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Tunable ponderomotive squeezing in an optomechanical system with two coupled resonators*

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We investigate properties of the ponderomotive squeezing in an optomechanical system with two coupled resonators, where the tunable two-mode squeezing spectrum can be observed from the output field. It is realized that the squeezing orientation can be controlled by the detuning between the left cavity and pump laser. Especially, both cavity decay and environment temperature play a positive role in generating better pondermotive squeezing light. Strong squeezing spectra with a wide squeezing frequency range can be obtained by appropriate choice of parameters present in our optomechanical system.

Keywords: ponderomotive squeezing light, optomechanical system, coupled resonators

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1. Introduction

The optomechanical system has sparkled the interest of a vast scientific community due to its distinct applications, ranging from sensing of masses, forces, and displacements at the ultimate quantum limits to tests of the validity of quantum mechanics at a macroscopic level, up to the realization of quantum interfaces for quantum information networks.[1–6] In a typical optomechanical system, a movable mirror couples the cavity field by the radiation pressure, and this coupling can lead to many remarkable effects, for example, the optomechanically induced transparency (OMIT),[7–13] the quantum ground state cooling of the nanomechanical resonators,[14–17] entanglement between cavity modes and/or mechanical modes,[18–25] and squeezed states of the light or the mechanical modes.[26–34]

On the other hand, the squeezed light, with quantum fluctuations below the level of vacuum noise or the standard quantum limit, is particularly useful for ultra-sensitive force measurements,[35] gravitational wave detection,[36] and is also an important resource in quantum information science for continuous-variable information processing.[37] Over twenty years ago, an optomechanical cavity was proposed to generate the so-called ponderomotive squeezing of light[38,39] due to its similarity to a Kerr nonlinear medium. The mechanical element is shifted by the radiation pressure, proportionally to the intracavity intensity of optical field, which in turn modulates the phase of light leading to the correlations between the amplitude and phase quadratures of light. Therefore the optomechanical system behaves similarly to a cavity filled with a nonlinear Kerr medium. It is known that when a cavity containing a Kerr medium is driven by an intense laser, one can obtain appreciable squeezing in the spectrum of quadrature fluctuations at the cavity output. The above analogy therefore suggests that a strongly driven optomechanical cavity will also be able to produce quadrature squeezing at its output. So far, various schemes have been put forward for generating squeezed states either in the light[40,41] or the mechanical system[42,43] and, quite recently, squeezing of the optical field has been experimentally realized in optomechanical systems.[44–47] Experimental realization of the squeezed states is also extended to create and stabilize entangled states,[48] which are important for quantum information processing and precision measurements.

In this paper, we demonstrate a tunable two-mode squeezed light observable in a coupled-resonator optomechanical system, where the two resonators are coupled through photon hopping. One cavity is coherently driven by the pump laser field and the other cavity couples to a high-Q mechanical oscillator via radiation pressure. In such a system, the introduction of two coupled resonators owning tunable coupling strength will present distinct features. One of the features is double-optomechanical induced transparency.[49] Here we show the feature of the squeezing spectrum and discuss in detail the dependence of squeezing properties (including optimal squeezing position, squeezing degree and squeezing range) on system parameters, such as the detuning between the cavity and pump laser, the coupling strength between the two resonators, the power of the pump laser, the environment temperature, and the decay of cavity. We find that the single peak in squeezing spectrum turns to two peaks due to the stronger tun-

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neling strength. Compared to previous studies,[50–53] where the ponderomotive squeezing is generated in an optomechanical cavity assisted by an atomic ensemble or two movable mirrors, our scheme owns some favorable features. First, the two-mode squeezed light generated in our scheme is created by the introduction of the two tunneling resonators. In addition, the tunneling strength is tunable by adjusting the positions of the resonators. Second, strong squeezing spectrum with wide squeezing frequency range can be obtained by appropriate choice of parameters present in our optomechanical system. Finally, both the cavity decay and environment temperature play a positive role in generating better pondermotive squeezing light, which means that the squeezing light in our model is more easily to be observed in experiments.

