# Suppression of fluorescence in a squeezed vacuum

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A scheme which results in the simultaneous narrowing of all three spectral peaks of the Mollow triplet is proposed. In particular, we investigate the effect on resonance fluorescence of processing the squeezed output of an optical parametric oscillator by propagation through a linear dielectric. The frequency-dependent squeezing phase induced by this process is found to result in subnatural spectral widths for appropriate parameter choices. The line-narrowing mechanisms at the Rabi sidebands manifest themselves in the second-order correlation function of the fluorescent field which is found to decay at a subnatural rate. The potential for experimental realization of these phenomena is discussed. [S1050-2947(96)08912-3]

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## I. INTRODUCTION

It is now well known that squeezed states of light [1] can have a profound effect on atomic damping processes. In particular, it has been established [2] that when a two-level atom interacts with a broadband squeezed vacuum, the transverse atomic polarization quadratures exhibit different decay rates. One of these decay rates is slower than found in the normal vacuum and becomes arbitrarily slow in the strong squeezing limit. Subnatural linewidths have been predicted [3] when an intense monochromatic pump resonant with the natural atomic transition frequency is added to the model of an atom damped by a broadband squeezed vacuum. The emission spectrum is found to be highly sensitive to the relative phase of the pump and the squeezed vacuum. For appropriate choices of pump phase the central peak of the familiar Mollow triplet is found to display a subnatural width which becomes arbitrarily narrow in the strong squeezing limit.

More recently, various studies have been carried out on the influence of squeezing bandwidths on resonance fluorescence. Parkins and Gardiner [4] and Parkins [5,6] studied the influence of colored squeezed light on resonance fluorescence by stochastic simulation of the wave function. They found that the inclusion of squeezing of bandwidths into the model results in new line-narrowing mechanisms at the Rabi sidebands. Furthermore, it was established that the degeneracy of the squeezing source strongly affects the relative widths of the three spectral peaks of the fluorescent spectrum which occur in the strong driving regime. These qualitative results were also found by Yeoman and Barnett [7], who presented an analytic technique to investigate the behavior of a coherently driven atom damped by a squeezed vacuum with finite squeezing bandwidth. By deriving the master equation in the dressed atom picture it is possible to relax the restriction that the squeezing bandwidth is much greater than the Rabi frequency. However, the theory still requires the squeezing bandwidth to be much greater than the atomic linewidth in order to ensure that the field is $\delta$ correlated on time scales comparable to natural atomic lifetimes. In a strong driving field, this naturally allows the analytical study of squeezed fields similar to those of Parkins and Gardiner; that is, squeezed fields with different squeezing properties at the frequencies corresponding to the three peaks of the fluorescent spectrum.

The line-narrowing mechanisms found at the Rabi sidebands in the presence of colored squeezed light suggest that quantum effects will manifest themselves in the second-order correlation function [8–10]. Carmichael and Walls [11] demonstrated that the decay rate of the second-order correlation function is equal to the widths of the Rabi sidebands. Thus, subnatural decay rates will occur for the conditions which result in subnatural widths of the Rabi sidebands [12]. The sensitivity of photocounting techniques suggests that it might be desirable to consider a measurement of the second-order correlation function rather than the emission spectrum.

In this paper we shall propose and theoretically investigate an experiment which should result in significant narrowing of all three spectral peaks even for modest squeezing strengths. We shall show that by processing the output of a parametric oscillator by propagation through a linear dielectric [13] or some other suitably dispersive element, the three spectral peaks occurring in the strong driving regime can all be made to exhibit subnatural widths. For modest squeezing, we shall show that these widths may be less than those predicted in the strong squeezing limit when the squeezed beam has colored noise properties. Furthermore, these quantum effects manifest themselves in the second-order correlation function. We establish that this particular statistical measure decays at a rate forbidden by semiclassical theory thus providing a signature of the nonclassical behavior arising from the quantum correlations of the squeezed field.

