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Effects of a squeezed vacuum on absorptive optical bistability

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**EFFECTS OF A SQUEEZED VACUUM ON
ABSORPTIVE OPTICAL BISTABILITY**

by

Steven Frederick Haas

A Thesis Submitted to
COMMITTEE ON OPTICAL SCIENCES (Graduate)
In Partial Fulfillment of the Requirements
For the Degree of
MASTER OF SCIENCE
In the Graduate College
THE UNIVERSITY OF ARIZONA

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ABSTRACT

We calculate the effects of a squeezed vacuum on absorptive optical bistability (AOB) using the different relaxation rate approximation for in-quadrature and in-phase components of a two-level system developed by Gardiner.¹ An expression for the complex absorption coefficient is developed, and the result applied to the AOB equation for the unidirectional ring cavity. We find a significant degradation of bistability for values of the in-quadrature decay constant less than or equal to approximately .5 of the in-phase decay constant. Effects of detuning and relative phase of the pump field to the squeezed vacuum field are also examined.

1. INTRODUCTION

Recent work by Gardiner¹ has shown that the effects of squeezed vacuum on the phase decay of a two-level system can be approximated by having different relaxation rates of the in-quadrature and in-phase components. Ritsch and Zoller² have derived the probe absorption spectra of a two-level system in a squeezed vacuum with a reduced probe field linewidth as a result. Additional work by An *et al.*,⁵ has shown that results similar to Ritsch and Zoller may be derived using Fourier series methods to solve the Schrödinger equations of motion. In this paper we examine the effects of squeezed vacuum on absorptive optical bistability (AOB) in the unidirectional ring cavity using the population density matrix equations of motion. We derive a formula predicting the onset of AOB and illustrate the onset numerically.

2. EQUATIONS OF MOTION

A unidirectional ring cavity of length L with an input field E_i , output field E_t , and a squeezing field, contains a nonlinear two-level medium of length l in one portion of the cavity (Fig. 1).

Figure 1. Ring cavity schematic

Szoke *et al.*,³ were the first to develop the equations for optical bistability. We use the form of the absorptive optical bistability given in Meystre *et al.*,⁴ where the complex input field E_i is related to the complex transmitted field E_t by the absorptive optical bistability equation

$$E_i \cong E_t \left[1 + \frac{\alpha \ell}{T} \right], \quad (1)$$

where T is the transmittance of the cavity input and output mirrors, and α is the complex absorption coefficient of the medium. Since α is a function of the field within the cavity, we assume a monochromatic, single-mode field of the form

$$E(z, t) = \frac{1}{2} \mathcal{E}(t) e^{i(Kz - \nu t)} + \text{c.c.}, \quad (2)$$

where $\mathcal{E}(t)$ is a complex amplitude that we assume varies little in atomic lifetimes. This field induces a polarization in the medium of a similar form

$$P(z, t) = \frac{1}{2} \mathcal{P}(t) e^{i(Kz - \nu t)} + \text{c.c.}, \quad (3)$$

where $\mathcal{P}(t)$ is a slowly varying complex polarization. The absorption coefficient α is given in terms of the induced polarization $\mathcal{P}$ by

$$\alpha = \frac{-iK}{2\epsilon} \cdot \frac{\mathcal{P}}{\mathcal{E}}, \quad (4)$$

where ϵ is the permittivity of the medium. The complex polarization is $P(z, t)$ of Eq. (3) with $\mathcal{P}$ given as

$$\mathcal{P} = 2\wp \rho_{ab}, \quad (5)$$

with $\wp$ being the dipole interaction matrix element and ρ_{ab} the off-diagonal element of the 2-level population matrix. We assume both the dipole and rotating wave approxi-

mations and will do so throughout the remainder of this paper.