The rest of this paper is organized as follows. In Section 2 we describe the model studied, derive the analytical expressions and the quantum Langevin equations of the optomechanical system. In Section 3, we give explicit relations for the ponderomotive squeezing spectrum of the transmitted field, and discuss the squeezing properties of the system. A conclusion of our findings is given in Section 4.

2. Model and solution

For an optomechancial system consisting of two coupled optical resonator as shown in Fig. 1, two cavity fields are coupled via the photon tunneling interaction. The cavity on the left is driven by the pump laser field, and the cavity on the right couples to the motion of oscillator through the radiation pressure. In the rotating frame of pump laser frequency ω_0, the Hamilton is given by

$$\begin{aligned} H =& \sum_{j=1}^{2}\hbar\Delta_j a_j^\dagger a_j + \frac{\hbar\omega_{\rm m}}{2}(p^2+q^2) + \hbar J(a_1^\dagger a_2 + a_2^\dagger a_1) \\ & - \hbar g a_2^\dagger a_2 q + {\rm i}\hbar\varepsilon_{\rm p}(a_1^\dagger - a_1), \end{aligned} \tag{1}$$

where a_j ($a_j^\dagger$) is annihilation (creation) operator for the cavity mode, and $\Delta_j = \omega_j - \omega_0$ is the detuning of cavity field with respect to the driving field. The pump laser is input from the left cavity with the amplitude $\varepsilon_{\rm p}$, which is related to the input power P and the damping rate κ_1 of the left cavity by $\varepsilon_{\rm p} = \sqrt{2P\kappa_1/\hbar\omega_0}$; $g = \frac{\omega_0}{L}\sqrt{\hbar/m\omega_{\rm m}}$ denotes the coupling between the oscillator and the right cavity, the operator $q(p)$ is the dimensionless position (momentum) of the mechanical oscillator with mass m and frequency $\omega_{\rm m}$. The photon tunneling amplitude between the cavities is J.

Using the Heisenberg equations of motion and taking into account the corresponding damping and noise terms, we can reach the quantum Langevin equations for the operators of the mechanical and optical modes

$$\begin{aligned} \dot{q} &= \omega_{\rm m} p, \\ \dot{p} &= -\omega_{\rm m} q + g a_2^\dagger a_2 - \gamma_{\rm m} p + \xi, \\ \dot{a}_1 &= -{\rm i}\Delta_1 a_1 - {\rm i}J a_2 + \varepsilon_{\rm p} - \kappa_1 a_1 + \sqrt{2\kappa_1}a_{1{\rm in}}, \\ \dot{a}_2 &= -{\rm i}\Delta_2 a_2 - {\rm i}J a_1 + {\rm i}g a_2 q - \kappa_2 a_2 + \sqrt{2\kappa_2}a_{2{\rm in}}, \end{aligned} \tag{2}$$

where κ_j denotes the decay rate of the two cavities. The mechanical oscillator is connected to a thermal bath at a damping rate $\gamma_{\rm m}$ with a mean thermal excitation number $n = 1/(e^{\frac{\hbar\omega_{\rm m}}{k_{\rm B}T}} - 1)$ with $k_{\rm B}$ being the Boltzmann constant and T the temperature of the mechanical bath. The mechanical mode is also affected by a random Brownian force ξ with correlation function $\langle\xi(t)\xi(t')\rangle \simeq \gamma_{\rm m}(2n+1)\delta(t-t')$. Furthermore, $a_{j{\rm in}}$ represents the input vacuum noise operator and its nonzero correlation function is $\langle a_{j{\rm in}}(t)a_{j{\rm in}}^\dagger(t')\rangle = \delta(t-t')$.



Fig. 1. Coupled-resonator optomechanical system consisting of two tunneling-coupled optical cavity modes a_1 and a_2, and a movable mirror coupled to the mode a_2 via radiation pressure. The system is driven by a pump filed $\varepsilon_{\rm p}$ from the left side with frequency ω_0.

