Time-dependent spectrum of a strongly driven two-level atom in the squeezed vacuum

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The time-dependent fluorescence spectrum from a strongly driven two-level atom damped by a broadband-squeezed vacuum is obtained analytically using the counting-rate definition of Eberly and Wodkiewicz [J. Opt. Soc. Am. **67**, 1252 (1977)] and employing the high-field approximation. The spectrum shows considerable modifications compared to the corresponding time-dependent spectrum in a normal vacuum. It is now dependent on the relative phase between the squeezed vacuum field and the coherent driving field. In particular, the time development of the central peak of the Mollow triplet (which is known to show a phase-dependent linewidth in the steady state) is discussed. The effects of detuning on the spectrum are also presented.

I. INTRODUCTION

The time-dependent physical spectrum of resonance fluorescence has been studied recently, for both two-level[1] and three-level[2–4] atoms. Such studies not only reveal the idea of time dependence in a spectral measurement but also take into account the existence of physical measurement devices such as the Fabry-Pérot interferometer, and thus give deeper insight into the understanding of the fluorescence process itself.[5] Eberly, Kunasz, and Wodkiewicz[1] have studied the time-dependent spectrum in considerable detail for a strongly driven two-level atom and observed changes in the spectrum with time and found that the total spectral intensity oscillates according to the Rabi flopping frequency. The calculated transient spectrum is symmetric and has a three-peaked structure for a resonant field but is quite asymmetric under off-resonance conditions before reaching the steady-state values. The effect of laser noise[1] in the time-dependent spectrum is also accounted for. More recently, the time dependence of resonance fluorescence from a two-level atom has been studied both experimentally[6] and theoretically[7–17] by various groups. Some of these studies[9,10,14–17] include also the effect of the slowly varying pulse shape of the irradiating laser field on the time-dependent spectrum.

The recent successful generation[18–20] of squeezed light has opened new avenues in the studies of atom-field interactions. A broadband squeezed light can modify the spectroscopic properties of an atom. Gardiner[21] has considered the interaction of a two-level atom with a broadband-squeezed vacuum and shown that the atomic dipole decay is phase dependent. That is, one quadrature of the atomic polarization decays with an enhanced rate and the other with a reduced rate compared with the normal atomic decay. Carmichael, Lane, and Walls[22] have considered the interaction of a coherently driven two-level atom with a broadband-squeezed vacuum and have obtained the steady-state spectrum of the fluorescent radiation which shows very interesting features as compared with the Mollow spectrum in the normal vacuum. It is now dependent on the relative phase between the driving field and the squeezed vacuum. If squeezing is large, coherent scattering is negligible at low driving fields, whereas it dominates in normal resonance fluorescence; the incoherent scattering is composed of components one of which is much broader and the other narrower than the natural linewidth. At very high driving field intensities, the fluorescence spectrum is a phase-sensitive triplet. In particular, it is possible to change a spectrum with a broad central peak into one with a subnatural linewidth by changing the relative phase between the squeezed vacuum and the driving field by π.[22] The possibility of obtaining a subnatural linewidth in the absorption spectrum of a two-level atom in a broadband-squeezed vacuum has been discussed by Ritsch and Zoller.[23]

The above theories describing the interaction of an atom with squeezed input fields are essentially based on a white noise or broadband assumption. For a realistic source of light such as the output of a degenerate parametric amplifier, this assumption is not strictly valid. More recently, the spectroscopic properties of an atom embedded in a finite bandwidth squeezed vacuum have been reexamined.[24–26]

In this paper we discuss the transient spectra of a two-level atom driven by a strong coherent field and interacting with a broadband-squeezed vacuum field. In Sec. II we formulate the problem and obtain the equation of motion of atomic variables in the high-field approximation. In Sec. III we obtain the analytical solution for the time-dependent spectrum in the high-field approximation and discuss the results.

