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Electromagnetically Induced Transparency with Quantized Fields in Optocavity Mechanics

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We report electromagnetically induced transparency using quantized fields in optomechanical systems. The weak probe field is a narrow band 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 of the order of 100 mK paving the way for using optomechanical systems as memory elements.

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I. INTRODUCTION

The interaction of a nano mechanical system via the 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 the nano mechanical systems [4–7]. In a previous paper [8] we showed how such upconversion processes can lead to electromagnetically induced transparency 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 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 the squeezed light in atomic vapors. They essentially reported that squeezing of the probe is not graded much by the quantum noise of the medium under EIT conditions. Following this a number of other experiments [16, 17] on EIT using quantized fields were reported. The EIT with quantized fields is very significant in storage of fields at single photon level [18–21].

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

The organization of the paper is as follows. In Sec. II, we describe the model, derive the equations of the motion for the system, and obtain the steady-state mean values.

In Sec. III, we show how to detect EIT with quantized fields, 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.

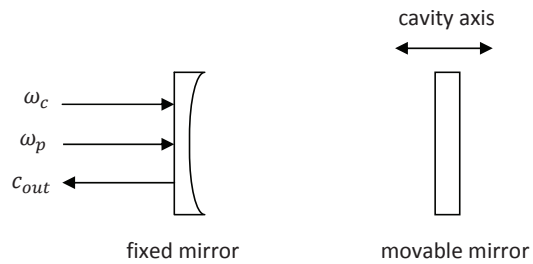


FIG. 1: Sketch of the studied system. A coherent coupling field at frequency ω_c and a squeezed vacuum at frequency ω_p enter the cavity through the partially transmitting mirror.

II. MODEL

The model we are going to consider has been discussed in detail before [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 partial 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 the 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 that the fre-

quency $\omega_0(q)$ 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, 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

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

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 Hamiltonian (1) by deriving the Heisenberg equations of motion and adding the damping and noise terms. The basic equations are given by

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

Here we have introduced the thermal Langevin force ξ with 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:

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

where K_B is the Boltzmann constant. Throughout this paper the following Fourier relations are used

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

where $f^\dagger(-\omega) = [f(-\omega)]^\dagger$. The c_{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,

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

where N is the photon number in the squeezed vacuum, and $M = \sqrt{N(N+1)}$. The antinormally ordered term has a broad band contribution coming from vacuum noise. Note that by setting $M = 0$, we would obtain standard phase independent quantum field with 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 time derivatives to zero. These are found to be

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

where

$$\Delta = \omega_0 - \omega_c - gQ_s \quad (7)$$

is the effective cavity detuning.

III. THE OUTPUT FIELD AND ITS MEASUREMENT

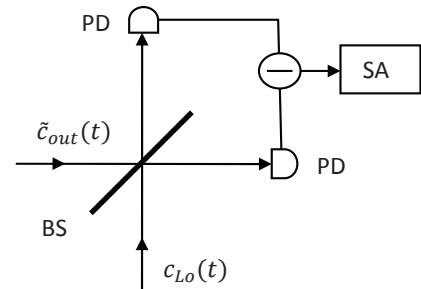


FIG. 2: Sketch of the measurement of the output field. The output field $\tilde{c}_{out}(t)$ is mixed with a strong local field $c_{Lo}(t)$ centered around the probe frequency ω_p at a beam splitter, where $\tilde{c}_{out}(t)$ is defined as the sum of the output field $c_{out}(t)$ from the cavity and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

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 $\tilde{c}_{out}(t)$ with a strong local field $c_{Lo}(t)$ 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 , $c_{Lo}(t) = c_{Lo}e^{-i\delta_0 t}$, 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_{Lo} . If c_{Lo} is real, the homodyne spectrum $X(\omega)$ of the output field measured by the spectrum analyzer is given by

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

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 the photon number in the cavity is large enough so that each operator can be written as a linear sum of the steady-state mean value and a small fluctuation, which yields

$$Q = Q_s + \delta Q, \quad P = P_s + \delta P, \quad c = c_s + \delta c, \quad (9)$$

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 $c_{out}(\omega) = \sqrt{2\kappa}c(\omega) - c_{in}(\omega)$, we can find the fluctuations $\delta c_{out}(\omega)$ of the output field. For the purpose of Fig. (2) we define the output field as $\tilde{c}_{out}(\omega) = c_{out}(\omega) + c_{in}(\omega)$, then we find the result

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

in which

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

where

$$d(\omega) = -4\omega_m \Delta g^2 |c_s|^2 + (\omega_m^2 - \omega^2 - i\gamma_m \omega) [(\kappa - i\omega)^2 + \Delta^2]. \quad (12)$$

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. $c_{in}(t) = C e^{-i(\omega_p - \omega_c)t}$, then $c_{in}(\omega) = 2\pi C \delta(\omega - \omega_p + \omega_c)$ and $c_{in}^\dagger(-\omega) = 2\pi C^\dagger \delta(\omega + \omega_p - \omega_c)$. Thus the fluctuations of the output field $\delta\tilde{c}_{out}(t) = \frac{1}{2\pi} \int_{-\infty}^{+\infty} V(\omega)\xi(\omega)e^{-i\omega t}d\omega + C E(\omega_p - \omega_c)e^{-i(\omega_p - \omega_c)t} + C^\dagger F(\omega_c - \omega_p)e^{-i(\omega_c - \omega_p)t}$. Therefore, $E(\omega_p - \omega_c)$ is the component at the probe frequency ω_p [which in 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 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 $c_{in}(\omega)$ 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 set up of Fig. 2

