Photon-counting statistics of resonance fluorescence in a squeezed vacuum

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We discuss the consequences of the interaction of a broadband squeezed vacuum on the photon statistics of a coherently driven two-level atom. The statistics is described in terms of the Mandel function $Q(T)$ and the probability $P(n,T)$ of counting n photons in a time interval T. The phase-dependent decay in the squeezed vacuum is highlighted with these functions.

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I. INTRODUCTION

Resonance fluorescence from a two- and three-level atom placed in a squeezed vacuum has received considerable attention in recent years [1–12]. A particular case of interest here is a broadband squeezed vacuum, which essentially describes a time-stationary squeezed broadband field with Gaussian statistics [1–3]. In the case of a two-level atom interacting with a broadband squeezed vacuum, the component of the atomic polarization that is in phase with the low-noise quadrature phase of the field experiences reduced fluctuations [2] and therefore decays more slowly when compared with spontaneous emission in a normal vacuum. The other component of the atomic polarization, on the other hand, experiences increased fluctuations. Consequently, the fluorescence spectrum of a coherently driven two-level atom in a broadband squeezed vacuum exhibits a marked departure from the usual Mollow triplet in the normal vacuum [3]. It is dependent on the relative phase between the driving field and the squeezed vacuum. This opens possibilities of obtaining subnatural linewidths [2–8].

The phase-dependent decay of the atomic dipole in squeezed vacuum is also manifested in the statistics of photon emission [9,10,12]. This aspect of resonance fluorescence in a broadband squeezed vacuum was discussed recently in terms of the second-order intensity correlation function [9] $f(\tau)$, Mandel's response function [9,10] $Q(\tau)$, and the waiting-time distribution [12] $w(\tau)$. Here the functions $f(\tau)$ and $w(\tau)$ describe, respectively, the probabilities [13] of emission of any photon and the very next photon at time $t+\tau$ following the emission of a photon at time t. In a broadband squeezed vacuum the functions $f(\tau)$ and $w(\tau)$ evolve on a time scale slower than the spontaneous-emission time scale and exhibit a marked increase in the probability of detection of a subsequent photon for delay times much greater than the natural lifetime of the excited level [9,12]. This behavior is a consequence of the inhibition of fluorescence in a squeezed vacuum. The Mandel response function $Q(\tau)$, on the other hand, has been used to infer the sub-Poissonian statistics of the photon emission in a broadband squeezed vacuum [9,10].

In this paper we study the probability $P(n,T)$ that n photons are emitted in a given time interval T in steady state by a coherently driven two-level atom when placed in a broadband squeezed vacuum. Mandel [14] has calculated the distribution function $P(n,T)$ in a normal unsqueezed vacuum to highlight the antibunching and sub-Poissonian nature of the resonance fluorescence. Here we examine the consequences of phase-dependent decay in a squeezed vacuum in terms of the distribution $P(n,T)$. We also study the Mandel function $Q(T)$ to describe the statistics of photon emission. In Sec. II we define various quantities that are required to describe the statistics of photon emission. The problem of resonance fluorescence in a broadband squeezed vacuum is dealt with in Sec. III. Analytical results in the intense-field limit are discussed in Sec. IV. Finally, some important conclusions from this study are summarized in Sec. V.

II. BASIC DEFINITIONS

The probability that n photons are detected in some finite time interval T when fluorescence light falls on a photodetector is given by [14,15]

$$P(n,T)=\left\langle T_N{:}\frac{U^n e^{-U}}{n!}{:}\right\rangle , \tag{2.1}$$

where T_N represents the normal-ordering time-ordering operator and U is the integrated intensity

$$U=q\int_0^T I(\tau)d\tau . \tag{2.2}$$

Here $I(t)$ is the instantaneous intensity of fluorescence and q is a measure of the quantum efficiency of the detector. We now introduce the first-order generating function $G(\lambda,T)$ such that [15]

$$\begin{aligned} G(\lambda,T)&=\sum_{n=0}^{\infty}(1-\lambda)^n P(n,T) \\ &=\langle T_N{:}e^{-\lambda U}{:}\rangle=\sum_{k=0}^{\infty}\frac{(-\lambda)^k}{k!}\langle T_N{:}U^k{:}\rangle . \end{aligned} \tag{2.3}$$