II. FORMULATION OF THE PROBLEM

A. Master equation for a two-level atom damped by a squeezed vacuum

The Hamiltonian describing the radiative decay of a two-level atom of transition frequency ω_A via its interaction with the multimode field is given by[22] [under the rotating-wave approximation (RWA) and $\hbar=c=1$]

$$H_s=\omega_A S_z+H_0+(S_+\Gamma_0+S_-\Gamma_0^\dagger)\ , \tag{1}$$

where $S_z=\frac{1}{2}(|2\rangle\langle 2|-|1\rangle\langle 1|)$, $S_+=|2\rangle\langle 1|$, $S_-=|1\rangle\langle 2|$ with the ground state and excited state denoted by $|1\rangle$ and $|2\rangle$, respectively. H_0 is the term for the free radiation field and Γ_0 and $\Gamma_0^\dagger$ are bath operators corresponding to the positive and negative frequency component of the radiation field. We assume (1) a broadband-squeezed vacuum centered around ω_L (which will be the frequency of the coherent driving field also) and (2) all the modes coupled to the atom are squeezed so there is no spontaneous emission to the unsqueezed vacuum modes. As large squeezing modifies the polarization decay rate and one polarization quadrature decays faster than the natural lifetime, this assumption requires a squeezing bandwidth much larger than natural linewidth. In other words, the squeezing bandwidth should be much broader than the spectral width of fluorescence so that the fluorescent spectrum falls well within the bands of squeezed-vacuum modes. Finally, the bandwidth of squeezing is assumed to be sufficiently broad so that it appears as δ-correlated squeezed white noise to the atom. The correlation functions can be expressed as follows:

$$\begin{aligned}
&\langle \Gamma_0^\dagger(t)\Gamma_0(t')\rangle=\gamma N\delta(t-t')\ ,\\
&\langle \Gamma_0(t)\Gamma_0^\dagger(t')\rangle=\gamma(N+1)\delta(t-t')\ ,\\
&\langle \Gamma_0(t)\Gamma_0(t')\rangle=\gamma M\exp(-2i\omega_L t)\delta(t-t')\ ,\\
&\langle \Gamma_0^\dagger(t)\Gamma_0^\dagger(t')\rangle=\gamma M^*\exp(2i\omega_L t)\delta(t-t')\ .
\end{aligned} \tag{2}$$

Here γ is the radiative damping constant into the unsqueezed vacuum. The parameters N and M describe squeezing and obey the following relationship:

$$|M|^2\le N(N+1)\ ,$$
$$M=|M|\exp(i\varphi)\ ,$$

where phase ϕ will depend on the specific details of the scheme used to generate the squeezed vacuum.

The master equation for the reduced density operator can be easily obtained using Eqs. (1) and (2) as follows:

$$\begin{aligned}
\frac{\partial\rho}{\partial t}=&-i\omega_A[S_z,\rho]+\frac{\gamma}{2}(N+1)(2S_-\rho S_+-S_+S_-\rho-\rho S_+S_-)+\frac{\gamma}{2}N(2S_+\rho S_--S_-S_+\rho-\rho S_-S_+)\\
&-\gamma M\exp(-2i\omega_L t)S_+\rho S_+-\gamma M^*\exp(2i\omega_L t)S_-\rho S_-\ .
\end{aligned} \tag{3}$$

B. Equation of motion for atomic operators

The Hamiltonian describing a two-level atom driven by a coherent field and interacting with a squeezed vacuum can be written as

$$H=-[\mu E\exp(-i\omega_L t)S_++\mu^*E^*\exp(i\omega_L t)S_-]+H_s\ , \tag{4}$$

where E is the amplitude of the driving field, ω_L the frequency of the driving field. μ is the atomic dipole moment and H_s is given by (1). The master equation with this Hamiltonian can be easily obtained as follows:[22]

$$\begin{aligned}
\frac{\partial\rho}{\partial t}=&\,i[\mu E\exp(-i\omega_L t)S_++(\mu E)^*\exp(i\omega_L t)S_-,\rho]-i\omega_A[S_z,\rho]+\frac{\gamma}{2}(N+1)(2S_-\rho S_+-S_+S_-\rho-\rho S_+S_-)\\
&+\frac{\gamma}{2}N(2S_+\rho S_--S_-S_+\rho-\rho S_-S_+)-\gamma M^*\exp(-2i\omega_L t)S_+\rho S_+-\gamma M^*\exp(2i\omega_L t)S_-\rho S_-\ .
\end{aligned} \tag{5}$$