The key to all of this lies in the fact that the output of an optical parametric oscillator is a unidirectional beam of small solid angle. Consequently, it is possible to perform conventional optics with practical sources of squeezed light [14]. It has recently been shown that propagating a squeezed beam with a Gaussian profile through a dispersive medium results in a modulation of the squeezing phase [13]. By tuning the parameters representing the dielectric it is possible to engineer a situation where the squeezed field incident upon the two-level atom is maximally squeezed (unsqueezed) at a fre-

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quency corresponding to the central atomic transition and maximally unsqueezed (squeezed) at frequencies corresponding to the Rabi sidebands. It is this situation which we shall investigate in this paper. To do this we use the master equation derived by Yeoman and Barnett [7]. We note that although this master equation assumes that the two-level atom interacts isotropically with squeezed modes over a $4\pi$ solid angle, it represents, to a good approximation, the current experimental proposal of a squeezed beam injected into a Fabry-Pérot resonator. This ensures that the atom predominantly couples to squeezed modes [15,16]. Indeed, it has recently been shown that the dispersive profiles which occur in a weak to modest driving regime [17] may be reconstructed in a Fabry-Pérot cavity [18].

## II. THE BLOCH EQUATIONS AND FLUORESCENT SPECTRUM

We consider a coherently, resonantly, and strongly driven two-level atom, damped by its interaction with a squeezed field which has frequency-dependent squeezing parameters and squeezing phase. The Bloch equations describing such a system may be derived following the prescription of Ref. [7] where it is assumed that the pump frequency, the central frequency of squeezing, and the bare atomic transition are all exactly resonant. The result is

$$\langle \dot{\sigma_X} \rangle = -\gamma_\Omega \langle \sigma_X \rangle - 2G(\omega_A)\langle \sigma_Z \rangle + \gamma |M(\omega_A+\Omega)| \sin\Phi(\omega_A+\Omega)\langle \sigma_Y \rangle, \tag{1}$$

$$\langle \dot{\sigma_Z} \rangle = -(\gamma_0+\gamma_\Omega)\langle \sigma_Z \rangle - 2\Omega \langle \sigma_Y \rangle - \gamma, \tag{2}$$

$$\langle \dot{\sigma_Y} \rangle = \gamma |M(\omega_A)| \sin\Phi(\omega_A)\langle \sigma_X \rangle + \left( \frac{\Omega}{2} + \delta(\omega_A) \right) \langle \sigma_Z \rangle - \gamma_0 \langle \sigma_Y \rangle, \tag{3}$$

where

$$\gamma_0 = \gamma \left[ N(\omega_A) - |M(\omega_A)| \cos\Phi(\omega_A) + \frac{1}{2} \right], \tag{4}$$

$$\gamma_\Omega = \gamma \left[ N(\omega_A+\Omega) + |M(\omega_A+\Omega)| \cos\Phi(\omega_A+\Omega) + \frac{1}{2} \right]. \tag{5}$$

$G(\omega_A)$ and $\delta(\omega_A)$ are principal part integrals

$$\delta(\omega_A) = \frac{\gamma}{2\pi} \wp \int_{-\infty}^{\infty} \frac{d\Delta}{\Delta+\Omega} [|M(\Delta+\omega_A)| \cos\Phi(\Delta) + N(\Delta+\omega_A)], \tag{6}$$

$$G(\omega_A) = -\wp \int_{-\infty}^{\infty} \frac{d\Delta}{\Delta+\Omega} \frac{\gamma |M(\Delta+\omega_A)| \sin\Phi(\Delta)}{4\pi}, \tag{7}$$

where $\Delta$ is the deviation from the natural atomic transition frequency,

$$\Delta = \omega - \omega_A. \tag{8}$$

$N(\omega)$ and $|M(\omega)|$ characterize the squeezed vacuum field autocorrelation functions

$$\langle b^\dagger(\omega) b(\omega') \rangle = N(\omega)\delta(\omega-\omega'), \tag{9}$$

$$\langle b(\omega) b(\omega') \rangle = |M(\omega)| \exp[i\phi_S(\omega)]\delta(2\omega_A - \omega - \omega'), \tag{10}$$

while $\Phi(\omega)$, the squeezing phase, describes the phase difference between the phase of the squeezed field, $\phi_S(\omega)$, and twice the phase of the driving field, $\phi_L$, that is,