It then follows that

$$P(n,T)=\frac{(-1)^n}{n!}\frac{\partial^n}{\partial\lambda^n}G(\lambda,T)|_{\lambda=1} . \quad (2.4)$$

Note here that the quantity $\langle T_N{:}U^k{:}\rangle$ is kth factorial moment of n. Mandel's response function $Q(T)$ is the normalized second-order factorial moment [14]

$$Q(T)=\frac{2I(\infty)}{T}\int_0^T dt_2\int_0^{t_2} dt_1 f(t_1)-I(\infty)T , \quad (2.5)$$

where $I(\infty)$ is the steady-state fluorescence intensity and $f(\tau)$ is the normalized second-order intensity correlation function. The parameter Q is a natural measure of the departure of the variance of the photon number n from the variance of a Poisson distribution. The evaluation of Eq. (2.3) requires knowledge of the intensity correlations of arbitrary higher order. Using the factorization property of the higher-order intensity correlations [14,16].

$$\left\langle T_N{:}I(t)I(t+\tau_1)\cdots I\left[t+\sum_{i=1}^{k-1}\tau_i\right]{:}\right\rangle =k![I(\infty)]^k\prod_{i=1}^{k-1}f(\tau_i) , \quad (2.6)$$

we may rewrite Eq. (2.3) as

$$G(\lambda,T)=1-\lambda qI(\infty)T+\sum_{k=2}^{\infty}[-\lambda qI(\infty)]^k\times\int_0^T dt_k\cdots\int_0^T dt_1\prod_{i=2}^{k}f(t_i-t_{i-1}) . \quad (2.7)$$

Now defining $\tilde{G}(\lambda,Z)$ to be the Laplace transform of $G(\lambda,T)$, it follows that [17]

$$\tilde{G}(\lambda,Z)=\frac{1}{Z}-\frac{\lambda qI(\infty)}{Z^2[1+\lambda qI(\infty)\tilde{f}(Z)]} . \quad (2.8)$$

The Laplace-transformed distribution function $P(n,Z)$ then reads as [17]

$$\tilde{P}(n,Z)=\frac{1}{Z}-\frac{qI(\infty)}{Z^2[1+\tilde{C}(Z)]}, \quad n=0$$
$$\tilde{P}(n,Z)=\frac{qI(\infty)[q\tilde{C}(Z)]^{n-1}}{Z^2[1+q\tilde{C}(Z)]^{n+1}}, \quad n=1,2,3,\ldots \quad (2.9)$$

where $\tilde{C}(Z)$ is the Laplace transform of

$$C(\tau)=I(\infty)f(\tau) . \quad (2.10)$$

We mention here that $C(\tau)d\tau$ is the conditional probability of an atom emitting any photon between τ and $\tau+d\tau$ after it has emitted one at time $\tau=0$, where $d\tau$ is an infinitesimal time delay [13]. The counting distribution $P(n,T)$ then can be obtained by taking the inverse Laplace transform of Eq. (2.9). In the analysis that follows, we assume that the detector is an ideal one (with quantum efficiency of unity).

III. TWO-LEVEL ATOM IN SQUEEZED VACUUM

Consider a two-level atom of transition frequency ω_0 between an excited level $|1\rangle$ and ground level $|2\rangle$ interacting with a classical driving field of frequency ω_L. The master equation for the reduced atomic density operator $\rho(t)$ describing the interaction of this atom with a broadband squeezed vacuum has the form [2,3]

$$\frac{d\rho}{dt}=-i[H_0,\rho]-\gamma(N+1)(A_{11}\rho+\rho A_{11}-2A_{21}\rho A_{12})-\gamma N(A_{22}\rho+\rho A_{22}-2A_{12}\rho A_{21}) -2\gamma|M|\exp(-i\phi)A_{12}\rho A_{12}-2\gamma|M|\exp(i\phi)A_{21}\rho A_{21} , \quad (3.1)$$

