Light Amplification and Noise Reduction Based on Dressed-State Electromagnetically Induced Transparency

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Abstract It is suggested that the dressed-state electromagnetically induced transparency (EIT) can be used as an efficient way to produce light amplification and quantum noise reduction. The four-level tripod atoms are studied in the scheme by dressed-state atom and collective-mode approach. On two-photon resonance, two strong coherent fields induce the depopulation of a coherent superposition state of two lower states, which leads to quantum beat between the cavity modes. While the two strong coherent fields dress the atoms, the system is simply reduced to a standard EIT model, in which the laser transition, the coherent coupling, and the spontaneous decay constitute a successive population transfer channel to recycle the laser electron. It is for the very mechanism that the cavity modes oscillate with high intensities and exhibit squeezing.

Key words quantum optics; squeezed states; laser oscillation; dressed-state electromagnetically induced transparency; combination modes

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利用修饰态电磁感应透明产生光放大和噪声抑制

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摘要 提出了利用原子在修饰态下处于电磁感应透明(EIT)构型来产生光放大和量子噪声压缩。主要利用修饰态原子和组合模的方法讨论一个四能级 Tripod 原子系统。在双光子共振的条件下,两个强的相干场耦合原子产生两个叠加态,其中的一个叠加态与系统退耦合,这将导致两个腔场间产生量子拍。该四能级系统在修饰态下简化成了一个标准的三能级 EIT 模型,在这个 EIT 系统中,激光跃迁、原子相干耦合和自发衰减通道形成了一个连续的布居转移通道,激光电子有规律地循环。该机制提供了一种产生具有高强度振荡的压缩激光方式。

关键词 量子光学; 压缩态; 激光振荡; 修饰态电磁感应透明; 组合模

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1 Introdution

As is known, two-mode squeezing can lead to efficient distribution of entanglement and implementation of quantum channels[1-2]. The parametric oscillators have been demonstrated to be an effective way for the generation of two-mode squeezed light[3-6]. As a class of alternative and important schemes, four-wave mixing interactions[7-11] have been proposed for this purpose. In particular, for a two-level mixing system[7], one has the best achievable two-mode squeezing of 50%. However, for this class of schemes, two cavity modes operate below threshold and the steady state average amplitudes are zero. It is desirable to devise a scheme in which the two-mode squeezing occur well above threshold. So far, there have been several schemes been presented to obtain light amplification and two-mode squeezing[12-14]. For a closed Λ

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system, incoherent pumping from the lower to upper lasing level, and a successive decay from the upper lasing level to the auxiliary level add up to a large population in the auxiliary level. For a Raman system, a direct decay pathway from the lower lasing level to the auxiliary level makes the laser electron swiftly recycle to the auxiliary level. Therefore, only when coherent driving is combined with fast incoherent processes can good squeezing be achieved.

We suggest an efficient way to obtain light amplification and dynamical noise reduction. The dressed-state atom and combination mode approcach is employed in the calculations. It will be shown that two-mode squeezing with light amplification can be obtained via dressed-state electromagnetically induced transparency (EIT) in a four-level tripod atomic system. Three external coherent driving fields are coupled to the higher states to the three lower states, respectively. Two cavity fields are generated from two of the three transitions, respectively. On two-photon resonance, the two superposition states of the two lower states are induced by the coherent coupling of the two strong driving fields, and only one state is coupled to the system, while the other state is decoupled and has vanishing population. According to the combination modes, only one sum mode is mediated into the interaction, while the relative mode is decoupled from the system and keeps in vacuum state. It leads to the two original cavity fields in quantum beat. When two of the three driving fields induce the Stark splitting, the system is in dressed-state EIT, in which the laser electron is regularly recycled. On the basis of this, the two cavity modes run well above threshold and are in a two-mode squeezed state.