We can rewrite each Heisenberg operator as a steady-state value plus an additional fluctuation operator with zero-mean value, $a_j = a_{j{\rm s}} + \delta a_j (j = 1,2)$, $q = q_{\rm s} + \delta q$, $p = p_{\rm s} + \delta p$. Then we substitute these expressions into the Langevin equations and obtain the steady-state values

$$\begin{aligned} p_{\rm s} &= 0, \\ q_{\rm s} &= \frac{g|a_{2{\rm s}}|^2}{\omega_{\rm m}}, \\ a_{2{\rm s}} &= \frac{-{\rm i}J a_{1{\rm s}}}{\kappa_2 + {\rm i}\Delta_2'}, \\ a_{1{\rm s}} &= \frac{\varepsilon_{\rm p}}{\kappa_1 + {\rm i}\Delta_1 + \frac{J^2}{\kappa_2 + i\Delta_2'}}, \end{aligned} \tag{3}$$

where $\Delta_2' = \Delta_2 - g q_{\rm s}$ is the effective detuning of the right cavity.

Retaining only first-order terms in the small quantities $\delta o (o = a_1, a_2, q, p)$, the linearized Langevin equations are given as follows:

$$\begin{aligned} \delta\dot{q} &= \omega_{\rm m}\delta p, \\ \delta\dot{p} &= -\omega_{\rm m}\delta q + G_0(\delta a_2 + \delta a_2^\dagger) - \gamma_{\rm m}\delta p + \xi, \\ \delta\dot{a}_1 &= -(\kappa_1 + {\rm i}\Delta_1)\delta a_1 - {\rm i}J\delta a_2 + \sqrt{2\kappa_1}a_{1{\rm in}}, \\ \delta\dot{a}_2 &= -(\kappa_2 + {\rm i}\Delta_2')\delta a_2 - {\rm i}J\delta a_1 + {\rm i}G_0\delta q + \sqrt{2\kappa_2}a_{2{\rm in}}, \end{aligned} \tag{4}$$

where $G_0 = g a_{2{\rm s}}$ is the effective coupling strength.

It is more convenient to work in the frequency domain because, experimentally, fluctuations of the electric field are more convenient to be measured in the frequency domain than in the time domain. For calculation in frequency domain, we introduce the Fourier transform for an operator u

($u = \delta a_1,\ \delta a_2,\ \delta q,\ \delta p,\ a_{1\mathrm{in}},\ a_{2\mathrm{in}},\ \xi$) as

$$\begin{aligned} u(t) &= \frac{1}{\sqrt{2\pi}}\int_{-\infty}^{\infty} \mathrm{e}^{\mathrm{i}\omega t}u(\omega)\mathrm{d}\omega, \\ u^{\dagger}(t) &= \frac{1}{\sqrt{2\pi}}\int_{-\infty}^{\infty} \mathrm{e}^{-\mathrm{i}\omega t}u^{\dagger}(\omega)\mathrm{d}\omega. \end{aligned} \tag{5}$$

Then we rewrite Eqs. (4) in the frequency domain as

$$\begin{aligned} \mathrm{i}\omega\delta q(\omega) &= \omega_{\mathrm{m}}\delta p(\omega), \\ \mathrm{i}\omega\delta p(\omega) &= -\omega_{\mathrm{m}}\delta q(\omega) + G_0\left\{\delta a_2(\omega) + \delta a_2^{\dagger}(-\omega)\right\} \\ &\quad -\gamma_{\mathrm{m}}\delta p(\omega) + \xi(\omega), \\ \mathrm{i}\omega\delta a_1(\omega) &= -(\kappa_1 + \mathrm{i}\Delta_1)\delta a_1(\omega) - \mathrm{i}J\delta a_2(\omega) \\ &\quad +\sqrt{2\kappa_1}a_{1\mathrm{in}}(\omega), \\ \mathrm{i}\omega\delta a_1^{\dagger}(-\omega) &= -(\kappa_1 - \mathrm{i}\Delta_1)\delta a_1^{\dagger}(-\omega) + \mathrm{i}J\delta a_2^{\dagger}(-\omega) \\ &\quad +\sqrt{2\kappa_1}a_{1\mathrm{in}}^{\dagger}(-\omega), \\ \mathrm{i}\omega\delta a_2(\omega) &= -(\kappa_2 + \mathrm{i}\Delta_2')\delta a_2(\omega) - \mathrm{i}J\delta a_1(\omega) \\ &\quad +\mathrm{i}G_0\delta q(\omega) + \sqrt{2\kappa_2}a_{2\mathrm{in}}(\omega), \\ \mathrm{i}\omega\delta a_2^{\dagger}(-\omega) &= -(\kappa_2 - \mathrm{i}\Delta_2')\delta a_2^{\dagger}(-\omega) + \mathrm{i}J\delta a_1^{\dagger}(-\omega) \\ &\quad -\mathrm{i}G_0^{*}\delta q(\omega) + \sqrt{2\kappa_2}a_{2\mathrm{in}}^{\dagger}(-\omega). \end{aligned} \tag{6}$$

