Fluorescence spectrum of a two-level atom interacting with a quantized field in a Kerr-like medium

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A generalized Jaynes-Cummings model when an ideal cavity is filled with a Kerr-like medium is considered. The fluorescence spectrum produced by such a system using the infinity of transitions among the dressed states of the Jaynes-Cummings model is analyzed. A large number of resonances in the spectra are observed that are sensitive to the initial field statistics as well as the nonlinearity of the medium. Also, the fine structure of the spectra gets considerably modified by increasing the nonlinearity of the Kerr-like medium.

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The Jaynes-Cummings model (JCM) and its various extensions have been studied in great detail over the past few decades [1–7]. Many interesting features such as the collapse and revivals [6] of Rabi oscillations, quenching of spontaneous emissions [8], squeezing [9], and chaos [10] have been predicted. Recent developments in superconducting cavities [11,12] have made it possible to observe some of these phenomena. The JCM has been adopted also for investigation of emission spectra from two-level atoms [13,14]. It is well known that the infinite number of dressed states arise from the diagonalization of the JCM Hamiltonian. The transitions among the infinity of the dressed states are important in cavity QED where the radiation field consists of only a few number of photons [1]. The spectral features of the JCM have been examined and it has been shown how the fluorescence characteristic in a very-high-Q cavity are different from those in the free space [14]. The vacuum-field Rabi splitting is also predicted in the literature [15]. This splitting results due to transitions between the ground state and the first excited states of the interacting system comprised of a two-level atom and the cavity field mode. The vacuum-field Rabi splitting is reminiscent of similar effects observed in a more classical situation involving large photon or atoms numbers. Conversely, the cavity resonance splitting effect expected with a single atom in the mode is reminiscent of the cavity pulling effect observed when a macroscopic medium with an index of refraction different from unity is placed inside the cavity. Nevertheless, the vacuum-field Rabi splitting has been seen in absorption experiments at optical frequencies [16,17]. The transitions involving other dressed states can in principle be seen in nonlinear absorption and mixing experiments [18].

In this paper we deal with a modified JCM in which we consider the cavity filled with a Kerr-like medium (where the index of refraction is intensity dependent) and there is a passage in which an atom can travel through the cavity. Such a model has been considered previously and its dynamic properties have been analyzed in detail [19]. We study the influence of nonlinear coupling of a cavity mode to a Kerr-like medium on the dipole-dipole correlation function of an atom interacting with the cavity mode. The spectrum of the radiated field is determined by the Fourier transform of the dipole-dipole correlation function. We present the analytical results for the spectrum for an ideal cavity ($Q \cong \infty$). We demonstrate the effect of input field statistics on the spectrum by considering the field to be either in a coherent state or in a thermal state along with the changes in the spectrum due to the presence of the nonlinearity of the Kerr medium.

One can argue that it is purely academic to study the Kerr effect from a field with small photon numbers, as one would anticipate large photon numbers being necessary to produce significant nonlinear radiative shifts and thus related changes in the fluorescence spectrum. This argument is correct for most of the cases of atom-field interactions, but is quite incorrect for studies of nonlinearities in cavity QED where nonlinearities due to intensity-dependent Stark shifts caused by off-resonant levels have already been observed in micromasers [11,12,20,21] and have been proposed as ingredients for quantum-nondemolition measurements based on the Kerr effect, which is sufficiently large to produce significant radiative shifts for photon numbers down to vacuum level [21]. Alternatively, such radiative shifts are manifested in the cavity-fluorescence spectrum of a two-level atom. Although detecting a one-atom cavity spectrum is a challenging task, the spectrum contains interesting features such as a large number of transition among dressed states, its sensitivity to the input field statistics, etc., which will be discussed subsequently.

The model presented here consists of a single two-level atom surrounded by a nonlinear Kerr-like medium contained inside an infinite-Q cavity sustaining a single mode of the electromagnetic field. The cavity mode is coupled to the two-level atom as well as to the Kerr-like medium. The Kerr-like medium can be modeled as an anharmonic oscillator [19,22] and the total Hamiltonian of the system in the rotating-wave approximation takes the form

$$H=\omega_0 a^\dagger a+\omega_a S_z+\omega b^\dagger b+q b^{\dagger 2} b^2+\eta(a^\dagger b+b^\dagger a)+g(a^\dagger S_- + S_+ a) \qquad (1)$$

Here, a ($a^\dagger$) is the annihilation (creation) operator of the cavity field mode of frequency ω_0, b ($b^\dagger$) is the annihilation (creation) operator of the Kerr-like medium (transition frequency ω), and $S_{\pm}, S_z$ are the atomic operators. Next, we assume that the response time of the Kerr medium is small enough so that the medium follows the field in the adiabatic manner. In such a situation the Hamiltonian [Eq. (1)] can be transformed to an effective Hamiltonian containing only atomic and photon operators [19,22]:

$$H_{\text{eff}}=\omega_c a^\dagger a+\omega_a S_z+\chi a^{\dagger 2}a^2+g(a^\dagger S_- +S_+ a) . \quad (2)$$