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

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; the last two terms in Eq. (13)

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

$$X(0) = N[E(\omega_m) + E^*(\omega_m)]^2 + |E(\omega_m)|^2 + |F(\omega_m)|^2 + 4|V(\omega_m)|^2\gamma_m \coth\left[\frac{\hbar\omega_m}{2k_B T}\right]. \quad (14)$$

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, 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, $T = 20$ 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 contribution of the input vacuum noise is about 0.16, 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, $T = 20$ 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, dotdashed, and dashed 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 dotdashed curve ($\wp = 20$ mW and $M = \sqrt{N(N+1)}$) exhibit the EIT dip, which is the result of the destruc-

the EIT phenomenon in the homodyne spectrum of the output field. Since the original equations (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 $N \approx M$,

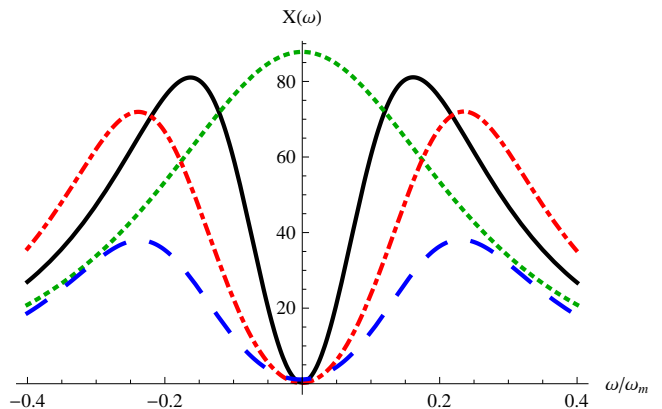


FIG. 3: (Color online) Plots of the homodyne spectrum $X(\omega)$ as a function of ω/ω_m for $N = 5$ in the absence (dotted curve) and the presence (solid, dotdashed, and dashed curves) of the coupling field for the temperature of the environment $T = 20$ mK. The solid curve is for $\wp = 10$ mW and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

tive 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 which 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], 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 $M = \sqrt{N(N+1)}$.

Next we increase the temperature to 100 mK. The 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 $N = 1, 5$ and $M = \sqrt{N(N+1)}$. In the presence of the coupling field ($\varphi = 10$ 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.

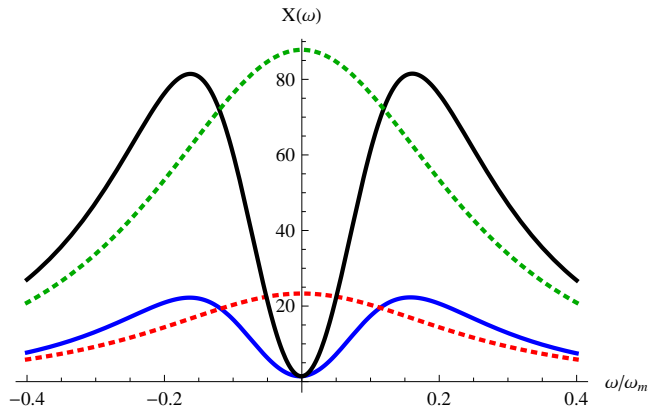


FIG. 4: (Color online) Plots of the homodyne spectrum $X(\omega)$ as a function of ω/ω_m for different values of the parameter N and $M = \sqrt{N(N+1)}$ in the absence (dotted curves) and the presence (solid curves) of the coupling field with power $\varphi = 10$ mW and the temperature of the environment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

The effects discussed above occur under 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, $Q' = \omega_m/\gamma_m = 6700$. The values for parameters T , φ , 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 $\varphi = 0, 10$ 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 rather weak dependence of the EIT curves on temperature. Therefore, current optomechanical designs can be used to realize quantum optical memory at 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_{in} so that it is a pulse field and use say a super Gaussian for the coupling field. The super Gaussian enables one to conveniently switch on and off the coupling field [28].

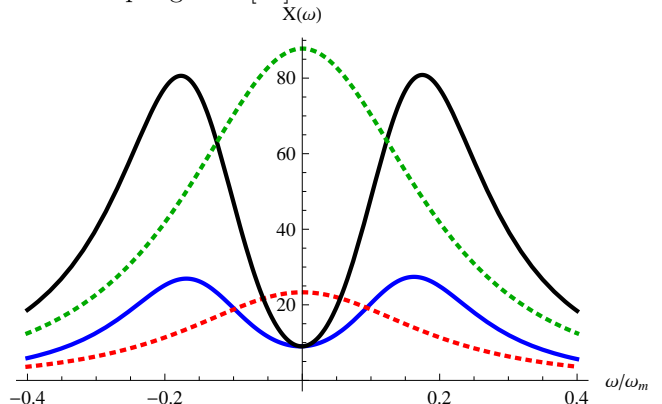


FIG. 5: (Color online) Same as in Fig. 4 but now the parameters used are from Ref. [22].

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 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 get be generalized to optomechanical systems working on the reactive coupling [23–25]. Our work suggest that optomechanical systems could be used as elements for quantum memory, explicit demonstration would be given elsewhere.

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