Derivation of the expression for α begins by considering the equation of motion for the population density matrix element coupling the two states. From An *et al.*⁵, this equation including squeezing effects is given as

$$\dot{\rho}_{ab} = -(\gamma_+ + i\delta)\rho_{ab} - \gamma_-\rho_{ba} + i\mathcal{V}_{ab} D \quad (6)$$

where

$$\gamma_{\pm} = \frac{\gamma_u \pm \gamma_v}{2}$$

$\delta = \omega - \nu$, detuning from line center with ω = line center frequency,

ν = pump field frequency.

γ_u = decay constant for the component of ρ_{ab} in-phase with
squeezed field

γ_v = decay constant for the component of ρ_{ab} in-quadrature with
squeezed field

$\mathcal{V}_{ab}$ = interaction energy matrix element

D = population density difference; $\rho_{aa} - \rho_{bb}$

Solving Eq. (6) in steady state ($\dot{\rho}_{ab} = 0$) and using the property $\rho_{ba} = \rho_{ab}^*$, we find

$$\rho_{ab} = \frac{-\gamma_-\rho_{ab}^* + i\mathcal{V}_{ab} D}{\gamma_+ + i\delta} \quad (7)$$

Substituting the complex conjugate of this expression back into Eq. (7), we find

$$(\gamma_+ + i\delta)\rho_{ab} = \frac{1}{\gamma_+ - i\delta} \cdot [\gamma_-^2 \rho_{ab} + i\gamma_- \mathcal{V}_{ab}^* D] + i\mathcal{V}_{ab} D . \quad (8)$$

$$\rho_{ab} = \frac{i(\gamma_+ - i\delta)\mathcal{V}_{ab} + i\gamma_- \mathcal{V}_{ab}^*}{\gamma_+^2 - \gamma_-^2 + \delta^2} D . \quad (9)$$

We can find an expression for the population difference, D , containing squeezing effects. The equation of motion for the population difference is given in An *et al.*⁵ as

$$\dot{D} = -\Gamma(N + D) - 2i [\mathcal{V}_{ab} \rho_{ab}^* + \text{c.c.}] . \quad (10)$$

Solving for D in the steady state gives

$$D = -\frac{N}{1 + I' \mathcal{L}'} , \quad (11)$$

where

N = number of dipoles per unit volume

$\mathcal{L}'$ = squeezed Lorentzian; $\frac{\gamma_u \gamma_v}{\gamma_u \gamma_v + \delta^2}$

$$I' = \frac{(\mathcal{E}_r / \hbar)^2}{\Gamma \gamma_v} + \frac{(\mathcal{E}_i / \hbar)^2}{\Gamma \gamma_u}$$

Γ = population difference decay constant

Substituting Eq. (11) and the interaction matrix element

$$\mathcal{V}_{ab} = -\frac{\wp \mathcal{E}}{2\hbar} . \quad (12)$$

into Eq. (9), we find

$$\rho_{ab} = i \wp \frac{N}{2\hbar(1+I'\mathcal{L}')} \cdot \frac{1}{\gamma_+^2 - \gamma_-^2 + \delta^2} \cdot [\gamma_- \mathcal{E}^* + (\gamma_+ - i\delta)\mathcal{E}] . \quad (13)$$

Now using $\mathcal{E} = |\mathcal{E}|e^{i\phi}$, where ϕ is the phase of the pump field relative to the squeezed vacuum field, and $\gamma_u \gamma_v = (\gamma_+^2 - \gamma_-^2)$, the equation for ρ_{ab} becomes

$$\rho_{ab} = i N \wp \frac{|\mathcal{E}|}{2\hbar} e^{i\phi} \cdot \frac{\gamma_- e^{-2i\phi} + \gamma_+ - i\delta}{(\gamma_u \gamma_v + \delta^2)(1+I'\mathcal{L}')} . \quad (14)$$

where the dimensionless intensity I' is now

$$I' = \frac{(\wp \mathcal{E}/\hbar)^2}{\Gamma} \cdot \left[\frac{\cos^2 \phi}{\gamma_v} + \frac{\sin^2 \phi}{\gamma_u} \right] . \quad (15)$$

