2} + 1\right). \tag{11}$$

The quantity $E_{\mathcal{N}}$ may be regarded as the amount of entanglement transferred from the field to the atom during the interaction. It has been shown that for the standard Jaynes–Cummings interaction the maximal entanglement which can be dynamically transferred from the field to the atoms is not monotonically increasing with the entanglement of the field (i.e. with the squeezing parameter r) [23]. There is an optimal value $r = 0.86$ for which the transferred entanglement has the maximum value 0.90 [23]. In figure 2, the entanglement (11) is plotted as a function of the squeezing parameter r. It is observed that for the intensity-dependent interaction (1) the transferred entanglement monotonically increases with increasing squeezing parameter r. For $r = 0.86$, we find the entanglement $E_{\mathcal{N}} = 0.95$. This is better than the optimally transferred value obtained in the case of standard

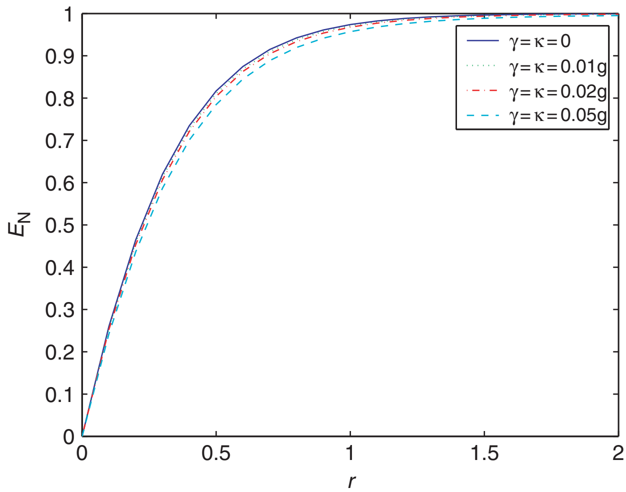

Figure 2. Logarithmic negativity vs. the squeezing parameter r. (The colour version of this figure is included in the online version of the journal.)

Jaynes–Cummings interaction. But, in many quantum information processes such as quantum teleportation [1], it is desired to have extremely entangled states for quantum information channels. However, as shown in figure 2, the states (9) could become extremely entangled states of two qubits only in the limit of infinite squeezing ($r \to \infty$). However, this case is not a physical situation because an infinite degree of squeezing means an infinite amount of excitation energy in the field [31]. Therefore, it is not a wise approach to generate extremely entangled states by continuously increasing the squeezing degree of the field. In the case of finite squeezing, we may (probabilistically) prepare extremely entangled states from the states (9) by means of a quantum jump technique [11, 32–35].

When the states (9) are obtained, two radio frequency pulses are, at the same time, applied to the transitions between $|g''\rangle$ and $|g\rangle$ of the atoms, respectively, as shown in figure 1(b). This action induces the unitary transformation: $|g\rangle \longrightarrow \cos\theta|g\rangle - i\sin\theta|g''\rangle$, and transforms the states (9) to the states

$$|\phi\rangle_{12} = \frac{1}{\sqrt{1+\lambda^2}}\Big[\cos^2\theta|gg\rangle_{12} - i\cos\theta\sin\theta(|gg''\rangle_{12} + |g''g\rangle_{12}) - \sin^2\theta|g''g''\rangle_{12} \pm \lambda|g'g'\rangle_{12}\Big]. \tag{12}$$

Then, the atoms in the states $|g''\rangle$ are monitored by use of the quantum jump technique [11]. If neither of the atoms is found in the state $|g''\rangle$, the states (9) are projected onto the ideal entangled states $(|gg\rangle_{12} \pm |g'g'\rangle_{12})/\sqrt{2}$ if one chooses $\cos^2\theta = \lambda$. The projection success probability is given by