By solving the above equations, we obtain

$$\begin{aligned} \delta a_1(\omega) &= \mu_1(\omega)a_{1\mathrm{in}}(\omega) + \mu_2(\omega)a_{1\mathrm{in}}^{\dagger}(-\omega) \\ &\quad +\mu_3(\omega)a_{2\mathrm{in}}(\omega) + \mu_4(\omega)a_{2\mathrm{in}}^{\dagger}(-\omega) \\ &\quad +\mu_5(\omega)\xi(\omega), \\ \delta a_1^{\dagger}(-\omega) &= \mu_1^{*}(-\omega)a_{1\mathrm{in}}^{\dagger}(-\omega) + \mu_2^{*}(-\omega)a_{1\mathrm{in}}(\omega) \\ &\quad +\mu_3^{*}(-\omega)a_{2\mathrm{in}}^{\dagger}(-\omega) + \mu_4^{*}(-\omega)a_{2\mathrm{in}}(\omega) \\ &\quad +\mu_5^{*}(-\omega)\xi(\omega), \end{aligned} \tag{7}$$

where the coefficients are given as follows:

$$\begin{aligned} \mu_1(\omega) &= -\frac{\sqrt{2\kappa_1}}{d}\{G^3\Lambda_{1-}\Lambda_{2+}\omega_{\mathrm{m}} \\ &\quad +\mathrm{i}G(J^2 + \Lambda_{1-}\Lambda_{2-})(\mathrm{i}G^2\omega_{\mathrm{m}} + \Lambda_{2+})\varepsilon\}, \\ \mu_2(\omega) &= -\frac{\sqrt{2\kappa_1}}{d}G^3J\omega_{\mathrm{m}}, \\ \mu_3(\omega) &= \frac{\mathrm{i}J}{d}\{\sqrt{2\kappa_2}G^3\Lambda_{1-}\omega_{\mathrm{m}} \\ &\quad +\mathrm{i}G\sqrt{2\kappa_2}(J^2 + \Lambda_{1-}\Lambda_{2-})\varepsilon\}, \\ \mu_4(\omega) &= \frac{\mathrm{i}\sqrt{2\kappa_2}}{d}G^3J\Lambda_{1-}\omega_{\mathrm{m}}, \\ \mu_5(\omega) &= \frac{\mathrm{i}G^2J}{d}(J^2 + \Lambda_{1-}\Lambda_{2-})\omega_{\mathrm{m}}. \end{aligned} \tag{8}$$

Here we have defined

$$\begin{aligned} \varepsilon &= \omega^2 - \omega_{\mathrm{m}}^2 - \mathrm{i}\gamma_{\mathrm{m}}\omega, \\ \Lambda_{1\pm} &= \kappa_1 + \mathrm{i}(\omega \pm \Delta_1), \\ \Lambda_{2\pm} &= \kappa_2 + \mathrm{i}(\omega \pm \Delta_2'), \\ d &= \mathrm{i}J\left\{\mathrm{i}G^3J\Lambda_{1-}\omega_{\mathrm{m}} + \mathrm{i}J(\mathrm{i}GJ^2 + \mathrm{i}G\Lambda_{1-}\Lambda_{2-})\varepsilon\right\}. \end{aligned} \tag{9}$$

In the following, we will study the ponderomotive squeezing numerically in the frequency domain.