where

$$H_0=\alpha(A_{12}+A_{21})+\Delta A_{11} . \quad (3.2)$$

Here A_{ij} is the atomic operator $|i\rangle\langle j|$, 2α is the laser Rabi frequency, 2γ is the Einstein coefficient for spontaneous emission in the usual unsqueezed vacuum, and $\Delta=(\omega_0-\omega_L)$ is the frequency detuning of the laser from the atomic resonance frequency. The quantity $\phi=2\phi_L-\phi_s$, where ϕ_L and ϕ_s are, respectively, the phases for the laser field and the squeezed vacuum. The parameters $N,M=|M|\exp(-i\phi_s)$ characterize the squeezed vacuum, with N being proportional to the number of photons in squeezed vacuum. Also $|M|^2\leq N(N+1)$, where the equality sign holds for a minimum uncertainty squeezed state. In arriving at Eq. (3.1), it is assumed that the bandwidth of the squeezing is sufficiently broad so that the squeezed vacuum appears to the atom as a δ-correlated white noise [1–3]. We further assume that there is a window for observation such that only fluorescence photons can be detected by the photodetector.

The normalized second-order intensity correlation function $f(\tau)$ may be expressed in terms of the atomic operators A_{ij} as

$$f(\tau)=\lim_{t\to\infty}\frac{\langle A_{12}(t)A_{12}(t+\tau)A_{21}(t+\tau)A_{21}(t)\rangle}{\langle A_{11}(t)\rangle^2} . \quad (3.3)$$

Similarly, the steady-state intensity is given by

$$I(\infty)=2\gamma\langle A_{11}(\infty)\rangle . \quad (3.4)$$

Using the master equation (3.1) and the definitions of Eqs. (2.10), (3.3), and (3.4), it is straightforward to obtain an expression for $\tilde{C}(Z)$ as

$$\tilde{C}(Z)=2\gamma\frac{N(Z)}{D_0(Z)}\ , \tag{3.5}$$

where

$$\begin{aligned}N(Z)&=2\gamma NZ^2+[2\gamma N\gamma_{us}+2\Gamma^2(1-r^2)]Z\\&\quad+2\gamma N\gamma_u\gamma_s+4\Gamma^2(\gamma_0-r^2\gamma)\ ,\\D_0(Z)&=Z^3+2\gamma_{us}Z^2+(\gamma_u\gamma_s+\gamma_{us}^2+4\Gamma^2)Z\\&\quad+\gamma_u\gamma_s\gamma_{us}+8\gamma_0\Gamma^2\ .\end{aligned} \tag{3.6}$$

In writing Eq. (3.6), we used the following definitions [3]:

$$\begin{aligned}\gamma_u&=2\gamma(N+|M|+\tfrac{1}{2})\ ,\\\gamma_s&=2\gamma(N-|M|+\tfrac{1}{2})\ ,\\\gamma_{us}&=\gamma_u+\gamma_s\ ,\end{aligned} \tag{3.7}$$

and

$$\gamma_0=\gamma[(N+\tfrac{1}{2})(1+r^2)+|M|(1-r^2)\cos\phi]\ . \tag{3.8}$$

The generalized Rabi frequency Γ and the parameter r have the following meaning:

$$\begin{aligned}\Gamma&=(\alpha^2+\tfrac{1}{4}\Delta^2)^{1/2}\ ,\\r&=\frac{\Delta}{2\Gamma}.\end{aligned} \tag{3.9}$$

The steady-state intensity is then simply given by

$$I(\infty)=C(\infty)=\lim_{Z\to 0}Z\tilde{C}(Z)\ . \tag{3.10}$$

Now substituting Eq. (3.5) into (2.9) we obtain

$$P(n,Z)=\begin{cases}\dfrac{-I(\infty)D_0(Z)}{ZD_1(Z)} & \text{for } n=0\\[2ex]\dfrac{(2\gamma)^{n-1}I(\infty)[N(Z)]^{n-1}[D_0(Z)]^2}{[D_1(Z)]^{n+1}} & \text{for } n>0\end{cases} \tag{3.11}$$

where

$$\begin{aligned}D_1(Z)&=Z^4+2\gamma_{us}Z^3+(\gamma_u\gamma_s+\gamma_{us}^2+4\gamma^2N+4\Gamma^2)Z^2\\&\quad+[\gamma_u\gamma_s\gamma_{us}+4\gamma^2N\gamma_{us}\\&\quad\quad+8\gamma_0\Gamma^2+4\gamma\Gamma^2(1-r^2)]Z\\&\quad+4\gamma^2N\gamma_u\gamma_s+8\gamma\Gamma^2(\gamma_0-r^2\gamma)\ .\end{aligned} \tag{3.12}$$

