Interaction of a two-level atom with a squeezed vacuum: Photon statistics and spectra

Perry R. Rice and Christopher A. Baird
Department of Physics, Miami University, Oxford, Ohio 45056
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We consider the interaction of a two-level atom with a squeezed vacuum, both in free space and in a cavity of moderate Q. In the latter case, only vacuum modes coupled to the cavity are squeezed. In both cases we calculate the following quantities for the fluorescent light fields: the second-order intensity correlation function $g^{(2)}(\tau)$, the spectrum of squeezing, the coherent spectrum, and the spectrum obtained in a pump-probe absorption measurement. Nonclassical behavior is discussed and comparison to an ordinary vacuum and thermal fields is made. [S1050-2947(96)05105-0]

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I. INTRODUCTION

In recent years, squeezed light sources have become available in the laboratory and attention has turned to their interaction with optical systems. In particular, much attention has been directed at modifying the radiative properties of an atom via interaction with a squeezed light. This began with the seminal work of Gardiner [1], who showed that the decay rate of the atomic polarization quadratures was phase dependent. Carmichael, Lane, and Walls [2] (hereafter referred to as CLW) considered resonance fluorescence when the atom is immersed in a squeezed vacuum. They predicted that for weak driving fields, independent of the relative phase between the driving field and the squeezed vacuum, the incoherent spectrum would narrow as the amount of squeezing was increased. In the limit of strong squeezing, a δ-function spectrum would be obtained. For stronger driving fields, the central peak of the Mollow spectrum could be broadened or narrowed, depending on the relative phase between the strong driving field and the squeezed vacuum. The photon number distribution $P(n)$ has been calculated by Jagatap and Lawande [3], showing phase-sensitive behavior for strong fields.

It was realized early on that experiments would probably require some sort of cavity system, as it is impractical to squeeze all of the vacuum modes that interact with an atom. Several theoretical calculations having to do with squeezing only the cavity modes have been presented. Savage [4] has calculated that for large Jaynes-Cummings coupling g and strong excitation, the width of the Rabi sidebands could be narrowed, but not below the natural linewidth. In a cavity of moderate Q, Courty and Reynaud [5] found that one of the Rabi sidebands could be suppressed for the proper detuning, essentially turning off spontaneous emission from one of the dressed states. Kennedy [6] has found similar behavior in the many-atom case. Rice and Pedrotti [7] have considered an extension of the work of CLW, again for an atom in a cavity of moderate Q. They found that it was possible to squeeze away the cavity enhancement part of the linewidth, but that to obtain measurably subnatural linewidths, the fraction of 4π sr the cavity mode subtends must be significant. This system has also been considered by Cirac [8], who investigated both the fluorescent spectrum and the steady-state population inversion, as discussed by Savage and Lindberg [9]. For very strong driving fields and finite-bandwidth squeezed light centered on the Rabi sidebands, Parkins [10], and Cirac and Sanchez-Soto [11] have found narrowing of one of the Rabi sidebands. Parkins, Zoller, and Carmichael [12] have calculated the fluorescent spectrum of a strongly coupled atom-cavity system, where the driving field is tuned to the one-photon dressed state resonance and have predicted narrowing. One notable example of a calculation in which the atom interacts with only one mode that is squeezed is the work of Vyas and Singh [13], who considered resonance fluorescence in the weak-field limit when the usual coherent driving field was replaced by the squeezed output of an optical parametric oscillator.

It is the purpose of this paper to discuss various aspects of this problem that have not been addressed to date: the second-order intensity correlation function $g^{(2)}(\tau)$, the spectrum of squeezing, the coherent spectrum, and the spectrum obtained in a pump-probe absorption measurement. The latter has been considered in free space by Ritsch and Zoller [14], who have also included finite-bandwidth effects. We calculate these quantities both for an atom completely embedded in a squeezed vacuum, and also for an atom inside a cavity, where the cavity modes are driven by a squeezed vacuum, but the vacuum modes the atom couples directly to (out the side of the cavity, for example) are ordinary vacuums. We use this phrase fully realizing that there is nothing ordinary about the vacuum in quantum mechanics, as it contains a wealth of interesting phenomena. We use the term simply to describe the case where none of the modes the atom interacts with are squeezed. Comparisons to the case when a thermal field replaces the squeezed field are made. One purpose we have is to discuss when application of a squeezed field does something interesting (i.e., nonclassical), and when the changes resulting from replacing the ordinary vacuum with a squeezed field depend solely on the fact that the input field has a nonzero photon occupation number, and not on the special phase dependent properties of squeezed light. In Sec. II we describe the two physical models under study, and discuss the Bloch equations for each system. In Secs. III and IV, we present the results for the second-order intensity correlation function $g^{(2)}(\tau)$ and coherent spectrum, respectively. In Secs. V and VI, we discuss the pump-probe

