Tunable ponderomotive squeezing induced by Coulomb interaction in an optomechanical system

Qin Wu(吴琴)†

School of Information Engineering, Guangdong Medical University, Dongguan 523808, *China*

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We investigate the properties of the ponderomotive squeezing in an optomechanical system coupled to a charged nanomechanical oscillator (NMO) nearby via Coulomb force. We find that the introduction of Coulomb interaction allows the generation of squeezed output light from this system. Our numerical results show that the degree of squeezing can be tuned by the Coulomb coupling strength, the power of laser, and the frequencies of NMOs. Furthermore, the squeezing generated in our approach can be used to measure the Coulomb coupling strength.

Keywords: ponderomotive squeezing, Coulomb interaction, optomechanical system

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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.[1–4] In a typical optomechanical system, the coupling between a movable mirror and the cavity field can lead to many remakable effects, such as the optomechanically induced transparency (OMIT),[5–11] the quantum ground state cooling of the nanomechanical resonators,[12–14] entanglement between cavity modes and/or mechanical modes,[15–22] and squeezed states of the light or the mechanical modes.[23–25]

On the other hand, the squeezed light, where fluctuations in one quadrature of the optical amplitude drop below the level of vacuum noise, is interesting from both fundamental and practical points of view. Such light can be used to improve the measurement sensitivity in applications ranging from gravitational wave detection[26] to even biological measurements.[27] Squeezed states of light are also a key ingredient for continuous-variable information processing.[28] It is well known that the standard method for generating optical squeezing is to drive a nonlinear optical medium,[29] which has been realized in Ref. [30]. Recently, the squeezing of the transmitted field can also be achievable in optomechanical systems,[31,32] where cavity photons are coupled to mechanical motion by radiation pressure. The optomechanical method of manipulating the quantum fluctuations has been termed as "ponderomotive squeezing",[30] which relies on the nonlinearity induced by the optomechanical interaction between the position of the oscillator and the intensity of the cavity mode.[31,33] Moreover, achieving this kind of ponderomotive squeezing experimentally is an important goal due to its applications in ultrahigh precision measurements.[34–36]

In this work, we demonstrate a tunable squeezed light observable in an optomechanical system, in which the two nanomechanical oscillators (NMO) are coupled through Coulomb force. We find that the single peak in squeezing spectrum is split to two peaks due to the Coulomb interaction. Specifically, the separation of the two optimal squeezing points can be tuned by the Coulomb interaction.

Compared to previous studies,[37,38] where the ponderomotive squeezing is generated in a cavity assisted by an atomic ensemble or in a ring cavity consisting of two mirrors without Coulomb force, our scheme owns some favorable features. First, the squeezed light generated in our scheme is created by Coulomb interaction. In addition, the Coulomb coupling has important advantages, for instance, the range of the interaction distance is from nanometer to meter,[39–41] the strength can be controlled by the bias voltage,[39,40] and it can also be applied to the different kinds of charged objects at different frequecies.[41] Second, the point with optimal squeezing can be tuned by the frequencies of the two NMOs. Finally, we find that the two optimal squeezing points are apart near linearly with respect to the Coulomb coupling strength. This feature reminds us of a practical application of the ponderomotive squeezing spectrum for detecting the Coulomb coupling strength. In this context, we have to emphasize that our proposal is essentially different from previous idea,[38] where the squeezed light is caused by the frequency difference between the two movable mirrors. In contrast, our studied squeezed light can be observed even for two identical movable mirrors.

The remainder of this paper is organized as follows. In Section 2, we present the model under study and analytical expressions of the optomechanical system, derive the quan-

†Corresponding author. E-mail: 905374532@qq.com

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tum Langevin equations, and the steady state of the system. In Section 3, we quantify the squeezing properties of the system. Finally we draw our conclusions in Section 4.