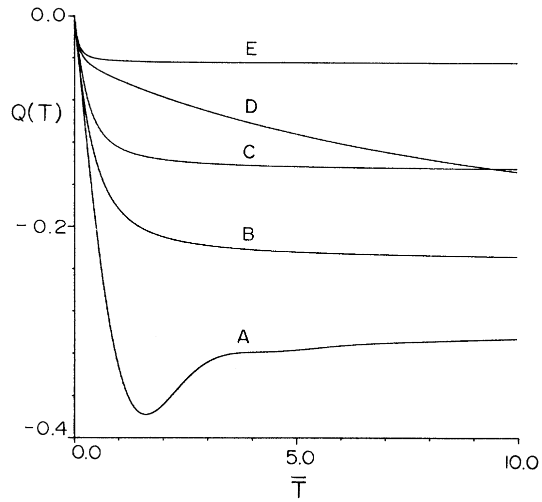

FIG. 1. Plot of $Q(T)$ vs $\bar{T}=2\gamma T$ for $\alpha/2\gamma=1$ in a normal vacuum (curve A) and in a squeezed vacuum (curves B–E). The parameters (N,ϕ) for B–E are (2,0), $(2,\pi)$, (10,0), and $(10,\pi)$, respectively.

The expression for $P(n,T)$ then depends on the roots of the biquadratic equation (3.12). These roots are generally difficult to obtain. In such cases we resort to the numerical evaluation of Eq. (3.11).

Before we discuss the counting distribution $P(n,T)$ in a squeezed vacuum, it is instructive to study Mandel's function $Q(T)$. For Poisson statistics, $Q=0$, whereas $Q<0$ and $Q>0$ describe, respectively, the sub- and super-Poissonian statistics [14]. Analytical expression for $Q(T)$ is generally difficult to obtain. We display in Fig. 1 the behavior of $Q(T)$ in a squeezed vacuum compared with that in normal vacuum for $\alpha/2\gamma=1$. This figure is plotted by evaluating $Q(T)$ numerically. It is clear from this figure that the squeezed vacuum tends to smoothen the behavior of $Q(T)$ and that this behavior is phase dependent. Specifically, $Q(T)$ is more negative for $\phi=0$ as compared with the $\phi=\pi$ case. When squeezing increases, the sub-Poissonian nature of the emission decreases.

The function $Q(T)$ for large counting times ($T\to\infty$) can be obtained analytically. Using Eqs. (2.5), (2.10), and (3.5) it is straightforward to show that

$$\begin{aligned}Q(\infty)=-\frac{8\gamma}{(\gamma_u\gamma_s\gamma_{us}+8\gamma_0\Gamma^2)^2}\{&8\Gamma^4r^2(\gamma_0-\gamma)+\gamma N\gamma_u^2\gamma_s^2\\&+\Gamma^2[2(\gamma_0-r^2\gamma)(\gamma_u\gamma_s\gamma_{us}^2)-4\gamma N(2\gamma_{us}\gamma_0-\gamma_u\gamma_s)-(1-r^2)\gamma_u\gamma_s\gamma_{us}]\}\ .\end{aligned} \tag{3.13}$$