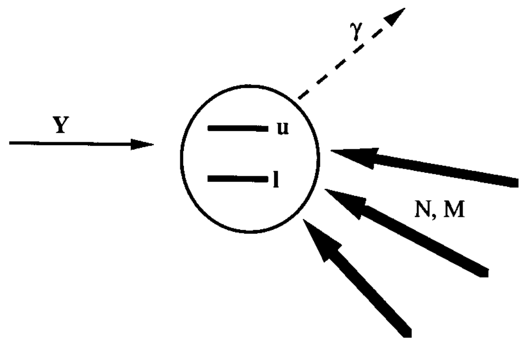

FIG. 1. Two-level atom driven by a coherent field Y, spontaneous emission rate γ, immersed in a squeezed vacuum characterized by M,N. Hereafter referred to as the SV case.

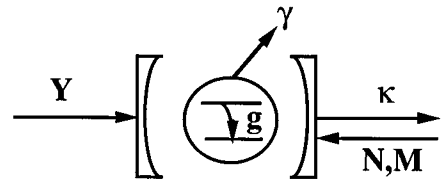

FIG. 2. Two-level atom in an optical cavity, cavity loss rate κ, spontaneous emission rate Γ into noncavity modes, driven by a coherent field Y, with a squeezed vacuum characterized by an M,N incident on the output mirror. Hereafter referred to as the CS case.

absorption spectrum and the spectrum of squeezing. Throughout, comparisons to coherent and thermal fields will be discussed, and we conclude in Sec. VII.

II. PHYSICAL MODELS

The first model, which we refer to as the squeezed-vacuum (SV) model, is shown in Fig. 1. It consists of a two-level atom in free space, embedded in a squeezed vacuum. The atom has a free-space lifetime of γ^{-1}, and interacts with a driving field, characterized by the dimensionless parameter Y. The squeezed field is characterized as usual by the parameters N and M. N of course is just the mean photon number of the squeezed vacua, whereas M describes the phase-sensitive properties of the squeezed state. For a coherent squeezed state, such as that produced by an optical parametric oscillator, N and M satisfy the following relation:

$$|M|=\sqrt{N(N+1)}. \tag{1}$$

If one were to set $M=0$ while retaining a nonzero N, the resulting model would be that of an atom embedded in a thermal background. If a result is independent of M, then one would obtain the same result for a thermal field with the same photon occupation number.

It was shown by CLW that this system was described by the following Bloch equations on resonance:

$$\langle\dot{\sigma}_x\rangle=-\gamma(N-|M|\cos\phi+\tfrac{1}{2})\langle\sigma_x\rangle+\gamma|M|\sin\phi\langle\sigma_y\rangle, \tag{2}$$

$$\langle\dot{\sigma}_y\rangle=-\gamma(N+|M|\cos\phi+\tfrac{1}{2})\langle\sigma_y\rangle+\gamma|M|\sin\phi\langle\sigma_x\rangle -(Y/\sqrt{2})\langle\sigma_z\rangle, \tag{3}$$

$$\langle\dot{\sigma}_z\rangle=-\gamma(2N+1)\left(\langle\sigma_z\rangle+\frac{1}{2N+1}\right)+(Y/\sqrt{2})\langle\sigma_y\rangle, \tag{4}$$

with $Y=4\sqrt{2}\mu E_{ext}/\hbar$. Here μ is the (scalar) transition dipole matrix element for the atom, and ϕ is the relative phase between the driving field and the squeezed vacua. Please note that these equations differ from those of CLW in two ways: our definition for $\sigma_\pm$ is half as large, and the phase of our driving field is shifted by $\pi/2$. The conventions that we use are consistent with those of Rice and Pedrotti. Implicit in all of our calculations (as in those of CLW) is the assumption that a small window of unsqueezed vacuum modes exists to view the fluorescence from the atom, uncontaminated by a strong squeezed field, which could swamp the fluorescence. The second model, which we refer to as the cavity-squeezed (CS) model, is shown in Fig. 2. It is essentially an extension of the SV model to make explicit the ''window'' through which the fluorescence is viewed. This is provided for in the unsqueezed modes out the side of the cavity, which we take to be an ordinary vacuum. The cavity is characterized by its field decay rate κ, and the coupling of the atom to the cavity mode is described by the Jaynes-Cummings coupling parameter g. In terms of cavity parameters that can be changed experimentally, these are defined in the following manner:

$$g=\mu(\omega_0/\hbar\epsilon_0 AL)^{1/2}, \tag{5}$$

$$\kappa=\pi/\mathscr{F}\tau_C, \tag{6}$$

where A is the transverse area of the field mode, L is the length of the cavity, $\mathscr{F}$ is the cavity finesse, and τ_C is the round-trip cavity time. Rice and Pedrotti have shown that if all of the atomic dynamics occur within the bandwidth of the cavity [$\Gamma(1+2C)$ and Rabi frequency $\Omega\ll\kappa$, the bad-cavity limit], the system is described by the following Bloch equations:

$$\langle\dot{\sigma}_x\rangle=-\frac{\Gamma}{2}[1+2C(1+2N-2|M|\cos\phi)]\langle\sigma_x\rangle +\Gamma(2C|M|\sin\phi)\langle\sigma_y\rangle, \tag{7}$$

$$\langle\dot{\sigma}_y\rangle=-\frac{\Gamma}{2}[1+2C(1+2N+2|M|\cos\phi)]\langle\sigma_y\rangle +\Gamma(2C|M|\sin\phi)\langle\sigma_x\rangle-(Y/\sqrt{2})\langle\sigma_z\rangle, \tag{8}$$

$$\langle\dot{\sigma}_z\rangle=-\Gamma[1+2C(2N+1)]\left(\langle\sigma_z\rangle+\frac{1}{1+2C(2N+1)}\right) +(Y/\sqrt{2})\langle\sigma_y\rangle, \tag{9}$$

where Γ is the decay rate of the atom into modes other than the privileged cavity mode (approximately γ unless the cavity subtends a large fraction of 4π sr), and Y and C are defined by

$$Y=2\sqrt{2}g\mathscr{E}/\kappa\Gamma, \tag{10}$$

$$C=g^2/\kappa\Gamma. \tag{11}$$

Here, $\mathscr{E}$ is the dimensionless intracavity driving field

$$\mathscr{E}=-\kappa(\mathscr{F}/\pi)(\omega_0/\hbar\epsilon_0 AL)^{1/2}\sqrt{T_1}e^{i\theta_1}E_{ext}, \tag{12}$$

where T_1 and θ_1 are, respectively, the transmission coefficient and phase change at the mirror through which the coherent field E_{ext} is inserted. We consider an essentially single-ended cavity, letting T_1 tend to zero and E_{ext} tend to infinity. Again, Y is a scaled driving field, and C is just the single-atom cooperativity parameter familiar from the optical bistability literature. It is assumed here that the other field injected into the cavity is squeezed, and that the bandwidth of the squeezing is large compared to the cavity bandwidth. We note that both systems are essentially the same, with three different decay rates γ_x, γ_y, and γ_z. The nonsqueezed vacuum modes in the CS model lead to a limit on the size of nonclassical effects. For example, the linewidth of the incoherent spectrum can only be made as narrow as Γ, which is significantly less than the free-space rate γ only when the cavity mode encloses a large fraction of 4π sr.

These equations simplify greatly for two choices of ϕ: 0 and $\pi/2$. These are the cases of most interest to us, and then both systems obey the following Bloch equations for $\phi=0$, where we have scaled time in units of γ in the SV case, and Γ in the CS case:

$$\langle\dot{\sigma}_x\rangle=-\gamma_x\langle\sigma_x\rangle, \tag{13}$$

$$\langle\dot{\sigma}_y\rangle=-\gamma_y\langle\sigma_y\rangle-(Y/\sqrt{2})\langle\sigma_z\rangle, \tag{14}$$

$$\langle\dot{\sigma}_z\rangle=-\gamma_z(\langle\sigma_z\rangle+\delta)+(Y/\sqrt{2})\langle\sigma_y\rangle, \tag{15}$$

