Electromagnetically induced transparency with quantized fields in optocavity mechanics

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We report electromagnetically induced transparency (EIT) using quantized fields in optomechanical systems. The weak probe field is a narrowband squeezed field. We present a homodyne detection of EIT in the output quantum field. We find that the EIT dip exists even though the photon number in the squeezed vacuum is at the single-photon level. The EIT with quantized fields can be seen even at temperatures on the order of 100 mK, thus paving the way for using optomechanical systems as memory elements.

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I. INTRODUCTION

The interaction of a nanomechanical system via radiation pressure [1,2] is like a three-wave interaction in nonlinear optics [3]. This interaction can lead to processes like upconversion; for example, a photon of frequency ω_c can be converted into a photon of frequency $\omega_p = \omega_c + \omega_m$, where ω_m is the frequency of the mechanical oscillator. Such upconversion processes have been useful in cooling nanomechanical systems [4–7]. In a previous article [8], we showed how such upconversion processes can lead to electromagnetically induced transparency (EIT) in optomechanical systems. The EIT in such systems turned out to share many of the features of EIT in atomic vapors. The EIT in optomechanical systems has been seen experimentally [9–11]. Traditionally, almost all of the EIT experiments in atomic systems and other systems have been done with coherent pump and probe fields [12–14]. Akamatsu *et al.* [15] did the very first experiment on EIT using squeezed light in atomic vapors. They essentially reported that squeezing of the probe is not degraded much by the quantum noise of the medium under EIT conditions. Subsequently, a number of other experiments [16,17] on EIT using quantized fields were reported. The EIT with quantized fields is very significant in the storage of fields at the single-photon level [18–21].

In this paper, we examine EIT in optomechanical systems using quantized fields. In optomechanical systems, noise is added by both the resonator and the mechanical system. We find the conditions when the perfect EIT of the quantized field results. We study how the temperature of the mechanical system can degrade EIT. We present detailed results for the designs of nanomechanical systems as used in Refs. [9,22]. We find that certain designs of nanomechanical systems are good even at temperatures on the order of 100 mK. Thus, such systems would be quite useful as optical memories at the single-photon level. The results that we present can be extended to the reactive case [23–25].

The organization of the paper is as follows. In Sec. II, we describe the model, derive the equations of motion for the system, and obtain the steady-state mean values. In Sec. III, we show how to detect the EIT with quantized fields, and we present a homodyne detection and obtain the relevant spectrum. In Sec. IV, we discuss the impact of the coupling field on the homodyne spectrum of the output field and show the existence of the EIT in the homodyne spectrum of the quantized field at the output.

II. MODEL

The model that we are going to consider has been discussed in detail previously [26,27] and is sketched in Fig. 1. The cavity consists of a fixed mirror and a movable mirror separated by a distance L. The fixed mirror is partially transmitting, while the movable mirror is 100% reflecting. The cavity is driven by a strong coupling field at frequency ω_c. A quantized weak probe field in a squeezed vacuum state at frequency ω_p is injected into the cavity through the fixed mirror. The movable mirror interacts with the cavity field through radiation pressure. The movable mirror is modeled as a harmonic oscillator with mass m, frequency ω_m, and decay rate γ_m. Moreover, the movable mirror and its environment are in thermal equilibrium at a low temperature T.

In such a system, the coupling between the movable mirror and the cavity field is dispersive, so the frequency $\omega_0(q)$ of the cavity field depends on the displacement q of the movable mirror: $\omega_0(q) = n\pi c/(L+q)$, where c is the light speed in vacuum and n is the mode number in the cavity. For $q \ll L$, we can expand $\omega_0(q)$ to the first order of q; thus, we have $\omega_0(q) \approx \omega_0(0) + \frac{\partial\omega_0(q)}{\partial q}q \approx \omega_0 - \frac{\omega_0}{L}q$, where we write $\omega_0(0)$ as ω_0.