In deriving Eq. (2) the adiabatic limit, i.e., frequencies ω and ω_0 are far from each other, has been used (see the Appendix). In the expression of H_{eff}, the frequency ω_c and the third-order nonlinear susceptibility χ are related to η and q by [19,22].

$$\chi=q\eta^4/(\omega-\omega_0)^4 ,\qquad \omega_c=\omega_0-\eta^2/(\omega-\omega_0) . \quad (3)$$

Note that the coupling constant χ is the dispersive part of the third-order nonlinearity of the Kerr-like medium. The eigenstates of H_{eff} [Eq. (2)] are the dressed states given by

$$\begin{aligned}
&H|0,g\rangle=-\tfrac{1}{2}\omega_a|0,g\rangle,\\
&H|\psi_n^{\pm}\rangle=\omega_n^{\pm}|\psi_n^{\pm}\rangle ,\\
&\omega_n^{\pm}=(n+\tfrac{1}{2})\omega_c+n^2\chi\pm\Omega_{n\Delta} ,\\
&\Omega_{n\Delta}=[(\Delta/2+n\chi)^2+g^2(n+1)]^{1/2} ,\\
&\Delta=\omega_c-\omega_a,\quad n=0,1,2,3,\ldots,\infty \qquad (4)\\
&|\psi_n^{\pm}\rangle=\begin{bmatrix}\cos(\Phi_n)\\ \sin(\Phi_n)\end{bmatrix}|n+1,g\rangle+\begin{bmatrix}\sin(\Phi_n)\\ \cos(\Phi_n)\end{bmatrix}|n,e\rangle ,\\
&\tan(\Phi_n)=g\sqrt{(n+1)}^{1/2}/[\Omega_{n\Delta}+(\Delta/2+n\chi)] .
\end{aligned}$$

The time evolution of the states can be expressed as

$$\begin{aligned}
&U|n,e\rangle=C_n|n,e\rangle+D_n|n+1,g\rangle ,\\
&U|n,g\rangle=A_n|n,g\rangle+B_n|n,e\rangle ,\qquad (5)\\
&U=\exp(-iHt) .
\end{aligned}$$

The dipole-dipole correlation function which is proportional to the field radiated sideways is defined as [14]

$$G(t,\tau)=\langle S_+(t+\tau)S_-(t)\rangle . \quad (6)$$

We suppose that at the initial moment ($t=0$) the atom is in the excited state $|e\rangle$ and the cavity mode is prepared in the state $|\zeta\rangle$,

$$|\zeta\rangle=\sum_{n=0}^{\infty}Q_n|n\rangle , \quad (7)$$

with the photon-number distribution $P_n=|Q_n|^2$. The initial-state vector of the system under consideration can be written as

$$|\psi\rangle_0=|\zeta\rangle * |e\rangle=\sum_n Q_n|n,e\rangle . \quad (8)$$

The correlation function $G(t,\tau)$ thus reads as

$$G(t,\tau)=\sum_n P_n C_n(t)C_n^*(t-\tau)A_n(\tau) , \quad (9)$$

where

$$\begin{aligned}
A_n(\tau)&=\cos^2(\Phi_{n-1})\exp(-i\omega_{n-1}^{+}\tau)\\
&\quad+\sin^2(\Phi_{n-1})\exp(-i\omega_{n-1}^{-}\tau) ,\\
C_n(\tau)&=\sin^2(\Phi_n)\exp(-i\omega_n^{+}\tau)\\
&\quad+\cos^2(\Phi_n)\exp(-i\omega_n^{-}\tau) ,\quad n\ge 1
\end{aligned} \quad (10)$$

and

$$A_0(\tau)=\exp(-i\omega_a\tau/2) . \quad (11)$$

Next, we calculate the Fourier transform of the time-averaged dipole-dipole correlation function

$$S(\nu)=\operatorname{Re}\int_0^{\infty}d\tau\exp(-i\nu\tau-\gamma\tau)\overline{G(t,\tau)} . \quad (12)$$