$$\Phi(\omega) = 2\phi_L - \phi_S(\omega). \tag{11}$$

Finally, we note that $\Omega$ denotes the Rabi frequency and that the atomic polarization quadratures and population inversion have been defined as

$$\langle \sigma_X \rangle = \frac{1}{2}[\langle \sigma_- \rangle e^{i(\phi_L - \omega_A t)} + \langle \sigma_+ \rangle e^{-i(\phi_L - \omega_A t)}], \tag{12}$$

$$\langle \sigma_Y \rangle = \frac{1}{2i}[\langle \sigma_- \rangle e^{i(\phi_L - \omega_A t)} - \langle \sigma_+ \rangle e^{-i(\phi_L - \omega_A t)}], \tag{13}$$

$$\langle \sigma_Z \rangle = \langle [\sigma_+, \sigma_-] \rangle, \tag{14}$$

where $\sigma_\pm$ represents the Pauli pseudospin operators for the atom and $\phi_L$ and $\omega_A$, respectively, denote the phase of the pump and the central resonance frequency of the atomic transition.

For the remainder of this discussion we consider the parameters which maximize the potential for subnatural linewidths and also permit analytic expressions to be found for the fluorescent spectrum [19]. To do this we make two principal assumptions. First, we consider a squeezing phase with frequency dependence given by

$$\Phi(\Delta) = \Phi(\omega_A) + \frac{\pi\Delta}{\Omega} \tag{15}$$

with $\Phi(\omega_A)=0$ or $\pi$. The result of this assumption is that the terms $\sin\Phi(\omega_A)$ and $\sin\Phi(\omega_A+\Omega)$ appearing in the Bloch equations (1)–(3) are identically zero. This permits simple analytic expressions for the eigenvalues of the Bloch matrix. As we have already explained, a squeezed field with squeezing phase characteristics described by Eq. (15) may be generated by propagating a squeezed beam through dispersive optical elements such as certain types of glass [13].

The second assumption which we make is that the unprocessed squeezed beam is sufficiently broadband so that the squeezing parameters $N(\omega)$ and $|M(\omega)|$ may be replaced by constants $N$ and $|M|$. Within this second assumption we note that we may substitute $\gamma_0 - \gamma_\Omega = \delta(\omega_A) = 0$ into the Bloch matrix shown in Eq. (17).

A natural consequence of these two assumptions is to forbid the study of strongly squeezed fields with frequency-dependent squeezing phase. The reason for this is that in order for the Markoff approximation to be valid we require the squeezing phase to be slowly varying compared to the natural atomic linewidth. Thus, if we wish to study the situation where the squeezing phase at the Rabi sideband frequencies differs significantly from the squeezing phase at the central atomic resonance, it is necessary to consider a strong driving regime in which the three spectral peaks of the fluorescent spectrum are well separated. However, in this strong driving regime the squeezing parameters $N(\omega)$ and $|M(\omega)|$ may be approximated as constants only for weakly or modestly squeezed fields. This is due to the fact that in practice there exists a correspondence between the squeezing bandwidth and the squeezing strength of realizable squeezed beams. For example, the squeezing parameters associated with the beam produced by an optical parametric oscillator are slowly varying in the frequency domain only when operating significantly below threshold [20]. In this parameter regime the squeezed output exhibits only modest squeezing. Alternative sources of broadband squeezed light, such as the squeezed output of optical parametric amplifiers [21], give only modest squeezing. In particular, it has been demonstrated that the output of an optical parametric amplifier [21] has noise fluctuations 0.5 dB below the shot-noise limit with no noticeable loss in noise reduction over a bandwidth of the order of GHz. Either way, we may only consider modestly and weakly squeezed fields if we wish to study the situation where the squeezing phase is significantly different at the Rabi sideband and central atomic frequencies. However, this limit does allow us to invoke the approximation

$$\Omega \gg \gamma\left(N+|M|+\frac{1}{2}\right) \tag{16}$$

in Eq. (3).