Let c ($c^\dagger$) be the annihilation (creation) operators for the cavity field and Q and P be the dimensionless operators for the position and momentum of the movable mirror with $Q = \sqrt{\frac{2m\omega_m}{\hbar}}q$ and $P = \sqrt{\frac{2}{m\hbar\omega_m}}p$. Note that the commutation relation for Q and P is $[Q,P] = 2i$. In a frame rotating at the frequency ω_c of the coupling field, the Hamiltonian for the system is

$$H = \hbar(\omega_0 - \omega_c)c^\dagger c - \hbar g c^\dagger c Q + \frac{\hbar\omega_m}{4}(Q^2 + P^2) + i\hbar\varepsilon(c^\dagger - c). \tag{1}$$

In the above equation, the parameter $g = (\omega_c/L)\sqrt{\hbar/(2m\omega_m)}$ is the coupling strength between the cavity field and the movable mirror, where we assume $\omega_0 \simeq \omega_c$. The parameter ε is the real amplitude of the coupling field, depending on its power $\wp$ by $\varepsilon = \sqrt{\frac{2\kappa\wp}{\hbar\omega_c}}$, where κ is the photon loss rate due to the transmission of the fixed mirror.

The time evolution of the total system is obtained from the Hamiltonian Eq. (1) by deriving the Heisenberg equations of

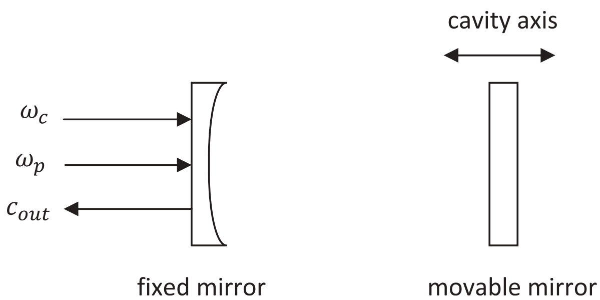

FIG. 1. Sketch of the studied system. A coherent coupling field at frequency ω_c and a squeezed vacuum at frequency ω_p enter the cavity through the partially transmitting mirror.

motion and adding the damping and noise terms. The basic equations are given by

$$\begin{aligned} \dot{Q} &= \omega_m P, \\ \dot{P} &= 2gn_c - \omega_m Q - \gamma_m P + \xi, \\ \dot{c} &= i(\omega_c - \omega_0 + gQ)c + \varepsilon - \kappa c + \sqrt{2\kappa}c_{\rm in}, \\ \dot{c}^\dagger &= -i(\omega_c - \omega_0 + gQ)c^\dagger + \varepsilon - \kappa c^\dagger + \sqrt{2\kappa}c_{\rm in}^\dagger. \end{aligned} \tag{2}$$

Here, we have introduced the thermal Langevin force ξ with a vanishing mean value, resulting from the coupling of the movable mirror to the environment. The Langevin force ξ has the correlation function in the frequency domain

$$\langle \xi(\omega)\xi(\Omega)\rangle = 4\pi\gamma_m \frac{\omega}{\omega_m}\left[1 + \coth\left(\frac{\hbar\omega}{2k_B T}\right)\right]\delta(\omega+\Omega), \tag{3}$$

where k_B is the Boltzmann constant. Throughout this paper, the following Fourier relations are used:

$$\begin{aligned} f(t) &= \frac{1}{2\pi}\int_{-\infty}^{+\infty} f(\omega)e^{-i\omega t}d\omega, \\ f^\dagger(t) &= \frac{1}{2\pi}\int_{-\infty}^{+\infty} f^\dagger(-\omega)e^{-i\omega t}d\omega, \end{aligned} \tag{4}$$

where $f^\dagger(-\omega) = [f(-\omega)]^\dagger$. $c_{\rm in}$ represents the input quantum field, which is centered around the frequency $\omega_p = \omega_c + \omega_m$ with a finite bandwidth Γ. The quantized field has the following nonvanishing correlation functions:

$$\begin{aligned} \langle c_{\rm in}(\omega)c_{\rm in}(\Omega)\rangle &= 2\pi\frac{M\Gamma^2}{\Gamma^2+(\omega-\omega_m)^2}\delta(\omega+\Omega-2\omega_m), \\ \langle c_{\rm in}(\omega)c_{\rm in}^\dagger(-\Omega)\rangle &= 2\pi\left[\frac{N\Gamma^2}{\Gamma^2+(\omega-\omega_m)^2}+1\right]\delta(\omega+\Omega), \end{aligned} \tag{5}$$