2. Model and solution

As sketched in Fig. 1, our system consists of a high-quality Fabry–Pérot cavity with a fixed mirror and a nanomechanical oscillator NMO_1. NMO_1 with the bias gate voltage V_1 is coupled to another charged NMO_2 with the bias gate voltage $-V_2$ through the Coulomb force. The optomechanical cavity of the length L is driven by one strong pump field ε_p from the left side of the cavity. The output field is represented by ε_{out}. q_1 and q_2 represent the small displacements of NMO_1 and NMO_2 from their equilibrium positions, and r_0 is the equilibrium distance between the two NMOs. Then, the Hamiltonian in the rotating frame at frequency of the pump field ω_p can be written as

$$\begin{aligned} H &= \hbar\Delta_c c^\dagger c + \frac{\hbar\omega_1}{2}(p_1^2+q_1^2) + \frac{\hbar\omega_2}{2}(p_2^2+q_2^2) \\ &\quad - \hbar g c^\dagger c q_1 + \hbar\lambda q_1 q_2 + \mathrm{i}\hbar\varepsilon_p(c^\dagger - c), \end{aligned} \tag{1}$$

where $\Delta_c = \omega_c - \omega_p$ is the detuning of the pump field from the bare cavity. The first term in Eq. (1) is for the single-mode cavity field with frequency ω_c and annihilation (creation) operator c ($c^\dagger$). The second (third) term describes the vibration of the charged NMO_1 (NMO_2) with frequency ω_1 (ω_2), effective mass m_1 (m_2), the position q_1 (q_2), and the momentum operator p_1 (p_2).[42] NMO_1 couples to the cavity field due to the radiation pressure with the coupling strength $g = (\omega_c/L)\sqrt{\hbar/m\omega_1}$. The fifth term in Eq. (1) presents the Coulomb coupling between the charged NMO_1 and NMO_2,[41] where the NMO_1 and NMO_2 take the charges $Q_1 = C_1V_1$ and $Q_2 = -C_2V_2$, with C_1 (C_2) and V_1 ($-V_2$) being the capacitance and the voltage of the bias gate, respectively. The last term in Eq. (1) describes the interaction between the cavity field and the input field. The pump field strength ε_p depends on the power $\wp_p$ of coupling field, $\varepsilon_p = \sqrt{2\kappa\wp_p/\omega_p}$, where κ is the cavity decay rate.

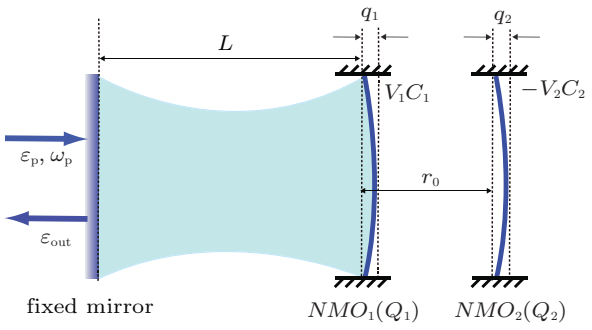


Fig. 1. (color online) Schematic diagram of the system.

Using the Heisenberg equations of motion, and considering photon losses from the cavity and the Brownian noise from the environment, we can obtain the following nonlinear quantum Langevin equations for the operators of the mechanical and optical modes:[42]

$$\begin{aligned} \dot{q_1} &= \omega_1 p_1, \\ \dot{p_1} &= -\omega_1 q_1 - \gamma_1 p_1 + g c^\dagger c - \lambda q_2 + \xi_1, \\ \dot{q_2} &= \omega_2 p_2, \\ \dot{p_2} &= -\omega_2 q_2 - \gamma_2 p_2 - \lambda q_1 + \xi_2, \\ \dot{c} &= -(\kappa + \mathrm{i}\Delta_c)c + \mathrm{i}gcq_1 + \varepsilon_p + \sqrt{2\kappa}c_{in}, \end{aligned} \tag{2}$$