3. Ponderomotive squeezing

It was shown that the optomechanical coupling can lead to squeezing of the nanomechanical motion, which can be inferred by measuring the squeezing of the transmitted microwave field. Here we analyze the squeezing properties of the transmitted field, which is accessible to experiment and useful for practical applications.[44–46] The squeezing spectrum of the transmitted field is given by[39]

$$S(\omega) = \int \mathrm{d}\omega' \left\langle \delta X_{\theta}^{\mathrm{out}}(\omega)\delta X_{\theta}^{\mathrm{out}}(\omega')\right\rangle, \tag{10}$$

where $\delta X_{\theta}^{\mathrm{out}}(\omega) = \mathrm{e}^{-\mathrm{i}\theta}\delta a_{\mathrm{out}}(\omega) + \mathrm{e}^{\mathrm{i}\theta}\delta a_{\mathrm{out}}^{\dagger}(\omega)$ and $\delta a_{\mathrm{out}} = \sqrt{2\kappa}\delta a - a_{\mathrm{in}}$.

Using the results of Eq. (7), and the following correlation functions in the frequency domain:

$$\begin{aligned} &\langle a_{\mathrm{in}}(\omega')a_{\mathrm{in}}^{\dagger}(\omega)\rangle = \delta(\omega' - \omega), \\ &\langle \xi(\omega')\xi(\omega)\rangle \simeq \gamma_{\mathrm{m}}(2\bar{n} + 1)\delta(\omega' + \omega), \end{aligned} \tag{11}$$

we can obtain

$$\begin{aligned} &\left\langle \delta X_{\theta}^{\mathrm{out}}(\omega)\delta X_{\theta}^{\mathrm{out}}(\omega')\right\rangle \\ &= \delta(\omega - \omega')\left[1 + 2B_{a_1^{\dagger}a_1} + (\mathrm{e}^{-2\mathrm{i}\theta}B_{a_1a_1} + \mathrm{c.c})\right], \end{aligned} \tag{12}$$

where

$$\begin{aligned} B_{a_1^{\dagger}a_1} &= 2\kappa_1\left\{|\mu_2(\omega)|^2 + |\mu_4(\omega)|^2 + |\mu_5(\omega)|^2\,\gamma_{\mathrm{m}}(2n+1)\right\}, \\ B_{a_1a_1} &= 2\kappa_1\{\mu_1(\omega)\mu_2(-\omega) - \frac{\mu_2(-\omega)}{\sqrt{2\kappa_1}} \\ &\quad +\mu_3(\omega)\mu_4(-\omega) + \mu_5(\omega)\mu_5(-\omega)\gamma_{\mathrm{m}}(2n+1)\}. \end{aligned} \tag{13}$$

Substituting Eq. (12) into Eq. (10) yields the squeezing spectrum in a simple form

$$S_{\theta}(\omega) = 1 + 2B_{a_1^{\dagger}a_1} + (\mathrm{e}^{-2\mathrm{i}\theta}B_{a_1a_1} + \mathrm{c.c.}). \tag{14}$$

As the phase angle θ is an adjustable parameter, we choose appropriate θ by solving $dS_{\theta}(\omega)/\mathrm{d}\theta = 0$, then we obtain

$$\mathrm{e}^{2\mathrm{i}\theta_{\mathrm{opt}}} = \pm\frac{B_{a_1a_1}(\omega)}{|B_{a_1a_1}(\omega)|}. \tag{15}$$

Here, we choose the negative value to optimize the degree of squeezing. Thus the squeezing spectrum can be rewritten as

$$S_{\mathrm{opt}}(\omega) = 1 + 2B_{a_1^{\dagger}a_1} - 2\,|B_{a_1a_1}|\,. \tag{16}$$

In order to obtain squeezing, the term $2B_{a_1^{\dagger}a_1} - 2|B_{a_1a_1}|$ on the right-hand side must be negative, i.e., $S_{\mathrm{opt}}(\omega) < 1$.

Now we turn to discuss the feasibility of the squeezing spectrum in the coupled-resonator optomechanical system, and the relationship of the squeezing properties to the system parameters, such as the detunings between cavities and the

pump laser, the tunneling strength of the two resonators, the power of pump laser, the environment temperature, and decay rate of the cavity. For simplicity, we assume that two cavity fields have the same decay rate, i.e., $\kappa_1 = \kappa_2 = \kappa$. To demonstrate the numerical results, we use the experimentally realizable parameters:[54,55] $\lambda_c = 1064$ nm, $\omega_m = 2\pi \times 10^7$ Hz, $\gamma_m = 628$ Hz, $L = 1$ mm, $m = 12$ ng.

