Effect of the dynamic Stark shift on dipole squeezing in two-photon processes

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Atomic-dipole squeezing in the presence of the Stark shift in two-photon Jaynes-Cummings model for an atom initially prepared in a coherent superposition of states, interacting with the coherent- or squeezed-vacuum-state field, is studied. Closed expressions for dispersive and absorptive components of the dipole for the coherent-field input under the large-$\bar{n}$ approximation are derived. It is shown that atomic-dipole squeezing is sensitive to both the phase of the initial field and the relative phase between the two atomic levels. Our results show that, for certain choices of these phases, permanent dipole squeezing can be obtained.

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I. INTRODUCTION

Squeezing of the radiation field due to its potential applications in optical communication [1], high-precision interferometry for gravitational-wave detection [2], and laser spectroscopy [3] has attracted a great deal of interest over the years. Although the Heisenberg uncertainty principle is inviolate, in squeezed states fluctuations in one of the quadratures are reduced below the standard quantum limit at the expense of increased fluctuations in the other quadrature. This idea of reduction of quantum fluctuations in a dynamic observable at the expense of enlarged fluctuations in its canonical conjugate variable can be extended to the atomic variables, termed atomic-dipole squeezing. Production of such states in resonance fluorescence has been predicted by Walls and Zoller [4]. Wódkiewicz [5], Wódkiewicz and Eberly [6], and Arvind [7] have considered squeezing in the elements of su(2) algebra. Recently, the Jaynes-Cummings model has been analyzed for dipole squeezing [8,9]. A relationship between the field squeezing and the dipole squeezing has also been established by some authors. The dipole squeezing is shown to produce reduced fluctuations in the radiated light [4,8,10].

The linear superposition principle, which is one of the most fundamental features of quantum mechanics, has led to a new class of squeezed states, called superposition squeezed states of field and atoms [8,11–13]. The idea of preparing atoms in the coherent superposition of states has become very popular, particularly due to its applications to the noise quenching by correlated emission [14], quantum beats [15], and noise-free amplifiers [16]. It has been shown that the superposition of atomic states has a dramatic effect on pure quantum-mechanical effects in atom-field interaction [17,18]. The phenomenon of coherent trapping for certain choices of the relative phase of the field and the atomic dipole is observed [18]. Phase-sensitive noise-free amplification has also been studied [16].

We devoted this paper to an investigation of the atomic-dipole squeezing in the two-photon Jaynes-Cummings model. We consider the interaction of both coherent and squeezed-vacuum field inputs with an atom initially prepared in coherent superposition of ground and excited states. To make the model closer to the experimental realization, we include the effect of the dynamic Stark shift. When the two atomic levels are coupled with comparable strength to the intermediate relay level, the Stark shift becomes significant and cannot be ignored [19-21]. We show that the atomic-dipole squeezing exhibits a phase sensitivity and undergoes dramatic changes with a variation in the phase of the field or the relative phase of the two atomic levels. The results for dipole squeezing incorporating Stark shift are radically different from the results obtained in the absence of Stark shift. Hence, the inclusion of Stark shift is very crucial for the validity of the effective Hamiltonian.

The present paper is organized as follows. In Sec. II we define the model and derive analytic expressions for atomic-dipole squeezing. In Sec. III numerical results are presented for both coherent and squeezed-vacuum input fields for various initial conditions.

II. TWO-PHOTON JAYNES-CUMMINGS MODEL IN THE PRESENCE OF STARK SHIFT

We consider on-resonance interaction of a single-mode field of frequency ω with an effective two-level atom of transition frequency ω_0 through two-photon transitions in a lossless cavity. The intermediate relay level is so located as to give rise to a significant Stark shift. The effective Hamiltonian describing such a system (with $\hbar=1$) is given by

$$H_{\text{eff}}=\omega a^\dagger a+\omega_0\sigma_3+a^\dagger a(\beta_2|a\rangle\langle a|+\beta_1|b\rangle\langle b|)+g(a^{\dagger 2}\sigma^- +a^2\sigma^+)\ , \tag{1}$$

where $\omega_0=2\omega$. a and $a^\dagger$ are the usual photon destruction and creation operators, and $\sigma_3=|a\rangle\langle a|-|b\rangle\langle b|$, $\sigma^+=|a\rangle\langle b|$, and $\sigma^-=|b\rangle\langle a|$ are the atomic flopping operators expressed in terms of the excited ($|a\rangle$) and ground ($|b\rangle$) states of the atom. β_1 and β_2 are the parameters describing the intensity-dependent Stark shifts of the two levels due to the virtual transitions to the intermediate relay level and g is the effective two-photon coupling constant.