where γ_1 (γ_2) and κ are the decay rates for NMO_1 (NMO_2) and the cavity, respectively. The quantum Brownian noise ξ_1 (ξ_2) comes from the coupling between NMO_1 (NMO_2) and its own environment with correlation function [43]

$$\langle \xi(t)\xi(t')\rangle \simeq \gamma_m(2\bar{n}+1)\delta(t-t'), \tag{3}$$

where $\bar{n} = [\exp(\hbar\omega_m/k_BT) - 1]^{-1}$ is the mean number of thermal excitation, with k_B being the Boltzmann constant. The c_{in} is the input vacuum noise operator with correlation function[43]

$$\langle c_{in}(t)c_{in}^\dagger(t')\rangle = \delta(t-t'). \tag{4}$$

We can rewrite each Heisenberg operator as a steady-state value plus an additional fluctuation operator with zero-mean value, $c = c_s + \delta c$, $q_i = q_{is} + \delta q_i$, $p_i = p_{is} + \delta p_i$ ($i = 1,2$). Then we insert these expressions into the Langevin equations and obtain the steady-state values

$$\begin{aligned} p_{1s} &= 0, \quad p_{2s} = 0, \\ q_{1s} &= \frac{g|c_s|^2}{\omega_1 - \lambda^2/\omega_2}, \quad q_{2s} = \frac{-\lambda}{\omega_2} q_{1s}, \\ c_s &= \frac{\varepsilon_p}{\kappa + \mathrm{i}\Delta}, \end{aligned} \tag{5}$$

where $\Delta = \Delta_c - gq_{1s}$ is the effective cavity detuning.

Retaining only first-order terms in the small quantities δo ($o = c, q_i, p_i$), the linearized Langevin equations for the fluctuations are given by

$$\begin{aligned} \delta\dot{q_1} &= \omega_1 \delta p_1, \\ \delta\dot{p_1} &= -\omega_1 \delta q_1 - \gamma_1 \delta p_1 + G(\delta c + \delta c^\dagger) - \lambda \delta q_2 + \xi_1, \\ \delta\dot{q_2} &= \omega_2 \delta p_2, \\ \delta\dot{p_2} &= -\omega_2 \delta q_2 - \gamma_2 \delta p_2 - \lambda \delta q_1 + \xi_2, \\ \delta\dot{c} &= -(\kappa + \mathrm{i}\Delta)\delta c + \mathrm{i}G\delta q_1 + \sqrt{2\kappa}c_{in}, \end{aligned} \tag{6}$$

where $G = gc_s$ 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. Using the following Fourier transform for an operator u ($u=\delta c,\delta p_i,\delta q_i,c_{\rm in}$, and ξ_i),

$$\begin{aligned} u(t) &= \frac{1}{\sqrt{2\pi}}\int_{-\infty}^{\infty} {\rm e}^{{\rm i}\omega t}u(\omega){\rm d}\omega,\\ u^{\dagger}(t) &= \frac{1}{\sqrt{2\pi}}\int_{-\infty}^{\infty} {\rm e}^{-{\rm i}\omega t}u^{\dagger}(\omega){\rm d}\omega, \end{aligned} \tag{7}$$

we rewrite Eq. (6) in the frequency domain in matrix form

$$\boldsymbol{AX}=\boldsymbol{B}, \tag{8}$$

where

$$\begin{aligned} \boldsymbol{B} &= (0,\xi_1,0,\xi_2,\sqrt{2\kappa}c_{\rm in},\sqrt{2\kappa}c_{\rm in}^{\dagger})^{\rm T},\\ \boldsymbol{X} &= (\delta q_1,\delta p_1,\delta q_2,\delta p_2,\delta c,\delta c^{\dagger})^{\rm T},\\ \boldsymbol{A} &= \begin{pmatrix} {\rm i}\omega & -\omega_1 & 0 & 0 & 0 & 0\\ -\omega_1 & \gamma_1+{\rm i}\omega & \lambda & 0 & -G & -G\\ 0 & 0 & {\rm i}\omega & -\omega_2 & 0 & 0\\ \lambda & 0 & \omega_2 & \gamma_2+{\rm i}\omega & 0 & 0\\ -{\rm i}G & 0 & 0 & 0 & \Theta_1 & 0\\ {\rm i}G & 0 & 0 & 0 & 0 & \Theta_2 \end{pmatrix}, \end{aligned}$$