$$P = \frac{2\lambda^2}{1+\lambda^2}. \tag{13}$$

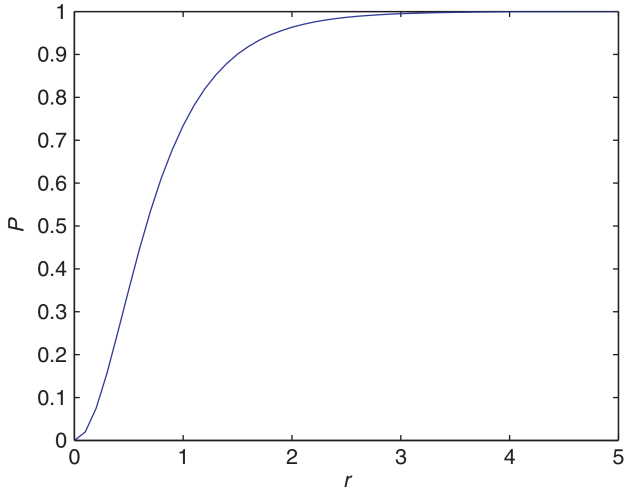

Figure 3. Success probability vs. squeezing parameter without dissipation. (The colour version of this figure is included in the online version of the journal.)

In figure 3, the dependence of the success probability (9) on the squeezing parameter is shown. One can observe that the success probability monotonically increases with increasing squeezing parameter and exceeds 50% when r is larger than 0.6.

3. Effects of spontaneous emission and photon leakage

In this section, we investigate the effects of atomic spontaneous emission and cavity photon leakage on the generation of the entangled states. According to the quantum jump theory [32–35], if no photons are lost through either spontaneous emission or cavity leakage, the evolution of the atom–field system is governed by the Schrödinger equation

$$i\frac{\partial}{\partial t}|\Psi(t)\rangle = H_c|\Psi(t)\rangle \tag{14}$$

with the conditional Hamiltonian

$$H_c = -i\frac{\gamma}{2}\sum_{j=1,2} a_j^\dagger a_j - i\frac{\kappa}{2}\sum_{j=1,2}|e_j\rangle\langle e_j| + \sum_{j=1,2} g_j\left[a_j\sqrt{a_j^\dagger a_j}\sigma_{j+} + \sqrt{a_j^\dagger a_j}a_j^\dagger\sigma_{j-}\right], \tag{15}$$

where γ is the atomic spontaneous emission rate and κ is the cavity decay rate.

Since each of the atoms interacts independently with one of the local cavity modes, we may solve equation (14) for each of the atoms and the corresponding local cavity mode, respectively. When atom j is in a state $c_j^g|g\rangle + c_j^e|e\rangle$ and mode j is in a Fock state of m photons at the initial moment, from equation (14), we find that at time t atom j and mode j are in the state

$$
\begin{aligned}
|\Psi(t)_m\rangle_j = {}& e^{-(\gamma m/2)t-(\Delta/4)t} c_j^g \Bigg[\Bigg[\left(\frac{\Delta}{8i\Omega_m}+\frac{1}{2}\right)e^{i\Omega_m t} + \left(-\frac{\Delta}{8i\Omega_m}+\frac{1}{2}\right)e^{-i\Omega_m t}\Bigg]|g\rangle_j|m\rangle_j \\
& + \Bigg[\left(-\frac{g_j m}{2\Omega_m}\right)e^{i\Omega_m t} + \left(\frac{g_j m}{2\Omega_m}\right)e^{-i\Omega_m t}\Bigg]|e\rangle_j|m-1\rangle_j\Bigg] \\
& + e^{-(\gamma(m+1))/2t-(\Delta/4)t} c_j^e \Bigg[\Bigg[\left(-\frac{\Delta}{8i\Omega_{m+1}}+\frac{1}{2}\right)e^{i\Omega_{m+1} t} + \left(\frac{\Delta}{8i\Omega_{m+1}}+\frac{1}{2}\right)e^{-i\Omega_{m+1} t}\Bigg] \\
& \times |e\rangle_j|m\rangle_j + \Bigg[\left(-\frac{g_j(m+1)}{2\Omega_{m+1}}\right)e^{i\Omega_{m+1} t} + \left(\frac{g_j(m+1)}{2\Omega_{m+1}}\right)e^{-i\Omega_{m+1} t}\Bigg]|g\rangle_j|m+1\rangle_j\Bigg],
\end{aligned}
\tag{16}
$$

where $\Delta = \kappa - \gamma, \Omega_m = \sqrt{g_j^2 m^2 - (\gamma-\kappa)^2/16}$. When the cavity field is initially in the squeezed vacuum state (4), at time t, the entire system will be in the state

$$
|\Psi(t)\rangle = \sqrt{1-\lambda^2}\sum_{m=0}^{\infty}|\Psi(t)_m\rangle_1 \otimes |\Psi(t)_m\rangle_2. \tag{17}
$$

If $\gamma \neq \kappa$, we find that at any time the state (17) cannot be separated into a product state of the atoms and the field modes. If $\gamma = \kappa$ and $g_1 = g_2 = g$, we obtain the following results.