2 Laser oscillation and two-mode squeezing

2.1 Model and equation

We consider N four-level tripod-type atoms are put in a two-mode optical cavity (as shown in Fig. 1). Three external coherent fields of frequencies ω_{kl} are applied to the three transitions $|k_\mu\rangle-|l_\mu\rangle$ ($k=4$; $l=1,2,3$) with complex Rabi frequencies Ω_{4l}, respectively. Two cavity modes $a_{1,2}$ of frequencies $\nu_{1,2}$ are generated from the transitions $|1_\mu,2_\mu\rangle-|4_\mu\rangle$, respectively. The transitions $|4_\mu\rangle-|1_\mu,2_\mu,3_\mu\rangle$ are dipole-allowed, while the transitions between the three lower states $|1_\mu,2_\mu,3_\mu\rangle$ are dipole forbidden. In the dipole and rotating wave approximation and in the interaction picture, the master equation of the atom-field density operator is given by[15]

$$\dot{\rho}=-\frac{\mathrm{i}}{\hbar}[H,\rho]+\ell_a\rho+\ell_f\rho, \tag{1}$$

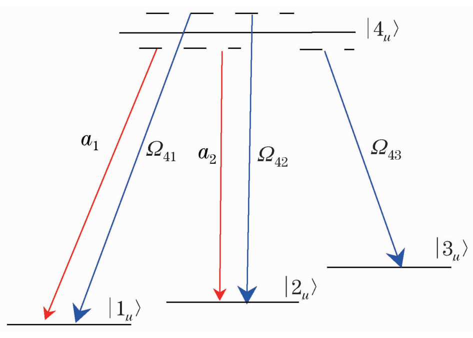

Fig. 1 Schematic energy level diagram for the four-level tripod atomic system

where $H=H_0+V$, H_0 represents the free term of the atoms and the interaction of the coherent fields with the atoms, and V denotes the interaction of the cavity fields with atoms,

$$H_0=\hbar\Delta_{41}|4_\mu\rangle\langle 4_\mu|+\sum_{\mu=1}^{N}\frac{\hbar}{2}\{\Omega_{41}|4_\mu\rangle\langle 1_\mu|+\Omega_{42}\exp[-\mathrm{i}(\Delta_{41}-\Delta_{42})t]|4_\mu\rangle\langle 2_\mu|+\Omega_{43}\exp[-\mathrm{i}(\Delta_{41}-\Delta_{43})t]|4_\mu\rangle\langle 3_\mu|\}+\mathrm{H.c.}, \tag{2}$$

$$V=\sum_{\mu=1}^{N}\hbar\left[g_1a_1|4_\mu\rangle\langle 1_\mu|\exp(\mathrm{i}\eta_1t)+g_2a_2|4_\mu\rangle\langle 2_\mu|\exp(\mathrm{i}\eta_2t)\right]+\mathrm{H.c.}, \tag{3}$$

H. c. represents Hermitian conjugates. $\hbar$ is the reduced Planck constant, a_l and $a_l^\dagger$ ($l=1,2$) are annihilation and creation operators, and $g_{1,2}$ are atom-field coupling constants. $\Delta_{kl}=\omega_{kl}-\bar{\omega}_{kl}$ are the detunings between the coherent field frequencies and atomic resonance frequencies, where $\bar{\omega}_{kl}$ are the atomic resonance frequencies, $k=4$, $l=1,2,3$. $\delta_1=\nu_1-\bar{\omega}_{41}$ and $\delta_2=\nu_2-\bar{\omega}_{42}$ are the detunings between the cavity frequencies and atomic resonance frequencies. $\eta_l=\omega_l-\nu_l$ ($l=1,2$), represents the detunings between the coherent driving field frequencies and the cavity field frequencies. $\ell_a\rho$ and $\ell_f\rho$ represent the atomic and field decay terms, respectively.

$$\ell_a\rho=\sum_{k=4;l=1,2,3}\gamma_{kl}\ell_{lk}\rho,\ \ell_{lk}\rho=\sum_{\mu=1}^{N}\frac{1}{2}\left[2\sigma_{lk}^{(\mu)}\rho\sigma_{kl}^{(\mu)}-\rho\sigma_{kl}^{(\mu)}\sigma_{lk}^{(\mu)}-\sigma_{kl}^{(\mu)}\sigma_{lk}^{(\mu)}\rho\right], \tag{4}$$

$$\ell_f\rho=\sum_{l=1,2}\kappa_l\ell_{al}\rho,\ \ell_{al}\rho=\frac{1}{2}(2a_l\rho a_l^\dagger-\rho a_l^\dagger a_l-a_l^\dagger a_l\rho), \tag{5}$$