where N is the photon number in the squeezed vacuum and $M = \sqrt{N(N+1)}$. The antinormally ordered term has a broadband contribution coming from vacuum noise. Note that by setting $M = 0$ we would obtain a standard phase-independent quantum field with a mean number of photons $\frac{N\Gamma^2}{\Gamma^2+(\omega-\omega_m)^2}$ around the frequency $\omega = \omega_m$.

The mean values at steady state can be obtained from Eq. (2) by setting all of the time derivatives to zero. These are found to be

$$P_s = 0,\quad Q_s = \frac{2g|c_s|^2}{\omega_m},\quad c_s = \frac{\varepsilon}{\kappa + i\Delta}, \tag{6}$$

where

$$\Delta = \omega_0 - \omega_c - gQ_s \tag{7}$$

is the effective cavity detuning.

III. THE OUTPUT FIELD AND ITS MEASUREMENT

The output field is a quantum field; it contains many Fourier components. Since the quantized input field is centered around $\omega_p = \omega_c + \omega_m$, the interesting component of the output field is near the probe frequency ω_p, so we mix the output field $\tilde{c}_{\rm out}(t)$ with a strong local field $c_{\rm lo}(t)$ centered around the probe frequency ω_p at a 50:50 beam splitter, as shown in Fig. 2. In a frame rotating at the frequency ω_c, $c_{\rm lo}(t) = c_{\rm lo}e^{-i\delta_0 t}$, where $\delta_0 = \omega_p - \omega_c$. The difference between the output signals from the two photodetectors is sent to the spectrum analyzer, and the output signal from the spectrum analyzer depends on the phase of $c_{\rm lo}$. If $c_{\rm lo}$ is real, then the homodyne spectrum $X(\omega)$ of the output field measured by the spectrum analyzer is given by

$$\begin{aligned} &\langle [c_{\rm lo}^*(t)\tilde{c}_{\rm out}(t) + {\rm c.c.}][c_{\rm lo}^*(t')\tilde{c}_{\rm out}(t') + {\rm c.c.}]\rangle \\ &\quad = \frac{c_{\rm lo}^2}{2\pi}\int d\omega e^{-i\omega(t-t')}X(\omega). \end{aligned} \tag{8}$$

Thus, in our investigations of EIT with quantized fields, $X(\omega)$ is the quantity of interest.

In order to study the EIT effect in the homodyne spectrum $X(\omega)$ of the output field, we will calculate the fluctuations of the output field. The steady-state part would not contribute as it is at the frequency of the coupling field. We assume that the photon number in the cavity is large enough so that each

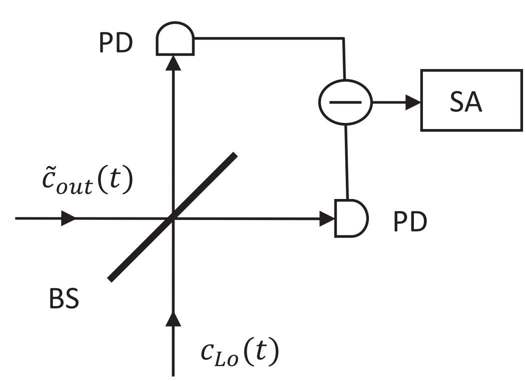

FIG. 2. Sketch of the measurement of the output field. The output field $\tilde{c}_{\rm out}(t)$ is mixed with a strong local field $c_{\rm lo}(t)$ centered around the probe frequency ω_p at a beam splitter, where $\tilde{c}_{\rm out}(t)$ is defined as the sum of the output field $c_{\rm out}(t)$ from the cavity and the input quantized field $c_{\rm in}(t)$. BS, 50:50 beam splitter; PD, photodetector; SA, spectrum analyzer.

