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ATOMS IN A NARROW-BANDWIDTH SQUEEZED VACUUM*

R. Tanaś[1]
Nonlinear Optics Division, Institute of Physics, Adam Mickiewicz University, Umultowska 85, 61-614 Poznań, Poland

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Two possible descriptions of evolution of a two-level atom driven by a strong laser field and subjected to a squeezed vacuum with finite bandwidth are discussed. One is the master equation approach in which the squeezed vacuum is treated as a Markovian reservoir to the atom, and the other is the coupled-systems (or cascaded-systems) approach in which the degenerate parametric oscillator (DPO) producing squeezed vacuum is a part of the system. Examples of optical spectra obtained using both approaches are given.

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1 Introduction

Radiative properties of atoms change dramatically when the normal vacuum fluctuations of the electromagnetic field reservoir are replaced by the anisotropic, phase-dependent quantum fluctuations of the squeezed vacuum. Gardiner [1] has shown that in the squeezed vacuum the atomic dipole moment can decay with two vastly different rates, one much longer and the other much shorter than that in the normal vacuum. In consequence, a subnatural linewidth has been predicted in the spontaneous emission spectrum. The addition of a coherent driving field to the problem introduces a strong dependence of the atom dynamics and the fluorescence and absorption spectra on the relative phase between the coherent field and the squeezed field. Carmichael et al. [2] have shown, for example, that the central peak of the Mollow [3] triplet, depending on the phase, can either be much narrower or much broader than the natural linewidth of the atom. Apart from these quantitative changes, the qualitative changes of the fluorescence spectrum have also been predicted.

Another spectroscopic feature accessible to experimental verification is the probe absorption spectrum. Mollow [4] has predicted that the absorption spectrum of a weak field probing a system of two-level atoms driven by an off-resonant laser field consists of one absorption and one emission component at the Rabi sidebands and a small dispersionlike component at the center of the spectrum. The probe field can be amplified due to the population inversion

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[1] E-mail address: tanas@kielich.amu.edu.pl

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between the dressed states of the atom despite the fact that there is no population inversion between the bare atomic states.

Most of the studies dealing with the problem of a two-level atom in a squeezed vacuum assumed that the squeezed vacuum is broadband, i.e., the bandwidth of the squeezed vacuum is much larger than the atomic linewidth and the Rabi frequency of the driving field. Experimental realizations of squeezed states, however, indicate that the bandwidth of the squeezed light is typically of the order of the atomic linewidth. The most popular schemes for generating squeezed light are those using parametric oscillator operating below threshold, the output of which is a squeezed beam with a bandwidth of the order of the cavity bandwidth [5, 6].

First studies of the finite-bandwidth effects have been performed by Gardiner et al. [5], Parkins and Gardiner [6] and Ritsch and Zoller [7]. The approaches were based on stochastic methods and numerical calculations, and were applied to analyze the narrowing of the spontaneous emission and absorption lines. The fundamental effect of narrowing has been confirmed, but the effect of finite bandwidth was to degrade the narrowing of the spectral lines rather than enhance it. Later, however, numerical simulations done by Parkins [8, 9] demonstrated that for strong driving fields a finite bandwidth of squeezing can have positive effect on the narrowing of the Rabi sidebands. A nice review on the subject has been written by Parkins [10].

In 1993 Gardiner [11] and Charmichael [12] proposed a new theoretical approach that can be used to describe the evolution of the atom interacting with the narrow-bandwidth squeezed vacuum — the coupled-systems (or cascaded-systems) approach. In this approach one considers a quantum system consisting of two subsystems: first is the degenerate parametric oscillator (DPO) producing the squeezed vacuum field and second is the atom driven by the output of the first system. It is assumed that there is no coupling back from the second system to the first. This approach leads to the master equation (in variables of both systems) describing the evolution of the coupled-systems. It, however, does not allow for analytical solutions and the resulting master equation must be solved numerically. Gardiner and Parkins [13] used this approach to study the effects of driving systems by various kinds of nonclassical light including squeezed light.

