Dynamical suppression of the Autler-Townes doublet in the presence of a cavity

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We study the Autler-Townes absorption spectrum of the upper and intermediate levels of a three-level cascade atom in which the intermediate and ground levels are strongly driven by a laser field as well as weakly coupled by a cavity mode [Phys. Rev. **100**, 703 (1955)]. Surprisingly, one of the Autler-Townes peaks may be greatly suppressed or enhanced by tuning the cavity resonant frequency. The larger the Rabi frequency, the more pronounced the cavity-induced suppression or enhancement. The Autler-Townes spectrum carries full information on the cavity-induced modification of the strongly driven levels. A physical interpretation is presented in the dressed-state basis. [S1050-2947(98)08205-5]

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I. INTRODUCTION

Over the past decade, considerable interest has been devoted to systems in which atoms interact with a modified vacuum, such as that presented by a cavity [1], where the electromagnetic modes are concentrated around the cavity resonant frequency. The coupling of the atoms to the modified electromagnetic vacuum is therefore frequency dependent. For an excited atom located inside such a cavity, the cavity mode is the only one available to the atom for emission. If the atomic transition is in resonance with the cavity, the spontaneous emission rate into the particular cavity mode is enhanced [2]; otherwise, it is inhibited [3], which respectively result in a broadening and narrowing of the spontaneous emission spectrum [4]. When the atom-cavity coupling is very strong, the spontaneously emitted photon may be repeatedly absorbed and emitted by the atom before it leaves the cavity and the spectrum correspondingly exhibits a vacuum Rabi splitting [5].

When the atom is strongly driven by a laser field, the atom-laser system may be considered to form a new, dressed atom [6] whose energy-level structure is intensity dependent and whose spontaneous emission dominates at the three frequencies $\omega_L, \omega_L \pm \bar{\Omega}$. For such a coherently driven two-level atom placed inside a cavity, theoretical studies have predicted a phenomenological richness not found in the absence of the strong driving, for example, dynamical suppression of the spontaneous emission rate [7,8], population inversion in both bare and dressed-state bases [8,9], distortion and narrowing [7–9] of the Mollow triplet, and multipeaked spectral profiles [9,10]. All these features are very sensitive to the cavity resonant frequency because of the cavity enhancement of the dressed atomic transitions at the frequencies $\omega_L, \omega_L \pm \bar{\Omega}$. Recently, Lange and Walther [11] have observed the dynamical suppression of spontaneous emission in a microwave cavity. In the optical-frequency regimes, Zhu *et al.* [12] have also reported experimental studies of the effects of the cavity detuning on the radiative properties of a coherently driven two-level atom. They have shown that the atomic fluorescence of a strongly driven two-level atom is enhanced when the cavity frequency is tuned to one of the sidebands of the Mollow fluorescence triplet, whereas it is inhibited by tuning to the other sideband. The enhancement of atomic fluorescence at one sideband is a direct demonstration of population inversion.

Here we study how a cavity coupled to a coherently driven two-level system affects transitions between one of the cavity-coupled levels and an auxiliary level that is not coupled to the cavity. This paper investigates the Autler-Townes absorption spectrum [13] of a three-level cascade atom with one pair of levels coherently driven and coupled to a cavity frequency tunable around the transition frequencies of the atomic dressed states. Somewhat surprisingly, we find that one of the Autler-Townes peaks may be greatly suppressed or enhanced depending on the cavity frequency detuning.

There have been a few previous studies dealing with cavity effects in a three-level atomic system with the two upper levels coupled to a fixed-frequency cavity, without a driving laser. Field [14] reported that the vacuum Rabi splitting may induce a probe transparency at the transition frequency of the ground and cavity-coupled upper levels. He also showed that lasing without inversion can occur when the cavity decay reservoir is warmed up. Without considering cavity loss, Zubairy [15] found that the spectral profile of spontaneous emission from the intermediate level to the ground level has the same characteristics as the photon statistics of the cavity field coupled to the upper and intermediate levels. He therefore proposed to measure the quantum state of the cavity mode by monitoring the Autler-Townes spectrum. The modification of quantum jumps of a three-level atom, with two upper levels weakly coupled to a cavity of low-Q value, has also been reported recently [16].

The present model, different from those in Refs. [14–16] and described in Sec. II, assumes that a *frequency-tunable* cavity couples to a pair of strongly driven levels. The strong, coherent driving field may dynamically modify the cavity-induced effects. Section III shows that the Autler-Townes spectrum of a cavity-uncoupled level and one of the two

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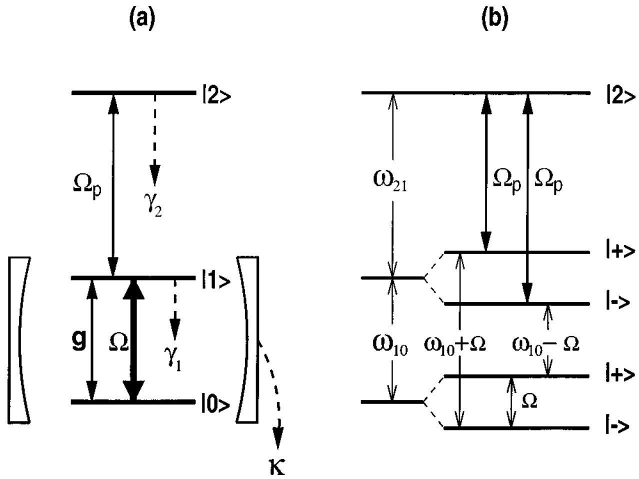

FIG. 1. (a) Atomic energy-level diagram for observation of the Autler-Townes spectrum, where the ground level $|0\rangle$ and the intermediate level $|1\rangle$ are coupled by a cavity mode and driven by a laser field. (b) Dressed-state level diagram for $\omega_L=\omega_{10}$.

cavity-coupled, coherently driven levels fully carries the information of the cavity effects on the two coherently driven levels. We also present a brief account of the theory in the dressed-state picture. Conclusions are presented in Sec. IV.