This transform is directly related to the fluorescence spectrum with the identification of γ as the width associated with the detector [13]. On substituting (9)–(11) in (12) and carrying out the integration and other necessary operations we get

$$\begin{aligned}
S(\nu)=P_0\{&\sin^4(\Phi_0)\gamma/[\gamma^2+(\nu-\omega_0^{+}-\omega_a/2)^2]\cos^4(\Phi_0)\gamma/[\gamma^2+(\nu-\omega_0^{-}-\omega_a/2)^2]\}\\
+\sum_n P_n\{&\sin^4(\Phi_n)\cos^2(\Phi_{n-1})\gamma/[\gamma^2+(\nu+\omega_{n-1}^{+}-\omega_n^{+})^2]+\sin^4(\Phi_n)\sin^2(\Phi_{n-1})\gamma/[\gamma^2+(\nu+\omega_{n-1}^{-}-\omega_n^{+})^2]\\
&+\cos^4(\Phi_n)\cos^2(\Phi_{n-1})\gamma/[\gamma^2+(\nu+\omega_{n-1}^{+}-\omega_n^{-})^2]\\
&+\cos^4(\Phi_n)\sin^2(\Phi_{n-1})\gamma/[\gamma^2+(\nu+\omega_{n-1}^{-}-\omega_n^{-})^2]\} .
\end{aligned} \quad (13)$$

It is easy to see that $S(\nu)$ consists of the resonant structures which arises from transitions among the various dressed states [14],

$$\nu-\omega=\begin{cases}\tfrac{1}{2}(\Delta\pm\Omega_{0\Delta}), & \psi_0^{\pm}\Longrightarrow|0,g\rangle,\\ \chi\{(n-1)^2-n^2\}\mp\tfrac{1}{2}(\Omega_{n-1,\Delta}-\Omega_{n,\Delta}), & \psi_n^{\pm}\Longrightarrow\psi_{n-1}^{\pm},\\ \chi\{(n-1)^2-n^2\}\pm\tfrac{1}{2}(\Omega_{n,\Delta}+\Omega_{n-1,\Delta}), & \psi_n^{\pm}\Longrightarrow\psi_{n-1}^{\mp} .\end{cases} \quad (14)$$

The intensity of these structures is dependent on the detuning Δ susceptibility χ, and the value of n. For $\Delta=0$, $\Omega_{n,\Delta}=\chi n^2+[\chi^2 n^2+g^2(n+1)]^{1/2}$ and the peaks occur depending on the weightage factor at

$$\nu-\omega=\pm g,$$

$$\chi[n^2-(n-1)^2)\pm\{[n^2\chi^2+g^2(n+1)]^{1/2}+[(n-1)^2\chi^2+g^2 n]^{1/2}\}^{1/2}, \quad (15)$$

$$\chi[n^2-(n-1)^2]\pm\{[n^2\chi^2+g^2(n+1)]^{1/2}-[(n-1)^2\chi^2+g^2 n]^{1/2}\}^{1/2},$$

$$\text{for } n>0 .$$

The peaks at $\pm g$ correspond to the usual vacuum-field Rabi splitting and have been considerably discussed elsewhere. It may be important here to note that in the spectrum (i) the structure for $n>0$ arises from the transitions among different dressed states, (ii) the form of the structure is dependent on the initial photon statistics of the field as well as on the parameter χ. In particular, we consider the cases where the cavity field is prepared in the coherent state with

$$P_n=\exp(-\bar{n})\bar{n}^n/n! \quad (16)$$

or in the thermal or chaotic state with

$$P_n=\bar{n}^n/(1+\bar{n})^{n+1} . \quad (17)$$

First we present numerical results for the coherent field by substituting (16) in (13) and carrying out the sum over n. Here we take the detector width $\gamma/g=0.1$ for all the cases and define $\delta=(\nu-\omega_c)/g$. Figures 1, 2, and 3 corre-