When the atoms are initially in the state $|gg\rangle_{12}$, at $t = \pi/2g$, the atom–field system is conditionally in the state

$$
|\Psi(t)\rangle = \sqrt{1-\lambda^2}\sum_{m=0}^{\infty}\lambda^{2m}e^{-\gamma m\pi/g}|2m\rangle_1|2m\rangle_2 \otimes \left(|gg\rangle_{12} - \lambda e^{-\gamma\pi/2g}|ee\rangle_{12}\right). \tag{18}
$$

The probability of the atom-field being in this state is given by

$$
p_1 = \frac{(1-\lambda^2)(1+\lambda^2 e^{-\gamma\pi/g})}{1-\lambda^4 e^{-2\gamma\pi/g}}. \tag{19}
$$

The state of the atoms is now separated from the state of the cavity field. The atoms are in the pure state

$$
|\phi\rangle_{12} = \frac{1}{\sqrt{1+\lambda^2 e^{-\gamma\pi/g}}}\left(|gg\rangle_{12} - \lambda e^{-\gamma\pi/2g}|ee\rangle_{12}\right). \tag{20}
$$

By use of the method mentioned in the preceding section, we can project the state (20) onto the extremely entangled state $(|gg\rangle_{12} - |ee\rangle_{12})/\sqrt{2}$ with success probability

$$
P_{01} = \frac{2\lambda^2 e^{-\gamma\pi/g}}{1+\lambda^2 e^{-\gamma\pi/g}}. \tag{21}
$$

On combining (19) and (21), we obtain the total success probability of preparing the entangled state from the state $\sqrt{1-\lambda^2}\sum_{m=0}^{\infty}|m\rangle_1|m\rangle_2 \otimes |gg\rangle_{12}$

$$
P_1 = \frac{2(1-\lambda^2)\lambda^2 e^{-\gamma\pi/g}}{1-\lambda^4 e^{-2\gamma\pi/g}}. \tag{22}
$$

For the initial state $|ee\rangle_{12}$, at $t = \pi/2g$, the atom–field coupled system is conditionally in the separable state

$$|\Psi(t)\rangle = -\sqrt{1-\lambda^2}\sum_{m=0}^{\infty}\lambda^{2m}e^{-((\gamma(2m+1)\pi)/2g)}|2m+1\rangle_1|2m+1\rangle_2 \otimes \left(|gg\rangle_{12} - \lambda e^{-(\gamma\pi/2g)}|ee\rangle_{12}\right). \tag{23}$$

The probability of the atom-field being in this state is given by

$$p_2 = \frac{(1-\lambda^2)(1+\lambda^2 e^{-\gamma\pi/g})e^{-\gamma\pi/g}}{1-\lambda^4 e^{-2\gamma\pi/g}}. \tag{24}$$

The atoms are now in the pure state

$$|\phi\rangle_{12} = \frac{1}{\sqrt{1+\lambda^2 e^{-\gamma\pi/g}}}\left(|gg\rangle_{12} - \lambda e^{-\gamma\pi/2g}|ee\rangle_{12}\right). \tag{25}$$

The success probability for getting the entangled state $(|gg\rangle_{12} - |ee\rangle_{12})/\sqrt{2}$ from (25) is

$$P_{02} = \frac{2\lambda^2 e^{-\gamma\pi/g}}{1+\lambda^2 e^{-\gamma\pi/g}}. \tag{26}$$

The total success probability of preparing the entangled state from the state $\sqrt{1-\lambda^2}\sum_{m=0}^{\infty}|m\rangle_1|m\rangle_2 \otimes |ee\rangle_{12}$ is given by

$$P_2 = \frac{2(1-\lambda^2)\lambda^2 e^{-2\gamma\pi/g}}{1-\lambda^4 e^{-2\gamma\pi/g}}. \tag{27}$$