The limit on squeezing strength described above is, however, not unduly restrictive. The reason for this is that the master equation, used to generate the Bloch equations (1)–(3), is derived within the assumption that the field dynamics are Markoffian. Specifically, this means that we require the field to be sufficiently slowly varying in the frequency domain so that it appears $\delta$ correlated over atomic time scales. This must be true for all values of Rabi frequency. However, in the absence of a driving field it is found that the transverse atomic polarization quadratures exhibit two distinct decay rates [2] given by $\gamma(N+|M|+\frac{1}{2})$ and $\gamma(N-|M|+\frac{1}{2})$. Thus, for the Markoff approximation to be valid we now require the field bandwidth to be very much greater than each of these distinct rates. For a field with frequency-dependent squeezing phase given by Eq. (15), the bandwidth is approximately equal to the Rabi frequency, $\Omega$. Consequently, for the theory to be valid it is necessary in any case that the inequality (16) is satisfied.

A final simplication of the Bloch equations (1)–(3) results from the assumptions we have made about the squeezing phase and strength. Within the constraint of Markoffian field dynamics the correction to broadband squeezing characteristics, $2G(\omega_A)$, will be negligibly small. This may be reasoned if we substitute (15) into (7) and impose the condition for Markoffian dynamics (16). As a result of this substitution we find that, to a first approximation, $G(\omega_A)$ may be neglected.

Substituting these assumptions of strong driving, modest squeezing, and Markoffian field dynamics into the Bloch equations leads to a simplified set of equations

$$\frac{\partial}{\partial t}\begin{bmatrix}\langle\sigma_X\rangle\\ \langle\sigma_Z\rangle\\ \langle\sigma_Y\rangle\end{bmatrix}=\begin{pmatrix}-\gamma_\Omega & 0 & 0\\ 0 & -2\gamma_\Omega & -2\Omega\\ 0 & \dfrac{\Omega}{2} & -\gamma_\Omega\end{pmatrix}\begin{bmatrix}\langle\sigma_X\rangle\\ \langle\sigma_Z\rangle\\ \langle\sigma_Y\rangle\end{bmatrix}-\begin{bmatrix}0\\ \gamma\\ 0\end{bmatrix}. \tag{17}$$

The fluorescent spectrum may now be calculated in the usual manner by applying the quantum regression theorem [22] to the formal solution of the Bloch equations and substituting these results into the Wiener-Khintchine theorem [23]. For the parameter regime resulting in (17), the result may be approximated by

$$S(\omega)=\frac{1}{4}\left(\frac{\gamma}{\Omega}\right)^2\delta(\omega-\omega_A)+\frac{\gamma_\Omega}{16\pi}\left[\frac{4}{\gamma_\Omega^2+(\omega-\omega_A)^2}+\frac{3}{\left(\dfrac{3\gamma_\Omega}{2}\right)^2+(\omega-\omega_A+\Omega)^2}+\frac{3}{\left(\dfrac{3\gamma_\Omega}{2}\right)^2+(\omega-\omega_A-\Omega)^2}\right]. \tag{18}$$

Inspection of Eq. (18) reveals that in the strong driving regime the fluorescent spectrum is composed of a $\delta$ function and a three-peaked structure comprising three well-separated Lorentzian peaks. These contributions arise, respectively, from elastic and inelastic scattering [24]. The incoherent contribution to the spectrum of scattered light has a central peak of width $\gamma_\Omega$ and two sidebands shifted in frequency by $\Omega$ with widths $\frac{3}{2}\gamma_\Omega$. Consequently, appropriate choices of $\Phi(\omega_A)$ result in simultaneous line narrowing or line broadening of all three spectral peaks of the incoherent contribution. These peaks may not, however, be made arbitrarily narrow even when the squeezed beam is a minimum uncertainty state obeying the equality

$$|M|=\sqrt{N(N+1)}. \tag{19}$$

This is due to the fact that the strong squeezing limit, as we have already explained, is not only forbidden in the theory but is also incompatible in practice with the required frequency dependence of the squeezing phase.

Further examination of Eq. (18) reveals two interesting features of the fluorescent spectrum. First, the integrated spectral intensity of the coherent component of the spectrum of scattered light is vanishingly small in the strong driving limit. This is in keeping with the behavior found in the absence of squeezed light [24]. Second, the ratio of the integrated spectral intensities of the central peak of the incoherent component to each of the Rabi sidebands is 2:1. This behavior is also found in the normal vacuum and has been explained in the dressed state basis [25].