From Eq. (5) we can get equations of motion for $\langle S_+\rangle$, $\langle S_-\rangle$, and $\langle S_z\rangle$ in a frame rotating with frequency ω_L as

$$\begin{aligned}
\langle \dot S_+\rangle=&\,[i\Delta-\gamma(N+\tfrac{1}{2})]\langle S_+\rangle-\gamma|M|e^{-i\phi}\langle S_-\rangle\\
&-i\Omega_0\langle S_z\rangle\ ,\\
\langle \dot S_-\rangle=&\,[-i\Delta-\gamma(N+\tfrac{1}{2})]\langle S_-\rangle-\gamma|M|e^{i\phi}\langle S_+\rangle\\
&+i\Omega_0\langle S_z\rangle\ ,\\
\langle \dot S_z\rangle=&-\gamma(2N+1)\left[\langle S_z\rangle+\frac{1}{2N+1}\right]-2i\Omega_0\langle S_+\rangle\\
&+2i\Omega_0\langle S_-\rangle\ ,
\end{aligned} \tag{6}$$

where $\Delta=\omega_A-\omega_L$, $\Omega_0=|\mu E|$ and $\phi=2\phi_L-\varphi$.

A great deal of simplification results if we restrict ourselves to the case of a strong field ($\Omega_0\gg\gamma$). For this we transform to operators defined by[27]

$$\begin{aligned}
&R_x=S_y\ ,\\
&R_y=-(\Delta/2\Omega)S_x+(\Omega_0/\Omega)S_z\ ,\\
&R_z=(\Omega_0/\Omega)S_x+(\Delta/2\Omega)S_z\ ,
\end{aligned} \tag{7}$$

where $\Omega^2=\Omega_0^2+\Delta^2/4$.

Next, we go to the interaction picture defined by

$$\bar\psi=e^{2i\Omega t R_z}\psi e^{-2i\Omega t R_z},\quad \psi\equiv\begin{bmatrix}\langle S_+\rangle\\ \langle S_-\rangle\\ \langle S_z\rangle\end{bmatrix}\ . \tag{8}$$

With this, the resulting equation splits into two parts: one containing no oscillatory terms and the other involving rapidly oscillating terms like $\exp(\pm 2i\Omega t)$ and $\exp(\pm 4i\Omega t)$. If we make the secular approximation, i.e., neglect oscillatory terms and revert back to the Schrödinger picture, we arrive at the following equations:

$$\langle \dot{R}_+(t)\rangle = (2i\Omega - \gamma_+)\langle R_+(t)\rangle \ , \tag{9a}$$

$$\langle \dot{R}_z(t)\rangle = -\gamma_0\langle R_z(t)\rangle - \sqrt{r}\,\gamma \ , \tag{9b}$$

where

$$\begin{aligned} \gamma_+ &= (\gamma/2)[(2N+1)(3-r)/2 - |M|(1-r)\cos\phi] \ , \\ \gamma_0 &= \gamma[(2N+1)(1+r)/2 + |M|(1-r)\cos\phi] \ , \\ r &= \Delta^2/4\Omega^2 \ . \end{aligned} \tag{10}$$

Equations (9a) and (9b) can be solved and the solution can be written in the compact form

$$\begin{aligned} \langle R_+(t+\tau)\rangle &= \langle R_+(t)\rangle e^{(2i\Omega-\gamma_+)\tau} \ , \\ \langle R_z(t+\tau)\rangle &= \langle R_z(t)\rangle e^{-\gamma_0\tau} - (\gamma\sqrt{r})/\gamma_0 \ . \end{aligned} \tag{11}$$

The steady-state solutions are

$$\begin{aligned} \langle R_+\rangle_{ss} &= 0 \ , \\ \langle R_z\rangle_{ss} &= -(\gamma/\gamma_0)\sqrt{r} \ . \end{aligned} \tag{12}$$

III. TIME-DEPENDENT FLUORESCENCE SPECTRUM

The time-dependent fluorescence spectrum can be written as a double convolution integral[5]

$$I(D,t,\Gamma) = 2\Gamma\,\mathrm{Re}\left[\int_0^t dt'\,e^{-\Gamma(t-t')}\int_0^{t-t'} d\tau\, e^{(\Gamma/2-iD)\tau}\langle S_+(t+\tau)S_-(t)\rangle\right] . \tag{13}$$

In Eq. (13), Γ is the full width at half maximum (FWHM) of the transmission peak of the Fabry-Pérot interferometer used to measure the frequency spectrum of the fluorescence, D is the detuning or the frequency offset of the Fabry-Pérot line center (ω) above the laser frequency, that is $D=\omega-\omega_L$. This scheme for calculating the spectrum assumes that there exists a small "window" of unsqueezed vacuum modes through which we can view the fluorescence.[22]