For the initial state $|eg\rangle_{12}$, at $t = \pi/2g$, the atom–field system is conditionally in the separable state

$$|\Psi(t)\rangle = -i\sqrt{1-\lambda^2}\sum_{m=0}^{\infty}\lambda^{2m}e^{-((\gamma(4m+1)\pi)/4g)}|2m+1\rangle_1|2m\rangle_2 \otimes \left(|gg\rangle_{12} + \lambda e^{-\gamma\pi/2g}|ee\rangle_{12}\right). \tag{28}$$

The probability of the atom–field being in this state is given by

$$p_3 = \frac{(1-\lambda^2)(1+\lambda^2 e^{-\gamma\pi/g})e^{-\gamma\pi/2g}}{1-\lambda^4 e^{-2\gamma\pi/g}}. \tag{29}$$

The atoms are in the pure state

$$|\phi\rangle_{12} = \frac{1}{\sqrt{1+\lambda^2 e^{-\gamma\pi/g}}}\left(|gg\rangle_{12} + \lambda e^{-\gamma\pi/2g}|ee\rangle_{12}\right). \tag{30}$$

The success probability for getting the entangled state $(|gg\rangle_{12} + |ee\rangle_{12})/\sqrt{2}$ from (30) is

$$P_{03} = \frac{2\lambda^2 e^{-\gamma\pi/g}}{1+\lambda^2 e^{-\gamma\pi/g}}. \tag{31}$$

The total success probability of preparing the entangled state from the state $\sqrt{1-\lambda^2}\sum_{m=0}^{\infty}|m,m\rangle_{12}\otimes|eg\rangle_{12}$ is given by

$$P_3 = \frac{2(1-\lambda^2)\lambda^2 e^{-3\gamma\pi/2g}}{1-\lambda^4 e^{-2\gamma\pi/g}}. \tag{32}$$

As shown in figure 2, the logarithmic negativity of the states (20), (25) and (30) monotonically increases with increasing squeezing parameter r and is not strongly affected by the dissipation processes. However, in figures 4–6, it is observed that the

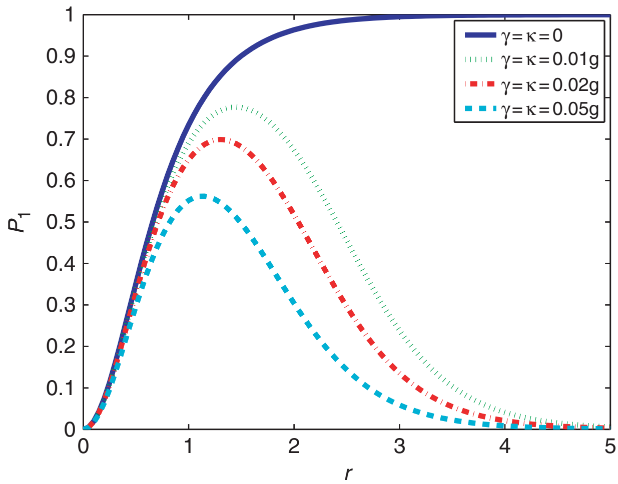

Figure 4. Total success probability P_1 vs. squeezing parameter r. (The colour version of this figure is included in the online version of the journal.)

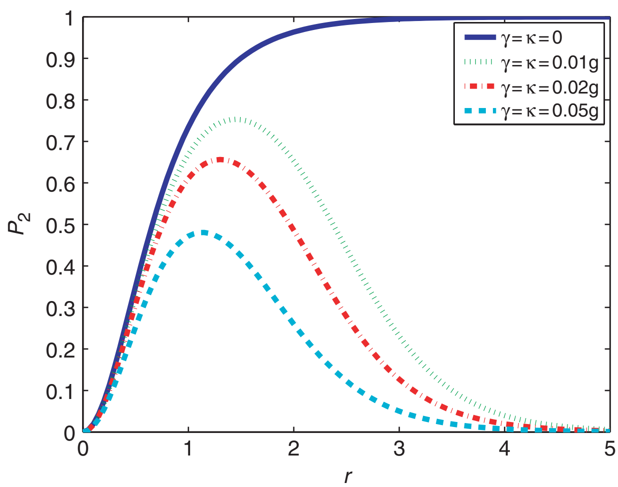

Figure 5. Total success probability P_2 vs. squeezing parameter r. (The colour version of this figure is included in the online version of the journal.)