where for the SV case we have

$$\gamma_x=(N+M+1/2), \tag{16}$$

$$\gamma_y=(N-M+1/2), \tag{17}$$

$$\gamma_z=(2N+1), \tag{18}$$

$$\delta=-1/2, \tag{19}$$

and in the CS case we have

$$\gamma_x=\tfrac{1}{2}[1+2C(2N+2M+1)], \tag{20}$$

$$\gamma_y=\tfrac{1}{2}[1+2C(2N-2M+1)], \tag{21}$$

$$\gamma_z=[1+2C(2N+1)], \tag{22}$$

$$\delta=-\tfrac{1}{2}(1+2C). \tag{23}$$

Results for $\phi=\pi/2$ are obtained simply by changing the sign of M. The SV case can be recovered from the CS case by letting $\Gamma\to\gamma/2$ and $C\to 1/2$. We note that in the strong squeezing limit, M is approximated by

$$M=N+\frac{1}{2}-\frac{1}{8N}. \tag{24}$$

It is instructive to examine the eigenvalues of the Bloch equations, as they play a large role in determining the results presented here. They are

$$\lambda_1=-\gamma_x, \tag{25}$$

$$\lambda_2=-\gamma_y+\gamma_z/2+\Omega, \tag{26}$$

$$\lambda_2=-\gamma_y+\gamma_z/2-\Omega, \tag{27}$$

where $\Omega=(1/2)[(\gamma_y-\gamma_z)^2-2Y^2]^{1/2}$. We now turn our attention to photon statistics and spectra.

III. SECOND-ORDER INTENSITY CORRELATION FUNCTION $g^{(2)}(\tau)$

The second-order intensity correlation function $g^{(2)}(\tau)$ is defined as

$$g^{(2)}(\tau)=\frac{\langle\sigma_+(0)\sigma_+(\tau)\sigma_-(\tau)\sigma_-(0)\rangle}{\langle\sigma_+(0)\sigma_-(0)\rangle^2}, \tag{28}$$

where we have defined

$$\sigma_\pm=\sigma_x\pm i\sigma_y, \tag{29}$$

which is the probability of detecting a photon at time τ given that one was detected at time 0, relative to that same probability for a field in a coherent state. In the steady state, we may use the quantum regression theorem to evaluate the necessary correlation functions, which then obey the Bloch equations but with different initial conditions. The result is

$$g^{(2)}(\tau)=1-\exp\{-(\gamma_y+\gamma_z)\tau\}\left[\cosh(\Omega\tau)+\frac{\Phi}{\Omega}\sinh(\Omega\tau)\right], \tag{30}$$

where

$$\Omega=\tfrac{1}{2}[(\gamma_y-\gamma_z)^2-2Y^2]^{1/2}, \tag{31}$$

$$\Phi=\left\{\frac{\gamma_y-\gamma_z}{2}+\frac{\sqrt{2}Y\langle\sigma_y\rangle}{1+2\langle\sigma_z\rangle}\right\}. \tag{32}$$

For weak fields ($Y\to 0$), in the SV case, this result simplifies to

$$g^{(2)}(\tau)\approx 1-\exp[-(2N+1)\tau], \tag{33}$$

independent of M, i.e., it is the same for thermal fields and squeezed fields. Similarly, in the CS case, we find

$$g^{(2)}(\tau)\approx 1-\exp\{-[2N(1+2C)+1]\tau\}. \tag{34}$$

Plots of these are shown in Fig. 3. We see that of course $g^{(2)}(0)=0$, as it must be for a single two-level atom, but that $g^{(2)}(\tau)$ rapidly approaches unity, approaching it at a faster rate as the squeezing is increased. This result makes sense if we recall that the squeezed vacuum state is a state of nonzero mean photon number N. As the squeezing parameter N is increased, the atom interacts with a larger amount of photons, and it takes a correspondingly shorter time to reexcite the atom after an emission event. In the weak-field limit, there is essentially no coherent field to provide a phase ref-