We note that[27]

$$\begin{aligned} S_+ &= i/2(1+\sqrt{r})R_+ + i/2(1-\sqrt{r})R_- + (1-r)^{1/2}R_z \ , \\ S_z &= -i/2(1-r)^{1/2}(R_+ - R_-) - \sqrt{r}R_z \ , \end{aligned} \tag{14}$$

so that for an atom initially in the ground state we get

$$\begin{aligned} \langle S_+(t+\tau)S_-(t)\rangle = {} & \tfrac{1}{8}(1-r)\sqrt{r}e^{-\gamma_0 t}(-e^{(2i\Omega-\gamma_+)t} + e^{(-2i\Omega-\gamma_+)t}) \\ & + \tfrac{1}{4}(1-r)[(\tfrac{1}{2}-\sqrt{r}\,\nu/\nu_0)e^{(2i\Omega-\gamma_+)t} + (\tfrac{1}{2}+\sqrt{r}\,\gamma/\gamma_0)e^{(-2i\Omega-\gamma_+)t}] \\ & - \tfrac{1}{8}(1-r)[(1+\sqrt{r})e^{(2i\Omega-\gamma_+)(t+\tau)} + (1-\sqrt{r})e^{(-2i\Omega-\gamma_+)(t+\tau)}] \\ & - \tfrac{1}{8}(1-r)[(1+\sqrt{r})e^{(2i\Omega-\gamma_+)t-\gamma_0\tau} + (1-\sqrt{r})e^{(-2i\Omega-\gamma_+)t-\gamma_0\tau}] \\ & + \tfrac{1}{4}\sqrt{r}(\gamma/\gamma_0)(1-r)[(1+\sqrt{r})e^{(2i\Omega-\gamma_+)t} - (1-\sqrt{r})e^{(-2i\Omega-\gamma_+)t}] \\ & + \tfrac{1}{4}(1-r)e^{-\gamma_0\tau} + \tfrac{1}{2}(1-r)r(\gamma/\gamma_0)e^{-\gamma_0\tau} + (1-r)r(\gamma/\gamma_0)^2 \ . \end{aligned} \tag{15}$$

After substituting the above expression in Eq. (13), and carrying out integration completely, we finally arrive at the following explicit expression for the time-dependent spectrum:

$$I(D,t,\Gamma) = \sum_{n=1}^{8} S_n \ , \tag{16}$$

where

$$\begin{aligned} S_1 = {} & -\tfrac{1}{8}(1-r)\sqrt{r}\,[(x_1\{e^{-\gamma_0 t} - e^{-(\Gamma/2+\gamma_+)t}\cos[(2\Omega-D)t]\} - \gamma_1 e^{-(\Gamma/2+\gamma_+)t}\sin[(2\Omega-D)t] \\ & - Z_1(e^{-\gamma_0 t} - e^{-rt})) - [\Omega\to-\Omega]] \ , \\ S_2 = {} & \tfrac{1}{4}(1-r)[(\tfrac{1}{2}-\sqrt{r}\,\gamma/\gamma_0)(x_2 - e^{-\gamma_+ t}\{X_2\cos[(2\Omega-D)t] - Y_2\sin[(2\Omega-D)t]\} - Z_2(1-e^{-\Gamma/t})) \\ & + (\tfrac{1}{2}+\sqrt{r}\,\gamma/\gamma_0)[\Omega\to-\Omega]] \ , \\ S_3 = {} & -\tfrac{1}{8}(1-r)[(1+\sqrt{r})(e^{-\gamma_+ t}[(X_3-Z_3)\cos(2\Omega t) - (Y_3-Z_4)\sin(2\Omega t)] \\ & - e^{-(\Gamma/2+\gamma_+)t}\{X_3\cos[(2\Omega-D)t] + Y_3\sin[(2\Omega-D)t]\} + Z_3 e^{-\Gamma t}) - (1-\sqrt{r})[\Omega\to-\Omega]] \ , \end{aligned}$$