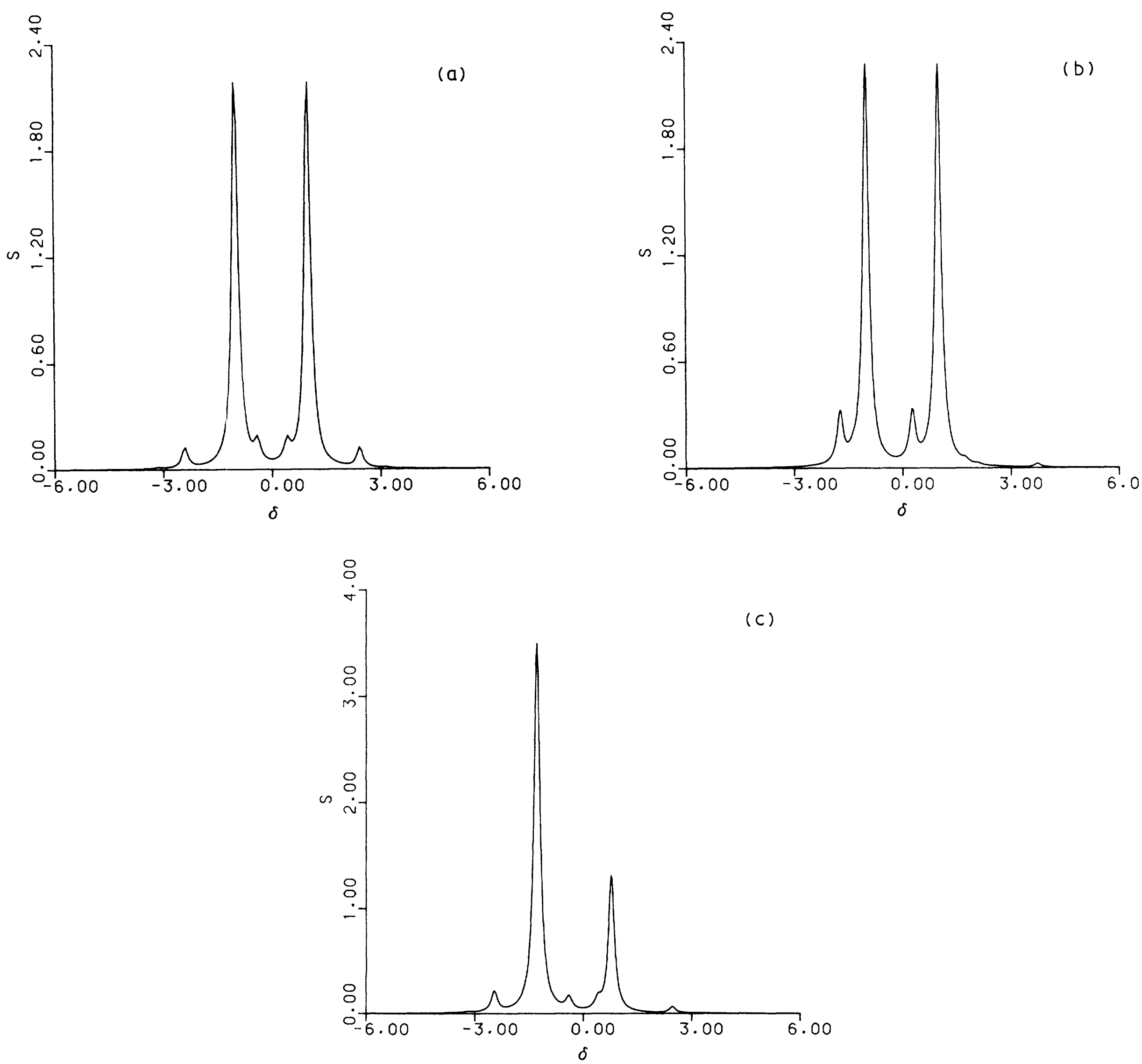

FIG. 1. The time-averaged spectrum $S(\nu)$ in an ideal cavity as a function of $\delta=(\nu-\omega_c)/g$ for an input coherent field with $\bar{n}=0.1$, $\gamma/g=0.1$ and (a) $\Delta=0$, $\chi=0$, (b) $\Delta=0$, $\chi=1.0$, (c) $\Delta=0.5$, $\chi=0.0$.

spond to the spectra for a coherent field with increasing values of average photon number $\bar{n}$. From Fig. 1(a) where $\Delta=0$ and $\chi=0, \bar{n}=0.1$ it is clear that for small value of $\bar{n}$, the predominant peaks in the spectrum are due to the vacuum-field Rabi oscillations which occur at $\delta=\pm 1$. The effect of $\chi \neq 0$ and $\Delta \neq 0$ has been shown in Figs. 1(b) and 1(c), respectively. The effects due to χ are less pronounced over those due to Δ when $\bar{n}$ is small. It is also seen from the expression (13) of the spectrum that for nonzero detunings the weight factors of the peaks get considerably modified and an asymmetry is introduced in the spectrum. This dependence arises from $\cos(\Phi_n)$ and $\sin(\Phi_n)$ in (13). However, as $\bar{n}$ is increased additional resonances which are of the comparable magnitude to the vacuum-field Rabi peaks start appearing in the spectrum. This is shown in Fig. 2(a) for $\bar{n}=1.0$ and $\chi=0$. We find a number of peaks corresponding to $\delta=[\sqrt{(n+1)}+\sqrt{n}\,]$ for $n=1,2,3,\ldots$, etc. towards the right of $\delta=1$ and $(\sqrt{2}-1)$ towards the left of $\delta=1$. The overall spectrum is symmetric around $\delta=0$. Introduction of nonlinearity [$\chi=0.5$, $\bar{n}=1$, Fig. 2(b)] causes drastic changes in the spectrum. The symmetry around $\delta=0$ is removed and most of the peaks are situated towards the right-hand side of the $\delta=0$. For simplicity we have considered $\Delta=0$ in all these cases. For larger $\bar{n}$ [Fig. 3(a), $\chi=0$, $\bar{n}=10$] the spectrum matches with the Mollow spectrum but along with a fine structure due to the dressed states. It is important to note that this fine structure persists for a very large value of $\bar{n}$ (which we have verified separately) and the separation between the successive fine-structure peaks is of the order of $1/\sqrt{n}$. The presence of nonzero χ [Fig. 3(c), $\bar{n}=10$, $\chi=0.2$] changes both the quantitative and qualitative nature of the spectrum. The relative magnitude of the peaks is greatly changed and their separa-