II. MODEL

We consider a three-level cascade atom in which the ground and intermediate levels are driven by a laser field and coupled by a cavity mode. The upper and intermediate levels decay into background modes at the rates γ_2 and γ_1, respectively, while the cavity mode is damped by a vacuum reservoir at the rate κ, as depicted in Fig. 1(a). In a frame rotating at the frequencies ω_L and ω_{21} [17] the master equation of the density matrix operator ρ for the combined atom-cavity system is of the form

$$\dot{\rho}=-i[H_A+H_C+H_I,\rho]+\mathcal{L}_A\rho+\mathcal{L}_C\rho, \tag{1}$$

where

$$H_A=-\Delta A_{00}+\frac{\Omega}{2}(A_{10}+A_{01}), \tag{2a}$$

$$H_C=\delta a^\dagger a, \tag{2b}$$

$$H_I=g(a^\dagger A_{01}+A_{10}a), \tag{2c}$$

$$\mathcal{L}_A\rho=\gamma_2(2A_{12}\rho A_{21}-A_{22}\rho-\rho A_{22}) + \gamma_1(2A_{01}\rho A_{10}-A_{11}\rho-\rho A_{11}), \tag{2d}$$

$$\mathcal{L}_C\rho=\kappa(2a\rho a^\dagger-a^\dagger a\rho-\rho a^\dagger a), \tag{2e}$$

with

$$\Delta=\omega_{10}-\omega_L, \quad \delta=\omega_C-\omega_L, \tag{3}$$

where $A_{lk}=|l\rangle\langle k|$ is an atomic operator, a and $a^\dagger$ are the annihilation and creation operators of the cavity mode, respectively, Ω is the Rabi frequency of the driving field, g is the atom-cavity coupling constant, ω_L and ω_C are, respectively, the frequencies of the driving laser and the cavity mode, and ω_{lk} is the level splitting between the levels $|l\rangle$ and $|k\rangle$ $(l,k=0,1,2)$.

We assume that the atom-cavity coupling is weak and the cavity has a low-Q value, so that $\kappa\gg g\gg\gamma$, that is, the cavity field decay dominates. The cavity field response to the continuum modes is much faster than that produced by its interaction with the atom, so that the atom always experiences the cavity mode in the state induced by the vacuum reservoir. Thus one can adiabatically eliminate the cavity-mode variables, giving rise to a master equation for the atomic variables only. As the derivation is tedious, we refer readers to [8,11] and here only outline the key points.

We temporarily disregard $\mathcal{L}_A\rho$ in the elimination of the cavity mode since it remains unchanged at the end of these operations. First we perform a canonical transformation to the atom-cavity interaction picture (1) by

$$\tilde{\rho}=e^{i(H_A+H_C)t}\rho e^{-i(H_A+H_C)t}. \tag{4}$$

The master equation then takes the form

$$\partial_t\tilde{\rho}=-i[\tilde{H}_I(t),\tilde{\rho}]+\mathcal{L}_C\tilde{\rho},$$

or

$$\partial_t(e^{-\mathcal{L}_Ct}\tilde{\rho})=-ie^{-\mathcal{L}_Ct}[\tilde{H}_I(t),\tilde{\rho}], \tag{5}$$

where $\tilde{H}_I(t)=g[\tilde{A}_{01}(t)a^\dagger\exp(i\delta t)+\text{H.c.}]$, with $\tilde{A}_{01}(t)=\exp(iH_At)A_{01}\exp(-iH_At)$. We next introduce the operator χ [8,11,18]

$$\chi=e^{-\mathcal{L}_Ct}\tilde{\rho}, \tag{6}$$

which, according to Eq. (5), obeys the equation

$$\begin{aligned}\dot{\chi}(t)=&-ige^{\kappa t}\{[a^\dagger,\tilde{A}_{01}(t)\chi(t)]e^{i\delta t}+[a,\chi(t)\tilde{A}_{10}(t)]e^{-i\delta t}\}\\&-ige^{-\kappa t}\{[\tilde{A}_{01}(t),\chi(t)a^\dagger]e^{i\delta t}\\&+[\tilde{A}_{10}(t),a\chi(t)]e^{-i\delta t}\}.\end{aligned} \tag{7}$$