operator can be written as a linear sum of the steady-state mean value and a small fluctuation, which yields

$$Q = Q_s + \delta Q, \quad P = P_s + \delta P, \quad c = c_s + \delta c, \tag{9}$$

where δQ, δP, and δc are the small fluctuations around the steady state. By substituting Eq. (9) into Eq. (2), one can arrive at the linearized equations for the fluctuation operators. Further, we transform the linearized equations into the frequency domain by Eq. (4) and solve it; we can obtain the fluctuations $\delta c(\omega)$ of the cavity field. Then, using the input-output relation $c_{\text{out}}(\omega) = \sqrt{2\kappa}c(\omega) - c_{\text{in}}(\omega)$, we can find the fluctuations $\delta c_{\text{out}}(\omega)$ of the output field. For the purpose of Fig. 2, we define the output field as $\tilde{c}_{\text{out}}(\omega) = c_{\text{out}}(\omega) + c_{\text{in}}(\omega)$; then we find the result

$$\delta\tilde{c}_{\text{out}}(\omega) = V(\omega)\xi(\omega) + E(\omega)c_{\text{in}}(\omega) + F(\omega)c_{\text{in}}^{\dagger}(-\omega), \tag{10}$$

in which

$$\begin{aligned}
V(\omega) &= \frac{\sqrt{2\kappa}\, g c_s \omega_m i}{d(\omega)}[\kappa - i(\omega + \Delta)],\\
E(\omega) &= \frac{2\kappa}{d(\omega)}\{2ig^2|c_s|^2\omega_m + (\omega_m^2 - \omega^2 - i\gamma_m\omega)\\
&\quad \times [\kappa - i(\omega+\Delta)]\},\\
F(\omega) &= \frac{4\kappa}{d(\omega)}\omega_m g^2 c_s^2 i,
\end{aligned} \tag{11}$$

where

$$\begin{aligned}
d(\omega) &= -4\omega_m\Delta g^2|c_s|^2 + (\omega_m^2 - \omega^2 - i\gamma_m\omega)\\
&\quad \times [(\kappa - i\omega)^2 + \Delta^2].
\end{aligned} \tag{12}$$

The first term on the right-hand side of Eq. (10) refers to the contribution of the thermal noise of the movable mirror, and the other two terms represent the contribution of the squeezed vacuum. To illustrate the meaning of the last two terms, let the squeezed vacuum be a single mode, i.e., $c_{\text{in}}(t) = Ce^{-i(\omega_p-\omega_c)t}$; then $c_{\text{in}}(\omega) = 2\pi C\delta(\omega - \omega_p + \omega_c)$ and $c_{\text{in}}^{\dagger}(-\omega) = 2\pi C^{\dagger}\delta(\omega + \omega_p - \omega_c)$. Thus, the fluctuations of the output field $\delta\tilde{c}_{\text{out}}(t) = \frac{1}{2\pi}\int_{-\infty}^{+\infty} V(\omega)\xi(\omega)e^{-i\omega t}d\omega + CE(\omega_p - \omega_c)e^{-i(\omega_p-\omega_c)t} + C^{\dagger}F(\omega_c - \omega_p)e^{-i(\omega_c-\omega_p)t}$. Therefore, $E(\omega_p - \omega_c)$ is the component at the probe frequency ω_p, which in the rotating frame is $\omega_p - \omega_c$, and $F(\omega_c - \omega_p)$ is the component at the new frequency $2\omega_c - \omega_p$, which in the rotating frame is $\omega_c - \omega_p$, due to the nonlinear interaction between the movable mirror and the cavity field.

