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Phase-dependent fluctuations of resonance fluorescence beyond the squeezing regime

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ABSTRACT

Phase-dependent quantum features of the light scattered by a two-level atom driven by a monochromatic laser were investigated theoretically using the method of conditional homodyne detection [Carmichael, Castro-Beltran, Foster, Orozco, Phys. Rev. Lett. 85 (2000) 1855]. The splitting of fluctuations into terms of second and third order correlations of the dipole noise is obtained analytically. For the out-of-phase quadrature and weak laser driving the former are known to be squeezed. The third order fluctuations, newly found in this paper, grow with the laser intensity, contaminate squeezing below saturation, and dominate above it. They are responsible for the non-classicality and non-Gaussianity of the fluorescence for moderate to strong driving. Conditional homodyne detection, in both time and frequency domains, illustrates more general phase-dependent effects than squeezing, and is much less restricted by finite collection and quantum detector efficiencies than standard homodyne detection schemes.

1. Introduction

Resonance fluorescence is a cornerstone of quantum optics for its displays of quantum signatures of light, such as antibunching [1] and sub-Poissonian statistics [2] (intensity fluctuations), and squeezing [3,4] (phase-dependent amplitude fluctuations). Squeezing in the light scattered by a single two-level atom is still an open problem. The requirement of low laser intensity meets serious obstacles in the homodyne detector's finite collection and quantum efficiencies, while for stronger driving saturation destroys squeezing [5]. Proposed remedies include increase of detection's solid angle [6] and enhancement of squeezing via bad cavities [7], feedback [8], and atomic coherence [6].

Recently, Carmichael, Orozco, and coworkers demonstrated for a cavity QED system [9,10] a robust scheme for the measurement of weak squeezed light: conditional homodyne detection (see also Refs. [11–17]). CHD features two inequalities that a quantum field, squeezed or not, violates, thus non-classicality of a quadrature does not depend on squeezing. In CHD, the dynamics of a quadrature (field amplitude) is recorded by balanced homodyning on the cue of photon (intensity) detections. The conditioning on photon detections greatly reduces counting noise and nearly eliminates the problem of detector efficiency by not degrading the signal, making this method a much better prospect for measurement of squeezing of resonance fluorescence than standard homodyne methods.

Progress in this direction was reported by Gerber et al. who measured the amplitude–intensity correlation of the resonance fluorescence of a three-level atom in a setup similar to CHD [18]. The regression of the dipole field upon a photon emission, which projects the atom in the ground state, was demonstrated. The squeezing conditions, however, were not achieved. One wonders if squeezing in resonance fluorescence is close to be seen. Here and in CHD, the weak field limit, due to the low photon flux from the atom, translates into impractical data collection times due to laser and setup instabilities. Thus, experiments deviate from this limit and squeezing, though maximal near saturation [4], gets mixed with barely explored additional fluctuations.

In this paper we investigate the character, mainly non-classical and non-Gaussian, of phase-dependent fluctuations in resonance fluorescence for arbitrary laser field strengths within the framework of CHD, in both the time and spectral domains, providing better understanding of amplitude noise than the restrictive squeezing. The amplitude–intensity correlation is of third order in the dipole's electric field amplitude. Splitting this field into a mean plus noise, second and third order correlations in the dipole fluctuations are obtained. In the weak field limit the out-of-phase quadrature's second order term is squeezed [3,4]. The third order term, the main analytical result of this paper, grows with the laser intensity and, still below saturation, *contaminates* the squeezed part, actually enhancing the characteristic negative peak in the spectrum. This regime is probably closer, experimentally, than the deep weak field limit. Above saturation, on the other hand, the spectrum has the Rabi sidebands

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(second order term) mixed with a large third order term. The latter not only signal a deviation from Gaussian (second order) noise, but account for the non-classicality of the out-of-phase quadrature in the strong field regime, as seen from the inequalities of CHD [9,10].