Firstly, we discuss the dependence of ponderomotive squeezing on the detunings between the cavities and the pump laser in Fig. 2, where we have set $T = 15$ mK, $P = 10$ mW, and $\kappa = J = 0.5\omega_m$. One can find from this figure that, in the case of weak tunneling strength of the two resonators ($J = 0.5\omega_m$), the squeezing spectrum presents only a single peak at the frequency $\omega = \omega_m$ (Figs. 2(a) and 2(c)) or at the frequency $\omega = -\omega_m$ (Figs. 2(b) and 2(d)). The squeezing of light is optimized at $\omega = \omega_m$ for $\Delta_1 = \omega_m$, i.e., the left cavity is resonant with the Stokes sideband of the laser, as shown in Figs. 2(a) and 2(c). However, it is seen from Figs. 2(b) and 2(d) that the optimal squeezing of light appears at $\omega = -\omega_m$ for $\Delta_1 = -\omega_m$, i.e., the left cavity is resonant with the anti-Stokes sideband of the laser. This implies that we can control the point where the optimal squeezing occurs through selecting different detuning Δ_1 between the left cavity and the pump laser. It comes from the fact that the optimal squeezing spectrum is emitted from the left cavity. Therefore, the optimal squeezing position is strongly dependent on detuning Δ_1, but independent of Δ_2. Comparing with Figs. 2(a) and 2(c), as well as Figs. 2(b) and 2(d), it is not difficult to see that the degree of squeezing can be controlled by appropriate choice of the effective detuning Δ_2' between the right cavity and the pump laser.

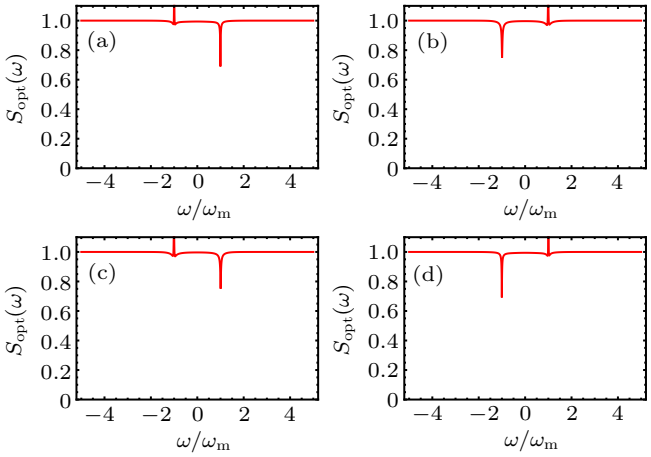


Fig. 2. The squeezing spectrum with different detunings between the cavities and the pump laser when $T = 15$ mK, $P = 10$ mW, $J = \kappa = 0.5\omega_m$: (a) $\Delta_1 = \Delta_2' = \omega_m$; (b) $\Delta_1 = -\omega_m$, $\Delta_2' = \omega_m$; (c) $\Delta_1 = \omega_m$, $\Delta_2' = -\omega_m$; (d) $\Delta_1 = \Delta_2' = -\omega_m$.

With the enhancement of the tunneling strength of the two resonators, two peaks appear in the transmitted field, which can be easily seen from Fig. 3. This is due to the fact that an additional hybrid energy level is effectively introduced into the original three-level system by the stronger tunneling coupling between the optomechanical resonator and the left resonator. It is worth mentioning that our proposal is quite different from the previous idea,[51,52] where two-mode squeezed light is caused by the frequency difference between the two movable mirrors[51] or the Coulomb coupling between the two movable mirrors.[52] The introduction of tunneling coupling between the two resonators also leads to a split in the OMIT window, as reported in Ref. [49]. Comparing the four curves in Fig. 3, we can find that the degree of squeezing at the point $\omega \simeq \pm\omega_m$ increases and the point where the maximal squeezing occurs shifts slightly to the right side with the tunneling strength between the two resonators.