$\sigma_{kl}^{(\mu)}=|k_\mu\rangle\langle l_\mu|$ ($k,l=1\sim4$) are the atomic projection operators for $k=l$ and the flip operators for $k\neq l$. γ_{kl} are the atomic decay rates from $|k_\mu\rangle$ to $|l_\mu\rangle$, and κ_l ($l=1,2$) are the decay rates of the two cavity fields. We focus on the case of two-photon resonance $\Delta_{41}=\Delta_{42}=\Delta$, and $\eta_1=\eta_2=\eta$, $g_1^\mu=g_2^\mu=g^\mu$ are assumed for simplicity. We define the two superposition states $|c_\mu\rangle=\cos\theta|1_\mu\rangle+\sin\theta|2_\mu\rangle$, $|d_\mu\rangle=-\sin\theta|1_\mu\rangle+\cos\theta|2_\mu\rangle$, with $\tan\theta=\left|\frac{\Omega_{42}}{\Omega_{41}}\right|$, and the combination modes $A=\cos\theta a_1\exp(\mathrm{i}\phi_1)+\sin\theta a_2\exp(\mathrm{i}\phi_2)$, $B=-\sin\theta a_1\exp(\mathrm{i}\phi_1)+\cos\theta a_2\exp(\mathrm{i}\phi_2)$. Then the Hamiltonians H_0 and V are respectively rewritten as

$$H_0=\sum_{\mu=1}^{N}\hbar\Delta\sigma_{44}^{(\mu)}+\frac{\hbar}{2}\{\Omega\sigma_{4c}^{(\mu)}+\Omega_3\sigma_{43}^{(\mu)}\exp[\mathrm{i}(\Delta-\Delta_{43})t]\}+\mathrm{H.c.}, \tag{6}$$

$$V=\sum_{\mu=1}^{N}\hbar g^\mu\left[A\sigma_{4c}^{(\mu)}+B\sigma_{4d}^{(\mu)}\right]\exp(\mathrm{i}\eta t)+\mathrm{H.c.}, \tag{7}$$

with $\Omega=\sqrt{|\Omega_{41}|^2+|\Omega_{42}|^2}$. It is seen from Hamiltonian Eq. (6) that only the superposition state $|c_\mu\rangle$ is coupled to the applied fields while the coherent superposition state $|d_\mu\rangle$ is decoupled. The state $|d_\mu\rangle$ is empty, which leads that the mode B always stays in its vacuum state. That is to say, there exists a quantum beat between the original modes $a_{1,2}$, which makes the relative mode B is independent of the sum mode A[16]. So the four-level tripod atomic system is reduced to a conventional EIT (states $|c_\mu\rangle$, $|3_\mu\rangle$, $|4_\mu\rangle$ form a Λ system) as shown in Fig. 2(a).

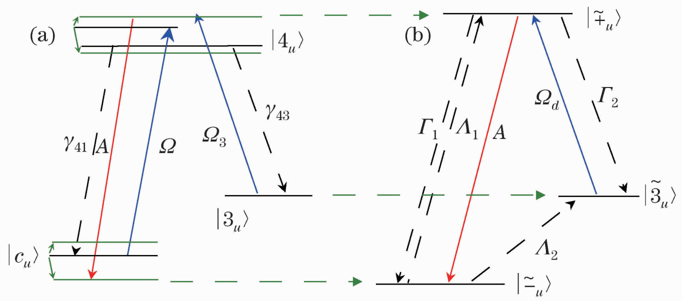

Fig. 2 (a) Transitions in the equivalent three-level Λ type atom in the bare tates basis; (b) dressed transitions for the equivalent EIT system

For clarity, we transform the bare atomic states to the dressed states[16]. We assume that the coherent driven fields are much stronger than the cavity fields $\bar{\Omega}=\sqrt{\Delta^2+\Omega^2}\gg(|g_j\langle a_j\rangle|,\gamma_{kl})$ ($j=1,2$; $k=4$, $l=1,2,3$). The μ-th atomic dressed states induced by the driving fields Ω are $|\tilde{+}_\mu\rangle=-\sin\phi|c_\mu\rangle+\cos\phi|4_\mu\rangle$, $|\tilde{-}_\mu\rangle=-\cos\phi|c_\mu\rangle+\sin\phi|4_\mu\rangle$, with $\tan(2\phi)=-\frac{1}{\delta}$, $\delta=\frac{\Delta}{\Omega}$, and the eigenvalues are $\lambda_\pm=\frac{1}{2}(-\Delta\pm\bar{\Omega})$.