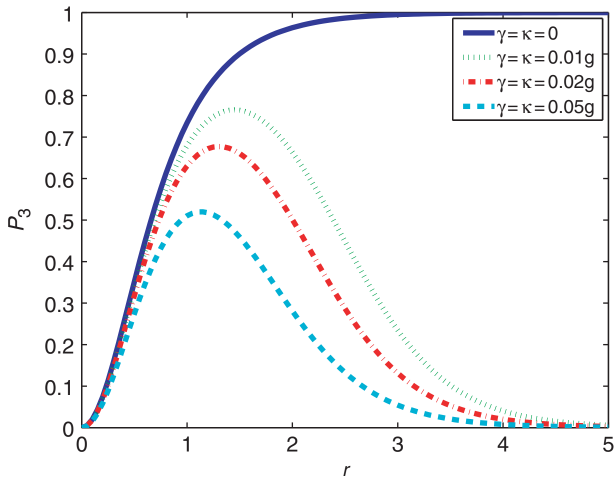

Figure 6. Total success probability P_3 vs. squeezing parameter r. (The colour version of this figure is included in the online version of the journal.)

total success probability first increases with increasing degree of squeezing, in contrast to the case without dissipation, but decreases with further increase of the degree of squeezing when the squeezing parameter is larger than a specific value, and finally approaches zero if squeezing is made stronger and stronger. This means that the high degree of squeezing in the initial state of the cavity field is not always helpful for a high probability of creating extremely entangled states when spontaneous emission and cavity decay exist. Given spontaneous emission and photon leakage rates, as shown in figures 4–6, there is an optimal squeezing degree for which the total success probability becomes a maximum. We also observe that the larger the rates are, the less the optimal squeezing degree is. The physical reason for this phenomenon is simple. Since photon leakage and spontaneous emission take place, at $\pi/2g$, the atom–field is in the states (18), (23) and (28) only under the condition that no photon escapes out of the cavities. As shown in (19), (24) and (29), the probabilities of no photon leakage contain the factor $(1-\lambda^2)$. It is this factor that results in the no photon escaping probability and thus the total success probability final decrease as the squeezing degree increases. The mean photon number of the field increases with increasing squeezing degree [31]. Therefore, the decrease in no photon leakage probability with increasing squeezing degree reflects the very simple fact that the more photons there are in the cavity, the larger the escape probability for one photon.

In a recent experiment [36], a squeezed vacuum with reduced noise level $-9.7\,\mathrm{dB}$ was generated. The corresponding squeezing parameter r was around 1.1. As shown in figures 4–6, for the order of decay rate $\gamma=\kappa\sim 10^{-2}g$, the optimal squeezing parameter is around 1.3. With present technology, the dissipation rates in a cavity can be controlled to the order of $\min(\gamma,\kappa)/g\sim 165$ [37]. Therefore, the scheme under consideration may be feasible in present experiments.

4. Summary

We consider an atom–field coupled system which consists of two four-level atoms and two spatially separated cavity fields driven by a two-mode squeezing light. Each of the atoms resonantly interacts with one local field. The interaction between the atom and the field is described by the intensity-dependent Jaynes–Cummings model. We find that at a certain time the atoms can be decoupled from the cavity field if both the atoms are in either the excited state or the ground state, or if one of the atoms is in the excited state but another is in the ground state, and the cavity field is prepared in two-mode squeezing vacuum state. From the decoupled atomic states, we can (probabilistically) create extremely entangled states of two atoms by a quantum jump technique. The total success probability of creating extremely entangled states monotonically increases with increasing squeezing degree of the field if the various dissipation processes are depressed. The condition that the rate of atomic spontaneous emission is equal to the decay rate of the cavity field is required to decouple the atoms from the cavity field at a certain time if atomic spontaneous emission and cavity photon leakage are considered. We also find that if dissipation processes take place, the total success probability first increases but then decreases with increasing degree of squeezing, and finally approaches zero if the squeezing parameter becomes very large. This means that strong squeezing is not always helpful for creating extremely entangled states of two atoms. We show that this phenomenon results from the fact that the probability of a photon leaking out of the cavity increases with increasing squeezing degree.

Acknowledgments

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