In order to obtain the essential features of $Q(\infty)$, we consider the two-level atom driven at exact resonance ($r=0$). Furthermore, we use the notation Q_{s0} and $Q_{s\pi}$ to describe the function Q in a squeezed vacuum for $\phi=0$ and π, respectively. Similarly, Q corresponding to the unsqueezed normal vacuum will be denoted by Q_u. In the normal vacuum Eq. (3.13) reduces to the familiar one given by Mandel [14]

$$Q_u(\infty)=-\frac{6(\alpha/\gamma)^2}{[2(\alpha/\gamma)^2+1]^2}\ , \tag{3.14}$$

which has a minimum value of $-\frac{3}{4}$ when $\alpha=\gamma/\sqrt{2}$. However, under off-resonance conditions the absolute value of $Q_u(\infty)$ can be larger than $\frac{3}{4}$. In a squeezed vacuum at exact resonance, $2\gamma_0=\gamma_{u,s}$ for $\phi=0,\pi$ and Eq. (3.13) yields

$$Q_s(\infty)=-\frac{2\Gamma^2(\gamma_x+2\gamma)+2\gamma N\gamma_x^2}{(\gamma_x\gamma_{us}+4\Gamma^2)^2}\ , \tag{3.15}$$

where $\gamma_x=\gamma_{s,u}$ for $\phi=0$ and π, respectively. When squeezing is high ($N\gg1$) and the squeezed vacuum is in a minimum uncertainty state [$M^2=N(N+1)$],

$$\begin{aligned}Q_{s\pi}(\infty)&=-\frac{1}{2N}\frac{1}{[(\alpha/4\gamma N)^2+1]}\ ,\\ Q_{s0}(\infty)&=-4\frac{(2\alpha/\gamma)^2+1/8N}{[(2\alpha/\gamma)^2+1]^2}\ .\end{aligned} \tag{3.16}$$

It is clear that $Q_{s\pi}(\infty)$ increases monotonically to zero as the squeezing parameter N is increased. On the other hand, $Q_{s0}(\infty)$ is weakly dependent on N for reasonable values of the field strength α. For $\alpha\to0$, $Q_{s\pi}=Q_{s0}$, as can be seen from Eq. (3.16). Furthermore, $Q_{s0}(\infty)$ shows an absolute maximum value of 1 for $\alpha\cong\gamma/2$ when the vacuum is highly squeezed. This indicates that the emission in this case is more sub-Poissonian than that in the normal vacuum.

Another interesting case is when the driving field is strong. In that case

$$Q(\infty)=-\frac{r^2\gamma(\gamma_0-\gamma)}{\gamma_0^2}\ , \tag{3.17}$$

which depends critically on the detuning parameter r. In the normal vacuum, Eq. (3.17) shows that $Q_u(\infty)\to r^2/2$, whereas for $N\gg1$, $M^2=N(N+1)$ and for small r, $Q_{s\pi}(\infty)\to16Nr^2$ and $Q_{s0}(\infty)\to-r^2/2N$. The large value of $Q_{s\pi}(\infty)$ is reminiscent of the one exhibited by a system undergoing a quantum jump [9]. Moreover, the emission is sub-Poissonian for $\phi=0$ and super-Poissonian for $\phi=\pi$ in the squeezed vacuum. The high-field situation is of particular interest as it exhibits the essential features of the interaction of a squeezed vacuum with a two-level atom. It is also amenable to the analytical treatment as shown in Sec. IV.

IV. ANALYTICAL RESULTS IN THE INTENSE-FIELD LIMIT

When the effective Rabi frequency $\Gamma\gg\gamma N$, we may invoke high-field approximation and obtain approximate analytical results for $P(n,T)$. In this limit, the steady-state intensity is given by

$$I(\infty)=\frac{\gamma}{\gamma_0}(\gamma_0-r^2\gamma)\ . \tag{4.1}$$

Moreover, $D_0(z)$ admits the following analytical roots:

$$Z_1=-2\gamma_0,\quad Z_{2,3}=-\gamma_1\pm2i\Gamma\ , \tag{4.2}$$

where

$$\gamma_1=\gamma[(N+\tfrac{1}{2})(3-r^2)-|M|(1-r^2)\cos\phi]\ . \tag{4.3}$$

The analytical expression for the second-order correlation function then reads

$$\begin{aligned}f(\tau)=1&+\frac{r^2(\gamma-\gamma_0)}{(\gamma_0-r^2\gamma)}e^{-2\gamma_0\tau}\\ &-\frac{(1-r^2)\gamma_0}{(\gamma_0-r^2\gamma)}e^{-\gamma_1\tau}\cos(2\Gamma\tau)\ .\end{aligned} \tag{4.4}$$