The bare state $|3_\mu\rangle$ keeps unchanged but is denoted by $|\tilde{3}_\mu\rangle$ for consistency. We tune the driving field Ω_{43} such $\Delta_{43}=\Delta-\lambda_+$ that it is resonant with the $|\tilde{3}_\mu\rangle-|\tilde{+}_\mu\rangle$ transition, and tune the cavity fields such $\eta=-\Delta-\overline{\Omega}$ that they are resonant with Rabi sideband transition $|\tilde{+}_\mu\rangle-|\tilde{-}_\mu\rangle$. After rotating transform and secular approximation (neglecting the fast oscillating terms), the interaction Hamiltonian is rewritten as

$$H=\sum_{\mu=1}^{N}\hbar[g_A A\tilde{\sigma}_{+-}^{(\mu)}+\Omega_d\tilde{\sigma}_{+3}^{(\mu)}]+\text{H.c.}, \tag{8}$$

with $g_A=g\cos^2\phi$, and $\Omega_d=\frac{1}{2}\Omega_{43}\cos\phi$. By separating the relative mode B, we derive the master equations for the density $\tilde{\rho}$ of the sum mode A after the above transformations as

$$\frac{\partial}{\partial t}\tilde{\rho}=-\frac{\mathrm{i}}{\hbar}[H,\tilde{\rho}]+\ell'_a\tilde{\rho}+\ell_A\tilde{\rho}, \tag{9}$$

where the damping terms $\ell'_a\tilde{\rho}$ has the form

$$\ell'_a\tilde{\rho}=\Lambda_1\ell_{+-}\tilde{\rho}+\Gamma_1\ell_{-+}\tilde{\rho}+\Lambda_2\ell_{3-}\tilde{\rho}+\Gamma_2\ell_{3+}\tilde{\rho}+\Gamma_p\ell_p\tilde{\rho}, \tag{10}$$

with the phase damping $l_p\tilde{\rho}=\frac{1}{2}(2\tilde{\sigma}_p\rho\tilde{\sigma}_p-\rho\tilde{\sigma}_p^2-\tilde{\sigma}_p^2\rho)$, $\tilde{\sigma}_p=\tilde{\sigma}_{++}-\tilde{\sigma}_{--}$. For the atomic damping we have taken $\gamma_{41}=\gamma_{42}$ for simplicity. $\ell_{kl}\tilde{\rho}$ has the same form as $\ell_{kl}\rho$ as in Eq. (4). For the damping of cavity fields we have assumed that $\kappa_1=\kappa_2=\kappa$. $\ell_A\tilde{\rho}$ has the same form as in Eq. (5) except for the substitutions of A for a_1. The parameters in Eq. (10) are $\Lambda_1=\gamma_{41}\sin^4\phi$, $\Gamma_1=\gamma_{41}\cos^4\phi$, $\Lambda_2=\gamma_{43}\sin^2\phi$, $\Gamma_2=\gamma_{43}\cos^2\phi$, and $\Gamma_p=\frac{1}{4}\gamma_{41}\sin^2 2\phi$. The dressed transitions described by Eqs. (8) and (10) are indicated in Fig. 2(b).

2.2 Steady-state intensities and two-mode squeezing

To obtain the quantum correlations for the cavity fields, we calculate the normally ordered part of the output fluctuation spectrum

$$S(\omega)=2\int_0^\infty \mathrm{d}\tau\cos(\omega\tau)\frac{\langle :i(t+\tau),i(t):\rangle}{\langle i(t)\rangle}=Q\left(\frac{4\kappa}{\lambda}\right)\frac{\lambda^2}{\omega^2+\lambda^2}, \tag{11}$$

where $i(t)=\kappa A^\dagger(t)A(t)$ represents the output intensity operator, λ is proportional to the differential gain and Q is the Mandel factor $Q=\frac{\langle :(\Delta I)^2:\rangle}{I}$. We know that $S(\omega)=0$ is corresponding to shot noise and $-1\leqslant S(\omega)\leqslant 0$ to sub-Poissonian statistics. Correspondingly, the normally ordered parts of the output spectra for the two original modes are indicated as $S_l(\omega)=\frac{S(\omega)}{2}$. Quadrature squeezing in the respective modes occurs when sub-Poissonian statistics in the sum mode is existent.

