SENSITIVE ABSORPTION MEASUREMENTS USING AMPLITUDE–SQUEEZED LIGHT

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The application of amplitude–squeezed light obtained from semiconductor diode lasers with different types of negative feedback to sensitive absorption measurements is considered. It is shown that the sensitivity is limited by the type of feedback loop and by the quality of the detector. It can be much better than the usual shot–noise limit. For existing photodetectors, a factor–of–three improvement in the sensitivity can be, obtained.

INTRODUCTION

Amplitude squeezing is the production of light whose photon–number fluctuations are less than those of light in the Glauber coherent state. Amplitude–squeezed light exhibits antibunching, where the photon arrivals are more regularly spaced than the purely random arrivals of a coherent state. It also exhibits a distribution of photon number that is narrower than the Poisson distribution of a coherent state. These properties were first observed in experiments on resonance fluorescence.[1,2]

Amplitude squeezing has also been observed in a semiconductor diode laser with negative feedback.[3,4] In these experiments, the laser output was detected using a photodiode, and the, photodiode output was inverted, amplified, and applied to the laser drive current in a negative feedback loop. Because the laser output is detected to produce squeezing, squeezed light is not available for other applications. If a beamsplitter is used to deflect a portion of the light before detection, the deflected portion will actually exhibit greater fluctuations than if feedback is not used. The first solution is to replace the feedback by a high-impedance constant-current laser power supply. This technique has been demonstrated to produce squeezing at high frequencies of 350 – 450 MHz, (Ref. 5) but it is difficult to implement at low frequencies because of the so–called "1/f" noise processes.

The second potential solution to the problem of using the squeezed light generated by negative feedback is to place an absorber between the laser and the detector and use the squeezed light for measurement of absorption. This possibility is analyzed in this paper and found to have no advantage over using coherent light.

The third solution to the problem is to use a quantum nondemolition measurement of photon number in the feedback loop.[6–8] In this case most of the light is available for making absorption measurements. In the most commonly discussed approach, the laser output is passed through a nonlinear Kerr medium. Such a measurement has been demonstrated experimentally,[9] although the sensitivity which was achieved in this first experiment was insufficient.

In this paper, we consider the use of amplitude-squeezed light for making very sensitive absorption measurements. The application is absorption spectroscopy of very weak lines. We show that the sensitivity is limited by the quality of the detector and can be much better than the usual shot–noise limit.

AMPLITUDE SQUEEZING

We analyze laser operation based on the geometry of Fig. 1. Light from the laser is detected by the photodiode. The photocurrent consists of a deterministic component and a random noise contribution. In quantum optics, the noise contribution is related to statistical properties of the light.[10] In the semiclassical interpretation that will be used here, the noise is added within the photodiode; it is the shot noise in the photodiode current. The photodiode output is inverted and combined with a bias current. The amplified result is used as the laser drive current.

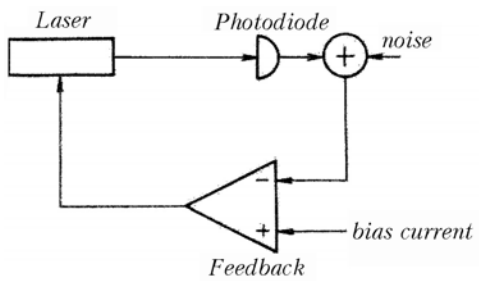

FIG. 1. Schematic diagram of feedback–generated optical squeezing configuration.

The number of electrons produced by the photodiode within the Nyquist interval of the system can be expressed as

$$n_d = \eta_d - n_n, \tag{1}$$

where η_d is the quantum efficiency of the detector, m_l is the average number of photons in the Nyquist interval, and n_n is the number of noise electrons in a particular interval. The quantity n_n is a zero–mean random variable with variance given by

$$\sigma_n^2 = \langle n_d \rangle, \tag{2}$$

which implies the ideal laser case.

The number of laser photons is given by

$$m_l = \eta_l\, G(n_b - n_d), \tag{3}$$

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where η_1 is the differential quantum efficiency of the laser, G is the amplifier gain, and n_b is the number of electrons per Nyquist interval in the bias current.