Similarly, Mandel's response function $Q(T)$ can be written as

$$\begin{aligned}Q(T)=&\frac{\gamma(\gamma-\gamma_0)r^2}{\gamma_0^2}\left[1+\frac{1}{2\gamma_0T}(e^{-2\gamma_0T}-1)\right]\\ &-\frac{2\gamma(1-r^2)}{(\gamma_1^2+4\Gamma^2)}\left[\gamma_1+\frac{1}{(4\Gamma^2+\gamma_1^2)}[(\gamma_1^2-4\Gamma^2)(e^{-\gamma_1T}\cos2\Gamma T-1)-4\Gamma^2\gamma_1e^{-\gamma_1T}\sin2\Gamma T]\right]\ .\end{aligned} \tag{4.5}$$

The second term in Eq. (4.5) contributes negligibly to $Q(T)$. Moreover, $Q(T)$ exhibits sub-Poissonian character for a normal as well as a squeezed vacuum ($\phi=0,\pi$) in a narrow initial time domain [9]. Finally, approximate roots of $D_1(z)$ may be obtained as

$$\begin{aligned}Z_{1,2}&=-\lambda_\mp=-(\beta+\gamma\mp\tilde{\gamma})\ ,\\ Z_{3,4}&=-\gamma_1+\tfrac{1}{2}\gamma(1-r^2)\pm2i\Gamma\ ,\end{aligned} \tag{4.6}$$

where

$$\tilde{\gamma}=(\beta^2+r^2\gamma^2)^{1/2},\quad \beta=\gamma_0-\tfrac{1}{2}\gamma(1+r^2)\ . \tag{4.7}$$

From the structure of $P(n,Z)$ of Eqs. (3.11) and (3.12) it is clear that the contribution to $P(n,T)$ by the terms $\exp(Z_{3,4}T)$ is $O(\gamma N/\Gamma^2)$ or less as compared with that by $\exp(Z_{1,2}T)$. Therefore we can conveniently neglect their contribution to $P(n,T)$ in the intense-field limit. The analytical expression for $P(n,T)$ in this limit then reads as

$$P(n,T)=F_+^{(n)}e^{-\lambda_-T}+F_-^{(n)}e^{-\lambda_+T}\ , \tag{4.8}$$

where

$$F_+^{(0)}=\frac{\lambda_+\gamma_2}{4\gamma_0\tilde{\gamma}}\ , \tag{4.9}$$

$$F_+^{(n)}=\frac{\gamma^{n-1}I(\infty)\gamma_2^2\gamma_3^{n-1}}{(2\tilde{\gamma})^{n+1}}\frac{2}{n}\sum_{k_1=0}^{2}\frac{(\gamma_2)^{-k_1}}{k_1!(2-k_1)!}\sum_{k_2=0}^{n-k_1}\left[\frac{\gamma_3}{1-r^2}\right]^{-k_2}\frac{1}{k_2!(n-1-k_2)!}$$
$$\times\sum_{k_3=0}^{n-k_1-k_2}(2\tilde{\gamma})^{-k_3}\frac{(n+k_3)!}{k_3!}\frac{T^{n-k_1-k_2-k_3}}{(n-k_1-k_2-k_3)!}\quad \text{for } n>0 \tag{4.10}$$

and

$$F_-^{(n)}=F_+^{(n)}(\tilde{\gamma}\rightarrow-\tilde{\gamma})\ . \tag{4.11}$$

Furthermore,

$$\gamma_2=\beta+\gamma r^2+\tilde{\gamma},\quad \gamma_3=\beta(1+r^2)+\tilde{\gamma}(1-r^2)\ . \tag{4.12}$$

This analytical form of $P(n,T)$ agrees very well with the numerical results obtained from Eqs. (3.11) and (3.12).