in which we have defined $\Theta_1=\kappa+{\rm i}\omega+{\rm i}\Delta$ and $\Theta_2=\kappa+{\rm i}\omega-{\rm i}\Delta$.

By solving Eq. (8), we obtain

$$\delta c=\mu_1(\omega)c_{\rm in}+\mu_2(\omega)c_{\rm in}^{\dagger}+\mu_3(\omega)\xi_1+\mu_4(\omega)\xi_2, \tag{9}$$

where the coefficients are given by

$$\begin{aligned} \mu_1(\omega) &= \frac{\sqrt{2\kappa}}{\Lambda_+}+\frac{\sqrt{2\kappa}}{d}G^2\Lambda_-\omega\omega_1\varepsilon_2,\\ \mu_2(\omega) &= \frac{\sqrt{2\kappa}}{d}G^2\Lambda_+\omega\omega_1\varepsilon_2,\\ \mu_3(\omega) &= \frac{G\Lambda_+\Lambda_-\omega\omega_1\varepsilon_2}{d},\\ \mu_4(\omega) &= -\frac{G\lambda\Lambda_+\Lambda_-\omega\omega_1\omega_2}{d}, \end{aligned} \tag{10}$$

in which we have defined

$$\begin{aligned} \varepsilon_1 &= \omega_1^2-\omega^2+{\rm i}\gamma_1\omega,\\ \varepsilon_2 &= \omega_2^2-\omega^2+{\rm i}\gamma_2\omega,\\ \Lambda_{\pm} &= \kappa+{\rm i}(\omega\pm\Delta),\\ d &= \Lambda_+\{G^2\omega\Lambda_-\omega_1\varepsilon_2+\Lambda_+[-G^2\omega\omega_1\varepsilon_2\\ &\quad +\Lambda_-(-{\rm i}\lambda^2\omega\omega_1\omega_2+{\rm i}\omega\varepsilon_1\varepsilon_2)]\}. \end{aligned} \tag{11}$$

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.[34–36] The squeezing spectrum of the transmitted field is given by[32]

$$S_{\theta}(\omega)=\int {\rm d}\omega'\langle\delta X_{\theta}^{\rm out}(\omega)\delta X_{\theta}^{\rm out}(\omega')\rangle, \tag{12}$$

where $\delta X_{\theta}^{\rm out}(\omega)={\rm e}^{-{\rm i}\theta}\delta a_{\rm out}(\omega)+{\rm e}^{{\rm i}\theta}\delta a_{\rm out}^{\dagger}(\omega)$ and $\delta a_{\rm out}=\sqrt{2\kappa}\delta a-a_{\rm in}$. By using the results of Eqs. (9) and the following correlation functions in the frequency domain:

$$\begin{aligned} &\langle c_{\rm in}(\omega')c_{\rm in}^{\dagger}(\omega)\rangle=\delta(\omega'+\omega),\\ &\langle\xi(\omega')\xi(\omega)\rangle\simeq\gamma_{\rm m}(2\bar{n}+1)\delta(\omega'+\omega), \end{aligned} \tag{13}$$

we can obtain

$$\begin{aligned} &\langle\delta X_{\theta}^{\rm out}(\omega)\delta X_{\theta}^{\rm out}(\omega')\rangle\\ &=[1+2B_{c^{\dagger}c}+({\rm e}^{-2{\rm i}\theta}B_{cc}+{\rm e}^{2{\rm i}\theta}B_{c^{\dagger}c^{\dagger}})]\delta(\omega+\omega'), \end{aligned} \tag{14}$$