The Langevin equations are derived from the master equation by means of the generalized P representation of Drummond and Gardiner[17]. The atomic variables are described by collective operators $\tilde{\sigma}_{kl}=\frac{1}{N}\sum_{\mu=1}^{N}\tilde{\sigma}_{kl}^{\mu}$. We choose the normal ordering $A^\dagger$, $\tilde{\sigma}_{+3}$, $\tilde{\sigma}_{+-}$, $\tilde{\sigma}_{3-}$, $\tilde{\sigma}_{++}$, $\tilde{\sigma}_{33}$, $\tilde{\sigma}_{--}$, $\tilde{\sigma}_{-3}$, $\tilde{\sigma}_{-+}$, $\tilde{\sigma}_{3+}$, A and define the correspondence between the c-numbers and operators as $\alpha\leftrightarrow A(\alpha\leftrightarrow A)$, $v_1\leftrightarrow\tilde{\sigma}_{3+}(v_1^\dagger\leftrightarrow\tilde{\sigma}_{+3})$, $v_2\leftrightarrow\tilde{\sigma}_{-+}(v_2^\dagger\leftrightarrow\tilde{\sigma}_{+-})$, $v_3\leftrightarrow\tilde{\sigma}_{-3}(v_3^\dagger\leftrightarrow\tilde{\sigma}_{3-})$, $z_l\leftrightarrow\tilde{\sigma}_{ll}(l=\pm,3)$. The set of equations for the c-numbers are derived as

$$\dot{\alpha}=-\frac{1}{2}k\alpha-\mathrm{i}g_A N v_2+F_\alpha, \tag{12}$$

$$\dot{v}_1=-\gamma_1 v_1+\mathrm{i}\Omega_d(z_+-z_3)+\mathrm{i}g_A\alpha v_3^\dagger+F_{v_1}, \tag{13}$$

$$\dot{v}_2=-\gamma_2 v_2+\mathrm{i}g_A\alpha(z_+-z_-)-\mathrm{i}\Omega_d v_3+F_{v_2}, \tag{14}$$

$$\dot{v}_3=-\gamma_3 v_3+\mathrm{i}g_A\alpha v_1^\dagger-\mathrm{i}\Omega_d v_2+F_{v_3}, \tag{15}$$

$$\dot{z}_-=-(\Lambda_1+\Lambda_2)z_-+\Gamma_1 z_++\mathrm{i}g_A(\alpha v_2^\dagger-\alpha^* v_2)+F_{z_-}, \tag{16}$$

$$\dot{z}_3=\Lambda_2 z_-+\Gamma_2 z_++\mathrm{i}\Omega_d(v_1^\dagger-v_1)+F_{z_3}. \tag{17}$$

Populations follow the closure relation $z_- + z_+ + z_3 = 1$. The parameters in Eqs. (12)～(17) are $\gamma_1 = \frac{1}{2}(\Gamma_1 + \Gamma_2 + \Gamma_{ph})$, $\gamma_2 = \frac{1}{2}(\Gamma_1 + \Gamma_2 + \Lambda_1 + \Lambda_2) + 2\Gamma_{ph}$, $\gamma_3 = \frac{1}{2}(\Lambda_1 + \Lambda_2 + \Gamma_{ph})$, and $F_x(t)$ are the noise forces. The noise correlations can be easily calculated from the generalized Einstein relations[1]. We assume the atomic variables change much more rapidly than the cavity fields ($\gamma_{41}, \gamma_{42} \gg \kappa$). The atomic variables can be adiabatically eliminated. By solving the steady state Eqs. (12)～(17), we can easily obtain the linear gain G and the stable intensities $I = g_A^2 \langle \alpha^* \alpha \rangle$ for sum mode A. It is not hard to find that $G > \kappa$, which means that the cavity modes operate above threshold. Using the mode transform relation, we obtain the respective intensities for the two original cavity modes $\langle I_l \rangle = g^2 \langle a_l^\dagger a_l \rangle = \frac{1}{2} g^2 I$, ($l = 1, 2$).

In the numerical calculations, we scale Rabi frequencies, coupling constants, detunings, and decay rates in units of $2\gamma_{43}$. The cooperativity parameter is defined as $C = \frac{2g^2 N}{\kappa \gamma_{41}}$. The intensities $\langle I_1 \rangle = \langle I_2 \rangle$ are in units of $4g^2\gamma_{43}^2$. The zero-frequency output spectrum and the respective intensities for the two cavity modes are plotted in Fig. 3. It is seen from Fig. 3(a) that for a wide range of parameters, the squeezing is existent and the respective intensities are large. Both the variances and the intensities are strongly dependent on the cooperativity parameter C. As C is small, the laser intensities are also relatively small and the range of squeezing is relatively narrow. As C increases, the laser intensities rise with relatively wide range of squeezing. When the cooperativity parameter $C = 800$, the minimal output spectrum $S(0) \approx -0.49$ is obtained. That corresponds to 49% squeezing for the respective mode, and the intensities are also strongly intensified $\langle I_1 \rangle = \langle I_2 \rangle \approx 148$ at the same time. In a word, for a large range of parameters, two cavity fields operate well-above threshold and display sub-Poissonian statistics.