Only the atom-cavity interaction is involved. Due to the smallness of the coupling constant g, we can perform a second-order perturbation calculation with respect to g by means of standard projection operator techniques. Noting that

$$\text{Tr}_C\chi(t)\equiv\text{Tr}_C\tilde{\rho}(t)\equiv\tilde{\rho}_A(t), \tag{8}$$

we trace out the cavity variables to obtain the master equation for the reduced density matrix operator $\tilde{\rho}_A$ of the atom. Under the Born-Markovian approximation, the resulting master equation is of the form

$$\tilde{\rho}_A(t) = -g^2 \int_0^\infty e^{-(\kappa - i\delta)\tau}[\tilde{\rho}_A(t)\tilde{A}_{10}(t-\tau)\tilde{A}_{01}(t) - \tilde{A}_{01}(t)\tilde{\rho}_A(t)\tilde{A}_{10}(t-\tau)]d\tau - g^2 \int_0^\infty e^{-(\kappa + i\delta)\tau}[\tilde{A}_{10}(t)\tilde{A}_{01}(t-\tau)\tilde{\rho}_A(t) - \tilde{A}_{01}(t-\tau)\tilde{\rho}_A(t)\tilde{A}_{10}(t)]d\tau. \quad (9)$$

Finally transforming $\tilde{\rho}_A$ back to the original picture via $\rho_A = \exp(-iH_At)\tilde{\rho}_A \exp(iH_At)$ and restoring the $\mathcal{L}_A\rho_A$ contribution, we obtain the atomic master equation to be

$$\dot{\rho}_A = -i[H_A,\rho_A] + \mathcal{L}_A\rho_A + \gamma_c(A_{01}\rho_A S_+ + S_-\rho_A A_{10} - A_{10}S_-\rho_A - \rho_A S_+ A_{01}), \quad (10)$$

where $\gamma_c = g^2/\kappa$ specifies the emission rate of the atom into the cavity mode and

$$S_- = \kappa \int_0^\infty d\tau\, e^{-(\kappa + i\delta)\tau}\tilde{A}_{01}(-\tau) = \alpha_0(A_{11} - A_{00}) + \alpha_1 A_{10} + \alpha_2 A_{01}, \quad (11)$$

$$S_+ = (S_-)^\dagger, \quad (12)$$

with

$$\alpha_0 = \frac{\kappa\Omega}{4\bar{\Omega}^2}\left[\frac{2\Delta}{\kappa + i\delta} - \frac{\bar{\Omega} + \Delta}{\kappa + i(\delta - \bar{\Omega})} + \frac{\bar{\Omega} - \Delta}{\kappa + i(\delta + \bar{\Omega})}\right],$$

$$\alpha_1 = \frac{\kappa\Omega^2}{4\bar{\Omega}^2}\left[\frac{2}{\kappa + i\delta} - \frac{1}{\kappa + i(\delta - \bar{\Omega})} - \frac{1}{\kappa + i(\delta + \bar{\Omega})}\right], \quad (13)$$

$$\alpha_2 = \frac{\kappa}{4\bar{\Omega}^2}\left[\frac{2\Omega^2}{\kappa + i\delta} + \frac{(\bar{\Omega} + \Delta)^2}{\kappa + i(\delta - \bar{\Omega})} + \frac{(\bar{\Omega} - \Delta)^2}{\kappa + i(\delta + \bar{\Omega})}\right],$$

where $\bar{\Omega} = \sqrt{\Omega^2 + \Delta^2}$ is a generalized Rabi frequency. Obviously, the coefficients $\alpha_0, \alpha_1, \alpha_2$ are Rabi frequency dependent and resonant when the cavity frequency is tuned to $\delta = 0, \pm\bar{\Omega}$.

The first term of Eq. (10) describes the coherent evolution of the driven atom, whereas the final term represents the cavity-induced atomic decay into the cavity mode. The $\mathcal{L}_A\rho_A$ term describes spontaneous emission of photons into the background modes.

From this cavity-modified master equation, the equations of motion for the reduced density matrix elements are found to be

$$\dot{\rho}_{00} = (2\gamma_1 + \gamma_c\alpha_2 + \gamma_c\alpha_2^*)\rho_{11} - \left(\gamma_c\alpha_0 - i\frac{\Omega}{2}\right)\rho_{01} - \left(\gamma_c\alpha_0^* + i\frac{\Omega}{2}\right)\rho_{10}, \quad (14a)$$

$$\dot{\rho}_{11} = 2\gamma_2\rho_{22} - (2\gamma_1 + \gamma_c\alpha_2 + \gamma_c\alpha_2^*)\rho_{11} + \left(\gamma_c\alpha_0 - i\frac{\Omega}{2}\right)\rho_{01} + \left(\gamma_c\alpha_0^* + i\frac{\Omega}{2}\right)\rho_{10}, \quad (14b)$$

$$\dot{\rho}_{10} = -(\gamma_1 + i\Delta + \gamma_c\alpha_2)\rho_{10} + \gamma_c\alpha_1\rho_{01} + \left(\gamma_c\alpha_0 - i\frac{\Omega}{2}\right)\rho_{00} + \left(\gamma_c\alpha_0 + i\frac{\Omega}{2}\right)\rho_{11}, \quad (14c)$$

$$\dot{\rho}_{21} = -(\gamma_1 + \gamma_2 + \gamma_c\alpha_2^*)\rho_{21} + \left(\gamma_c\alpha_0^* + i\frac{\Omega}{2}\right)\rho_{20}, \quad (14d)$$

$$\dot{\rho}_{20} = -(\gamma_2 + i\Delta)\rho_{20} + i\frac{\Omega}{2}\rho_{21}, \quad (14e)$$

with $\rho_{00} + \rho_{11} + \rho_{22} = 1$ and $\rho_{lk} = \rho_{kl}^*$, where ρ_{00}, ρ_{11}, and ρ_{22} are the atomic populations in the levels $|0\rangle$, $|1\rangle$, and $|2\rangle$, respectively, and ρ_{lk} $(l \neq k)$ describes the atomic coherences. (Here we have dropped the subscript A from the atomic density matrix operator ρ_A for brevity.)