$$
\begin{aligned}
S_4 &= -\tfrac{1}{8}(1-r)((1+\sqrt{r}\,)\{ e^{(\Gamma-\gamma_+)t}[(X_4-Z_5)\cos(2\Omega t)-(Y_4-Z_6)\sin(2\Omega t)] \\
&\qquad -[X_4\cos(Dt)-Y_4\sin(Dt)]e^{(\Gamma/2-\gamma_0)t}\}-(1-\sqrt{r}\,)[\Omega\to-\Omega]) , \\
S_5 &= \tfrac{1}{4}\sqrt{r}\,(\gamma/\gamma_0)(1-r)((1+\sqrt{r}\,)\{ e^{-\gamma_+t}[(X_5-Z_7)\cos(2\Omega t)-(Y_5-Z_8)\sin(2\Omega t)] \\
&\qquad -e^{-\Gamma t/2}[X_5\cos(Dt)-Y_5\sin(Dt)]+Z_7e^{-\Gamma t}\}-(1-\sqrt{r}\,)[\Omega\to-\Omega]) , \\
S_6 &= -\tfrac{1}{2}(1-r)r(\gamma/\gamma_0)\{(X_6-Z_9)e^{-\gamma_0 t}-e^{\Gamma t/2}[X_6\cos(Dt)-Y_6\sin(Dt)]+Z_9e^{-\Gamma t}\} , \\
S_7 &= \frac{2\Gamma(1-r)}{(\Gamma/2+\gamma_0)^2+D^2}\left[\tfrac{1}{2}+\tfrac{1}{2}e^{-\Gamma t}+\frac{\gamma_0}{\Gamma}(1-e^{-\Gamma t})-e^{-(\Gamma/2+\nu_0)t}\cos(Dt)\right] , \\
S_8 &= r(1-r)(\gamma/\gamma_0)^2\frac{2\Gamma}{(\Gamma/2)^2+D^2}[\tfrac{1}{2}+\tfrac{1}{2}e^{-\Gamma t}-e^{-\Gamma t/2}\cos(Dt)] ,
\end{aligned}
\tag{17}
$$

where $[\Omega\to-\Omega]$ represents preceding terms with Ω interchanged with $-\Omega$. Here, coefficients X_i's are defined as follows:

$$
\begin{aligned}
X_1 &= 2\Gamma[(\Gamma/2)^2-\gamma_+^2+(4\Omega-D)^2]/\{[(\Gamma/2-\gamma_+)^2+(2\Omega-D)^2][(\Gamma/2+\gamma_+-\gamma_0)^2+(2\Omega-D)^2]\} , \\
X_2 &= 2\Gamma[(\Gamma/2)^2-\gamma_+^2+(2\Omega-D)^2]/\{[(\Gamma/2-\gamma_+)^2+(2\Omega-D)^2][(\Gamma/2+\gamma_+)^2+(2\Omega-D)^2]\} , \\
X_3 &= \Gamma[\Gamma(\Gamma/2-\gamma_+)-2D(2\Omega-D)]/(\{[(\Gamma/2)^2+D^2][(\Gamma/2-\gamma_+)^2+(2\Omega-D)^2]\}) , \\
X_4 &= \frac{2\Gamma[(\Gamma/2-\gamma_++\gamma_0)(\Gamma/2-\gamma_0)+D(2\Omega+D)]}{[(\Gamma/2-\gamma_++\gamma_0)^2+(2\Omega+D)^2][(\Gamma/2-\gamma_0)^2+D^2]} , \\
X_5 &= \frac{\Gamma[\Gamma(\Gamma/2-\gamma_+)+2D(D+2\Omega)]}{[(\Gamma/2)^2+D^2][(\Gamma/2-\gamma_+)^2+(D+2\Omega)^2]} , \\
X_6 &= \frac{\Gamma[(\Gamma/2-\gamma_0)+2D^2]}{[(\Gamma/2)^2+D^2][(\Gamma/2-\gamma_0)^2+D^2]} ,
\end{aligned}
\tag{18}
$$

and Y_i's are defined as follows:

$$
\begin{aligned}
Y_1 &= 4\Gamma(2\Omega\gamma_+-\gamma_+D)/\{[(\Gamma/2-\gamma_+)^2+(2\Omega-D)^2][(\Gamma/2+\gamma_+-\gamma_0)^2+(2\Omega-D)^2]\} , \\
Y_2 &= 2\Gamma[2\gamma_+(2\Omega-D)]/\{[(\Gamma/2-\gamma_+)^2+(2\Omega-D)^2][(\Gamma/2+\gamma_+)^2+(2\Omega-D)^2]\} , \\
Y_3 &= 2\Gamma[D(\Gamma/2-\gamma_+)+\Gamma/2(2\Omega-D)]/\{[(\Gamma/2)^2+D^2][(\Gamma/2-\gamma_+)^2+(2\Omega-D)^2]\} , \\
Y_4 &= 2\Gamma[2\Omega(\Gamma/2-\gamma_0)-D(\Gamma/2+\gamma_+-2\gamma_0)]/\{[(\Gamma/2-\gamma_++\gamma_0)^2+(2\Omega-D)^2][(\Gamma/2-\gamma_0)^2+D^2]\} , \\
Y_5 &= 2\Gamma(D\gamma_++\Gamma\Omega)/\{[(\Gamma/2)^2+D^2][(\Gamma/2-\gamma_+)^2+(D+2\Omega)^2]\} , \\
Y_6 &= 2\Gamma D\gamma_0/\{[(\Gamma/2)^2+D^2][(\Gamma/2-\gamma_0)^2+D^2]\} ,
\end{aligned}
\tag{19}
$$