Recently, Yeoman and Barnett [14] have proposed an analytical technique for investigating the behavior of a coherently driven atom damped by a squeezed vacuum with finite bandwidth. In the approach, they have derived a master equation and analytic expressions for the fluorescent spectrum for the simple case of a two-level atom exactly resonant with the frequencies of both the squeezed field and the driving field. Their analytical results agree with that of Parkins [8, 9] and show explicitly that the width of the central peak of the fluorescent spectrum depends solely on the squeezing present at the Rabi sideband frequencies. The advantage of their dressed atom method over the more complex treatments based on adjoint equation or stochastic methods [8, 9, 11–13, 15] is that simple analytical expressions for the spectra can be obtained, thus displaying explicitly the factors that determine the intensities of the spectral features and their widths. Recently, we have followed the idea of Yeoman and Barnett, and we have rederived in different way and extended, to include the detuning of the driving field from atomic resonance, the master equation describing the evolution of an atom in a squeezed vacuum with finite bandwidth [16] (our master equation differs slightly from that of Yeoman and Barnett). In this paper I would like to address the questions: What are the limits of applicability of our Markovian master equation tested by reference to the Gardiner-Parkins [13] coupled-systems approach? and What new, unexpected effects can be obtained for narrow bandwidth squeezed vacuum?

2 Two approaches to the problem

2.1 Coupled-systems approach

One of the simplest and convenient means of calculating the effects of finite squeezing bandwidth is provided by the coupled-systems approach [11, 12]. In this approach one considers a quantum system consisting of two subsystems. A field $b_{in}(1,t)$ drives the first system, and give rise to an output $b_{out}(1,t)$ which, after a propagation delay τ, becomes the input field $b_{in}(2,t)$ to the second system. In the coupled-systems approach it is assumed that the output from the first system drives the second system without there being any coupling back from the second system to the first, which experimentally can be achieved by appropriate isolation techniques.

In our case the first system is a degenerate parametric oscillator (DPO) the output of which drives a two-level atom. Using the quantum noise methods [17] one can arrive at the following master equation [13], which in the frame rotating with the frequency $\omega_L = \omega_s$, with ω_L being the laser driving field frequency and ω_s being the squeezing carrier frequency, has the form

$$\begin{aligned}\dot{\rho} &= \frac{1}{2} i\left[(\Delta\sigma_z - \Omega\left(\sigma_+ + \sigma_-\right) + (\epsilon a^{\dagger 2} - \epsilon^* a^2)), \rho\right] + \frac{\kappa}{2}\left\{2\, a\rho\, a^\dagger - \rho\, a^\dagger\, a - a^\dagger\, a\rho\right\} \\ &- \sqrt{\eta\kappa\gamma}\left\{[\sigma_+, a\,\rho] + [\rho\, a^\dagger, \sigma_-]\right\} + \frac{\gamma}{2}\left\{2\,\sigma_-\rho\,\sigma_+ - \rho\,\sigma_+\,\sigma_- - \sigma_+\,\sigma_-\rho\right\}, \qquad (1)\end{aligned}$$

where $\Delta = \omega_L - \omega_A = \omega_s - \omega_A$ is the detuning of the laser (and squeezed) field frequency from the atomic resonance, κ is the DPO cavity bandwidth, γ is the natural atomic linewidth and Ω is the Rabi frequency of the laser field driving the atom. The parameter η $(0 < \eta \leq 1)$ describes the matching of the incident squeezed vacuum to the modes surrounding the atom. For perfect matching $\eta = 1$, whereas $\eta < 1$ for an imperfect matching, which is always the case in experimental situations [18–21]. In order to observe the effects of the squeezed vacuum on the atom the parameter η should be as close to unity as possible. This requirement could be difficult to achieve in experiments, although some schemes involving optical cavities have been proposed [22, 23]. On the other hand, if the fluorescent field radiated by the atom to the non-squeezed modes is to be observed, η cannot be exactly unity because the radiation rate to the non-squeezed modes, which is $(1-\eta)\gamma$, would be zero and no fluorescence would be observed. In the coupled-systems approach one has a choice of detecting either transmitted light or the fluorescent light radiated by the atom to the modes of the ordinary vacuum. The transmitted light is a superposition of the squeezed vacuum coming from the DPO and the field radiated by the atom to the modes occupied by the squeezed vacuum.