In Fig. 1 we plot the incoherent component of the Mollow spectrum and the incoherent component of the fluorescent spectrum described in Eq. (18) for parameters corresponding

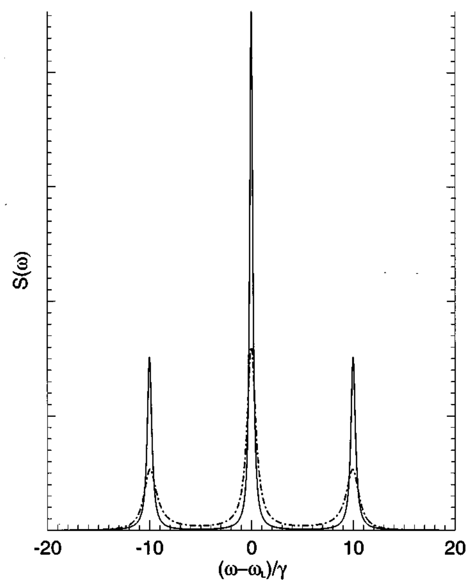


FIG. 1. Mollow triplet when $\Omega=10\gamma$, $N=0$ (dash dot line); fluorescent spectrum in the presence of squeezed vacuum with frequency-dependent squeezing phase when $\Omega=10\gamma$, $N=0.3$, $|M|=[N(N+1)]^{1/2}$, $\Phi(\omega_A)=0$, $\Phi(\omega_A+\Omega)=\pi$, $G(\omega_A)=0$ (solid line).

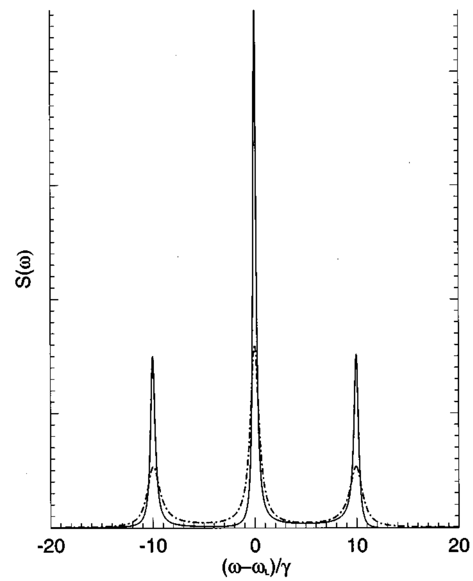


FIG. 2. Mollow triplet when $\Omega=10\gamma$ (dash dot line), $N=0$; fluorescent spectrum in the presence of squeezed vacuum with frequency-dependent squeezing phase when $\Omega=10\gamma$, $N=0.3$, $|M|=[N(N+1)]^{1/2}$, $\Phi(\omega_A)=0$, $\Phi(\omega_A+\Omega)=\pi$, $G(\omega_A)\neq 0$ (solid line).

to the modestly squeezed output of an optical parametric oscillator obeying the equality shown in Eq. (19). The squeezing phase at the central atomic frequency, $\Phi(\omega_A)$, is chosen to be 0 so that all three spectral peaks simultaneously exhibit subnatural linewidths. This figure shows that even for moderate squeezing strengths significant line narrowing of the Rabi sidebands is possible when the squeezed beam is processed by propagation through a linear dielectric or suitable dispersive element. Indeed, an 80% reduction in linewidth occurs for all three spectral peaks when $N=0.3$ as depicted in Fig. 1. This line narrowing significantly exceeds the Rabi sideband line narrowing possible when the squeezed field is the finite-bandwidth output of a nondegenerate or degenerate parametric oscillator [7].

In plotting Fig. 1 we have made the approximation that $G(\omega_A)=0$. The exact form of the fluorescent spectrum for nonzero $G(\omega_A)$ is rather complicated and we do not detail the results of the calculation here. Instead, we simply plot the result. This is shown in Fig. 2, where we plot the fluorescent spectrum for nonzero $G(\omega_A)$ along with the corresponding Mollow spectrum. This figure shows that the line narrowing predicted in Fig. 1 is maintained when $G(\omega_A)$ is finite but small due to the restrictions imposed by the Markoff approximation. The principal difference between Fig. 1 and Fig. 2 is the slight asymmetry of the central peak in Fig. 2. As we shall now show, this is due to a population inversion in the dressed state basis [25–27].