III. SUPPRESSION OF THE AUTLER-TOWNES DOUBLET AND ITS ORIGIN

Our aim is to investigate the influence of the cavity coupled to the two coherently driven levels, $|0\rangle$ and $|1\rangle$ on the Autler-Townes doublet resulting from transitions of one of the two driven levels, for example, $|1\rangle$ to a third level $|2\rangle$. Therefore, we introduce a very weak, frequency-tunable laser beam, indicated by Ω_p in Fig. 1, to probe the transitions between $|1\rangle$ and $|2\rangle$, after the system reaches the steady state.

According to linear-response theory, the probe absorption spectrum can be expressed as the Fourier transformation of the stationary mean value of the two-time commutator of the atomic transition operators: $\lim_{t\to\infty}\langle A_{12}(t)A_{21}(t+\tau)\rangle$, which can be calculated from Eqs. (14d) and (14e) by means of the quantum regression theorem. We find

$$\Lambda_{12}(\omega_p) = \mathrm{Re}\left[\frac{(z + \gamma_2 - i\Delta)(\overline{\rho_{11}} - \overline{\rho_{22}}) + (\gamma_c\alpha_0 - i(\Omega/2))\overline{\rho_{01}}}{(z + \gamma_1 + \gamma_2 + \gamma_c\alpha_2)(z + \gamma_2 - i\Delta) + i(\Omega/2)(\gamma_c\alpha_0 - i(\Omega/2))}\right]_{z = i(\omega_p - \omega_{21})}, \quad (15)$$

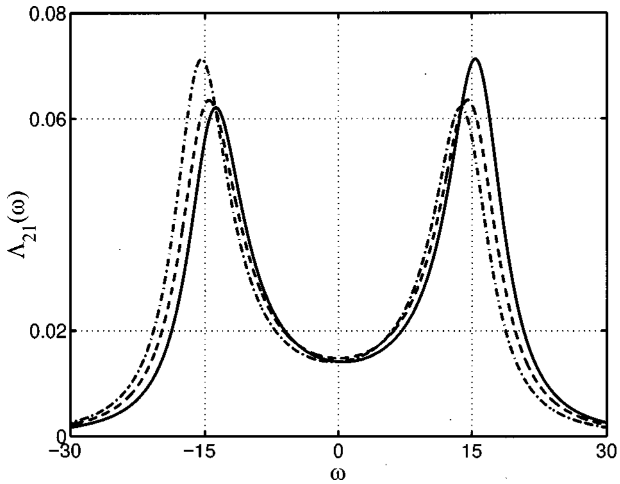

FIG. 2. Autler-Townes spectrum $\Lambda_{12}(\omega)$ as a function of the relative probe frequency $\omega=\omega_p-\omega_{21}$ for $\Omega=30$ and different cavity frequencies $\delta=-\Omega$ (dash-dotted curve), $\delta=0$ (dashed curve), and $\delta=\Omega$ (solid curve). We take the other parameters $\kappa=100$, $g=25$, $\gamma_1=1$, $\gamma_2=0.5$, and $\omega_L=\omega_{10}$ throughout these graphs. (Note that all variables are scaled by γ_1.)

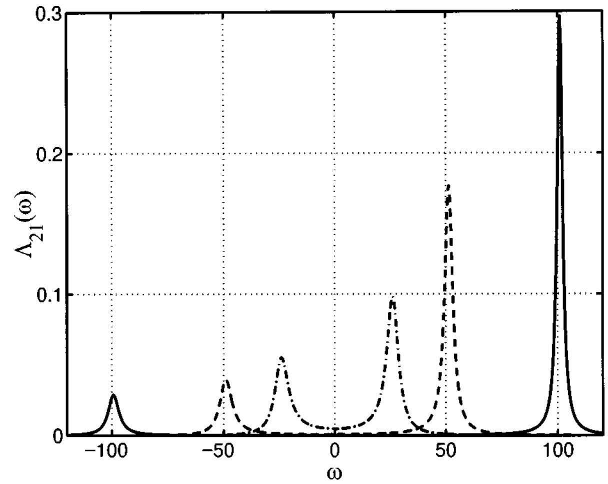

FIG. 4. Same as Fig. 2, but with $\delta=\Omega$ and various Rabi frequencies: $\Omega=50$ (dash-dotted curve), $\Omega=100$ (dashed curve), and $\Omega=200$ (solid curve).

where ω_p is the probe frequency, $\overline{\rho_{01}}$ is the steady-state coherence of the driven levels, and $\overline{\rho_{11}},\overline{\rho_{22}}$ are the steady-state populations of the states $|1\rangle$ and $|2\rangle$, respectively. It is not difficult to see that $\overline{\rho_{22}}=0$ in the steady state.