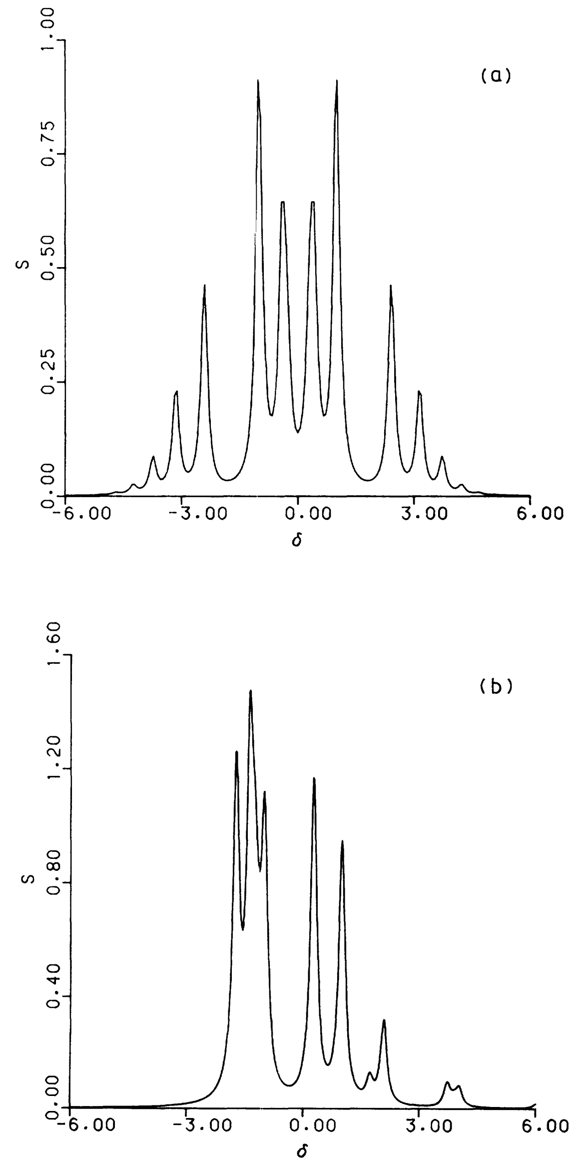

FIG. 2. The same as Fig. 1. but now $\bar{n}=1.0$, $\Delta=0$ and (a) $\chi=0$, (b) $\chi=0.5$.

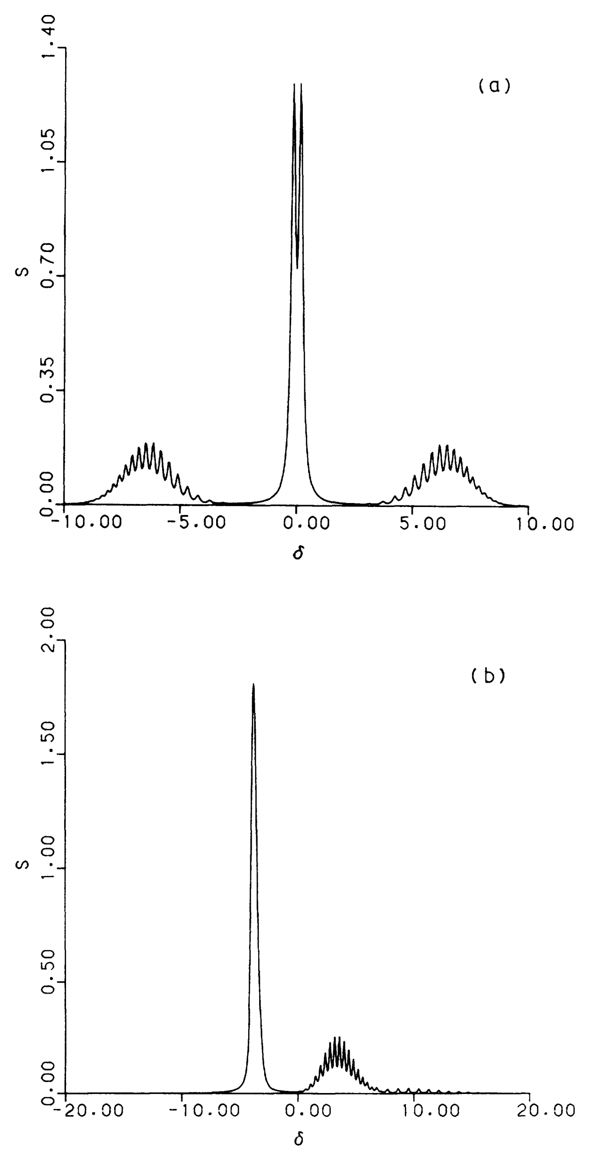

FIG. 3. The same as Fig. 1 but now $\bar{n}=10$, $\Delta=0$ and (a) $\chi=0$, (b) $\chi=0.2$.