## III. DRESSED STATE ANALYSIS

The nature of the atomic system under consideration suggests that an analysis of the evolution of the dressed states might be instructive. To do this we introduce two dressed states $|\alpha\rangle$ and $|\beta\rangle$,

$$|\alpha\rangle=\frac{1}{\sqrt{2}}[e^{-i\phi_L}|1\rangle+|2\rangle], \tag{20}$$

$$|\beta\rangle=\frac{1}{\sqrt{2}}[-e^{-i\phi_L}|1\rangle+|2\rangle], \tag{21}$$

defined to be the eigenstates of the time-independent interaction Hamiltonian [7] with eigenergies $\hbar\Omega/2$ and $-\hbar\Omega/2$, respectively. The Bloch equations in the dressed state basis are then found to be

$$\frac{\partial}{\partial t}(\rho_{\alpha\alpha}-\rho_{\beta\beta})=-\gamma_\Omega(\rho_{\alpha\alpha}-\rho_{\beta\beta})-4G(\omega_A)(\rho_{\alpha\beta}+\rho_{\beta\alpha}), \tag{22}$$

$$\frac{\partial\rho_{\alpha\beta}}{\partial t}=-\frac{\gamma}{2}(\rho_{\alpha\alpha}+\rho_{\beta\beta})-\frac{\gamma_\Omega}{2}\rho_{\beta\alpha}-\left(\frac{3\gamma_\Omega}{2}+i\Omega\right)\rho_{\alpha\beta}, \tag{23}$$

$$\frac{\partial\rho_{\beta\alpha}}{\partial t}=\left(\frac{\partial\rho_{\alpha\beta}}{\partial t}\right)^*, \tag{24}$$

where $\rho_{ij}=\langle i|\rho|j\rangle$ with $i,j$ denoting the dressed states $\alpha$ and $\beta$.

Inspection of (22)–(24) reveals that in the dressed state basis the Rabi frequency appears as a classical driving term in the equations of motion for the off-diagonal matrix elements. This feature corresponds to the frequency shift of the Rabi sidebands in the spectrum of scattered light. In addition to this we may identify the damping rate of the off-diagonal elements, $3\gamma_\Omega/2$, as the width of each Rabi sideband. The equations of motion for the diagonal elements, on the other hand, do not contain a driving term as they correspond to transitions with frequency $\omega_A$. In the absence of $G(\omega_A)$ the damping of the dressed state population inversion, and hence the width of the central resonance, is simply $\gamma_\Omega$. In this instance, the steady state of the dressed state population inversion is zero leading to an entirely symmetric fluorescent spectrum (see Fig. 1). However, when $G(\omega_A)$ is nonzero (see Fig. 2) we find a finite dressed state population inversion. This inversion naturally leads to an asymmetry in the fluorescent spectrum. However, this asymmetry differs from that present when $\Phi\neq 0$ or $\pi$ [26] due to the fact that the asymmetry occurs only at the central resonance and not at the sideband frequencies. This feature is predicted by the evolution of the dressed states. When $\Phi\neq 0$ or $\pi$ we find that the diagonal and off-diagonal matrix elements are all coupled via terms proportional to $\gamma|M|\sin\Phi$ [26]. A consequence of the fact that the matrix elements corresponding to each of the dressed state transitions are affected is an asymmetry in the entire fluorescent spectrum. However, in the regime of nonzero $G(\omega_A)$ we find that the equations of motion for the off-diagonal matrix elements are unchanged and that only the equation of motion for the population inversion is influenced by this quantity. Consequently, the dressed state population inversion in this special case leads to an asymmetry only of the central fluorescent peak.

## IV. THE SECOND-ORDER CORRELATION FUNCTION

In this section we calculate the second-order correlation function for the fluorescent field emitted by a strongly driven, two-level atom damped by a broadband squeezed bath with frequency-dependent squeezing phase.