The time evolution of a Heisenberg operator in terms of the well-known time-dependent unitary transformation operator $U(t)$ can be written as

$$\Theta(t)=U^{\dagger}(t)\Theta(0)U(t)\ , \tag{2}$$

where

$$U(t)=e^{-iH_{\text{eff}}t}\ . \tag{3}$$

By proceeding in the standard fashion we obtain the following expression for $U(t)$:

$$U(t)=\begin{bmatrix} \sin^2\theta_n e^{-iE_n^+t}+\cos^2\theta_n e^{-iE_n^-t} & \sin\theta_n\cos\theta_n(e^{-iE_n^+t}-e^{-iE_n^-t}) \\ \sin\theta_n\cos\theta_n(e^{-iE_n^+t}-e^{-iE_n^-t}) & \sin^2\theta_n e^{-iE_n^-t}+\cos^2\theta_n e^{-iE_n^+t} \end{bmatrix}\ , \tag{4}$$

where

$$\begin{aligned} &\sin\theta_n=(1/\sqrt{2})(1+\eta_n/\Omega_n)^{1/2}\ ,\\ &E_n^{\pm}=\omega(n+1)+\lambda_n^{\pm}\ ,\\ &\lambda_n^{\pm}=\left[g\frac{n(1+r^2)+2r^2}{2r}\pm\Omega_n\right]\ ,\\ &\Omega_n=[g^2(n+1)(n+2)+\eta_n^2]^{1/2}\ ,\\ &\eta_n=\frac{g}{2r}[n(1-r^2)-2r^2]\ , \end{aligned} \tag{5}$$

and

$$r=(\beta_1/\beta_2)^{1/2}\ . \tag{6}$$

We consider the situation in which the atom is initially in a coherent superposition of excited and ground states and that the initial field is an arbitrary superposition of Fock states, so that

$$\rho(0)=\sum_{n,m}C_{nm}[\cos(\theta/2)|a;n\rangle+\sin(\theta/2)e^{-i\phi}|b;n\rangle]\times[\cos(\theta/2)\langle n;a|+\sin(\theta/2)e^{i\phi}\langle n;b|]\ , \tag{7}$$

where $C_{nm}=C_nC_m^*$ and C_n are the coefficients in the Fock-state expansion of the initial field. θ is the degree of excitation and ϕ is the relative phase of the two atomic levels.

From Eqs. (2)–(7) the combined atom-field density matrix can be calculated in a straightforward manner, which when traced over the field variables gives rise to the following atomic density matrix:

$$\begin{aligned} \rho_A(t)=\sum_n\{&\cos^2(\theta/2)[C_{nn}|U_{11}^{(n)}|^2|a\rangle\langle a|+C_{n,n-2}U_{11}^{(n)}U_{21}^{(n-2)*}|a\rangle\langle b|+C_{n-2,n}U_{21}^{(n-2)}U_{11}^{(n)*}|b\rangle\langle a|\\ &+C_{n-2,n-2}|U_{21}^{(n-2)}|^2|b\rangle\langle b|]\\ &+\cos(\theta/2)\sin(\theta/2)e^{i\phi}[C_{n,n+2}U_{11}^{(n)}U_{12}^{(n)*}|a\rangle\langle a|+C_{nn}U_{11}^{(n)}U_{22}^{(n-2)*}|a\rangle\langle b|\\ &\qquad+C_{n-2,n+2}U_{21}^{(n-2)}U_{12}^{(n)*}|b\rangle\langle a|+C_{n-2,n}U_{21}^{(n-2)}U_{22}^{(n-2)*}|b\rangle\langle b|]\\ &+\cos(\theta/2)\sin(\theta/2)e^{-i\phi}[C_{n+2,n}U_{12}^{(n)}U_{11}^{(n)*}|a\rangle\langle a|+C_{n+2,n-2}U_{12}^{(n)}U_{21}^{(n-2)*}|a\rangle\langle b|\\ &\qquad+C_{nn}U_{22}^{(n-2)}U_{11}^{(n)*}|b\rangle\langle a|+C_{n,n-2}U_{22}^{(n-2)}U_{21}^{(n-2)*}|b\rangle\langle b]\\ &+\sin^2(\theta/2)[C_{n+2,n+2}|U_{12}^{(n)}|^2|a\rangle\langle a|+C_{n+2,n}U_{12}^{(n)}U_{22}^{(n-2)*}|a\rangle\langle b|\\ &\qquad+C_{n,n+2}U_{22}^{(n-2)}U_{12}^{(n)*}|b\rangle\langle a|+C_{n,n}|U_{22}^{(n-2)}|^2|b\rangle\langle b|]\}\ . \end{aligned} \tag{8}$$