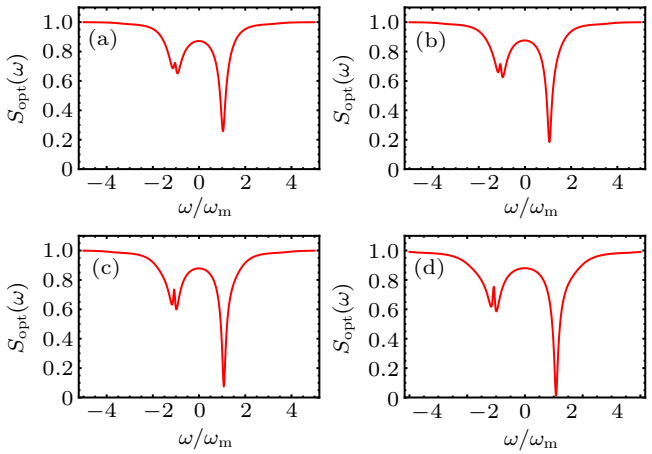


Fig. 3. The squeezing spectrum with different tunneling strength J of two resonators when $T = 15$ mK, $P = 10$ mW, $\kappa = 0.5\omega_m$, $\Delta_1 = \Delta_2' = \omega_m$: (a) $J = 2.3\omega_m$; (b) $J = 2.5\omega_m$; (c) $J = 2.7\omega_m$; (d) $J = 2.8\omega_m$.

Then we show the effect of the pump laser power on the squeezing spectrum in Fig. 4. It can be easily seen that the maximal squeezing degree remains unchanged with the increase of the pump laser power. However, the spectral range of the squeezed light becomes wider and wider by increasing the power of the pump laser. The half height width of the squeezing spectrum at frequency $\omega \simeq \omega_m$ is about $0.35\omega_m$ in the case of $P = 10$ mW (Fig. 4(a)). With the increase of the pump laser power, the half height width of the squeezing spectrum at frequency $\omega \simeq \omega_m$ changes to $0.9\omega_m$ ($P = 40$ mW), $1.4\omega_m$ ($P = 80$ mW), and $2.1\omega_m$ ($P = 200$ mW), which means that the frequency range of squeezing becomes wider with the increase of pump laser power. Especially, the frequency range of squeezing is about $-5 \leq \omega/\omega_m \leq 5$ in the case of $P = 200$ mW. This is due to the fact that a larger power will input amount photons which will effectively overcome the effects of thermal noise.

The robustness of squeezing dependent on the environment temperature T is shown in Fig. 5. Here we choose the power of pump laser is 1 mW and the other parameters are same as those used for Fig. 3. Figure 5 displays perfect two-mode squeezing spectrum at frequencies $\omega = \pm\omega_m$ with

different squeezing degree. In the case of $T = 15$ mK and $T = 55$ mK, the squeezing degree at $\omega = -\omega_m$ is smaller than that at $\omega = \omega_m$. However, the squeezing degree at $\omega = -\omega_m$ is larger than that at $\omega = \omega_m$ in the case of $T = 225$ mK, and the squeezing degrees at these two points $\omega = \pm\omega_m$ are equal to each other in the case of $T = 125$ mK. In other words, the squeezing degree at $\omega = -\omega_m$ increases with the temperature, while the squeezing degree at $\omega = \omega_m$ decreases with the temperature.

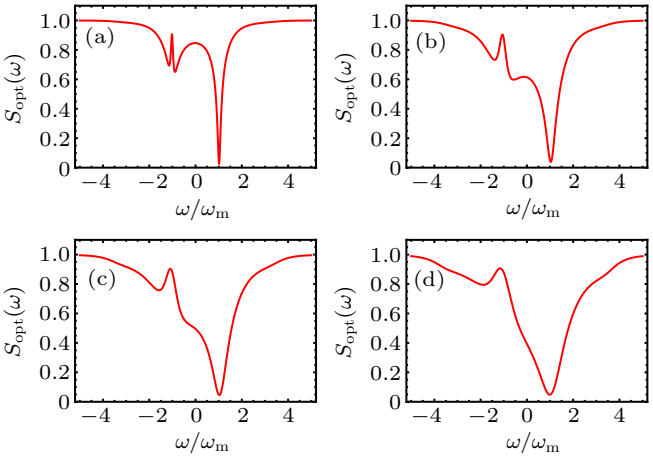


Fig. 4. The squeezing spectrum with different power of pump laser when $T = 15$ mK, J=2.3ω_m, κ=0.7ω_m, $\Delta_1 = \Delta_2' = \omega_m$: (a) $P = 10$ mW; (b) $P = 40$ mW; (c) $P = 80$ mW; (d) $P = 200$ mW.