In the absence of the cavity, the Autler-Townes spectrum is symmetric when $\Delta=0$. In order to see the modifications produced by the cavity, we concentrate on the $\Delta=0$ case in the following. We numerically investigate the Autler-Townes spectrum in Figs. 2–5, where we take the parameters $\kappa=100$, $g=25$, $\gamma_1=1$, $\gamma_2=0.5$, and $\Delta=0$ $(\omega_L=\omega_{10})$ throughout, but vary the Rabi frequency Ω and the cavity resonant frequency ω_C. (Note that all parameters are measured in units of γ_1.) Figure 2, for $\Omega=30$, clearly demonstrates asymmetries in the doublet when the cavity is detuned from the atomic resonance transition frequency ω_{10} of the levels $|0\rangle$ and $|1\rangle$. Specifically, for $\omega_C=\omega_{10}-\Omega$ (the dash-dotted curve), the lower-frequency (left-hand) sideband of the doublet is enhanced and the higher-frequency (right-hand) sideband is suppressed. For $\omega_C=\omega_{10}+\Omega$ (the solid curve), however, enhancement occurs at the higher-frequency sideband and suppression at the lower-frequency sideband. The larger the Rabi frequency, the more significant the enhancement and suppression of the Autler-Townes sidebands. See, for example, in Fig. 3 for $\Omega=200$, where the dash-dotted, dashed, and solid lines represent the doublet for the cavity frequency ω_C tuned to $\omega_{10}-\Omega$, ω_{10}, and $\omega_{10}+\Omega$, respectively. For clarity, we separately plot the lower-frequency sideband of the doublet in Fig. 3(a) and the higher-frequency sideband in Fig. 3(b). In this case, one sideband of the doublet is dramatically suppressed while the other is greatly enhanced, depending upon the cavity resonant frequency. The doublet is symmetric only when $\omega_C=\omega_{10}$, as shown in the dashed curves in Figs. 2 and 3.

Figures 2 and 3 also show that the linewidths and positions of the Autler-Townes doublet vary slightly with the cavity and Rabi frequencies. For large Rabi frequencies (see, for instance, Fig. 3), the lower-frequency sideband is narrower than the higher-frequency one and both sidebands are

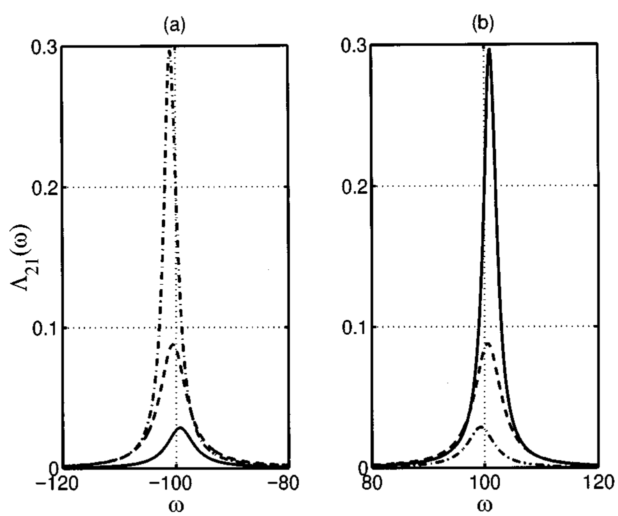

FIG. 3. Same as Fig. 2, but with $\Omega=200$: (a) lower-frequency sideband and (b) higher-frequency sideband.

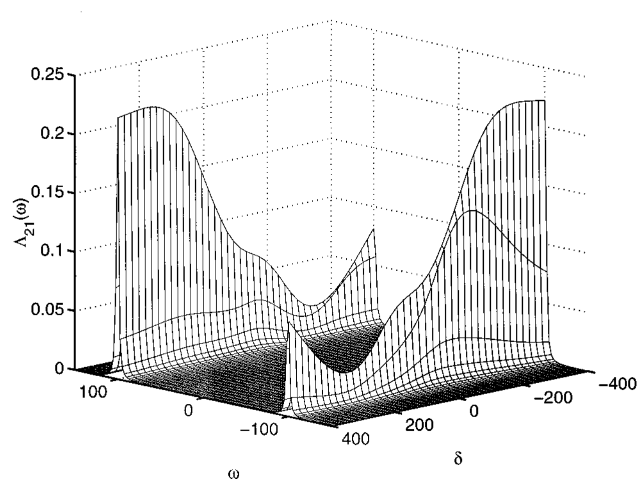

FIG. 5. Three-dimensional Autler-Townes spectrum as a function of the relative probe and cavity frequencies ω and δ for $\Omega=200$.

shifted slightly towards lower frequencies, when $\omega_C=\omega_{10}-\Omega$. The opposite shifts are obtained if $\omega_C=\omega_{10}+\Omega$.

Figure 4, in which the cavity frequency ω_C is fixed at $\omega_{10}+\Omega$ with the Rabi frequency Ω taking the values $\Omega=50$ (the dash-dotted line), 100 (the dashed line), and 200 (the solid line), respectively, clearly exhibits dynamical suppression and enhancement within the Autler-Townes doublet. Increasing the Rabi frequency can lead to significant suppression of the lower-frequency sideband and great enhancement of the higher-frequency sideband. In addition, increasing the Rabi frequency can give rise to narrowing of the enhanced sideband and broadening of the suppressed sideband.