The second-order correlation function [8,9], $g^{(2)}(t,t+\tau)$, is defined to be

$$g^{(2)}(t,t+\tau)=\frac{\langle E^{(-)}(t)E^{(-)}(t+\tau)E^{(+)}(t+\tau)E^{(+)}(t)\rangle}{\langle E^{(-)}(t)E^{(+)}(t)\rangle^2}, \tag{25}$$

where $E^{(-)}(t)$ and $E^{(+)}(t)$ represent the continuous-mode creation and annihilation operators of the field. It is possible to recast the second-order correlation function in terms of the atomic operators at retarded times using the source-field expressions [28]

$$E_S^{(-)}(\mathbf{r},t)=\mu\sigma_+\left(t-\frac{|\mathbf{r}-\mathbf{R}|}{c}\right), \tag{26}$$

$$E_S^{(+)}(\mathbf{r},t)=\mu\sigma_-\left(t-\frac{|\mathbf{r}-\mathbf{R}|}{c}\right), \tag{27}$$

with $\mathbf{R}$ denoting the atomic position. The free or vacuum field operators make no contribution to $g^{(2)}(t,t+\tau)$ due to the normal ordering in the definition (25) so we may obtain $g^{(2)}(t,t+\tau)$ by replacing the electric field operators by the source operators. For notational simplicity, we let $t$ denote the retarded time argument in (26) and (27). The second-order correlation function may now be defined solely in terms of the atomic operators

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

For the initial conditions corresponding to a photodetection occurring at time $t=0$,

$$\langle\sigma_X(0)\rangle=\langle\sigma_Y(0)\rangle=0, \tag{29}$$

$$\langle\sigma_Z(0)\rangle=-1, \tag{30}$$

the second-order correlation function may be calculated by integrating the Bloch equations (17) and substituting the resulting function $\langle\sigma_Z(\tau)\rangle$ along with the quantum regression into (28). Again we assume that the squeezed field is broadband to maximize the potential for new quantum effects. We also assume that the squeezing phase has the frequency dependence described in Eq. (15) and that we are in a strong driving regime where (16) is valid. The result of the calculation, to a good approximation, is

$$g^{(2)}(\tau)=1-\exp\left(-\frac{3\gamma_\Omega}{2}\right)\cos\Omega\tau, \tag{31}$$

where $\gamma_\Omega$ is defined in Eq. (5). In the absence of squeezing $(N=|M|=0)$ this reduces to the strong driving limit of the familiar expression derived by Carmichael and Walls [11],

$$g^{(2)}(\tau)=1-\exp\left(-\frac{3\gamma}{4}\right)\cos\Omega\tau. \tag{32}$$

From Eq. (31) we may deduce that the damping rate is equal to the widths of the Rabi sidebands of the fluorescent spectrum and hence, for appropriate choices of phase, will be significantly slower than that found in the absence of squeezing. Consequently, this statistical measure may be used as a signature of nonclassical effects.

In Fig. 3 we plot the second-order correlation function for a strongly driven two-level atom damped by a modestly squeezed bath with $N=0.3$. The squeezing phase at the central atomic frequency, $\Phi(\omega_A)$, is chosen to be zero. These parameters correspond to subnatural linewidths of the Rabi sidebands with widths approximately equal to $\gamma/5$ as shown in Figs. 1 and 2. Also plotted in Fig. 3 is the second-order correlation function for an atom damped by its interaction with the zero point motion of the normal vacuum as shown in Eq. (32). We note that the minimum which occurs after the first Rabi oscillation suggests a higher level of antibunching in the presence of the squeezed light considered here than that found in the normal vacuum. It is this phenomenon which may be sought in experiment.