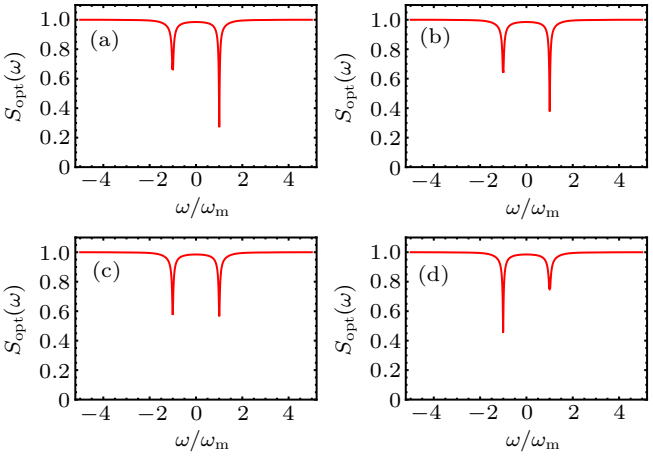


Fig. 5. The squeezing spectrum at different temperatures when $P =$ 1 mW, $J = 2.3\omega_m$, $\kappa = 0.5\omega_m$, $\Delta_1 = \Delta_2' = \omega_m$: (a) $T = 15$ mK; (b) $T = 55$ mK; (c) $T = 125$ mK; (d) $T = 225$ mK.

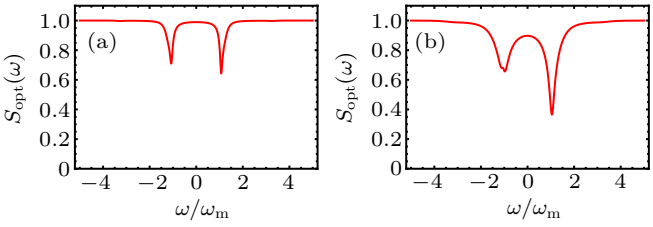


Fig. 6. The squeezing spectrum with different cavity decay rates of the cavities when $T = 15$ mK, $P = 10$ mW, $J = 2.3\omega_m$, $\Delta_1 = \Delta_2' = \omega_m$: (a) $\kappa = 0.1\omega_m$; (b) $\kappa = 0.4\omega_m$.

Finally, we show the squeezing spectrum with different decay rate κ of the cavities in Fig. 6. Comparing Fig. 6(a) ($\kappa = 0.1\omega_m$), Fig. 6(b) ($\kappa = 0.4\omega_m$) and Fig. 4(a) ($\kappa = 0.7\omega_m$), we can find that not only the squeezing degree is enhanced but also the squeezing range is widened with the decay of the cavity. This indicates the squeezing robustness with respect to the decay of the cavity. The optimal squeezed light can be even observable in the bad cavities. Therefore, the squeezed light spectrum in our model is easy to be achieved experimentally.

4. Conclusion

We have analyzed the properties of the ponderomotive squeezed light observable in an optomechanical system with two coupled resonators. We show that a two-mode tunable squeezed light can be achievable in this system due to the large tunneling strength between the two resonators. In the case of small tunneling strength where the squeezing spectrum presents only a single peak, the squeezing orientation is influenced by the detuning between the left cavity and the pump laser. We also find that the squeezing degree can be flexibly controlled by adjusting the tunneling strength, the decay of the cavities and the environment temperature. Especially, the larger the cavity dissipation, the higher the squeezing degree. In addition, the frequency range of squeezing becomes wider with the increase of laser power and cavity decay. In other words, even the cavity dissipation and environment temperature play a positive role in the generation of good squeezing light. Thus, the squeezed light spectrum in our model is easy to be observed experimentally.

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