By the aid of the correlation functions of the noise operators $c_{\text{in}}(\omega)$ and $\xi(\omega)$ and neglecting fast oscillating terms at frequency $\pm 2\omega_m$, we obtain the homodyne spectrum $X(\omega)$ of the output field as measured by the setup of Fig. 2,

$$\begin{aligned}
X(\omega) &= E(\omega+\omega_m)E(-\omega+\omega_m)\frac{M\Gamma^2}{\Gamma^2+\omega^2} + |E(\omega+\omega_m)|^2\\
&\quad \times \frac{N\Gamma^2}{\Gamma^2+\omega^2} + E^*(-\omega+\omega_m)E^*(\omega+\omega_m)\frac{M\Gamma^2}{\Gamma^2+\omega^2}\\
&\quad + |E(-\omega+\omega_m)|^2\frac{N\Gamma^2}{\Gamma^2+\omega^2} + |E(\omega+\omega_m)|^2\\
&\quad + |F(-\omega+\omega_m)|^2 + |V(\omega+\omega_m)|^2 2\gamma_m\frac{\omega+\omega_m}{\omega_m}\\
&\quad \times \left\{1 + \coth\left[\frac{\hbar(\omega+\omega_m)}{2k_BT}\right]\right\} + |V(-\omega+\omega_m)|^2 2\gamma_m\\
&\quad \times \frac{\omega-\omega_m}{\omega_m}\left\{1 + \coth\left[\frac{\hbar(\omega-\omega_m)}{2k_BT}\right]\right\},
\end{aligned} \tag{13}$$

where the first four terms in Eq. (13) originate from the squeezed vacuum, the next two terms not involving N and M are the contributions of the spontaneous emission of the input vacuum noise, and the last two terms result from the thermal noise of the movable mirror.

IV. EIT IN THE HOMODYNE SPECTRUM OF THE OUTPUT QUANTIZED FIELD

After having derived the homodyne spectrum of the output field, we next examine it numerically to explore the EIT phenomenon in the homodyne spectrum of the output field. Since the original Eqs. (2) are nonlinear, these can have instabilities. Thus, in the following, we work in the stable regime of the system. We first examine the frequency at which we expect transparency. This is $\omega = 0$. For $N \approx M$,

$$\begin{aligned}
X(0) &= N[E(\omega_m) + E^*(\omega_m)]^2 + |E(\omega_m)|^2 + |F(\omega_m)|^2\\
&\quad + 4|V(\omega_m)|^2\gamma_m\coth\left[\frac{\hbar\omega_m}{2k_BT}\right].
\end{aligned} \tag{14}$$

We use the parameters from the experimental paper [9] focusing on the EIT in the optomechanical system: the wavelength of the coupling field $\lambda = 2\pi c/\omega_c = 775$ nm, the coupling constant $g = 2\pi \times 12$ GHz/nm$\sqrt{\hbar/(2m\omega_m)}$, the mass of the movable mirror $m = 20$ ng, the frequency of the movable mirror $\omega_m = 2\pi \times 51.8$ MHz, the cavity decay rate $\kappa = 2\pi \times 15$ MHz, $\kappa/\omega_m = 0.289$, the mechanical damping rate $\gamma_m = 2\pi \times 41$ kHz, and the mechanical quality factor $Q' = \omega_m/\gamma_m = 1263$. In addition, we choose the linewidth of the squeezed vacuum $\Gamma = 2\kappa$ and consider the resonant case $\Delta = \omega_m$.

For $N = 10$ and $M = \sqrt{N(N+1)} \approx 10$, $\wp = 20$ mW, and $T = 20$ mK, the first term in Eq. (14), which is the contribution of the squeezed vacuum, is about 6.5×10^{-4}, the sum of the second and third terms in Eq. (14), which are the contributions of the input vacuum noise, is about 0.16, and the last term arising from the thermal noise of the movable mirror is about 0.14. The contribution of the input quantum field in principle can be obtained by doing the experiment with and without the quantized field and by subtracting the data, i.e., by studying $X(0) - X(0)|_{N=0}$. The squeezed field part in a sense exhibits perfect EIT. If $M = 0$, i.e., the input quantized field is phase insensitive, then such a field leads to a term $2N|E(\omega_m)|^2$, which is equal to 1.6 for the above-mentioned parameters, and hence there is no perfect EIT. The squeezed field changes $2N|E(\omega_m)|^2$ to $N[E(\omega_m) + E^*(\omega_m)]^2$, and for the above parameters, the number changes from 1.6 to 6.5×10^{-4}.