It is clear from the expression for $P(n,T)$ that in squeezed vacuum the counting statistics is dependent on the parameters N and M and the phase ϕ. Further $P(n,T)$ in the normal unsqueezed vacuum ($N=0$) may be obtained from expressions (4.11) and (4.12) by setting $\beta=0$. In order to discuss the behavior of $P(n,T)$ in a squeezed vacuum, we assume that the squeezed vacuum is in a highly squeezed ($N\gg1$) minimum uncertainty state [$M=(N^2+N)^{1/2}$]. Again for convenience, we use $P_u(n,T)$ to describe $P(n,T)$ in the normal unsqueezed vacuum. For a squeezed vacuum, we use the notation $P_{s\phi}(n,T)$, where ϕ is the specified phase.

First note the time scales involved in the expression for $P(n,T)$. Consider, for example, the situation where the two-level atom is driven near resonance ($r\cong0$). In the normal unsqueezed vacuum, $\lambda_\pm\rightarrow\gamma(1\pm r)$, whereas in the broadband squeezed vacuum these time constants depend on phase ϕ. Specifically for $\phi=\pi$, $\lambda_+\rightarrow\gamma/4N$ and $\lambda_-\rightarrow\gamma$, whereas for $\phi=0$, $\lambda_+\rightarrow2\gamma(2N+1)$ and $\lambda_-\rightarrow\gamma$, respectively. The slow time scale $\gamma/4N$ for $\phi=\pi$ is reflected in the counting statistics when the counting times are large. Interestingly at exact resonance ($r=0$), $F_-^{(n)}$ vanishes in a squeezed vacuum. Therefore the effect of the slow time scale on $P(n,T)$ can be realized only when $r\neq0$. Moreover, the distribution $P(n,T)$ is critically dependent on the value of detuning chosen. Due to widely separated time scales in the squeezed vacuum, the counting probability $P(n,T)$ depends on the length of the counting interval T. This may be seen from the analytical expression (4.9) for $P(0,T)$, which describes the probability of emitting no photons in time T. For $T>\gamma^{-1}$, $P_{s\pi}(0,T)$ is substantially higher than $P_u(0,T)$ or $P_{s0}(0,T)$. This is a consequence of the inhibition of the atomic decay [2–12] in the squeezed vacuum when $\phi=\pi$. In Figs. 2(a) and 2(b), we have plotted the temporal behavior of $P(n,T)$ for $n=1$ and 4, in a squeezed and normal vacuum. These plots show that $P(n,T)$ reaches a maximum for some time T_n, which is characteristic of the value of n chosen. For counting times $T<T_n$, $P_{s\pi}<P_u$, whereas $P_{s0}\geq P_u$. On the other hand, for $T>T_n$, $P_{s\pi}>P_u$ and $P_{s0}\leq P_u$. This contrasting behavior of $P_{s\pi}$ and P_{s0} is due to the difference in the fluctuations experienced by the atomic dipole [2,3]. For $\phi=\pi$, the re-