Because of the coupling between the operators of the two subsystems the master equation (1) does not allow for analytical solution. Since the DPO enters to the master equation as part of the system, and the ordinary vacuum, which is Markovian, plays the role of the reservoir to the total system, there are no requirements for the bandwidth of the squeezed vacuum to be broad. The cavity bandwidth κ can be comparable or even smaller than the atomic linewidth γ. This is an advantage of the approach with respect to approaches that treat squeezed vacuum as a reservoir to the atom. If the mean number of photons in the cavity is small, effective numerical solutions of the master equation (1) are possible [24]. For $\kappa \gg \gamma$ the solutions can be compared with the approximate Markovian master equation testing the validity of the latter. For squeezing bandwidths smaller than the natural atomic linewidth, for which standard Born-Markov master equations are not valid, one can find some new and unexpected features in the atomic spectra.

2.2 Squeezed vacuum as a reservoir

When the bandwidth of the DPO cavity, even though being finite, is much larger than the atomic linewidth, one can treat the squeezed vacuum as a reservoir to the atom and derive the master equation, in the Born-Markov approximation, that describes the dynamics of the atomic variables only. If, moreover, the atom is driven by a classical coherent laser field one can follow the Yeoman and Barnett [14] idea to perform the dressing transformation first and next apply the standard perturbation procedure to derive the master equation. The master equation obtained in this way has the structure which is known for the broadband squeezed vacuum, but with the new rate coefficients and some new terms that do not appear for broadband squeezing. We have derived such a master equation [16], which in the frame rotating with the laser frequency ω_L can be written as

$$\begin{aligned}
\dot{\rho} &= \frac{1}{2} i\left[(\gamma\,\delta\,\sigma_z - \Omega\,(\sigma_+ + \sigma_-)), \rho\right] \\
&+ \frac{1}{2}\gamma \tilde{N}\,(2\,\sigma_+\rho\,\sigma_- - \sigma_-\sigma_+\rho - \rho\,\sigma_-\sigma_+) \\
&+ \frac{1}{2}\gamma\,(\tilde{N}+1)\,(2\,\sigma_-\rho\,\sigma_+ - \sigma_+\sigma_-\rho - \rho\,\sigma_+\sigma_-) \\
&- \gamma\,\tilde{M}\,\sigma_+\rho\,\sigma_+ - \gamma\,\tilde{M}^*\,\sigma_-\rho\,\sigma_- \\
&+ \frac{1}{4} i\,\left(\beta\,[\sigma_+,[\sigma_z,\rho]] - \beta^*[\sigma_-,[\sigma_z,\rho]]\right) ,
\end{aligned} \tag{2}$$

where γ is the natural atomic linewidth,

$$\tilde{N} = N(\omega_L+\Omega') + \frac{1}{2}(1-\tilde{\Delta}^2)\,\gamma_n , \tag{3}$$

$$\tilde{M} = \left(|M(\omega_L+\Omega')| + i\,\tilde{\Delta}\,\delta_M\right) \mathrm{e}^{i\phi} - \frac{1}{2}(1-\tilde{\Delta}^2)\,(\gamma_n - i\,\delta_n) , \tag{4}$$

$$\delta = \frac{\Delta}{\gamma} + \tilde{\Delta}\,\delta_N + \frac{1}{2}(1-\tilde{\Delta}^2)\,\delta_n , \tag{5}$$

$$\beta = \gamma\tilde{\Omega}\left[\delta_N + \delta_M\,\mathrm{e}^{i\phi} - i\,\tilde{\Delta}\,(\gamma_n - i\,\delta_n)\right] , \tag{6}$$

$$\gamma_n = N(\omega_L) - N(\omega_L+\Omega') - (|M(\omega_L)| - |M(\omega_L+\Omega')|)\,\cos\phi , \tag{7}$$

$$\delta_n = (\,|M(\omega_L)| - |M(\omega_L+\Omega')|\,)\,\sin\phi , \tag{8}$$

$$\delta_N = \frac{1}{\pi}\mathcal{P}\int_{-\infty}^{\infty}\frac{N(x)}{x+\Omega'}\,\mathrm{d}x , \qquad \delta_M = \frac{1}{\pi}\mathcal{P}\int_{-\infty}^{\infty}\frac{|M(x)|}{x+\Omega'}\,\mathrm{d}x , \tag{9}$$