Non-Gaussian (non-zero odd-order) fluctuations in the field's amplitude are still poorly explored in quantum optics. Although third order fluctuations were also studied in phase-sensitive intensity correlations in the stationary state of resonance fluorescence [5,19], their role emphasized the relation between squeezing and intensity fluctuations instead. More recently, Shchukin and Vogel [20] proposed a hierarchy of conditions for the non-classicality of a field to all orders, but specific results for resonance fluorescence and many other optical systems are non-existent.

An additional feature of CHD is that it reveals non-Gaussian noise when the amplitude–intensity correlation is asymmetric, i.e., the light's intensity and amplitude have different noise properties, as observed in cavity QED both numerically [9,12] and experimentally [10], and numerically for the resonance fluorescence of a three-level atom [17]. For a two-level atom the small Hilbert space imposes a symmetric correlation [12,17], but the calculations for positive and negative intervals have to be interpreted differently.

This paper is organized as follows: a brief review of the theory of CHD, the splitting of fluctuations in resonance fluorescence and their (non-) classicality are the subject of Section 2. The spectral representation is presented in Section 3, and the conclusions are given in Section 4. Appendix A deals with the correlation for negative intervals, and Appendix B presents the analytical solutions to the equations for the correlations.

2. Theory

A setup of CHD is sketched in Fig. 1 (see also Ref.[9]). A quadrature of the field, $E_\phi \propto \sigma_\phi$, is measured in balance homodyne detection for a time τ whenever a photon is detected, $I \propto \sigma_+\sigma_-$, giving a third order correlation in the field amplitude [17], which in normalized form is

$$h_\phi(\tau) = \frac{\langle : \sigma_+(0)\sigma_-(0)\sigma_\phi(\tau) : \rangle}{\langle \sigma_+\sigma_- \rangle_{st} \langle \sigma_\phi \rangle_{st}}, \tag{1}$$

where $\sigma_\pm$, σ_z are Pauli pseudo spin operators,

$$\sigma_\phi = \left(\sigma_- e^{i\phi} + \sigma_+ e^{-i\phi} \right) / 2 \tag{2}$$

is the dipole quadrature operator, ϕ is the phase between the strong local oscillator and the driving field, and the dots : : indicate time and normal operator ordering. The brackets indicate averaging via a density operator, and the subindex *st* denotes steady state values.

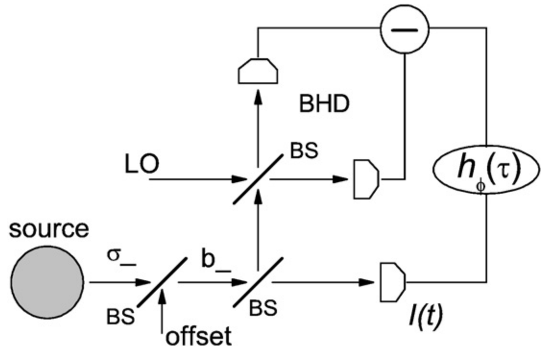

Fig. 1. Sketch for conditional homodyne detection. Balanced homodyne detection (BHD) of a quadrature $E_\phi(\tau)$ is made on the cue of photon detections $I(t)$. LO is the local oscillator beam and BS are beam splitters.

Ideally, conditioning and normalization make the correlation independent of detector efficiencies.

We work in a frame rotating at the laser frequency ν, $\tilde{\sigma}_\pm(t) = \sigma_\pm(t) e^{\mp i\nu t}$, and split the dipole dynamics into a mean $\alpha_\pm = \langle \tilde{\sigma}_\pm \rangle_{st}$ plus fluctuations $\Delta\tilde{\sigma}_\pm(t)$. The intensity at the start detector is proportional to

$$\langle \sigma_+\sigma_- \rangle = \alpha_+\alpha_- + \langle \Delta\sigma_+\Delta\sigma_- \rangle, \tag{3}$$

where $\alpha_+\alpha_-$ and $\langle \Delta\sigma_+\Delta\sigma_- \rangle$ are the coherent and incoherent parts, respectively, of the emission. The input field at the BHD port is

$$\tilde{\sigma}_\phi(t) = \alpha_\phi + \Delta\tilde{\sigma}_\phi(t), \tag{4}$$

where $\alpha_\phi = \langle \tilde{\sigma}_\phi \rangle_{st}$, and

$$\Delta\tilde{\sigma}_\phi = \frac{1}{2}\left(\Delta\tilde{\sigma}_- e^{i\phi} + \Delta\tilde{\sigma}_+ e^{-i\phi} \right). \tag{5}$$