in which

$$\begin{aligned} B_{c^{\dagger}c} &= 2\kappa(|\mu_2(-\omega)|^2+\gamma_1(2\bar{n}_1+1)|\mu_3(-\omega)|^2\\ &\quad +\gamma_2(2\bar{n}_2+1)|\mu_4(-\omega)|^2),\\ B_{cc} &= 2\kappa[\mu_1(\omega)\mu_2(-\omega)-\frac{\mu_2(-\omega)}{\sqrt{2\kappa}}\\ &\quad +\gamma_1(2\bar{n}_1+1)\mu_3(\omega)\mu_3(-\omega)\\ &\quad +\gamma_2(2\bar{n}_2+1)\mu_4(\omega)\mu_4(-\omega)]. \end{aligned} \tag{15}$$

Taking Eq. (14) into Eq. (12), we can obtain the squeezing spectrum in a simple form

$$S_{\theta}(\omega)=1+2B_{c^{\dagger}c}+{\rm e}^{-2{\rm i}\theta}B_{cc}+{\rm e}^{2{\rm i}\theta}B_{c^{\dagger}c^{\dagger}}. \tag{16}$$

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

$${\rm e}^{2{\rm i}\theta_{\rm opt}}=\pm\frac{B_{cc}(\omega)}{|B_{cc}(\omega)|}. \tag{17}$$

Here, we choose the negative value to optimize the degree of squeezing. Hence, equation (16) can be rewritten as

$$S_{\rm opt}(\omega)=1+2B_{c^{\dagger}c}-2|B_{cc}|. \tag{18}$$

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

For simplicity, we first assume two identical charged NMOs $\omega_1=\omega_2=\omega_{\rm m}$ and the detuning between the pump field and the cavity mode satisfying $\Delta\simeq\omega_{\rm m}$. To demonstrate the numerical results, we use the experimentally realizable parameters from Refs. [44] and [45]: $\lambda_{\rm c}=1064$ nm, $\kappa=\pi\times10^7$ Hz, $\omega_{\rm m}=2\pi\times10^7$ Hz, $\gamma_1=\gamma_2=400\pi$ Hz, $L=6$ mm, $m_1=m_2=12$ ng.

Firstly, we discuss the dependence of ponderomotive squeezing on the Coulomb interaction strength (units of $\lambda_0=\pi\times10^7$ Hz/m^2) in Fig. 2, where we have set $T_1=2$ mK, $T_2=3$ mK, and $\wp_{\rm p}=10$ mW. One can find from this figure that the squeezing spectrum presents only a single peak