$$\tilde{\Omega} = \frac{\Omega}{\Omega'} , \quad \tilde{\Delta} = \frac{\Delta}{\Omega'} , \quad \Omega' = \sqrt{\Omega^2+\Delta^2} , \tag{10}$$

and ϕ is the phase of squeezing ($M(\omega) = |M(\omega)| \exp(i\phi)$). In the derivation of Eq. (2) we have included the divergent frequency shifts (the Lamb shift) to the redefinition of the atomic transition frequency. Moreover, we have assumed that the squeezed vacuum is symmetric about the central frequency ω_L, so that $N(\omega_L - \Omega') = N(\omega_L + \Omega')$, and a similar relation holds for $M(\omega)$.

The Cauchy principal values of the integrals in Eq. (9) can be evaluated using the contour integration which for the degenerate parametric oscillator gives

$$\delta_N = \gamma\Omega' \frac{\lambda^2 - \mu^2}{4} \left[\frac{1}{\mu(\Omega'^2 + \mu^2)} - \frac{1}{\lambda(\Omega'^2 + \lambda^2)} \right], \tag{11}$$

$$\delta_M = \gamma\Omega' \frac{\lambda^2 - \mu^2}{4} \left[\frac{1}{\mu(\Omega'^2 + \mu^2)} + \frac{1}{\lambda(\Omega'^2 + \lambda^2)} \right], \tag{12}$$

$$\lambda = \frac{1}{2}\kappa + \epsilon, \quad \mu = \frac{1}{2}\kappa - \epsilon, \tag{13}$$

where κ is the bandwidth of the DPO cavity and ϵ the amplitude of the pump field.

The master equation (2) has the standard form known from the broadband squeezing approaches with the new effective squeezing parameters $\tilde{N}$ and $\tilde{M}$ given by (3) and (4). There are also new terms, proportional to β which are essentially narrow bandwidth modifications to the master equation. All the narrow bandwidth modifications are determined by the parameters γ_n, δ_n and the shifts δ_N and δ_M defined in (7) – (8) and (11)– (12). These parameters become zero when the squeezing bandwidth goes to infinity.

A clear advantage of the master equation (2) over the master equation (1) is that it leads to the optical Bloch equations for the atomic variables which allow for analytical solutions to the atomic dynamics and optical spectra. A number of atomic radiative features such as steady state solutions, resonance fluorescence spectra and probe absorption spectra derived from this master equation have already been discussed [16,25]. One has to remember, however, that the master equation (2) is valid for $\kappa \gg \gamma$ because of the Markov approximation. When κ becomes comparable to γ the Markov approximation breaks down and master equation (2) can even lead to unphysical results. In contrast, the master equation (1) is valid for any bandwidth κ of the DPO cavity and can be applied equally well for $\kappa < \gamma$. In this paper I will compare some of the results obtained for optical spectra using both methods.

3 Optical spectra

The stationary spectrum of the resonance fluorescence from a two-level atom is given by the Fourier transform of the two-time atomic correlation function as [3,26]

$$\mathcal{F}(\omega) = \frac{1}{\pi}\mathrm{Re}\left\{ \int_0^\infty \langle \sigma_+(0)\sigma_-(\tau)\rangle_{ss}\, \mathrm{e}^{i(\omega-\omega_L)\tau}\, \mathrm{d}\tau \right\}, \tag{14}$$

where Re denotes the real part of the integral. The two-time correlation function appearing in (14) can be found by applying the quantum regression theorem [27]. In the case of the coupled-systems approach this can be done merely numerically using the master equation (1) while from the master equation (2) we get the following equations of motion for the two-time correlation functions

$$\frac{\partial}{\partial \tau}\begin{pmatrix} \langle\sigma_+(0)\sigma_-(\tau)\rangle_{ss} \\ \langle\sigma_+(0)\sigma_+(\tau)\rangle_{ss} \\ \langle\sigma_+(0)\sigma_z(\tau)\rangle_{ss} \end{pmatrix} = \mathbf{A}\begin{pmatrix} \langle\sigma_+(0)\sigma_-(\tau)\rangle_{ss} \\ \langle\sigma_+(0)\sigma_+(\tau)\rangle_{ss} \\ \langle\sigma_+(0)\sigma_z(\tau)\rangle_{ss} \end{pmatrix} + \langle\sigma_+\rangle_{ss}\begin{pmatrix} 0 \\ 0 \\ -\gamma \end{pmatrix}, \quad (15)$$