For some sources the mean field of a quadrature is zero. Such is the case of the $\phi = 0$ quadrature in resonance fluorescence when the atom and laser frequencies are equal, where $\sigma_0 = 0$, leading to $h_0(\tau) = 0$. In order to observe this quadrature the source field has to be mixed with a coherent offset with real amplitude A in phase with the local oscillator [9]. Hence, we replace $\sigma_\pm$ with $b_\pm = \sigma_\pm + Ae^{\pm i\phi}$, and σ_ϕ with $b_\phi = \sigma_\phi + A$. Thus, their means are $\beta_\pm = \langle b_\pm \rangle_{st}$ and $\beta_\phi = \langle b_\phi \rangle_{st}$. The intensity at the start detector is $\langle b_+ b_- \rangle_{st}$.

With $\tilde{\sigma}_\pm(t) = \alpha_\pm + \Delta\tilde{\sigma}_\pm(t)$ and the offset, Eq. (1) is decomposed into second and third order correlations for the dipole fluctuations, $h_\phi(\tau) = h_\phi^{(2)}(\tau) + h_\phi^{(3)}(\tau)$, where

$$h_\phi^{(2)}(\tau) = 1 + \frac{\langle : [\beta_+ \Delta\tilde{\sigma}_-(0) + \beta_- \Delta\tilde{\sigma}_+(0)] \Delta\tilde{\sigma}_\phi(\tau) : \rangle}{\beta_\phi \langle b_+ b_- \rangle_{st}}, \tag{6a}$$

$$h_\phi^{(3)}(\tau) = \frac{\langle : \Delta\tilde{\sigma}_+(0) \Delta\tilde{\sigma}_-(0) \Delta\tilde{\sigma}_\phi(\tau) : \rangle}{\beta_\phi \langle b_+ b_- \rangle_{st}}. \tag{6b}$$

The amplitude–intensity correlation is time-symmetric for two-level atom resonance fluorescence [17]. For positive intervals, $h_\phi(\tau)$ is the measurement of a quadrature conditioned on a photon detection, the main subject of this paper. For negative intervals, however, the correlation is that of a measurement of the intensity conditioned on a quadrature detection. For completeness, this is included in Appendix A. In the following, we restrict the treatment to the exact resonance case, for which analytic solutions to the equations for the correlations of fluctuation operators are obtained in Appendix B.

2.1. Second and third order correlations

First, we quote the steady state values of the dipole operators from the Bloch equations [17,21,22]:

$$\alpha_\mp = \langle \tilde{\sigma}_\mp \rangle_{st} = \pm i \frac{Y/\sqrt{2}}{1 + Y^2}, \tag{7a}$$

$$\langle \sigma_+\sigma_- \rangle_{st} = \frac{1}{2}(1 + \langle \sigma_z \rangle_{st}) = \frac{Y^2/2}{1 + Y^2}, \tag{7b}$$

where

$$Y = \sqrt{2}\Omega / \gamma, \tag{8}$$

Ω is the atom-laser coupling and γ is the spontaneous emission rate.

The numerator in the second term of Eq. (6a) gives the second order fluctuations,

$$H_\phi^{(2)}(\tau) = 2Re\left[(\alpha_- + A) \langle \Delta\tilde{\sigma}_+(0) \Delta\tilde{\sigma}_\phi(\tau) \rangle \right], \tag{9}$$