The above result can be used to obtain expectation values of the atomic operators

$$\begin{aligned} \langle\sigma^-\rangle e^{2i\omega t}=\sum_n\,&[\cos^2(\theta/2)C_{n,n-2}U_{11}^{(n)}U_{21}^{(n-2)*}+\cos(\theta/2)\sin(\theta/2)e^{i\phi}C_{nn}U_{11}^{(n)}U_{22}^{(n-2)*}\\ &+\cos(\theta/2)\sin(\theta/2)e^{-i\phi}C_{n+2,n-2}U_{12}^{(n)}U_{21}^{(n-2)*}+\sin^2(\theta/2)C_{n+2,n}U_{12}^{(n)}U_{22}^{(n-2)*}]\ , \end{aligned} \tag{9}$$

$$\begin{aligned} \langle\sigma_3\rangle=2\sum_n\,&[\cos^2(\theta/2)C_{nn}|U_{11}^{(n)}|^2+\cos(\theta/2)\sin(\theta/2)e^{i\phi}C_{n,n+2}U_{11}^{(n)}U_{12}^{(n)*}\\ &+\cos(\theta/2)\sin(\theta/2)e^{-i\phi}C_{n+2,n}U_{12}^{(n)}U_{11}^{(n)*}+\sin^2(\theta/2)C_{n+2,n+2}|U_{12}^{(n)}|^2]\ . \end{aligned} \tag{10}$$

If the field is in a coherent state $|\alpha\rangle$, then

$$C_{nm}=\alpha^{n}\alpha^{*m}\exp(-|\alpha|^{2})/(n!m!)^{1/2}\ , \tag{11}$$

with $\alpha=|\alpha|e^{i\xi}$ and $|\alpha|^{2}=\bar{n}$, where ξ is the initial phase of the input field and $\bar{n}$ is the mean number of photons in the coherent state.

In general the sums in Eqs. (9) and (10) are computed numerically, however, these summations can be performed analytically for $\bar{n}\gg 1$. This is because for large $\bar{n}$, the dispersion in the Poissonian distribution of the coherent state in the $|n\rangle$ representations is much smaller than the mean number of photons $\bar{n}$. Hence, for $\bar{n}\gg 1$, only large n values contribute to the sum over n. Consequently, one can get closed-form expressions for (9) and (10) by expanding Ω_n and $\sin\theta_n$ in powers of n^{-1}. Retaining terms of order n^{0}

$$\begin{aligned}&\Omega_{n-2}=n\psi-\chi,\quad \chi=g/r(1+r^{2})\ ,\\&\psi=g(1+r^{2})/2r,\quad \sin(2\theta_{n-2})\simeq 2r/(1+r^{2})\ .\end{aligned} \tag{12}$$

According to Ref. [20] the corresponding results in the absence of the Stark shift for the large-n approximation follow by setting $r=1$ in Eq. (12).