where $\mathbf{A}$ is the 3×3 matrix

$$\mathbf{A} = \begin{pmatrix} -\gamma(\frac{1}{2}+\tilde{N}-i\delta) & -\gamma\tilde{M} & \frac{i}{2}\Omega \\ -\gamma\tilde{M}^* & -\gamma(\frac{1}{2}+\tilde{N}+i\delta) & -\frac{i}{2}\Omega \\ i(\Omega+\beta^*) & -i(\Omega+\beta) & -\gamma(1+2\tilde{N}) \end{pmatrix}, \quad (16)$$

and the initial values for the correlation functions are

$$\langle\sigma_+\sigma_-\rangle_{ss} = \frac{1}{2}(1+\langle\sigma_z\rangle_{ss}), \quad \langle\sigma_+\sigma_+\rangle_{ss} = 0, \quad \langle\sigma_+\sigma_z\rangle_{ss} = -\langle\sigma_+\rangle_{ss}. \quad (17)$$

Taking the Laplace transforms of (15) one gets the system of algebraic equations that can be easily solved giving the analytical expression for the resonance fluorescence spectrum [16].

To interpret some of the spectral features it is convenient to relate the resonance fluorescence spectrum to the quadrature noise spectrum (squeezing spectrum) as [28]

$$\mathcal{F}_{inc}(\omega+\omega_L) = S_X(\omega) + S_Y(\omega) + S_A(\omega), \quad (18)$$

where

$$S_X(\omega) = \frac{1}{2\pi}\mathrm{Re}\int_0^\infty \cos(\omega\tau)\left[\langle\sigma_+(0),\sigma_-(\tau)\rangle_{ss} + \langle\sigma_+(0),\sigma_+(\tau)\rangle_{ss}\right]\mathrm{d}\tau, \quad (19)$$

$$S_Y(\omega) = \frac{1}{2\pi}\mathrm{Re}\int_0^\infty \cos(\omega\tau)\left[\langle\sigma_+(0),\sigma_-(\tau)\rangle_{ss} - \langle\sigma_+(0),\sigma_+(\tau)\rangle_{ss}\right]\mathrm{d}\tau, \quad (20)$$

are, respectively, in-phase and out-of-phase quadrature components of the noise spectrum, and

$$S_A(\omega) = -\frac{1}{\pi}\int_0^\infty \sin(\omega\tau)\mathrm{Im}\,\langle\sigma_+(0),\sigma_-(\tau)\rangle_{ss}\,\mathrm{d}\tau \quad (21)$$

is the asymmetric contribution to the spectrum. In (19–(21), $\langle a,b\rangle \equiv \langle ab\rangle - \langle a\rangle\langle b\rangle$ denotes the covariance.

Effective numerical solutions of the master equation (1) are possible when the mean number of photons $\langle a^\dagger a\rangle$ in the cavity is small ($\langle a^\dagger a\rangle < 1$) [13, 24]. In this case it is sufficient to take about ten lowest photon states as a basis of the photon Hilbert space and the two atomic states that form the atomic Hilbert space. Steady state solutions of the master equation (1) together with the quantum regression theorem allow us to find optical spectra for the transmitted and fluorescent fields as well as atomic quadrature noise spectra for the cases when squeezing bandwidth is smaller or comparable to the natural atomic linewidth.

Since the coupled-systems approach is valid for any value of κ, numerical solutions for the resonance fluorescence spectra obtained from the master equation (1) can serve as a reference for the spectra obtained by solving Eqs. (15). It is particularly interesting to check how large must be κ to get reasonable agreement between the two approaches. In Fig. 1 I have plotted examples of the spectra for (a) weak and (b) strong fields. Figure 1(a) shows an example of the