Substituting Eq. (5) for $\mathcal{P}$ and $|\mathcal{E}|e^{i\phi}$ for $\mathcal{E}$ we find the absorption coefficient

$$\alpha = \alpha_0 \frac{\gamma(\gamma_+ - i\delta + \gamma_- e^{-2i\phi})}{(1 + I'\mathcal{L}')(\gamma_u \gamma_v + \delta^2)} , \quad (16)$$

where the linear absorption coefficient α_0 is defined as

$$\alpha_0 = NK \frac{\wp^2}{2\epsilon\hbar\gamma} . \quad (17)$$

The quantity I' must be expressed as a function of the transmitted field, E_t , instead of the internal pump field $\mathcal{E}$, if it is to be used in Eq. (1). For $\alpha \ll 1$, the transmitted field is

$$E_t = \sqrt{T} \mathcal{E} e^{(iK-\alpha)t} \cong \sqrt{T} \mathcal{E} , \quad (18)$$

or

$$\mathcal{E} \cong \frac{E_t}{\sqrt{T}} \rightarrow I \cong \frac{I_t}{T} . \quad (19)$$

with I and I_t being the pump field intensity and the transmitted field intensity, respectively. With squeezing, however, the pump field intensity is

$$I = \frac{(\varphi \mathcal{E} / \hbar)^2}{\Gamma \gamma_+} , \quad (20)$$

where γ_+ is defined in Eq. (6). By equating Eq. (19) and Eq. (20) we have

$$\frac{(\varphi \mathcal{E} / \hbar)^2}{\Gamma} = I_t \frac{\gamma_+}{T} , \quad (21)$$

and for I'

$$I' = I_t \frac{\gamma_+}{T} \left[\frac{\cos^2 \phi}{\gamma_v} + \frac{\sin^2 \phi}{\gamma_u} \right] . \quad (22)$$

Finally substituting Eq. (16) into Eq. (1), have

$$E_i \cong E_t \left[1 + \frac{2C\gamma(\gamma_+ - i\delta + \gamma_- e^{-2i\phi})}{(1 + I' \mathcal{L}')(\gamma_u \gamma_v + \delta^2)} \right] , \quad (23)$$

where if $\gamma_u = \gamma_v$ (representing the nonsqueezed case), the bistability parameter $2C = \alpha_0 l/T$ must be greater than or equal to 8 for bistability to occur.

For $\gamma_u \neq \gamma_v$ the situation is more complex since the $2C$ parameter becomes a function of squeezing and the relative phase ϕ . The effects of squeezing on $2C$ must now be determined.

We can rewrite Eq. (16) as

$$\alpha = \alpha_0 \frac{\gamma\gamma_+}{\gamma_u \gamma_v} \cdot \left[\frac{1 - \frac{i\delta}{\gamma_+} + \frac{\gamma_-}{\gamma_+} e^{-2i\phi}}{1 + \frac{\delta^2}{\gamma_u \gamma_v} + I'} \right] . \quad (24)$$

Or as,

$$\alpha = \alpha_0 \frac{\gamma\gamma_+}{\gamma_u \gamma_v} \cdot [\sigma_r - i\sigma_i] , \quad (25)$$

where

$$\sigma_r = \frac{1 + \frac{\gamma_-}{\gamma_+} \cos 2\phi}{F} \quad (26)$$

$$\sigma_i = \frac{\frac{\delta}{\gamma_+} + \frac{\gamma_-}{\gamma_+} \sin 2\phi}{F} \quad (27)$$

$$F = 1 + \frac{\delta^2}{\gamma_u \gamma_v} + Y \quad (28)$$

$$Y = I' \quad (29)$$

We introduce a new parameter $2C'$ such that $2C' = \alpha l/T$. The task now is to determine the values of $2C'$ for which optical bistability is possible. Using Eq. (2.5-3) from Gibbs,⁶ and ignoring cavity detuning