By substituting Eq. (11) and relations given by (12) in Eqs. (9) and (10) we get

$$\begin{aligned}\langle\sigma^{-}\rangle e^{2i\omega t}=e^{-2i\psi t}\Bigg\{&\frac{r\cos^{2}(\theta/2)e^{2i\xi}}{(1+r^{2})|\alpha|^{2}}\Bigg[-i\sin2\psi t(\bar{n}-\tfrac{1}{2})+ie^{-\delta(t)}[\bar{n}\sin\epsilon(t)-\tfrac{1}{2}\sin\alpha(t)]\\&\qquad+\frac{1-r^{2}}{1+r^{2}}\{-\cos2\psi t(\bar{n}-\tfrac{1}{2})+e^{-\delta(t)}[\bar{n}\cos\epsilon(t)-\tfrac{1}{2}\cos\alpha(t)]\}\Bigg]\\&+\cos(\theta/2)\sin(\theta/2)e^{i\phi}\left[\frac{2r^{2}}{(1+r^{2})^{2}}\cos2\psi t+e^{-\delta(t)}\left[\frac{1+r^{4}}{(1+r^{2})^{2}}\cos\alpha(t)+i\frac{1-r^{4}}{(1+r^{2})^{2}}\sin\alpha(t)\right]\right]\\&+\frac{r^{2}\cos(\theta/2)\sin(\theta/2)}{(1+r^{2})^{2}|\alpha|^{4}}e^{i(4\xi-\phi)}\{2\cos2\psi t[\bar{n}^{2}-2\bar{n}+(\bar{n}-1)\bar{e}^{\bar{n}}+1]\\&\qquad-e^{-\delta(t)}[2\bar{n}^{2}\cos\alpha(t)-4\bar{n}\cos\beta(t)+2\cos\gamma(t)]\\&\qquad+2[\bar{n}\cos2\chi t-\cos2(\psi+\chi)t]e^{-\bar{n}}\}\\&+\frac{\sin^{2}(\theta/2)e^{2i\xi}}{|\alpha|^{2}}\Bigg[\frac{ir}{2(1+r)^{2}}\{-2\sin2\psi t[\bar{n}-\tfrac{1}{2}-\tfrac{1}{2}(\bar{n}-1)e^{-\bar{n}}]\\&\qquad+e^{-\delta(t)}[-2\bar{n}\sin\beta(t)-\sin\gamma(t)]\\&\qquad+[\sin2(\psi+\chi)t-\bar{n}\sin2\chi t]e^{-\bar{n}}\}\\&\qquad+\frac{r(1-r^{2})}{2(1+r^{2})^{2}}(2\cos2\psi t\{\bar{n}-\tfrac{1}{2}[1+(\bar{n}-1)e^{-\bar{n}}]\}\\&\qquad-e^{-\delta(t)}[2\bar{n}\cos\beta(t)-\cos\gamma(t)]\\&\qquad+[\bar{n}\cos2\chi t-\cos2(\psi+\chi)t]e^{-\bar{n}})\Bigg]\Bigg\}\ ,\end{aligned} \tag{13a}$$

$$\begin{aligned}\langle\sigma_{3}\rangle=2\Bigg[&\cos^{2}(\theta/2)\left[\frac{1+r^{4}}{(1+r^{2})^{2}}+\frac{2r^{2}}{(1+r^{2})^{2}}e^{-\delta(t)}\cos\alpha(t)\right]\\&+\sin^{2}(\theta/2)\frac{2r^{2}}{(1+r^{2})^{2}}\{-[1+\bar{n}-\cos2\chi t-\bar{n}\cos2(\psi-\chi)t]e^{-\bar{n}}+1-\bar{e}^{\delta(t)}\cos\beta(t)\}\\&+2\frac{\cos(\theta/2)\sin(\theta/2)}{|\alpha|^{2}}\Bigg[\frac{r(1-r^{2})}{(1+r^{2})^{2}}\cos(2\xi-\phi)(\bar{n}-\tfrac{1}{2}-\tfrac{1}{2}(\bar{n}-1)e^{-\bar{n}}\\&\qquad-\{\tfrac{1}{2}[-\bar{n}\cos2(\psi-\chi)t+\cos2\chi t]e^{-\bar{n}}\\&\qquad+e^{-\delta(t)}[\bar{n}\cos\alpha(t)-\tfrac{1}{2}\cos\beta(t)]\})\\&\qquad+\frac{r}{1+r^{2}}\sin(2\xi-\phi)(-\tfrac{1}{2}[\bar{n}\sin2(\psi-\chi)t+\sin2\chi t)]e^{-\bar{n}}\\&\qquad-e^{-\delta(t)}[\tfrac{1}{2}\sin\beta(t)-\bar{n}\sin\alpha(t)])\Bigg]\Bigg]-1\ ,\end{aligned} \tag{13b}$$