and finally the Z_i's are defined as follows:

$$
\begin{aligned}
Z_1 &= \frac{2\Gamma(\Gamma/2-\gamma_+)}{(\Gamma-\gamma_0)}\Big/[(\Gamma/2-\gamma_+)^2+(2\Omega-D)^2], \quad Z_2=\frac{(\Gamma-2\gamma_+)}{(\Gamma/2-\gamma_+)^2+(2\Omega-D)^2} , \\
Z_3 &= 2\Gamma[(\Gamma/2-\gamma_+)^2-2\Omega(2\Omega-D)]/\{[(\Gamma/2-\gamma_+)^2+4\Omega^2][(\Gamma/2-\gamma_+)^2+(2\Omega-D)^2]\} , \\
Z_4 &= (4\Omega-D)(\Gamma/2-\gamma_+)/\{[(\Gamma/2-\gamma_+)^2+4\Omega^2][(\Gamma/2-\gamma_+)^2+(2\Omega-D)^2]\} , \\
Z_5 &= (\Gamma/2-\gamma_0)[(\Gamma-\gamma_+)+2\Omega D]/\{[(\Gamma/2-\gamma_0)^2+D^2][(\Gamma-\gamma_+)^2+4\Omega^2]\} , \\
Z_6 &= [2\Omega(\Gamma/2-\gamma_0)-D(\Gamma-\gamma_+)]/\{[(\Gamma/2-\gamma_0)^2+D^2][(\Gamma-\gamma_+)^2+4\Omega^2]\} , \\
Z_7 &= \Gamma[\Gamma(\Gamma-\gamma_+)+2\Omega D]/\{[(\Gamma/2)^2+D^2][(\Gamma-\gamma_+)^2+4\Omega^2]\} , \\
Z_8 &= 2\Gamma[D(\Gamma-\gamma_+)-\Gamma\Omega]/\{[(\Gamma/2)^2+D^2][(\Gamma-\gamma_+)^2+4\Omega^2]\} , \\
Z_9 &= \Gamma^2/\{(\Gamma-\gamma_0)[(\Gamma/2)^2+D^2]\} .
\end{aligned}
\tag{20}
$$

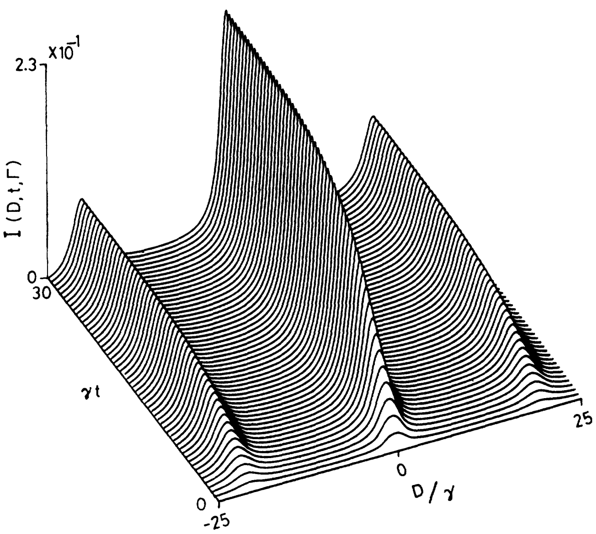

FIG. 1. Time development of physical spectrum (in arbitrary untis) of a two-level atom initially in its ground state and driving field in resonance with atomic transition, i.e., $\Delta/\gamma=0$ and $\Gamma/\gamma=0.1$, $2\Omega_0/\gamma=20$, and $N=0$, $|M|=0$ (unsqueezed vacuum).