For $N = 5$, $M = \sqrt{N(N+1)}$ and 0, and $T = 20$ mK, we plot the homodyne spectrum $X(\omega)$ of the output field as a function of the normalized frequency ω/ω_m in the absence (dotted curve) and presence (solid, dot-dashed, and dashed

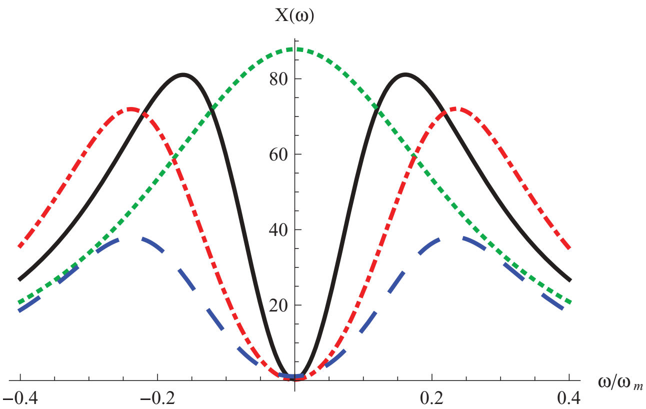

FIG. 3. (Color online) Homodyne spectrum $X(\omega)$ as a function of ω/ω_m for $N = 5$ in the absence (dotted curve) and the presence (solid, dot-dashed, and dashed curves) of the coupling field for the temperature of the environment $T = 20$ mK. The solid curve is for $\wp = 10$ mW and $M = \sqrt{N(N+1)}$, the dot-dashed curve is for $\wp = 20$ mW and $M = \sqrt{N(N+1)}$, and the dashed curve is for $\wp = 20$ mW and $M = 0$.

curves) of the coupling field in Fig. 3. First, let us look at the case that the input quantum field is phase dependent [$M = \sqrt{N(N+1)}$]. In the absence of the coupling field, one can note that the homodyne spectrum of the output field has a Lorentzian line shape. However, in the presence of the coupling field at different power levels, the solid curve [$\wp = 10$ mW and $M = \sqrt{N(N+1)}$] and the dot-dashed curve [$\wp = 20$ mW and $M = \sqrt{N(N+1)}$] exhibit the EIT dip, which is the result of the destructive interference between the squeezed vacuum and the scattering quantum field at the probe frequency ω_p generated by the interaction of the coupling field with the movable mirror. For $\wp = 20$ mW and $M = \sqrt{N(N+1)}$, the minimum value of $X(\omega)$ is about 0.22. Moreover, the linewidth of the dip for $\wp = 20$ mW is larger than that for $\wp = 10$ mW due to power broadening. Generally, the EIT dip has a contribution to its width that is proportional to the power of the coupling field. We indeed find that the width for $\wp = 20$ mW is $0.26\omega_m$, which is about twice the width for $\wp = 10$ mW. If the input quantum field is phase independent ($M = 0$) (the dashed curve), then we can see that the maximum value of $X(\omega)$ for $\wp = 20$ mW and $M = 0$ is about half that for $\wp = 20$ mW and $M = \sqrt{N(N+1)}$.

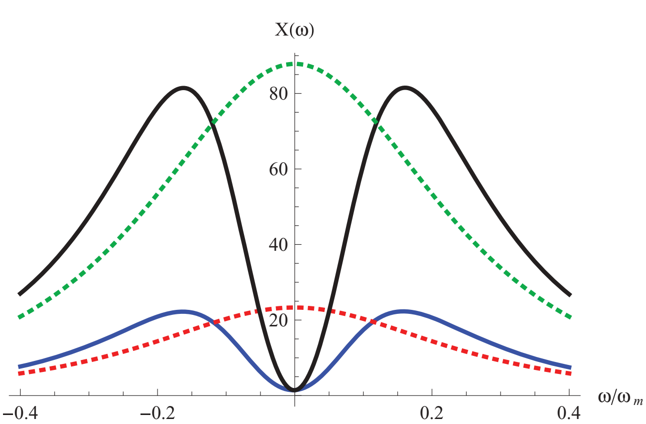

FIG. 4. (Color online) Homodyne spectrum $X(\omega)$ as a function of ω/ω_m for different values of the parameter N and $M = \sqrt{N(N+1)}$ in the absence (dotted curves) and the presence (solid curves) of a coupling field with power $\wp = 10$ mW and temperature of the environment $T = 100$ mK. The upper two curves are for $N = 5$, and the lower two curves are for $N = 1$.