$$\frac{1}{2C'} < \text{Sup} \left\{ -\sigma_r - Y \frac{d\sigma_r}{dY} + Y \left[\left(\frac{d\sigma_r}{dY} \right)^2 + \left(\frac{d\sigma_i}{dY} \right)^2 \right]^{1/2} \right\}, \quad (30)$$

where the suprema is the largest value of the LHS of Eq. (30) when Y is allowed to vary over all positive values. From Eqs. (26) and (27)

$$\frac{d\sigma_r}{dY} = - \frac{1 + \frac{\gamma_-}{\gamma_+} \cos 2\phi}{F^2} \quad (31)$$

$$\frac{d\sigma_i}{dY} = \frac{\frac{\delta}{\gamma_+} + \frac{\gamma_-}{\gamma_+} \sin 2\phi}{F^2}. \quad (32)$$

Then Eq. (30) becomes

$$\frac{1}{2C'} < \text{Sup} \left\{ \frac{1}{F^2} \cdot \left[- \left[1 + \frac{\delta^2}{\gamma_u \gamma_v} \right] \left[1 + \frac{\gamma_-}{\gamma_+} \cos 2\phi \right] + Y \sqrt{A} \right] \right\}, \quad (33)$$

where

$$A = 1 + \frac{2\gamma_-}{\gamma_+} \left[\cos 2\phi + \frac{\delta}{\gamma_+} \sin 2\phi \right] + \frac{\gamma_-^2}{\gamma_+^2} + \frac{\delta^2}{\gamma_+^2} . \quad (34)$$

Since we want the suprema, we find the maximum value of Y on the RHS of Eq. (33) by setting the differential with respect to Y of Eq. (33) to zero and solving for Y . This gives

$$\frac{\sqrt{A}}{F^2} - \frac{2}{F^2} \cdot \left[- \left[1 + \frac{\delta^2}{\gamma_u \gamma_v} \right] \left[1 + \frac{\gamma_-}{\gamma_+} \cos 2\phi \right] + Y\sqrt{A} \right] = 0 , \quad (35)$$

Or,

$$Y = \frac{\left[1 + \frac{\delta^2}{\gamma_u \gamma_v} \right] \left[\sqrt{A} + 2 \left[1 + \frac{\gamma_-}{\gamma_+} \cos 2\phi \right] \right]}{\sqrt{A}} . \quad (36)$$

Substituting into Eq. (33)

$$\frac{1}{2C'} < \frac{\left[1 + \frac{\delta^2}{\gamma_u \gamma_v} \right] \left[\sqrt{A} + 1 + \frac{\gamma_-}{\gamma_+} \cos 2\phi \right]}{F^2} . \quad (37)$$

Finally, solving for $2C$ and substituting for F

$$2C' > \frac{4 \left[1 + \frac{\gamma_-}{\gamma_+} \cos 2\phi + \sqrt{A} \right] \left[1 + \frac{\delta^2}{\gamma_u \gamma_v} \right]}{A} . \quad (38)$$

Now we determine the relationship between $2C'$ and $2C$. From Gardiner¹ the decay constants γ_u and γ_v are

$$\gamma_u = \gamma(2N + 2M + 1) , \quad (39)$$

$$\gamma_v = \gamma(2N - 2M + 1) , \quad (40)$$

$$M \geq \sqrt{N(N + 1)} . \quad (41)$$

Here N is proportional to the number of photons in the squeezed vacuum and M measures the amount of squeezing. Maximal squeezing is given by the equality in Eq. (41). From our definition of $2C'$ and $2C$ we write

$$2C' = 2C \frac{\gamma}{\gamma_+} \cdot \frac{\gamma\gamma_+}{\gamma_u \gamma_v} . \quad (42)$$