tion is no longer $1/\sqrt{n}$. The quantitative explanation of these drastic changes in the spectrum can be attributed to the changes brought out by the χ in the weighting factors of the various peaks in Eq. (13).

In order to get physical insight in the system we rewrite the effective Hamiltonian [Eq. (2)] as

$$H=\omega_a S_z+g[a^{\dagger}S_-\exp(i\omega_c t+\text{c.c.}]+\chi a^{\dagger 2}a^2 . \quad (18)$$

Further we transform it by using the inverse of the Glauber coherent state generator

$$D(\beta)=\exp(\beta a^{\dagger}-\beta^* a) , \quad (19)$$

where $|\beta\rangle$ is a coherent state such that $D^{-1}(\beta)|\beta\rangle=|0\rangle$. With this transformation we obtain the following new Hamiltonian:

$$H^T=\omega_a S_z+g[(a^{\dagger}+\beta^*)S_+\exp(i\omega_c t)+\text{c.c.}] +\chi(a^{\dagger}+\beta^*)^2(a+\beta)^2 . \quad (20)$$

Here in this transformed representation the two-level atom is in the excited state and the cavity field mode is in its vacuum state initially. The transition from the ground state $|g\rangle$ to the excited state $|e\rangle$ is governed by the semiclassical Hamiltonian

$$H_1=g(\beta^* S_-\exp(i\omega_c t)+\text{c.c.}) \quad (21)$$

and the quantum electrodynamic Hamiltonian

$$H_2=g(a^{\dagger}S_-\exp(i\omega_c t)+\text{c.c.}) . \quad (22)$$

Note that the change in photon numbers in the cavity mode field is brought out by the Hamiltonian H_2. Due to the presence of the nonlinear interaction term we have an additional part in the form

$$H_3=2\chi[(a^{\dagger}a+|\beta|^2)\beta^* a+\text{c.c.}]+\chi[a^{\dagger 2}\beta^2+\text{c.c.}] . \quad (23)$$

This term is responsible for the change of photons in the

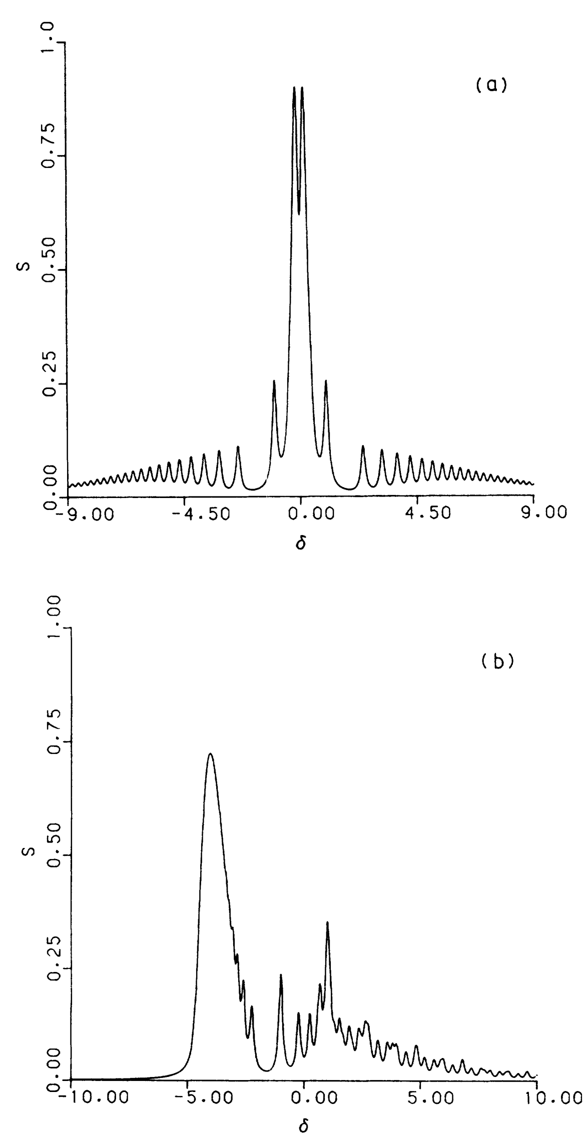

FIG. 4. The same as Fig. 1 but now for thermal field with $\bar{n}=10$, $\Delta=0$ and (a) $\chi=0$, (b) $\chi=0.2$.