Substituting Eq. (3) into Eq. (1) and solving for n_d produces

$$n_d = \frac{\eta_1\eta_d G n_b + n_n}{1 + \eta_1\eta_d G}, \tag{4}$$

The mean value is easily found to be

$$\langle n_d\rangle = \frac{\eta_1\eta_d G n_b}{1 + \eta_1\eta_d G}. \tag{5}$$

The variance is given by

$$\sigma^2_{n_d} = \frac{\langle n_d\rangle}{(1 + \eta_1\eta_d G)^2}. \tag{6}$$

which is less than the corresponding Poisson variance by the factor $(1 + \eta_1 \eta_d G)^{-2}$.

From Eq. (6), it is clear that we would like to operate with very large gain. In this case, we have $\langle n_d\rangle = n_b$ and $\sigma_n^2 = n_b/(\eta_1^2\eta_d^2 G^2)$. The noise can be very small. The corresponding relation for the open–loop, coherent–state cases are $\langle n_d\rangle = \eta_1\eta_d G n_b$ and $\sigma_{n_d}^2 = \eta_1\eta_d G n_b$.

INTRA–LOOP MEASUREMENT

We consider the case in which a weakly absorbing material is placed between the laser and the detector. The photodetector output becomes

$$n_d = \frac{\eta_1\eta_d G e^{-\delta} n_b + n_n}{1 + \eta_1\eta_d G e^{-\delta}}, \tag{7}$$

where δ is the absorption of the material. We are interested in making a very sensitive measurement of very small differences in absorption. We can therefore assume, δ is small and expand Eq. (7) in a Taylor series. Keeping only the first two terms

$$n_d \approx \frac{\eta_1\eta_d G n_b + n_n}{1 + \eta_1\eta_d G}\left[1 - \frac{\delta}{1 - \eta_1\eta_d G}\right]. \tag{8}$$

To construct an estimator of the absorption from the measured photocount, we set n_n to be equal zero in Eq. (8) and solve for δ. This suggests that we estimate the absorption by

$$\hat{\delta} = (1 + \eta_1\eta_d G)\left[1 - \frac{1 + \eta_1\eta_d G}{\eta_1\eta_d G n_b}\, n_d\right]. \tag{9}$$

The mean value of this estimator is the desired value of δ. The variance is

$$\sigma^2_{\hat{\delta}} = \frac{1 + \eta_1\eta_d G}{\eta_1\eta_d G n_b} + \frac{2\eta_1\eta_d G - 1}{\eta_1\eta_d G n_b}\,\delta, \tag{10}$$

which is approximately $1/n_b$ for large gain and small absorption.

For comparison, consider the open–loop case. Here the appropriate estimator is

$$\hat{\delta} = 1 - \frac{n_d}{\eta_1\eta_d G n_b}. \tag{11}$$

The variance in this case is

$$\sigma^2_{\hat{\delta}} = \frac{1 - \delta}{\eta_1\eta_d G n_b}, \tag{12}$$

which approaches for small values of δ. At first glance it appears that the open–loop configuration can achieve unlimited performance by using large enough gain. In practice, the laser power increases with gain in this case and limits the gain that can be used. For a more reasonable comparison of these two cases it should be noted that $\eta_1\eta_d G n_b$ is $\langle n_d\rangle$ in the open loop case. In the closed loop case with high gain $\langle n_d\rangle \approx n_b$. Therefore, the open–loop and closed–loop configurations are equivalent with

$$\sigma^2_{\hat{\delta}} \approx 1/\langle n_d\rangle \tag{13}$$

in either case.

EXTERNAL MEASUREMENT

To make a measurement of the absorption of a material placed outside of the loop, we must consider a quantum nondemolition measurement. For an ideal nondemolition measurement, the statistics of the output photons are found in the same way as the statistics of the photoelectrons in Section II. The mean value is

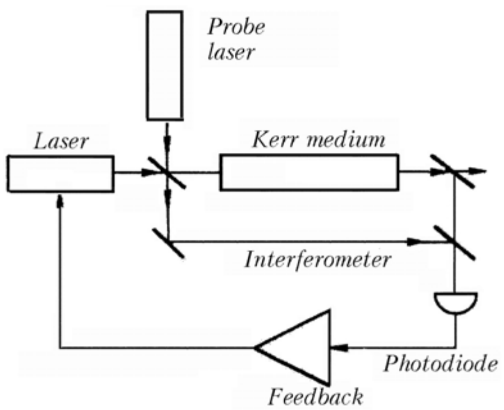

FIG. 2. Schematic diagram of quantum nondemolition measurement using nonlinear Kerr effect.

$$\langle m_0\rangle = \frac{\eta_1\eta_q G n_b}{1 + \eta_1\eta_q G}, \tag{14}$$

where η_q is the quantum efficiency of the nondemolition measurement.