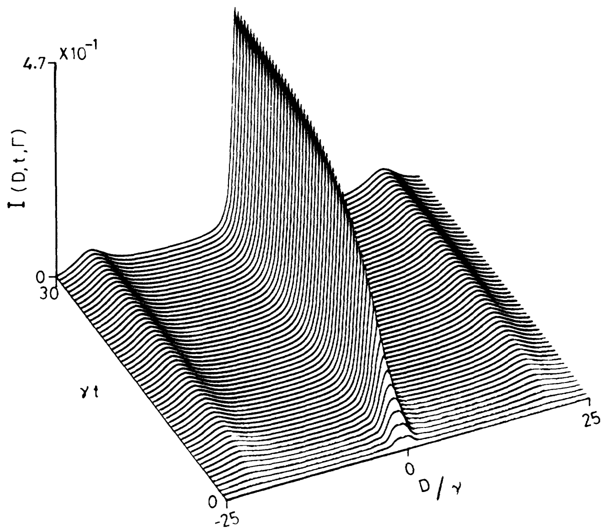

FIG. 3. The same as Fig. 2 but for $\phi=\pi$.

The expression (16) of the transient spectrum gets very much simplified under the exact resonance condition ($\Delta=0$). In Fig. 1 we have plotted the analytical time-dependent spectrum for the unsqueezed vacuum bath (for which $N=0$, $|M|=0$) for the atom initially in the ground state with $2\Omega_0/\gamma=20$ and $\Delta/\gamma=0$. We have also computed the spectrum with direct numerical integration of Eq. (5) and these are in excellent agreement with each other. In Figs. 2 and 3 we have plotted the transient spectrum with squeezed vacuum bath [with $|M|=\sqrt{N(N+1)}$] for the parameters $|M|=0.2$ and $\Delta/\gamma=0$, $2\Omega_0/\gamma=20$. In Fig. 2 we have kept $\psi=0$, $\Gamma=0.1$ and in Fig. 3 it is $\phi=\pi$, $\Gamma=0.1$. These spectra show some interesting behavior as compared to the spectra in the unsqueezed vacuum. We find that the central peak remains broadened (Fig. 2) or narrowed (Fig. 3) with respect to natural linewidth throughout its time development towards the steady state. However, it is interesting

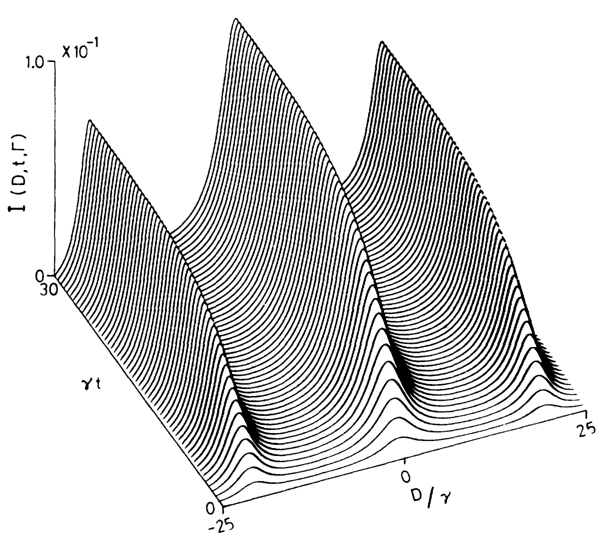

FIG. 2. The same as Fig. 1 but for squeezed vacuum $|M|=0.5$ and $\phi=0$.

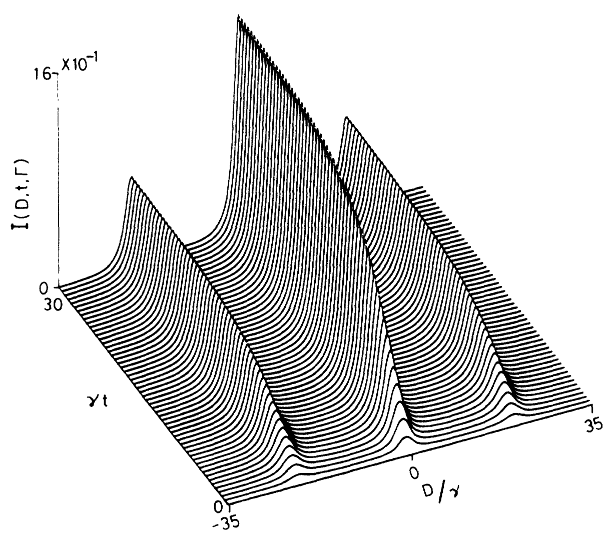

FIG. 4. Time development of physical spectrum (in arbitrary units) of a two-level atom initially in its ground state and driving field off resonance with atomic transition, i.e., $\Delta/\gamma=7$ and $\Gamma/\gamma=0.1$, $2\Omega_0/\gamma=20$, and $N=0$, $|M|=0$ (unsqueezed vacuum).