and

$$\alpha(t)=\bar{n}\sin 2\psi t+2(\psi-\chi)t\ , \tag{14a}$$

$$\beta(t)=\bar{n}\sin 2\psi t-2\chi t\ , \tag{14b}$$

$$\gamma(t)=\bar{n}\sin 2\psi t-2(\psi+\chi)t\ , \tag{14c}$$

$$\epsilon(t)=\bar{n}\sin 2\psi t+2(2\psi-\chi)t\ , \tag{14d}$$

$$\delta(t)=2\bar{n}\sin^2\psi t\ . \tag{14e}$$

The third term in both Eqs. (13a) and (13b) is the "interference" term and depends on the relative phase between the two atomic levels and the coherent state. In the case of $\theta=\pi$, i.e., when the atom is initially in the ground state, Eqs. (13) reduce to the results obtained in Ref. [9]. The corresponding expressions in the absence of the Stark shift are obtained by setting $\exp(2i\psi t)=1$ in Eq. (13a) and letting $r=1$ in the rest of the expression in Eqs. (13).

The dispersive and absorptive components of the slowly varying atomic-dipole operator can be written as [9]

$$\begin{aligned}\sigma_1&=\tfrac{1}{2}(\sigma^+e^{-i\omega_0 t}+\sigma^-e^{-i\omega_0 t})\ ,\\ \sigma_2&=\frac{1}{2i}(\sigma^+e^{-i\omega_0 t}-\sigma^-e^{-i\omega_0 t})\ .\end{aligned} \tag{15}$$

They obey the commutation relation

$$[\sigma_1,\sigma_2]=\frac{i}{2}\sigma_3\ . \tag{16}$$

The corresponding Heisenberg uncertainty relation is

$$(\Delta\sigma_1)^2(\Delta\sigma_2)^2\geq\tfrac{1}{16}\langle\sigma_3\rangle^2\ . \tag{17}$$

The atomic dipole is said to be squeezed if the variances in σ_1 and σ_2 satisfy the condition

$$\Delta\sigma_{1,2}^2\langle\tfrac{1}{4}|\langle\sigma_3\rangle|\ . \tag{18}$$

The above condition for squeezing in the dispersive and absorptive components may be written as

$$S_{1,2}<1\ , \tag{19}$$

with the corresponding squeezing parameters defined as

$$\begin{aligned}S_1&=\frac{1-4(\mathrm{Re}\langle\sigma^-\rangle e^{i\omega t})^2}{|\langle\sigma_3\rangle|}\ ,\\ S_2&=\frac{1-4(\mathrm{Im}\langle\sigma^-\rangle e^{i\omega t})^2}{|\langle\sigma_3\rangle|}\ .\end{aligned} \tag{20}$$

Substituting Eqs. (12)–(14) in Eq. (20) the closed expressions for S_1 and S_2 are obtained.

III. RESULTS AND DISCUSSION

A. Coherent-state input

In this section we present the results for $S_1(t)$ and $S_2(t)$ for $\theta=2\pi/3,\xi=0$ and for various choices of ϕ and the Stark-shift parameter r. Figures (1) and (2) show the time evolution of S_1 and S_2 for $\phi=0$ and for $r=0$ and 0.5. It is evident from the figures that only one quadratures squeezes in the absence of the Stark shift while both S_1 and S_2 squeeze alternatively in its presence. We also see that for the above-mentioned case the duration of squeezing reduces in the presence of the Stark shift. Figures (3) and (4) show how the dipole squeezing is affected by a variation in the relative phase of the two atomic levels ϕ from 0 to π. It is interesting to note that the dipole squeezing shows a weak phase dependence in the absence of Stark shift, but a strong phase dependence in its presence. This strong phase dependence in the presence of the Stark shift is exhibited in terms of permanent dipole squeezing as shown in Fig. (4). We have also studied the effect of varying the phase of the input field ξ on the dipole squeezing, which is seen to have almost the same effect as the relative phase of the two atomic levels. Here once again, we observe permanent dipole squeezing in the presence of the Stark shift.