We plot a three-dimensional Autler-Townes spectrum in Fig. 5 for $\Omega=200$, which shows how the doublet is modified when the cavity frequency varies. When $\delta=0$, the doublet is symmetric. When $\delta=200$, however, there is a minimum in the lower-frequency sideband and a maximum in the higher-frequency sideband. (When $\delta=-200$ the situation is reversed.)

Moreover, when $\delta\gg\Omega$, both sidebands tend to the same height and the symmetry in the doublet is recovered. (We have not presented the graph here.) This actually is the same as in free space because the cavity is so far off resonance with any atomic transition frequencies that no photon can be emitted into the cavity mode. In other words, the atom is effectively in free space.

We may gain a physical insight into the suppression and enhancement of the Autler-Townes doublet by working in the semiclassical dressed-state representation, $|\pm\rangle$, which linearly mixes the laser-driven levels $|0\rangle$ and $|1\rangle$. For a resonant driving field ($\omega_L=\omega_{10}$), the dressed states are

$$|+\rangle=\frac{1}{\sqrt{2}}(|0\rangle+|1\rangle), \quad |-\rangle=\frac{1}{\sqrt{2}}(|0\rangle-|1\rangle). \tag{16}$$

In the dressed-state picture, the weak probe laser couples the upper level $|2\rangle$ to a dressed state doublet. See, for instance, Fig. 1(b). The corresponding Autler-Townes spectrum, therefore, is expected to exhibit a doublet. In the secular approximation ($\Omega\gg\gamma_1,\gamma_2,\gamma_c$), the doublet takes the form

$$\Lambda_{12}(\omega_p)=\frac{1}{2}\mathrm{Re}\left[\frac{\overline{\rho}_{++}-\overline{\rho}_{22}}{z+(1/2)(\gamma_1+2\gamma_2+\gamma_c\alpha_2-\gamma_c\alpha_0+i\Omega)}+\frac{\overline{\rho}_{--}-\overline{\rho}_{22}}{z+(1/2)(\gamma_1+2\gamma_2+\gamma_c\alpha_2+\gamma_c\alpha_0-i\Omega)}\right]_{z=i(\omega_p-\omega_{21})}, \tag{17}$$

where

$$\overline{\rho}_{++}=\frac{\gamma_1+\gamma_c\ \mathrm{Re}(\alpha_2-\alpha_1+2\alpha_0)}{2\gamma_1+2\gamma_c\ \mathrm{Re}(\alpha_2-\alpha_1)},$$
$$\overline{\rho}_{--}=\frac{\gamma_1+\gamma_c\ \mathrm{Re}(\alpha_2-\alpha_1-2\alpha_0)}{2\gamma_1+2\gamma_c\mathrm{Re}(\alpha_2-\alpha_1)} \tag{18}$$

are the steady-state population of the dressed states $|\pm\rangle$, respectively, and $\overline{\rho}_{22}=0$.

Obviously, the heights of the Autler-Townes doublet depend strongly on the distribution of population within the dressed-state doublet. In the absence of the cavity (i.e., $\gamma_c=0$), the dressed-state populations are equal, $\overline{\rho}_{++}=\overline{\rho}_{--}=0.5$, and the Autler-Townes doublet is thus symmetric about the atomic resonance frequency ω_{21} of the levels $|1\rangle$ and $|2\rangle$ and separated by the dressed-state level splitting Ω.

In the presence of the cavity coupled to the two coherently driven levels $|0\rangle$ and $|1\rangle$, however, the dressed-state population is cavity-frequency dependent. For example, if the cavity frequency is tuned to ω_{10} ($\delta=0$), then both dressed-state populations are the same and the Autler-Townes components have the same height. However, if the cavity frequency is tuned to $\omega_C=\omega_{10}-\Omega$, the dressed-state populations in the case of large Rabi frequencies ($\Omega\gg\kappa$) are

$$\overline{\rho}_{++}\simeq\frac{\gamma_1+\gamma_c}{2\gamma_1+\gamma_c},$$
$$\overline{\rho}_{--}\simeq\frac{\gamma_1}{2\gamma_1+\gamma_c}. \tag{19}$$

It is obvious that $\overline{\rho}_{++}\neq\overline{\rho}_{--}$, so the Autler-Townes doublet is asymmetric. The amplitude of the lower-frequency sideband, proportional to the population of the dressed state $|+\rangle$, is larger than that of the higher-frequency one, which is proportional to the population of the dressed state $|-\rangle$. Furthermore, if $\gamma_c\gg\gamma_1$, the atomic population can be approximately trapped in the dressed state $|+\rangle$ and therefore the higher-frequency sideband of the Autler-Townes doublet is completely suppressed, while the lower-frequency one is greatly enhanced.

Otherwise, for $\delta=\Omega$, that is, when the cavity is in resonance with the frequency $\omega_{10}+\Omega$, the dressed-state populations are approximately given by

$$\overline{\rho}_{++}\simeq\frac{\gamma_1}{2\gamma_1+\gamma_c},$$
$$\overline{\rho}_{--}\simeq\frac{\gamma_1+\gamma_c}{2\gamma_1+\gamma_c}. \tag{20}$$

As a result, the lower-frequency sideband of the doublet is suppressed whereas the higher-frequency one is enhanced.