But from Eqs. (39), (40), (41) we find

$$\gamma_u \gamma_v \geq \gamma^2 . \quad (43)$$

Where equality indicates maximal squeezing. When substituted into Eq. (42) we have

$$2C' = 2C . \quad (44)$$

Therefore, Eq. (38) is

$$2C \geq 4 \cdot \frac{\left[1 + \frac{\gamma_-}{\gamma_+} \cos 2\phi + \sqrt{A}\right] \left[1 + \frac{\delta^2}{\gamma^2}\right]}{A}. \quad (45)$$

As a check, we remove the squeezing effects by letting $\gamma_u = \gamma_v = \gamma_+ = \gamma \rightarrow \gamma_- = 0$ so that Eq. (45) reduces to

$$C > 2 \left[1 + \left[1 + \frac{\delta^2}{\gamma^2} \right]^{1/2} \right], \quad (46)$$

which agrees with Gibbs⁶ provided we make the identification $\Delta_0^2 = \delta^2/\gamma^2$.

3. ANALYSIS

Examination of Eq. (23) shows that the bistability is dependent upon three variables: γ_v , δ , and ϕ . We now examine the effects of each of the three variables. The figures below show transmitted intensity as a function of the input intensity for the different variables. We take Eq. (23), multiply it by its complex conjugate and use Eq. (43) to get

$$\frac{I_i}{T} \cong \frac{I_t}{T} \left[\left[1 + \frac{2C\gamma(\gamma_+ + \gamma_- \cos 2\phi)}{(1 + I' \mathcal{L}')(\gamma^2 + \delta^2)} \right]^2 + \left[\frac{2C\gamma(\delta + \gamma_- \sin 2\phi)}{(1 + I' \mathcal{L}')(\gamma^2 + \delta^2)} \right]^2 \right]. \quad (47)$$

With the substitutions

$$X = \left[\frac{I_t}{T} \right] \cdot Q \quad (48)$$

$$Y = \left[\frac{I_t}{T} \right] \cdot Q \quad (49)$$

$$Q = \gamma_+ \left[\frac{\cos^2 \phi}{\gamma_v} + \frac{\sin^2 \phi}{\gamma_u} \right] \mathcal{L}, \quad (50)$$

$$S_1 = \frac{\gamma(\gamma_+ + \gamma_- \cos 2\phi)}{\gamma^2 + \delta^2} \quad (51)$$

$$S_2 = \frac{\gamma(\delta + \gamma_- \sin 2\phi)}{\gamma^2 + \delta^2}, \quad (52)$$

Eq. (47) may be written

$$X \cong Y \left[\left[1 + \frac{2C}{1+Y} \cdot S_1 \right]^2 + \left[\frac{2C}{1+Y} \cdot S_2 \right]^2 \right]. \quad (53)$$

The effects of squeezing are determined by the scaling factor Q along the X and Y axes and by two shape factors, S_1 and S_2 . Since Q is a function of γ_u^{-1} and γ_v^{-1} , as γ_v decreases (indicating greater squeezing), the effect is to amplify the pump saturation intensity through the $\cos^2 \phi$ term. Examining the $\phi = 0$ case and ignoring detuning, the denominator increases for small γ_v and therefore decreases the overall value of the function. This means that while the value of X for which bistability starts is lessened

through a smaller Q , the resulting Y for the onset of bistability is smaller as well. Figure 2 illustrates this effect with a plot of Y vs. X for zero relative phase and no detuning for different values of γ_v . γ_u is normalized to unity in all of the following figures.

Fig. 2. Y vs. X for, left to right, $\gamma_v = .01, .1, .5, .99$ with $\phi, \delta = 0$

Figure 3 shows a similar plot; however, here the relative phase is $\phi = \pi/2$ with $\delta = 0$. In this case, the effect of the relative phase is significant. Analysis of this effect is investigated further below.