We have also studied the effect of larger exciting

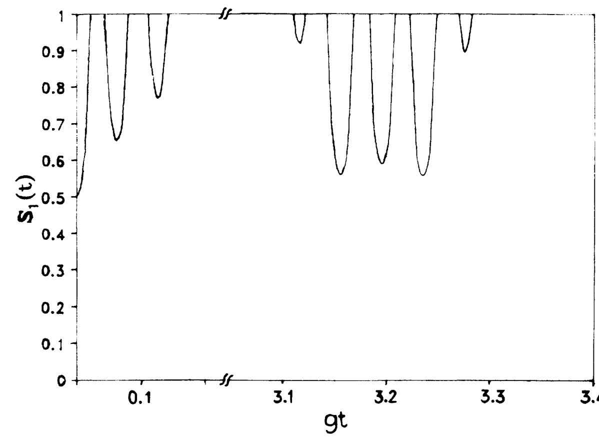

FIG. 1. Time evolution of the squeezing parameter $S_1(t)$ for coherent input in the absence of Stark shift for $\bar{n}=50$, $\theta=2\pi/3$, and $\phi=\xi=0$.

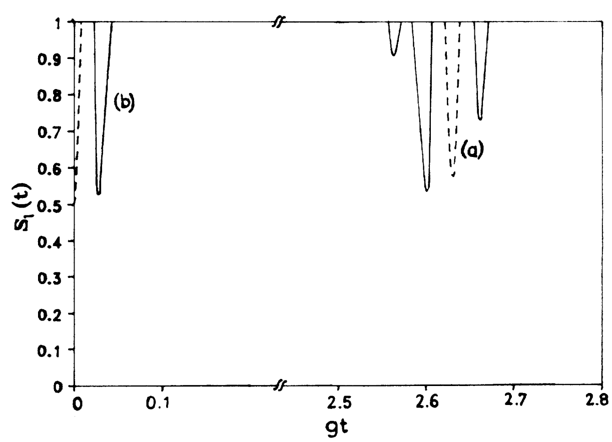

FIG. 2. Time evolution of the squeezing parameter for coherent input in the presence of Stark shift for $r=0.5$, $\bar{n}=50$, $\theta=2\pi/3$, and $\phi=\xi=0$. (a) $S_1(t)$ and (b) $S_2(t)$.

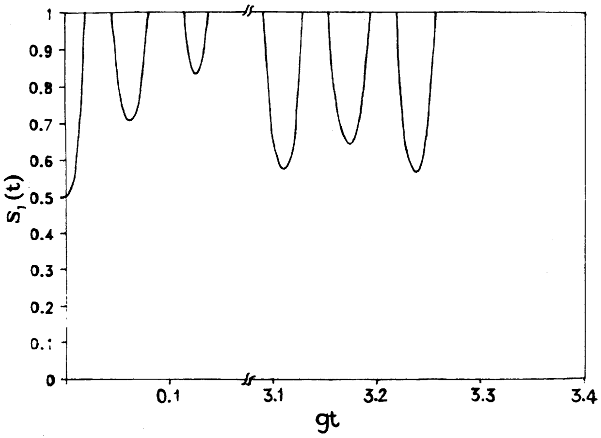

FIG. 3. Same as in Fig. 1, but for $\phi=\pi$.

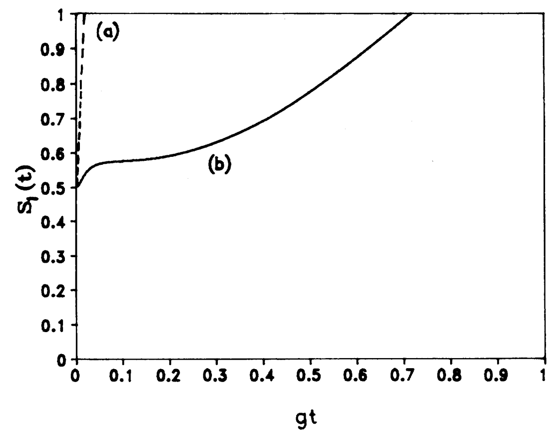

FIG. 5. $S_1(t)$ for the squeezed-vacuum input in the absence of Stark shift for $\bar{n}=10$, $\theta=2\pi/3$, and $\xi=0$. (a) $\phi=0$ and (b) π.

strength or large $\bar{n}$ on the dipole squeezing, which results in stronger squeezing and larger duration both in the presence and absence of the Stark shift.