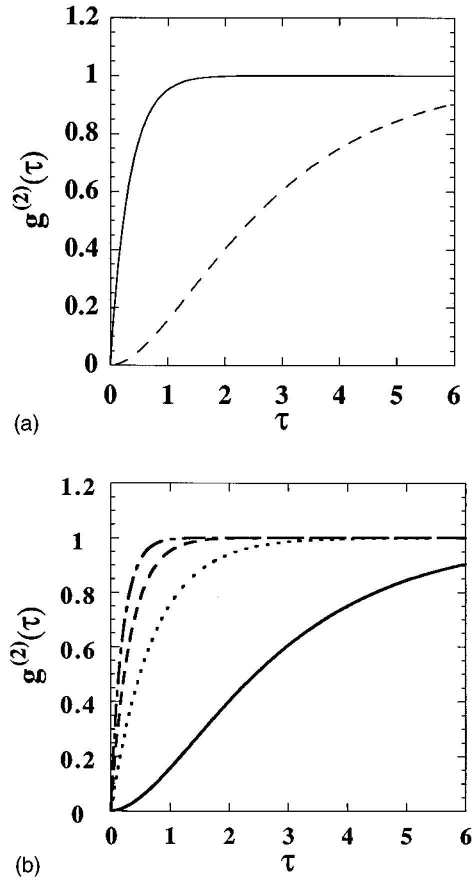

FIG. 3. Second-order intensity correlation function $g^{(2)}(\tau)$ for the SV case for weak fields. (a) For no squeezing (dash), thermal light (dot-dash), and squeezed vacuum (solid) with $N=1.0$. (b) For $N=0.0$ (solid), 0.2 (small dash), 1.0 (large dash), and 2.0 (dot-dash). Recall the scaling of the time axis $\tau=\gamma t$.

erence, hence the result is the same for all phases of the squeezed vacuum. The atom is predominantly in the ground state, and the interaction with the squeezed vacuum produces the same result as the thermal light of the same mean photon number, and there is no phase-sensitive behavior. This is similar to the behavior obtained by Vyas and Singh [13] for an atom interacting with one broadband squeezed mode.

In the limit of large squeezing, in the SV case we have

$$g^{(2)}(\tau)=1-\exp\{-N\tau\}\left[\cosh(\Omega_{LS}\tau)+\frac{1}{[1-Y^2/(2N^2)]^{1/2}}\sinh(\Omega_{LS}\tau)\right], \tag{35}$$

where here $\Omega_{LS}=N[1-Y^2/N^2]^{1/2}$. Notice that the large squeezing and weak coherent field limits do not commute. Once more $g^{(2)}(0)=0$ and the rise to unity is oscillatory in nature, due to the Rabi oscillations of the atom. The amplitude of the oscillations is reduced by the presence of a squeezed vacuum, just as there would be if the atom were interacting with a thermal, noisy reservoir. However, $g^{(2)}(\tau)$ is sensitive to the phase of the squeezed vacuum, or more precisely the relative phase between the squeezed vacuum and the coherent field. For $\phi=0$, when the quiet quadrature of the squeezed vacuum is in phase with the coherent driving field, and hence in phase with the Rabi oscillations of the atom, there is an increase in the amplitude of the oscillations in $g^{(2)}(\tau)$ relative to the thermal light of the same mean photon number. For the opposite phase $\phi=\pi/2$ there is an enhancement in the suppression of the oscillations relative to that of the thermal state. In both cases however, the oscillations are smaller than the case of an ordinary vacuum. This is seen in the eigenvalues of the Bloch equations, as the presence of the squeezed vacuum increases the magnitude of the real part, but reduces the magnitude of the imaginary part; faster damping and smaller oscillation frequency. Once again, increasing the squeezing washes out the oscillations, as can be seen in Fig. 4. Essentially similar behavior is observed for the CS case, although here the cavity enhancement of spontaneous emission increases the decay rate.

IV. COHERENT SPECTRUM

The fluorescent spectrum is given by

$$I_{inc}(\omega)=\langle\sigma_+(0)\rangle\langle\sigma_-(0)\rangle\delta(\omega-\omega_0)+\frac{1}{\pi}\,\mathrm{Re}\int_{-\infty}^{\infty}\langle\Delta\sigma_+(\tau)\Delta\sigma_-(0)\rangle\times\exp[i(\omega-\omega_0)\tau]d\tau. \tag{36}$$

The δ function component is known as the coherent, Rayleigh, or elastic scattering spectrum; the second component is the incoherent or inelastic scattering spectrum. We calculate the intensity of the coherent spectrum as the driving field is scanned across the central resonance frequency. This is determined using steady states obtained from the Bloch equations,

$$\langle\dot{\sigma}_x\rangle=-\gamma_x\langle\sigma_x\rangle-\Delta\langle\sigma_y\rangle, \tag{37}$$

$$\langle\dot{\sigma}_y\rangle=-\gamma_y\langle\sigma_y\rangle-(Y/\sqrt{2})\langle\sigma_z\rangle+\Delta\langle\sigma_x\rangle, \tag{38}$$

$$\langle\dot{\sigma}_z\rangle=-\gamma_z(\langle\sigma_z\rangle+\delta)+(Y/\sqrt{2})\langle\sigma_y\rangle, \tag{39}$$