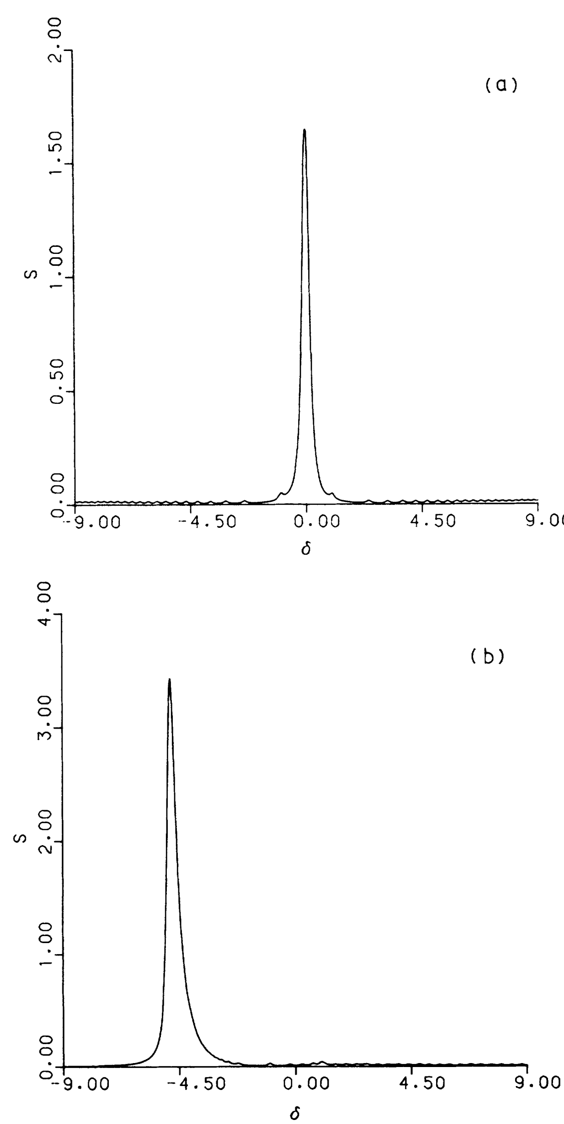

FIG. 5. The same as Fig. 4 but now $\bar{n}=100$, $\Delta=0$ and (a) $\chi=0$, (b) $\chi=0.2$.

mode while keeping the state of the atom intact. Hence in the case of strong nonlinear coupling, i.e., for $\chi\bar{n} > g\sqrt{\bar{n}}$, the term H_3 dominates and there is a reduction in the tendency of emission of photon from the excited state. This leads to a drastic change in the fluorescence spectrum of the system.

The fluorescence spectrum for an input thermal (chaotic) field are shown in Figs. 4 and 5, respectively, for $\bar{n}=10$ and 100, respectively. For small $\bar{n}$ the spectra for coherent as well as thermal input fields are alike. The effect of field statistics can be clearly observed by comparing Figs. 3 and 4 in the absence of the nonlinear medium as well as in the presence of the nonlinear medium. The transitions among the infinity of dressed states is present also in this case which give rise to the fine structure of the spectrum. However, increasing $\bar{n}$ up to 100 (Fig. 5) causes a further change in the quality of the spectrum.

Thus, these results show that the nonlinear coupling of the cavity field mode to the Kerr-like medium leads to drastic changes in the fluorescence spectrum for both coherent as well as thermal input fields. To this date the observation of the single atom-cavity spectrum remains yet to be achieved. The difficulty resides in the marginally small values of the Rabi frequency in the optical domain. The vacuum splitting effects, however, have been observed on systems made of a collection of N identical atoms coupled to a high-Q cavity mode. In the optical domain, it has been observed on the transmission spectrum of small Febry-Pérot cavities crossed by sodium or barium atomic beams. The minimum number of atoms required to see such splitting is about 30 in these experiments [16,17]. However, for Rydberg atoms coupled to microwave cavities, the number of atoms required to see such splitting would be much less. Under certain conditions, even a Rydberg atom may behave like an anharmonic oscillator [22], thus providing a Kerr-like dispersive coupling to cavity field. Hence with the present state of the art micromaser systems it may be possible to observe at least some of the features reported here. Detecting the effect of a single atom on the transmission spectrum of a cavity remains a challenging goal. However, it should be noted that the experiments in the cavity QED domain have completely changed the classical scenario involving a large number of photons to detect such effects. Even as low as one external photon on the average can affect significantly the spectrum produced by the excited atom in the cavity. In conclusion, these features reveal some noteworthy effects in the cavity QED of two-level atoms.