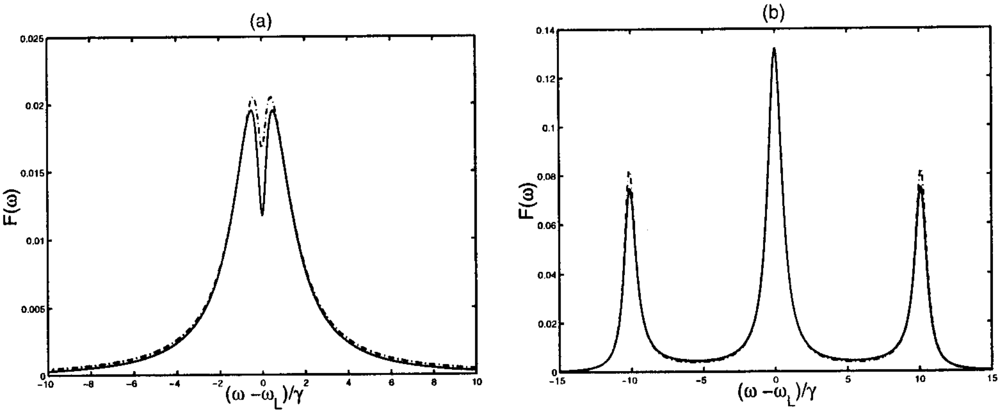

Fig. 1. Resonance fluorescence spectra plotted according to the coupled-systems approach (solid line) and the master equation (2) (dashed-dotted line). The parameters (in units of γ) are: (a) $\kappa = 40$, $\epsilon = 5$, $\Omega = 0.35$ and (b) $\kappa = 10$, $\epsilon = 1.25$, $\Omega = 10$.

spectrum for a special value of the Rabi frequency $\Omega = 0.35$ for which the spectrum shows a dip in the center. As it is seen for the squeezing bandwidth $\kappa = 40$ our master equation reproduces the dip and it is in quite a good agreement with the spectrum obtained from the coupled-systems approach. Figure 1(b) shows an example of the spectrum for $\Omega = 10$ and in this case $\kappa = 10$, which means that the squeezing bandwidth covers one third of the range of frequencies shown in the figure. This example convincingly shows that the bandwidth of squeezed light should be greater than the atomic linewidth but not necessarily greater than the Rabi frequency to make the Markovian approximation justified. The sidebands are slightly shifted with respect to the coupled-systems result, but the central line fits almost perfectly, for the squeezing bandwidth as small as $\kappa = 10$. For the above examples, we have assumed $\eta = 0.98$ in the master equation (1), and for better comparison the spectra are normalized to the same rate. Of course, as the values of the squeezing bandwidth κ become larger and larger the Markovian approximation works better and the analytical results based on our master equation are more reliable. A number of features concerning the resonance fluorescence spectra [16] as well as the probe absorption spectra [25] derived using the master equation (2) have already been discussed.

The situation is quite different when the squeezing bandwidth κ is comparable or smaller than the atomic linewidth. In this case the Markovian approximation breaks down and the master equation (2) becomes useless. One can still use the master equation (1) to calculate the atomic dynamics and optical spectra. Probably the most interesting and widely discussed problem related with modifications of the atomic radiation properties due to squeezed vacuum is the narrowing of the spectral lines predicted for the broadband squeezed vacuum. Gardiner and Parkins [13] have solved numerically the master equation (1) and have found that the squeezing induced line narrowing in the fluorescence spectrum appears only for the cavity linewidths κ sufficiently large with respect to the atomic natural linewidth γ. The narrowing decreases with decreasing κ and disappears for $\kappa \approx \gamma$. Can we find something interesting when κ becomes smaller than γ? Here, I will show an example of the structure that we have found recently [29], and which appears in the resonance fluorescence spectrum for $\kappa < \gamma$. Figure 2 shows a three-