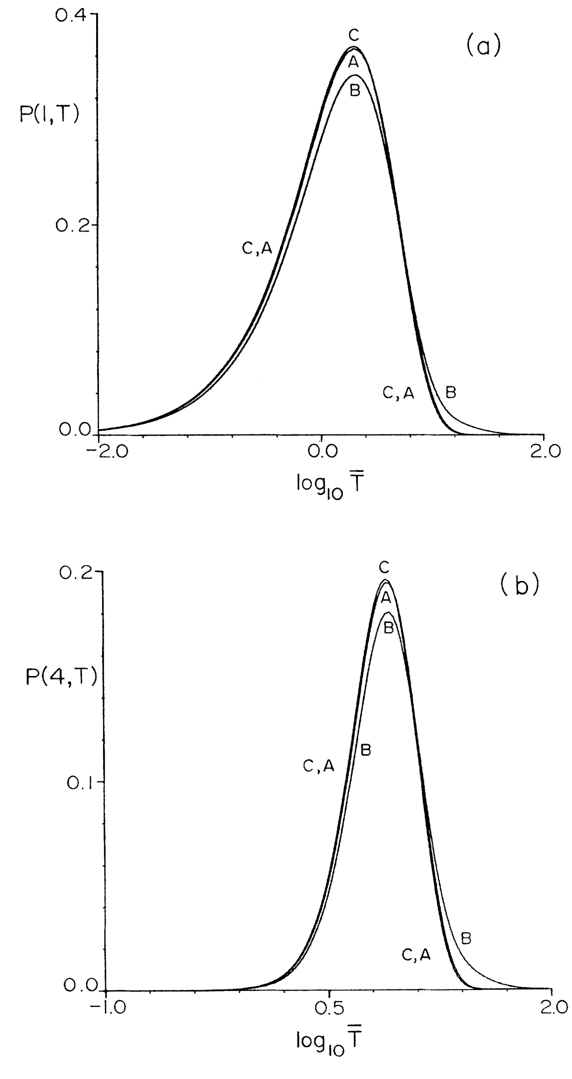

FIG. 2. Temporal behavior of $P(n,T)$ for (a) $n=1$ and (b) $n=4$, in a squeezed and normal vacuum with $\alpha/2\gamma=10$, $\Delta/2\gamma=2$, and $\bar{T}=2\gamma T$. Curves A–C correspond to $(N,\phi)=(0,0)$, $(2,\pi)$, and $(2,0)$, respectively.

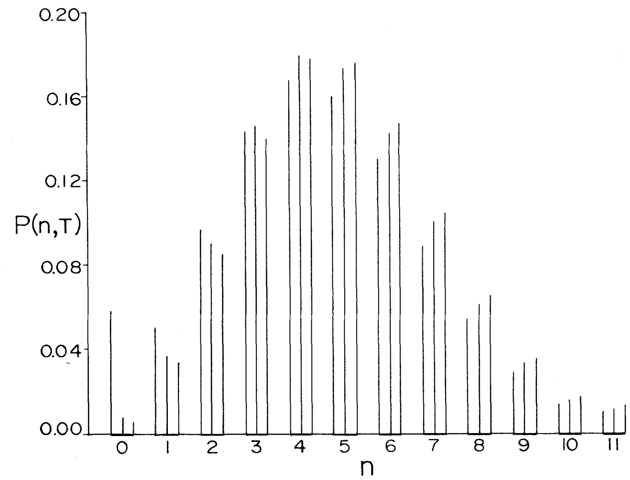

FIG. 3. Probability distribution $P(n,T)$ with $2\gamma T=10$. For each n, the central histogram corresponds to the normal vacuum. The left and right ones are for squeezed vacuum ($N=2$) with $\phi=\pi$ and 0, respectively. Other data are same as in Fig. 2. The Poisson distribution with the same mean coincides with the distribution in a normal vacuum.

duced fluctuations tend to inhibit the emission. As a consequece the probability of detection of photons is higher at longer times. This behavior may be further seen from Fig. 3, where we have plotted the distribution $P(n,T)$ vs n for counting time $2\gamma T=10$. This counting time is close to $T_n \cong T_3$. Therefore for $n<3$, $P_{s\pi}>P_u$, whereas $P_{s0}<P_u$. On the other hand, for $n>3$, $P_{s\pi}<P_u$ and $P_{s0}>P_u$, as can be seen from Fig. 3. This behavior is a manifestation of the phase-dependent decay in the squeezed vacuum.

V. CONCLUSIONS

In conclusion, we find that the statistics of resonance fluorescence in a broadband squeezed vacuum is markedly different from that in the normal vacuum. For large counting times, the analytical form of $Q(\infty)$ suggests that at exact resonance, the fluorescence in the squeezed vacuum has a sub-Poissonian character that depends on the phase ϕ. The case with $\phi=0$ is more sub-Poissonian than the case with $\phi=\pi$. Interestingly, for a Rabi frequency $\alpha\cong\gamma/2$, the distribution in a highly squeezed vacuum is narrower than that in the normal vacuum when a proper choice of the phase is made. In the high-field situation with nonzero detuning, $\phi=0$ and π exhibit sub- and super-Poissonian statistics, respectively. This phase-dependent decay in a squeezed vacuum is also reflected in the counting statistics $P(n,T)$. The analytical expression for $P(n,T)$ in the intense-field limit shows time scales that depend on the parameters N and ϕ. The inhibition of fluorescence in a squeezed vacuum is reflected in the probability of detecting photons at times longer than the spontaneous-emission time scale.

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