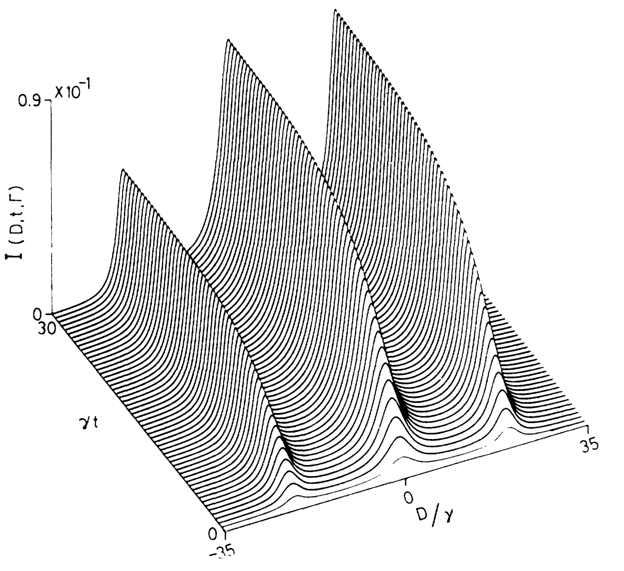

FIG. 5. The same as Fig. 4 but for squeezed vacuum $|M|=0.5$ and $\phi=0$.

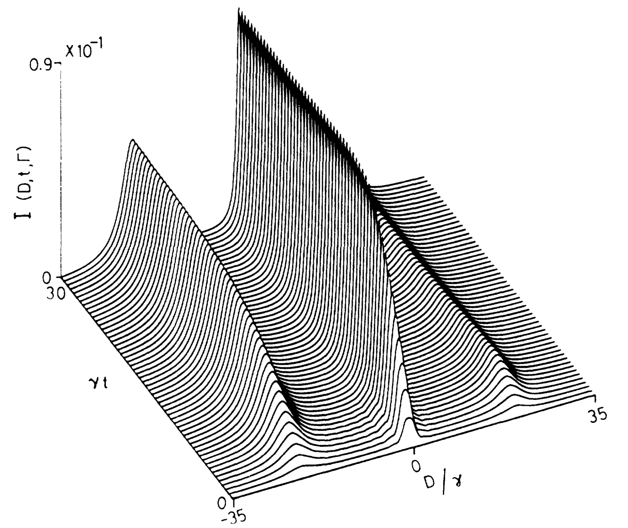

FIG. 6. The same as Fig. 5 but for $\phi=\pi$.

to note (in Fig. 3) a time lag between central peak and sidebands. In Figs. 4–6 we have introduced off-resonance conditions by fixing the detuning parameter $\Delta/\gamma=7$ (but all other conditions are the same as Figs. 1–3, respectively). In Fig. 4 we find, as expected for the unsqueezed bath, the symmetry of the spectrum is retained (i.e., peak heights of sidebands are equal) as it reaches steady state. But in the squeezed vacuum under the off-resonance condition of excitation there is asymmetry in the spectra both for $\phi=0$ and π (Figs. 5 and 6) throughout its time development towards the steady state. Now the peak heights of sidebands are unequal with respect to each other and one of the sideband peaks competes with the central peak. Nevertheless, it is interesting to note that the nature of the asymmetry in the two cases ($\phi=0$ and π) is opposite.

In conclusion the results presented here are valid for an atom embedded in a broadband-squeezed vacuum. In practice the light will not be perfectly δ correlated (say, in a four-wave-mixing process), but this analysis would be valid provided the input squeezing bandwidth is much larger than the fluorescence linewidth. However, the present results are expected to show further interesting modifications when finite bandwidth squeezed light is considered. For a colored squeezed vacuum, the atom not only interacts with the squeezed vacuum modes but also with the unsqueezed ones. A competition between these two interactions would considerably modify the nature of the steady state[25] as well as the time-dependent spectrum.

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