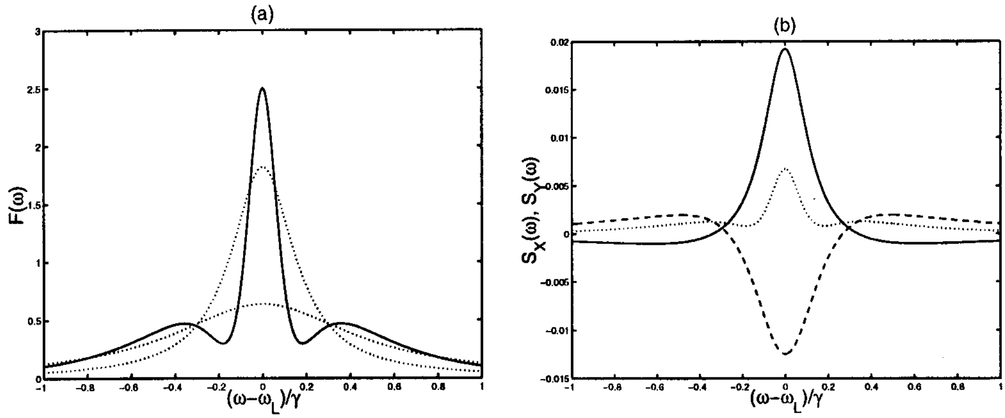

Fig. 2. (a) Resonance fluorescence spectrum for the narrow bandwidth squeezed vacuum plotted according to the coupled-systems approach: $\kappa = 0.35$, $\epsilon = 0.0087$, $\eta = 0.9$, $\Omega = 0$. Dotted lines are two Lorentzians with linewidths γ and κ. (b) Squeezing spectra $S_X(\omega)$ (solid line), $S_Y(\omega)$ (dashed line) and the resulting fluorescence spectrum (dotted line) for the same parameters.

peak structure of the resonance fluorescence spectrum for an atom driven by the squeezed vacuum field from DPO (no coherent pump) and the squeezing spectra explaining the structure of the fluorescence spectrum. The spectrum is obtained for $\Delta = 0$ (resonant case), $\kappa = 0.35$ (in units of γ) and the pump field amplitude $\epsilon = 0.0087$. The parameter $\eta = 0.9$ describes the fraction of the squeezed modes seen by the atom with respect to the fraction $(1 - \eta)$ of the ordinary vacuum modes. For reference, we plot the Lorentzian with the atomic linewidth (broader) and the Lorentzian the linewidth κ (narrower). One can see that the fluorescence spectrum has two components: the broad background with the natural linewidth at the wings and the narrow peak with the width narrower than the DPO cavity bandwidth at the center. The appearance of the unusual features in the spectra can be explained as arising from the squeezing produced by the atom. This is shown in Fig. 2(b) where we plot the squeezing spectra for the fluorescent field defined by (19) and (20). The $S_Y(\omega)$ quadrature is negative for frequencies near the carrier frequency ω_s, i.e., it shows squeezing near the center of the spectrum [30, 31]. This squeezing is responsible for the unusual shape of the resonance fluorescence spectrum. The $S_X(\omega)$ quadrature is positive, and adding the two squeezing spectra gives the fluorescence spectrum shown in Fig. 2(a) [or dotted line in figure Fig. 2(b)]. Clearly, the structure arises from squeezing in the fluorescence field and an experimental observation of the effect could be a manifestation of the quantum nature of squeezed light. I would like to emphasize the fact that the squeezing in the fluorescent field has been obtained by driving the atom with squeezed light only, without a coherent driving field.

The features discussed here depend crucially on the value of η, which should be as close to unity as possible to have the coupling between the two subsystems as high as possible. On the other hand, there is only a fraction $(1-\eta)\gamma$ of the radiation that goes to the modes different from the squeezed vacuum modes, and this rate must be non-zero to observe resonance fluorescence to the non-squeezed modes. In our calculations presented in Figs. 2 we have assumed $\eta = 0.9$. The features, however, degrade quickly as η decreases and disappear for $\eta \approx 0.6$.

4 Conclusion

Although it has been known for quite some time that radiative properties of the atom in a squeezed vacuum can be significantly different from that in ordinary vacuum, and a number of new effects have been predicted, their experimental verification appeared to be a demanding and difficult task. Recent experiments [21,32] with squeezed light give some hope for progress in this field of research. In this article I have compared two different theoretical descriptions of the atom interacting with a squeezed vacuum with finite bandwidth. I have shown that even for not very large squeezing bandwidth one can use a simple master equation derived under the Markovian approximation to describe atomic dynamics and optical spectra. I have also indicated a possibility to find new nonclassical features in a relatively unexplored area of research when atoms interact with a squeezed field with the bandwidth smaller than the natural atomic linewidth.

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