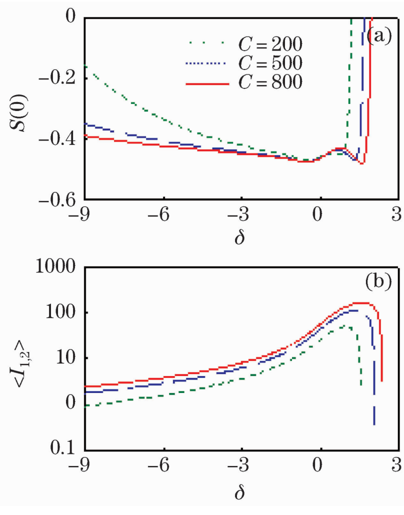

Fig. 3 (a) Zero-frequency output spectrum $S(0)$ and (b) steady intensities $\langle I_{1,2} \rangle$ versus the normalized detuning $\delta = \Delta/\Omega$ for different cooperativities $C = 200$, 500, 800 (the other parameters are chosen as $\Omega_{43} = 2$, $\Omega = 50$, $\gamma_{41} = 2$)

As for the mechanism of the squeezing, the following three aspects are important. 1) Quantum beat. Because the state $|d_\mu\rangle$ is not mediated into interaction with the cavity fields, the relative mode B is decoupled from the system, which leads to the two cavity modes in quantum beat. 2) Dressed state EIT. According to the dressed states induced by the effective field Ω, the sum mode A is tuned resonant with the transition $|\tilde{+}_\mu\rangle - |\tilde{-}_\mu\rangle$, and the coherent field Ω_3 is tuned resonant with the transition $|\tilde{3}_\mu\rangle - |\tilde{+}_\mu\rangle$. This leads to the system in a standard EIT configuration [as shown in Fig. 2(b)]. The coherent transition $|\tilde{3}_\mu\rangle - |\tilde{+}_\mu\rangle$ and the successive spontaneous decay $|\tilde{-}_\mu\rangle - |\tilde{3}_\mu\rangle$ transfer the atomic population from $|\tilde{-}_\mu\rangle$ to $|\tilde{+}_\mu\rangle$, which leads to a reduction of absorption. At the same time, atomic coherence between $|\tilde{3}_\mu\rangle$ and

$|\tilde{+}_{\mu}\rangle$ increases extra gain, as in lasing without inversion[12]. As a result, the sum mode runs above threshold. Laser oscillation with large intensities for the cavity modes can be obtained. 3) Intrinsic feedback. It is seen from Fig. 2(b) that there exist two intrinsic, incoherent channels $|\tilde{-}_{\mu}\rangle - |\tilde{+}_{\mu}\rangle$ and $|\tilde{-}_{\mu}\rangle - |\tilde{3}_{\mu}\rangle$, which is crucial for quantum correlations. By the two incoherent pathways and the coherent transitions $|\tilde{3}_{\mu}\rangle \xrightarrow{\Omega_d} |\tilde{+}_{\mu}\rangle$ and $|\tilde{+}_{\mu}\rangle \xrightarrow{A} |\tilde{-}_{\mu}\rangle$, the electron is regularly recycled. Such a recycling forms a deep intrinsic feedback[16]. As a result, the sum mode A has sub-shot noise when the system operates well above threshold. Since the system runs well above threshold, the saturation effect limits the degree of correlations. The present system can serve as an active device that provides high intensities and squeezed light.

3 Conclusion

It is possible to generate two-mode squeezing based on dressed-state EIT in a four-level tripod system. In terms of the combination modes and dressed states, the four-level tripod system can be reduced to a standard EIT system. Quantum beat occurs between the two cavity modes for the decoupling of the relative mode. In the dressed state EIT system, there exist electron recycling pathways to form an intrinsic feedback. The quantum beat and the deep feedback combine to make the sum mode operate well above threshold and have sub-shot noise. This leads to the two-mode squeezing with high intensities for the two cavity fields.

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