Fig. 3. Y vs. X for, left to right, $\gamma_v = .01, .1, .5, .99$ with $\phi = \pi/2$ and $\delta = 0$

The effect of detuning from line center is most easily seen by approximating the γ_+ and γ_- terms in Eq. (53) by noting that for small γ_v ($\cong .1$), both γ_+ and $\gamma_- \cong \gamma_u / 2$. Then the second term in the first squared term can be approximated as

$$\frac{2C}{1+Y} \cdot S_1 \cong \frac{2 \cdot 2C}{1+Y} \cdot \quad (54)$$

Similarly, the second squared term can be approximated as

$$\frac{2C}{1+Y} \cdot S_2 \cong \frac{2 \cdot 2C (\delta + \sin \phi \cos \phi)}{\cos^2 \phi (1+Y)} \cdot \quad (55)$$

Examination of these two terms and Eq. (50) shows that for $\delta \neq 0$ the primary effect of

detuning is through the $\mathcal{L}'$ term in Q . As δ becomes large Q becomes small, thereby decreasing the scaling. The overall ratio of the second squared term (see Eq. (52)) increases and therefore the shape of the function changes as well. The X necessary for bistability increases, but the resulting Y is increased too. The bistability becomes a mixture of absorptive and dispersive types; and may even disappear. Figure 4 illustrates the effects of detuning from line center for zero relative phase. Squeezing is present with $\gamma_v = .1$.

Fig. 4. Y vs. X for, left to right, $\delta = 0, 1, 2, 3$ with $\phi = 0$, $\gamma_v = .1$

The relative phase of the pump field to the squeezed vacuum field has very significant effects, as noted above. Returning to Eq. (53) and taking $\delta = 0$, we see that for $\phi = 0$ the shape factor S_2 in the second squared term disappears. The scale factor

Q is dominated by the γ_v^{-1} term from its $\cos^2\phi$ dependence (see Eq. (50)). The denominator becomes large from the amplifying effect of γ_v^{-1} and bistability is degraded. For $\phi = \pi/2$ the second squared term in Eq. (53) is again zero, while Q reduces to its nonsqueezed value, $1/2$, for $\gamma_u = 1$. If $0 \leq \phi \leq \pi/2$, both of the squared terms in Eq. (53) contribute through Q and the shape factors S_1 and S_2 . Figure 5 examines the effects of the pump field phase relative to the squeezed vacuum field. Squeezing is present with $\gamma_v = .5$; $\delta = 0$.

Fig. 5. Y vs. X for, left to right, $\phi = 0, 1.0, 1.9, \pi/2$ with $\delta = 0$, $\gamma_v = .5$

4. CONCLUSIONS

We have shown that absorptive optical bistability in the ring cavity is strongly affected by squeezing of the vacuum field. Squeezing will affect bistability through an increase of the atomic decay time and through the introduction of a relative phase term between the pump field and the squeezed vacuum field. The degree of squeezing, indicated by γ_v , necessary to affect bistability is on the order of $.5\gamma_u$. This implies that the squeezed field must increase the decay time of the two-level system by a factor of about 2 or more. Detuning of the pump field from atomic line center affects or eliminates optical bistability through a mixture of absorptive and dispersive types. We have also shown that the phase of the pump field relative to the squeezed vacuum field can play a significant role. For $\phi \cong \pi/2$ optical bistability can be affected, while for $\phi \cong 0$ the onset of optical bistability is dependent only on the degree of squeezing.

The phenomena investigated here will be difficult to observe experimentally. The problem is to force the atoms in an extended media to respond to the squeezed field with the same phase. Random motion of the atoms in the media introduce random phase fluctuations in the squeezed field. Hence the field will gradually lose its squeezed character. The key is to reduce the random motion of the atoms and to force the atoms to respond to the squeezed field with the same phase. This may be achieved, in principle, by cooling ($\cong 0K$) a very small number of atoms which are then subjected to the squeezed field.

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