APPENDIX

In this appendix we derive an expression for the effective Hamiltonian utilized for the Kerr-like medium in Eq. (2). For this purpose we start with the Hamiltonian

$$H=\omega_0 a^\dagger a+\omega b^\dagger b+q b^{\dagger 2}b^2+\eta(a^\dagger b+b^\dagger a) \ . \quad \text{(A1)}$$

The Heisenberg equation for the evolution of medium operator in the interaction picture is given by

$$\dot{B}=-2iq(B^\dagger B)B-i\eta e^{i\bar{\Delta}t}A \ , \quad \text{(A2)}$$

with

$$A=e^{i\omega_0 t}a \ , \quad B=e^{i\omega t}b \ , \quad \bar{\Delta}=\omega-\omega_0 \ . \quad \text{(A3)}$$

Next, we adiabatically eliminate [22,23] medium operators by formally integrating the Eq. (A2), assuming that the field operators in the interaction representation vary slowly with respect to optical frequencies, and the response time of the Kerr medium is very fast. Further, we make use of the conservation law [22]

$$a^\dagger a+b^\dagger b=N=\text{const}, \quad [H,N]=0 \ . \quad \text{(A4)}$$

For very large detuning $\bar{\Delta}\gg 0$, there is no exchange of photons with the medium. Therefore, $a^\dagger a$ remains constant, implying $b^\dagger b$ is a constant of motion in Eq. (A2). The formal integration of (A2) thus gives (under the assumption that $\bar{\Delta}$ is very large)

$$B=-(\eta/\bar{\Delta})e^{i\bar{\Delta}t}A \ . \quad \text{(A5)}$$

Reverting back from interaction picture and substituting for b in Eq. (A1) yields [22]

$$H_{\text{eff}}=\omega_c a^\dagger a+\chi a^{\dagger 2}a^2 \ , \quad \text{(A6)}$$

where ω_c and χ are as defined in Eq. (3).

[1] E. T. Jaynes and F. W. Cummings, Proc. IEEE **51**, 89 (1963).
[2] M. Tavis and F. W. Cummings, Phys. Rev. **118**, 692 (1969).
[3] P. Meystre, A. Quattropani, and H. P. Bates, Phys. Lett. **49A**, 856 (1974).
[4] T. Von Foerster, J. Phys. A **8**, 95 (1975).
[5] P. L. Knight and P. M. Radmore, Phys. Rev. A **26**, 676 (1982).
[6] Y. I. Yoo and J. H. Eberly, Phys. Rep. **118**, 289 (1985).
[7] R. R. Puri and G. S. Agarwal, Phys. Rev. A **35**, 3433 (1987).
[8] F. W. Cummings, Phys. Rev. A **33**, 1683 (1986).
[9] P. Meystre and M. S. Zubairy, Phys. Lett. **89A**, 390 (1982).
[10] R. F. Fox and J. Eidson, Phys. Rev. A **34**, 482 (1986).
[11] D. Meschede, H. Walther, and G. Muller, Phys. Rev. Lett. **54**, 551 (1985).
[12] G. Rempe, H. Walther, and N. Klein, Phys. Rev. Lett. **58**, 353 (1987).
[13] J. H. Eberly and K. Wodkiewicz, J. Opt. Soc. Am. **67**, 1252 (1977).
[14] J. J. Sanchez-Mondragon, N. B. Narozhny, and J. H. Eberly, Phys. Rev. Lett. **51**, 550 (1983); G. S. Agarwal, R. K. Bullough, and N. Nayak, Opt. Commun. **85**, 202 (1991).
[15] Y. Kaluzny, P. Goy. M. Gross, J. M. Raimond, and S.

Haroche, **51**, 1175 (1983).
[16] M. G. Raizen, R. J. Thomson, R. J. Brecha, H. J. Kimble and H. J. Carmichael, Phys. Rev. Lett. **63**, 240 (1989).
[17] Y. Zhu, D. J. Gauthier, S. E. Morin, Q. Wu, H. J. Carmichael, and T. W. Mossberg, Phys. Rev. Lett. **64**, 2499 (1990).
[18] G. V. Varda, M. Sanjay Kumar, and G. S. Agarwal, Opt. Commun. **62**, 328 (1987).
[19] V. Buzek and I. Jex, Opt. Commun. **78**, 425 (1990).
[20] H. Moya-Cessa, V. Buzek, and P. L. Knight, Opt. Commun. **85**, 267 (1991).
[21] M. Brune, S. Haroche, V. Lefevre, J. M. Raimond, and N. Zagury, Phys. Rev. Lett. **65**, 976 (1990).
[22] G. S. Agarwal and R. R. Puri, Phys. Rev. A **39**, 2969 (1990).
[23] L. M. Narducci, W. W. Eidson, P. Furcinitti, and D. S. Eteson, Phys. Rev. A **16**, 1665 (1977).