where we have included the detuning of the driving field from the atomic resonance by an amount Δ, normalized to the spontaneous emission rate γ. The steady-state solution to these equations is given by

$$\langle\sigma_x\rangle_{ss}=\frac{-\Delta\,\delta Y}{\sqrt{2}}\,\frac{1}{\gamma_x\gamma_y+\Delta^2+\dfrac{Y^2}{2}\dfrac{\gamma_x}{\gamma_z}}, \tag{40}$$

$$\langle\sigma_y\rangle_{ss}=\frac{-\gamma_x\delta Y}{\sqrt{2}}\,\frac{1}{\gamma_x\gamma_y+\Delta^2+\dfrac{Y^2}{2}\dfrac{\gamma_x}{\gamma_z}}, \tag{41}$$

$$\langle\sigma_z\rangle_{ss}=\frac{-\delta(\gamma_x\gamma_y+\Delta^2)}{\gamma_x\gamma_y+\Delta^2+\dfrac{Y^2}{2}\dfrac{\gamma_x}{\gamma_z}}. \tag{42}$$

Hence the coherent spectrum can be given as

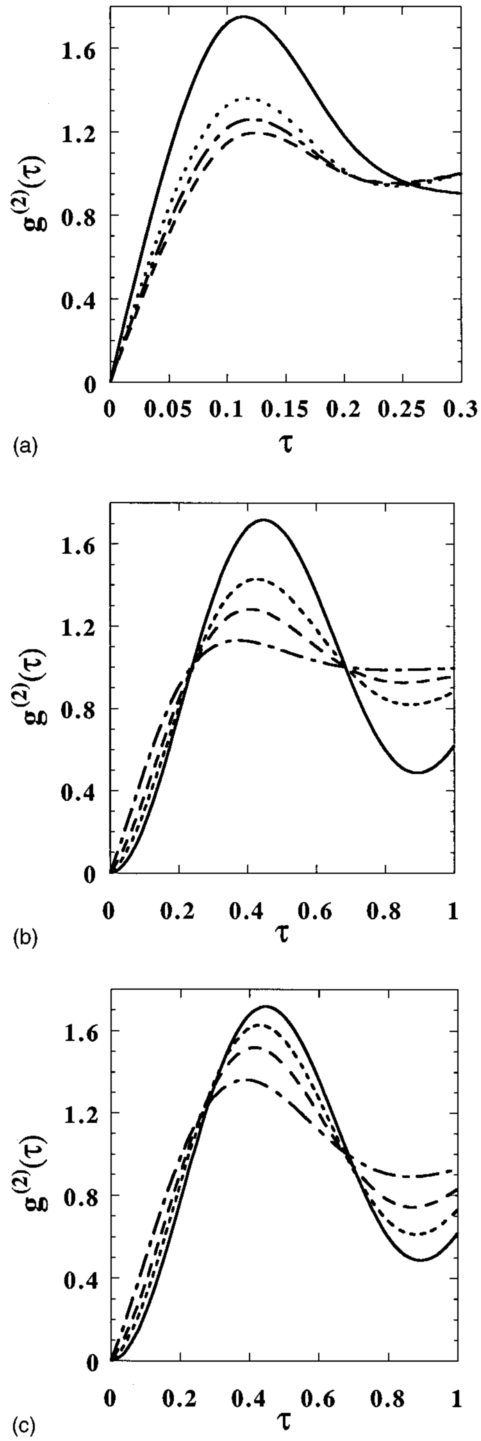

FIG. 4. Second-order intensity correlation function $g^{(2)}(\tau)$ for the SV case for strong fields, $Y=10.0$. (a) For no squeezing (solid), thermal light (dot-dash), and squeezed vacuum with $N=1.0$ and $\theta=0.0$ (small dash) and $\theta=\pi/2$ (large dash). (b) For $\phi=\pi/2$, $N=0.0$ (solid), 0.5 (small dash), 1.0 (large dash), and 2.0 (dot-dash). (c) For $\phi=0.0$, $N=0.0$ (solid), 0.5 (small dash), 1.0 (large dash), and 2.0 (dot-dash). Recall the scaling of the time axis $\tau=\gamma t$.

$$I_{coh} \propto \frac{4\gamma_x^2+\Delta^2}{\{4\gamma_x\gamma_y+\Delta^2+Y^2\}^2}. \tag{43}$$