As a final remark we note that the considerable reduction in damping rate of the second-order correlation function for realizable squeezing strengths is attributed to the processing of the squeezed beam to induce a frequency-dependent squeezing phase. In a broadband squeezed vacuum with constant squeezing phase the damping rate is *always* greater than that found in the normal vacuum irrespective of the value of the squeezing phase. In addition to this the damping rate depicted in Fig. 3 is markedly slower than that found in a squeezed vacuum with finite squeezing bandwidth [12] and constant squeezing phase. It has been established [12] that for a squeezed field with these properties, the damping rate of $g^{(2)}(\tau)$ is equal to the widths of the Rabi sidebands and has a value $\frac{1}{2}(2\gamma_0+\gamma_\Omega)$ with $\gamma_0$ and $\gamma_\Omega$ defined, respec-

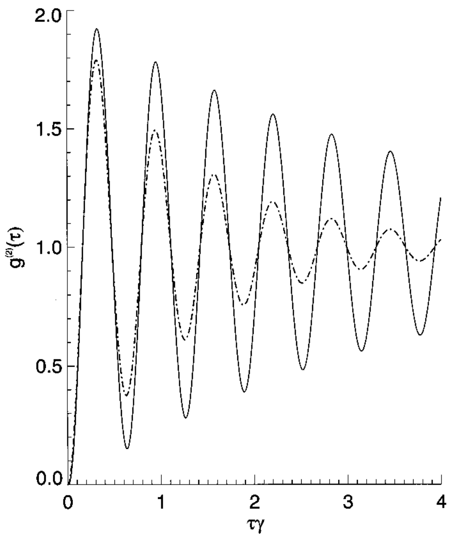


FIG. 3. Second-order correlation function when $\Omega=10\gamma$, $N=0$ (dash dot line); second-order correlation function in the presence of squeezed vacuum with frequency-dependent squeezing phase when $\Omega=10\gamma$, $N=0.3$, $|M|=[N(N+1)]^{1/2}$, $\Phi(\omega_A)=0$, $\Phi(\omega_A+\Omega)=\pi$, $G(\omega_A)\neq 0$ (solid line).

tively, in (4) and (5) and $\Phi(\omega_A)$ and $\Phi(\omega_A+\Omega)$ both equal to either 0 or $\pi$. When the squeezed field is the output of a degenerate parametric oscillator, the minimum damping rate permitted is $\gamma/4$. The parameters required to achieve this minimum correspond to the strong squeezing limit for the rather unrealistic situation where the squeezing is maximized at the central peak and is negligible at the Rabi sideband frequencies. When the squeezed field is the output of a nondegenerate parametric oscillator the minimum damping rate permitted approaches $\gamma/2$. This occurs for parameters corresponding to the strong squeezing limit with maximum squeezing occurring at the Rabi sideband frequencies and negligible squeezing at the central atomic resonance.

## V. DISCUSSION

It has recently been demonstrated that the two-photon absorption rate of a trapped atom illuminated with a squeezed beam deviates from the classical quadratic intensity dependence in the low intensity limit [29]. In this regime it is found that a linear intensity dominates. This is a direct consequence of the nonclassical field correlations and consequently provides a signature of quantum mechanical behavior. As experiments with squeezed light progress, it is pertinent to consider whether the line-narrowing mechanisms predicted for a two-level atom in a squeezed vacuum may be improved by processing the squeezed light. It is also topical to consider alternative experiments to measure quantum effects rather than the fluorescent spectrum with its inherent problems associated with frequency analyzing weak fluorescent beams.

In this paper we have investigated the effect of a squeezed beam with frequency-dependent squeezing phase on a coherently driven, two-level atom. Squeezed light with this characteristic may be generated by propagating the output of a parametric oscillator through a linear dielectric with dispersive properties. Appropriate parameter choices lead to markedly subnatural widths of all three peaks of the fluorescent spectrum in the strong driving regime. The line-narrowing mechanisms found at the sidebands manifest themselves in the second-order correlation function which is found to decay at a subnatural rate. This quantum statistical measure suggests that nonclassical effects in resonance fluorescence in a squeezed vacuum may be detected in photocounting experiments as well as in the fluorescent spectrum and the absorption spectrum [30,31]. Throughout, we have neglected the losses associated with propagation through the dielectric medium. This, however, does not present any great problem because the quantum effects noted in this paper occur at modest levels of squeezing.

As a final remark we note that the theory presented in this paper assumes that only one atom is present in the squeezed beam at any moment in time. When multiatom effects are taken into account it has been shown that the second-order correlation function contains terms arising from first-order and second-order correlation functions arising from the single atom theory [11,28]. Consequently, quantum effects beyond the Mollow regime will only register if all three spectral peaks exhibit subnatural widths. The theory presented here satisfies this criterion.

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