TABLE I. Linewidths and positions of the Autler-Townes doublet ($\omega_L=\omega_{10},\Omega\gg\kappa$).

Cavity frequency ($\delta=\omega_C-\omega_{21}$)	W_{LFS}	W_{HFS}	P_{LFS}	P_{HFS}
$\delta=-\Omega$	$\gamma_1+2\gamma_2$	$\gamma_1+2\gamma_2+\gamma_c/2$	$-\Omega/2-3g^2/8\Omega$	$\Omega/2-g^2/4\Omega$
$\delta=0$	$\gamma_1+2\gamma_2+\gamma_c/2$	$\gamma_1+2\gamma_2+\gamma_c/2$	$-\Omega/2-g^2/4\Omega$	$\Omega/2+g^2/4\Omega$
$\delta=\Omega$	$\gamma_1+2\gamma_2+\gamma_c/2$	$\gamma_1+2\gamma_2$	$-\Omega/2+g^2/4\Omega$	$\Omega/2+3g^2/8\Omega$

The asymmetric profile of the Autler-Townes doublet is the direct consequence of the unbalanced population distribution between the dressed states $|\pm\rangle$ induced by the cavity [8]. It is well known that spontaneous and stimulated transitions of a strongly driven atom in the bare state basis are equivalent to spontaneous emission between two adjacent dressed-state manifolds [6]. For the model at hand [see, for example, Fig. 1(b)], if the cavity frequency ω_C is tuned to ω_{10}, which is in resonance with spontaneous transitions of the same dressed state $|\pm\rangle$ from the upper manifold to the lower manifold, both transition rates are the same and the population distributes evenly within the dressed states $|\pm\rangle$. The Autler-Townes doublet is symmetric.

However, if the cavity frequency is tuned to $\omega_{10}-\Omega$, which is in resonance with the atomic transition from the state $|-\rangle$ of the upper dressed-state doublet to the state $|+\rangle$ of the lower dressed doublet and is far off resonance with the other downward transitions, the transition from $|-\rangle$ to $|+\rangle$ is enhanced, while the other is inhibited by the atom-cavity coupling. As a result, the population is most likely to be in the dressed state $|+\rangle$. Otherwise, if the cavity is in resonance with $\omega_{10}+\Omega$, the frequency of the atomic transition from $|+\rangle$ of the upper manifold to $|-\rangle$ of the lower manifold, the transition from $|+\rangle$ to $|-\rangle$ is enhanced while the other is suppressed: Therefore, there is more population in the dressed state $|-\rangle$. The corresponding Autler-Townes doublet is asymmetric.

In addition, Eq. (18) clearly shows that the linewidths and positions of the Autler-Townes doublet are also dependent on the cavity frequency and the driving intensity. For instance, the linewidth and position of the lower-frequency sideband of the doublet are, respectively,

$$W_{LFS}=\gamma_1+2\gamma_2+\gamma_c\,\mathrm{Re}(\alpha_2-\alpha_0),$$
$$P_{LFS}=-\frac{1}{2}[\Omega-\gamma_c\,\mathrm{Im}(\alpha_2-\alpha_0)]. \tag{21}$$

Those of the higher-frequency sideband are

$$W_{HFS}=\gamma_1+2\gamma_2+\gamma_c\,\mathrm{Re}(\alpha_2+\alpha_0),$$
$$P_{HFS}=\frac{1}{2}[\Omega+\gamma_c\,\mathrm{Im}(\alpha_2+\alpha_0)]. \tag{22}$$

The cavity-frequency dependence of the linewidths of the Autler-Townes doublet reflects the fact that the cavity can modify the transition rates of the dressed states $|\pm\rangle$ and the upper state $|2\rangle$, whereas the shifting of the positions is due to the cavity-induced level shifts of the dressed states $|\pm\rangle$. We list the linewidths and positions in the case of $\Omega\gg\kappa$ in Table I.

Both the linewidths and positions of the Autler-Townes doublet are different when the cavity frequency is detuned from the atomic transition frequency ω_{10}. One may be much broader than the other if $\gamma_c\gg\gamma_1,\gamma_2$, depending on the cavity frequency. Likewise, the positions of the doublet are also slightly shifted in the presence of the cavity.

IV. CONCLUSIONS

In this paper we investigate the absorption spectrum of the upper and intermediate levels of a three-level cascade atom in which the intermediate and ground levels are strongly driven by a laser field as well as weakly coupled by a cavity mode. We find that the Autler-Townes spectrum fully reflects the cavity-induced effects on the coherently driven levels. Specifically, the dynamical suppression and enhancement of the Autler-Townes doublet directly reflect the cavity-induced redistribution (polarization) of the dressed-state population since the lower- and higher-frequency sidebands are respectively proportional to the population of the dressed state $|\pm\rangle$, whereas the displacing of the positions of the doublet is the consequence of the cavity-induced level shifts of the dressed state $|\pm\rangle$. In addition, the cavity-dependent linewidths of the doublet are also due to the modification of the transition rates of the dressed state $|\pm\rangle$ and the upper state $|2\rangle$ due to the cavity.