A similar expression has been obtained by Ritsch and Zoller [14] for the SV case, where they found that the coherent spectrum was broadened by the presence of the squeezed vacuum. Essentially, the coherent spectrum is phase independent, and samples both the quiet and noisy quadratures. In the CS case, the coherent spectrum is given in the weak-field limit by

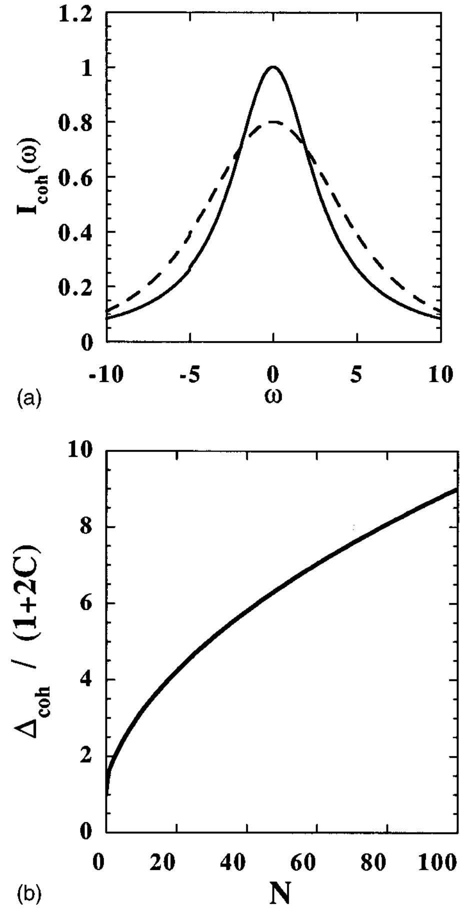

FIG. 5. Coherent spectrum for the CS case, for weak fields with $C=1.0$. (a) $N=0$ (solid) and $N=2.0$. (dash). (b) Linewidth of the coherent spectrum as a function of N. Frequencies are in units of Γ^{-1}.

$$I_{coh} \propto \frac{1+\Delta^2}{\{1+2C(1+4N)+\Delta^2\}^2} \tag{44}$$

which exhibits the same broadening as the SV case, as shown in Fig. 5.

V. PUMP-PROBE ABSORPTION SPECTRUM

Ritsch and Zoller [14] examined the absorption spectrum of a weak probe field as it is scanned across the central resonance, in the presence of a second pump field tuned to resonance. They obtained essentially the same result as Gardiner [1] and CLW [2] obtained for the incoherent spectrum. That is the separation of the usual Lorentzian spectrum into two distinct components as the amount of squeezing is increased; one that is narrowed and one which is broadened. In the limit of large squeezing the width of the narrow component can be made arbitrarily small. In the case of a two-level atom in a cavity, the CS case, the pump-probe absorption spectrum is given for weak pump fields by

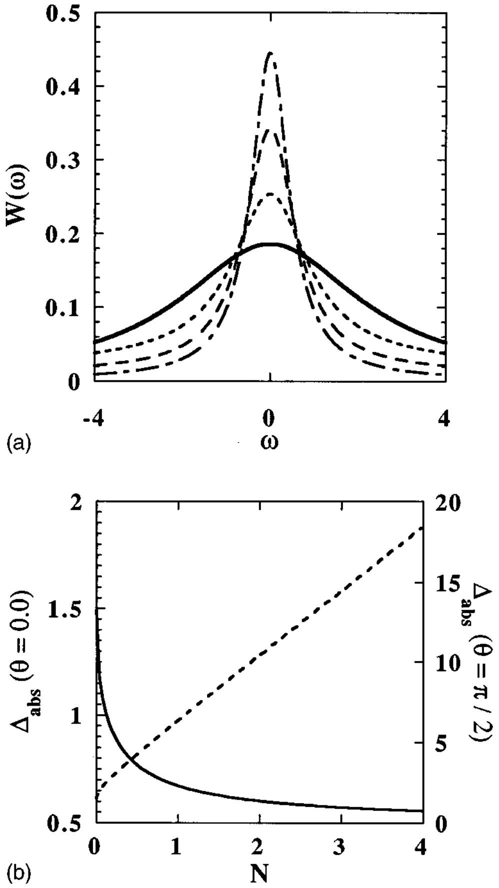

FIG. 6. Pump-probe absorption spectrum for the CS case for weak fields and $C=2.0$. (a) $N=0.0$ (solid), 0.5 (small dash), 2.0 (large dash), and 20.0 (dot-dash). (b) Linewidth of the two Lorentzian components of the pump-probe absorption spectrum as a function of N. Frequencies are in units of Γ^{-1}.