B. Squeezed-vacuum-state input

The probability amplitudes for the squeezed vacuum are given by

$$C_{2n}=(-1)(\mu)^{-1/2}\frac{\sqrt{(2n)!}}{n!}(e^{i\xi}\nu/2\mu)^n$$

and

$$C_{2n+1}=0,\quad n=0,1,2,3,\ldots\,. \tag{21}$$

Since the field is super-Poissonian, large-$\bar{n}$ approximation does not hold for this case; instead, we work directly with the expressions (9) and (10) and make use of Eq. (21).

The results for dipole squeezing are obtained by numerically computing the quantities in Eq. (20). Some of the results calculated from (20) for $\theta=2\pi/3$, $\xi=0$ and for various values of ϕ and Stark-shift parameter r are presented in Figs. 5 and 6. It is interesting to note that, irrespective of the presence of the Stark shift, both S_1 and S_2 do not squeeze when the atom is initially in the excited or ground state. Thus an atom initially in an incoherent state remains in it throughout its interaction with the squeezed-vacuum field. However, when the atom is initially prepared in the coherent superposition of states $(\theta=2\pi/2)$ one of the quadratures squeezes only for a small duration. Hence, initial coherence of the atomic states is destroyed during its interaction with the squeezed-vacuum input field. The time evolution of S_1 and S_2 for $\theta=2\pi/3,\xi=0$ and for two different values of ϕ is plotted in Figs. (5) and (6). An increase in ϕ from 0 to π results in increased duration of squeezing in the absence of the Stark shift while an opposite effect is seen in its presence. The phase of the input field ξ is seen to have the same effect on the dipole as the relative phase of the two atomic levels. The larger exciting strength has a desqueezing effect and with $\bar{n}=50$ and for certain choices of ϕ and ξ, squeezing is completely destroyed.

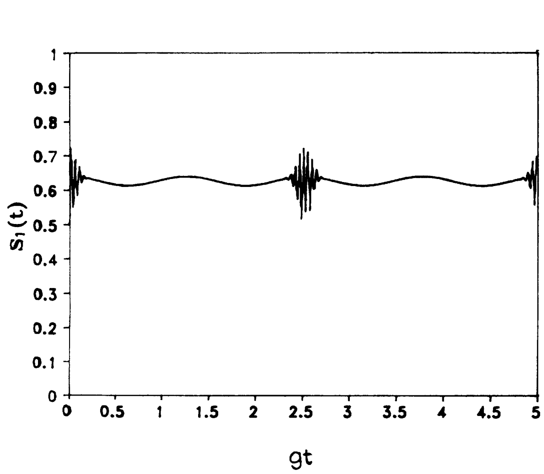

FIG. 4. $S_1(t)$ for the coherent input in the presence of Stark shift with all the parameters the same as in Fig. 2, but for $\phi=\pi$.

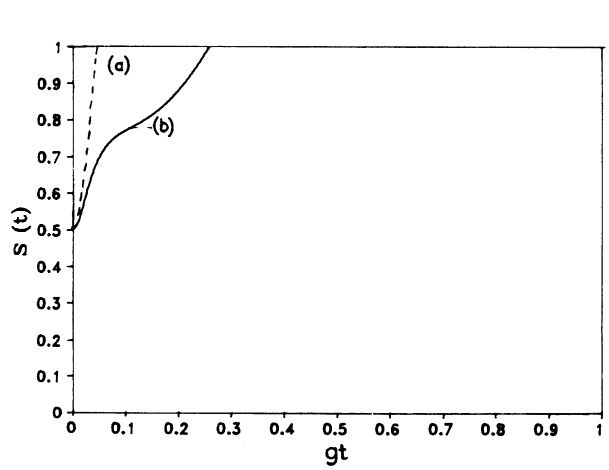

FIG. 6. $S_1(t)$ for the squeezed-vacuum input in the presence of Stark shift for $r=0.5$, $\bar{n}=10$, $\theta=2\pi/3$, and $\xi=0$. (a) $\phi=0$ and (b) π.

IV. CONCLUSION

It can be seen that the Stark shift plays an important role in the dynamics of the atom-field interaction. In particular, when the atom is considered initially in a coherent superposition of states, the duration of atomic-dipole squeezing is greatly enhanced for particular choices of the phase. The most significant manifestation of the Stark shift while taking the atom in coherent superposition of states is that for coherent field input and for a certain phase, permanent dipole squeezing is exhibited.

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