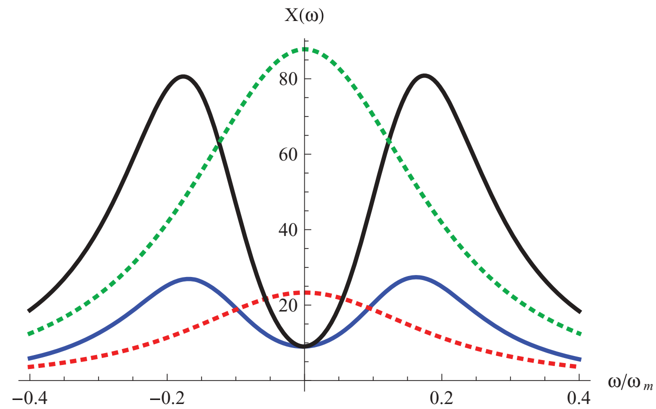

FIG. 5. (Color online) As in Fig. 4 but now the parameters used are from Ref. [22].

Next, we increase the temperature to 100 mK. Figure 4 displays the homodyne spectrum $X(\omega)$ of the output field against the normalized frequency ω/ω_m in the absence (dotted curves) and presence (solid curves) of the coupling field for $N = 1$ and 5 and $M = \sqrt{N(N+1)}$. In the presence of the coupling field ($\wp = 10$ mW), it is seen that the EIT dip still appears in the homodyne spectrum of the output field for $N = 1$ and 5. Note that the two dips almost have the same minimum values (about 1.43) and the same linewidth (about $0.15\omega_m$). Hence, the temperature of the environment is not detrimental to the EIT behavior.

The effects discussed above occur under a wide range of parameters. We demonstrate this by using the experimental parameters [22] $\lambda = 2\pi c/\omega_c = 1064$ nm, $L = 25$ mm, $g \approx 2\pi \times 11.28$ MHz/nm$\sqrt{\hbar/(2m\omega_m)}$, $m = 145$ ng, $\omega_m = 2\pi \times 947$ kHz, $\kappa = 2\pi \times 215$ kHz, $\kappa/\omega_m = 0.227$, $\gamma_m = 2\pi \times 141$ Hz, and $Q' = \omega_m/\gamma_m = 6700$. The values for the parameters T, $\wp$, N, M, Γ, and Δ are the same as those in Fig. 4. Shown in Fig. 5 is the homodyne spectrum $X(\omega)$ of the output field as the normalized frequency ω/ω_m is varied for $T = 100$ mK and $\wp = 0$ and 10 mW. Note that the EIT exists for $N = 1$ and 5 in the presence of the coupling field. The linewidth of the dip for $N = 5$ is about $0.2\omega_m$ and as expected gets broadened due to power. We have further studied the effect of temperature, and we find that there is a rather weak dependence of the EIT curves on temperature. Therefore, current optomechanical designs can be used to realize quantum optical memory at the single-photon level. This can be demonstrated using the numerical simulations and following the standard procedure as in Refs. [19–21]. One has to modulate the squeezed vacuum field $c_{\rm in}$ so that it is a pulse field and uses, say, a super-Gaussian for the coupling field. The super-Gaussian enables one to conveniently switch on and off the coupling field [28].

V. CONCLUSIONS

In conclusion, we have demonstrated EIT using quantum fields in optomechanical systems under a wide range of conditions. For squeezed quantum fields, we obtained the perfect EIT. The EIT gets degraded in phase-insensitive quantum fields. We have shown that even temperature is not critical for observations of EIT. The results can be generalized to optomechanical systems working on the reactive coupling [23–25]. Our work suggests that optomechanical systems could be used as elements for quantum memory, but explicit demonstration will be given elsewhere.

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