$$W(\omega)\propto\left\{\frac{\gamma_x}{\gamma_x^2+\omega^2}+\frac{\gamma_y}{\gamma_y^2+\omega^2}\right\}. \tag{45}$$

Here we see that the spectrum splits into two components as in the SV case, and we obtain essentially the same result as that of the incoherent spectrum in the CS case [7]. Note that the width of the spectrum cannot be made arbitrarily narrow, as it is limited by the value of Γ, the rate of spontaneous emission into the unsqueezed modes. The cavity enhancement of the linewidth, $\Gamma(1+2C)$ can be reduced using a squeezed vacuum, but is only subnatural if the cavity encloses a significant portion of 4π sr as was the case for the incoherent spectrum. Once again this weak-field result is independent of the phase of the squeezed vacuum. This is exhibited in Fig. 6. For thermal light of the same mean photon number, there is only a broadening of the spectrum.

VI. SPECTRUM OF SQUEEZING

It is known that for weak coherent excitation, the fluorescence emitted from a two-level atom is squeezed, reflecting nonclassical atomic polarization fluctuations [15,16]. The quadrature $\pi/2$ out of phase with the driving field is squeezed, and the quadrature in phase with the driving field is unsqueezed. Here we investigate the spectrum of squeezing for the SV and CS cases, in order to determine the effect of a squeezed vacuum on nonclassical effects obtained with a coherent field. The spectrum of squeezing is given by [16,17]

$$S(\omega,\theta)=8\eta\int_0^\infty d\tau\,\cos(\omega\tau)\mathrm{Re}\{\langle\Delta\sigma_+(0)\Delta\sigma_-(\tau)\rangle + e^{2i\theta}\langle\Delta\sigma_+(0)\Delta\sigma_+(\tau)\rangle\}. \tag{46}$$

$S(\omega,\theta)$ has a lower bound of -1 and 0 represents the shot-noise limit. Here η is the combined collection and detection efficiency. If one takes the limit of $Y\rightarrow 0$, the squeezing goes away entirely, and the spectra of squeezing are a positive narrow peak, and a positive broad peak, respectively. Changing the phase of the local oscillator relative to the phase of the coherent field merely changes which quadrature is broadened and which is narrowed. Summing these two spectra of squeezing results in the incoherent spectrum containing a narrowed and a broadened component [18]. These results can be understood in terms of the two quadratures of the atom as independent scatterers of squeezed noise (photons derived from the squeezed vacuum) where each quadrature acts as a bandpass filter of different linewidth. Changing the phase of the squeezed vacuum merely alters the bandwidths of the two quadratures. In the case of a thermal vacuum, the atom would still scatter noise, but the spectrum of squeezing would be independent of phase. Here, we see that the use of a squeezed vacuum results in the destruction of a nonclassical feature present when the system interacts with normal vacuum modes. In the CS case, one has the same qualitative behavior; with the linewidth of the two spectra of squeezing being given by the appropriate γ_x and γ_y.

VII. CONCLUSIONS

We have calculated several quantities for a single two-level atom immersed in a squeezed vacuum (SV model), and for an atom in a cavity that is driven by a coherent field as well as a squeezed vacuum, but interacts with an ordinary vacuum out the sides of the cavity (CS model). We find that the photon statistics of the fluorescent field exhibit photon antibunching $[g^{(2)}(0)<1]$, but that very rapidly $g^{(2)}(\tau)$ approaches the value of unity. Indeed, for weak fields, the correlation function is exactly the same as that of a thermal field of the same mean photon number. For strong fields, $g^{(2)}(\tau)$ exhibits phase-sensitive behavior as the phase of the squeezing is varied relative to the phase of the coherent driving field. Hence we see that in some respects a squeezed vacuum is similar to noise. This is due to the fact that although there are phase correlations for a squeezed vacuum, the average value of the field is 0, but the average value of the intensity is not. The coherent spectrum and pump-probe absorption spectrum for the CS case is very similar to those obtained by Ritsch and Zoller [14] for the SV case. The coherent spectrum is only broadened by the squeezed vacuum, and the pump-probe absorption spectrum splits into a narrow and a broad component. The width of this spectrum cannot be made arbitrarily narrow, but is limited by the spontaneous emission rate into unsqueezed modes. Finally, we have calculated the spectrum of squeezing for the fluorescent light, and have found that interaction with a squeezed vacuum can completely destroy the squeezing in the fluorescent light.

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