As an example, consider the Kerr–effect measurement shown in Fig. 2. The laser light is transmitted through a nonlinear Kerr medium. The refractive index of the medium is changed slightly by the laser power. A probe laser is used to measure the refractive index change interferometrically. The probe laser is operated at a different wavelength or different polarization so light from the primary laser passes through the beamsplitters and the Kerr medium with small loss. The interferometer output is detected, the photon number is inferred, and the result is amplified and used as the feedback signal. In this case, η_q would be the product of the following factors:

1) the transmission coefficient of the first beamsplitter and the front of the Kerr medium,

2) the change in the refractive index of the medium per signal photon,

3) the change in phase of the probe–laser light per unit change in refractive index,

4) the change in probe power out of the interferometer per unit change in phase, and

5) the interferometer detector quantum efficiency. Reflections of the signal laser from the second beamsplitter and the rear of the Kerr medium do not affect the measurement and are not included in η_q.

The variance of photon number is

$$\sigma_{m_0}^2 = \frac{\langle m_0 \rangle}{(1 + \eta_1 \eta_q G)^2} \tag{15}$$

by analogy with Eq. (6).

For a detector with less than unity quantum efficiency, the number of photoelectrons in any interval will be a binomial random variable with the number of incident photons as the number of samples and the quantum efficiency as the probability of a successful sample. The overall statistics can be found by averaging the binomial conditional statistics over the distribution of the number of incident photons. Thus

$$\eta_d \langle m_0 \rangle = \frac{\eta_1 \eta_q \eta_d G n_b}{1 + \eta_1 \eta_q G}, \tag{16}$$

and

$$\sigma_{n_d}^2 = \eta_d (1 - \eta_d) \langle m_0 \rangle + \eta_d^2 \sigma_{m_0}^2 =$$

$$= \eta_d (1 - \eta_d) \frac{\eta_1 \eta_q G n_b}{1 + \eta_1 \eta_q G} + \eta_d^2 \frac{\eta_1 \eta_q G n_b}{(1 + \eta_1 \eta_q G)^3}. \tag{17}$$

For large values of G, Eq. (17) reduces to

$$\sigma_{n_d}^2 = (1 - \eta_d) \langle n_d \rangle \tag{18}$$

which approaches the Poisson value of $\langle n_d \rangle$ for small values of quantum efficiency, but can become very small for near unit quantum efficiency. Here it should be noted that reflections from the rear of the Kerr medium must be included when calculating the effective value of η_d and can reduce squeezing significantly.

If we place an absorber in front of the detector, the photoelectron number becomes

$$n_d = \eta_d e^{-\delta} m_0 . \tag{19}$$

For small absorption, this suggests the estimator

$$\hat{\delta} = 1 - \frac{1 + \eta_1 \eta_q G}{\eta_1 \eta_q \eta_d G n_b} n_d . \tag{20}$$

The variance of this estimator is

$$\sigma_{\hat{\delta}}^2 = \frac{1 + \eta_1 \eta_q G}{\eta_1 \eta_q \eta_d G n_b} (1 - \eta_d + \delta) + \frac{1}{\eta_1 \eta_q G (1 + \eta_1 \eta_q G) n_b} \tag{21}$$

which approaches

$$\sigma_{\hat{\delta}}^2 = (1 - \eta_d + \delta) / \eta_d n_b \tag{22}$$

for large gain and small absorption.

From Eq. (22) we can see that squeezing reduces the uncertainty in the absorption measurement by a factor of $(1 - \eta_d + \delta)^{1/2}$ for small values of δ. Comparison of Eq. (22) with Eq. (13) shows that for the limit of small δ, the distribution width is $(1 - \eta_d)^{1/2}$ narrower than the distribution width obtained with an ideal laser. For high–efficiency photodetectors[11] $\eta_d = 0.9$, which results in three times the sensitivity.

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