As the dynamical modifications of the radiative properties of a two-level atom inside a cavity whose resonant frequency is tunable and the Autler-Townes doublet of a three-level atom system in free space have been observed in many laboratories [11–13,19], the scheme presented in this paper should provide an experimentally feasible way of investigating cavity effects and may lead to the probing of cavity quantum electrodynamics using Autler-Townes spectroscopic techniques.

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[1] See, for example, *Cavity Quantum Electrodynamics*, edited by P. R. Berman (Academic, London, 1994), and references therein.
[2] E. M. Purcell, Phys. Rev. **69**, 681 (1946).
[3] D. Kleppner, Phys. Rev. Lett. **47**, 233 (1981).
[4] P. Goy, J. M. Raimond, M. Gross, and S. Haroche, Phys. Rev. Lett. **50**, 1903 (1983); R. G. Hulet, E. S. Hilfer, and D. Kleppner, *ibid.* **55**, 2137 (1985); D. P. O'Brien, P. Meystre, and H. Walther, Adv. At. Mol. Phys. **21**, 1 (1985); W. Jhe, A. Anderson, E. A. Hinds, D. Meschede, L. Moi, and S. Haroche, Phys. Rev. Lett. **58**, 666 (1987); D. J. Heinzen, J. J. Childs, J. E. Thomas, and M. S. Feld, *ibid.* **58**, 1320 (1987).

[5] J. J. Sanchez-Mondragon, N. B. Narozhny, and J. H. Eberly, Phys. Rev. Lett. **51**, 550 (1983); G. S. Agarwal, *ibid.* **53**, 1732 (1984); M. G. Raizen, R. J. Thompson, R. J. Brecha, H. J. Kimble, and H. J. Carmichael, *ibid.* **63**, 240 (1989).

[6] C. Cohen-Tannoudji and S. Reynaud, J. Phys. B **10**, 345 (1977); in *Multiphoton Processes*, edited by J. H. Eberly and P. Lambropoulos (Wiley, New York, 1978).

[7] M. Lewenstein, T. W. Mossberg, and R. J. Glauber, Phys. Rev. Lett. **59**, 775 (1987); M. Lewenstein and T. W. Mossberg, Phys. Rev. A **37**, 2048 (1988).

[8] P. Zhou and S. Swain, Phys. Rev. A (to be published).

[9] T. Quang and H. Freedhoff, Phys. Rev. A **47**, 2285 (1993); H. Freedhoff and T. Quang, Phys. Rev. Lett. **72**, 474 (1994); J. Opt. Soc. Am. B **10**, 1337 (1993); **12**, 9 (1995).

[10] J. I. Cirac, H. Ritsch, and P. Zoller, Phys. Rev. A **44**, 4541 (1991); M. Loffler, G. M. Meyer, and H. Walther, *ibid.* **55**, 3923 (1997).

[11] W. Lange and H. Walther, Phys. Rev. A **48**, 4551 (1993); G. S. Agarwal, W. Lange, and H. Walther, *ibid.* **48**, 4555 (1993); **51**, 721 (1995); W. Lange, H. Walther, and G. S. Agarwal, *ibid.* **50**, R3593 (1994).

[12] Y. Zhu, A. Lezama, T. W. Mossberg, and M. Lewenstein, Phys. Rev. Lett. **61**, 1946 (1988); A. Lezama, Y. Zhu, S. Morin, and T. W. Mossberg, Phys. Rev. A **39**, R2754 (1989).

[13] S. H. Autler and C. H. Townes, Phys. Rev. **100**, 703 (1955).

[14] J. E. Field, Phys. Rev. A **47**, 5064 (1993).

[15] M. S. Zubairy, Phys. Lett. A **222**, 91 (1996).

[16] J. von Zanthier, G. S. Agarwal, and H. Walther, Phys. Rev. A **56**, 2242 (1997).

[17] An interaction picture transformed by a unitary operator $\exp(iH_0t)$, where $H_0=\omega_L(a^\dagger a-A_{00})+\omega_{21}A_{22}$, from the Schrödinger picture. Therefore, all quantities involved in the transition $|0\rangle\leftrightarrow|1\rangle$ have the referenced frequency ω_L, while those in the transition $|1\rangle\leftrightarrow|2\rangle$ are referenced by the frequency ω_{21}.

[18] S. M. Barnett and P. L. Knight, Phys. Rev. A **33**, 2444 (1986); J. I. Cirac, *ibid.* **46**, 4354 (1992).

[19] H. R. Gray and C. R. Stroud, Jr., Opt. Commun. **25**, 359 (1978); G. Orriols, Nuovo Cimento B **53**, 1 (1979); M. Kaivola, P. Thorsen, and O. Poulsen, Phys. Rev. A **32**, 207 (1985); Y. S. Bai, T. W. Mossberg, N. Lu, and P. R. Berman, Phys. Rev. Lett. **57**, 1692 (1986); K. J. Boller, A. Imamoglu, and S. E. Harris, *ibid.* **66**, 2593 (1991); M. Xiao, Y. Li, S. Jin, and J. Gea-Banacloche, *ibid.* **74**, 666 (1995); S. Papademetriou, M. F. Van Leeuwen, and C. R. Stroud, Jr., Phys. Rev. A **53**, 997 (1996).