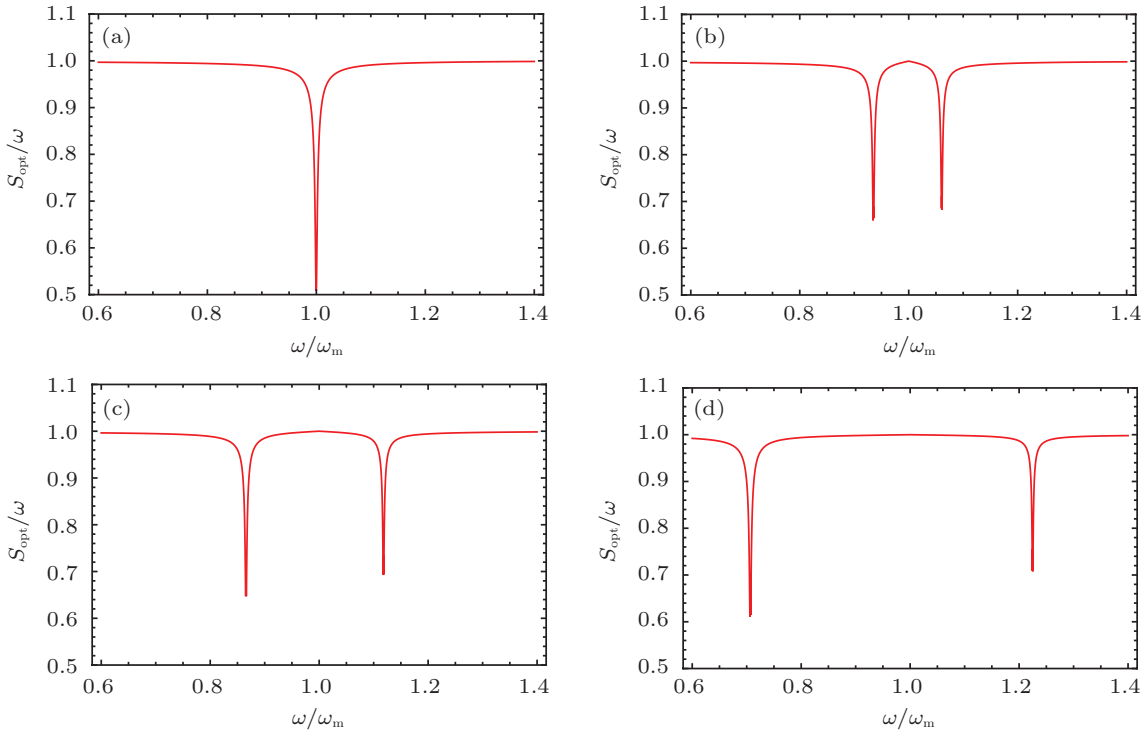


Fig. 2. (color online) The squeezing spectrum with different Coulomb coupling strength λ when $\wp_{\rm p} = 10$ mW, $T_1 = 2$ mK, $T_2 = 3$ mK. (a) $\lambda = 0$; (b) $\lambda = 0.25\lambda_0$; (c) $\lambda = 0.5\lambda_0$; (d) $\lambda = \lambda_0$ ($\lambda_0 = \pi \times 10^7$ Hz/m^2).

at the frequency $\omega = \omega_{\rm m}$ in the absence of the Coulomb interaction ($\lambda = 0$), and the single peak is split into two peaks in the presence of the Coulomb coupling. This is due to the fact that an additional hybrid energy level is introduced into the original three-level system by the Coulomb coupling between the external NMO and the optomechanical resonator, which is quite different from Ref. [38]. The introduction of Coulomb coupling also leads to a split in the OMIT window, as reported in Ref. [39]. Moreover, the two points $\omega_{\rm L,R}$ with optimal squeezing are further and further apart with the increase of λ. For example, $\omega_{\rm L(R)} = 0.93\omega_{\rm m}(1.06\omega_{\rm m}), 0.86\omega_{\rm m}(1.12\omega_{\rm m})$, and $0.71\omega_{\rm m}(1.23\omega_{\rm m})$ for the case of $\lambda = 0.25\lambda_0, 0.5\lambda_0$, and $\lambda = \lambda_0$, respectively. Obviously, the separation $d = \omega_{\rm R} - \omega_{\rm L}$ is $0.52\omega_{\rm m}$ when $\lambda = \lambda_0$, which is about twice of that in the case of $\lambda = 0.5\lambda_0$. In order to further explore the influence of Coulomb interaction on the separation d, we plot d as a function of the coupling strength λ (units of $\lambda_0 = \pi \times 10^7$ Hz/m^2) in Fig. 3, where the separation d almost linearly increases with λ within the regime $\lambda = \{0, 1.6\lambda_0\}$. This indicates the possibility to detect the Coulomb coupling strength between NMO$_1$ and NMO$_2$ by measuring the separation d in the squeezing spectrum of the output field. We also find that the degree of squeezing at point $\omega = \omega_{\rm L}(\omega_{\rm R})$ increases (decreases) with the Coulomb coupling strength.

Then we show the effect of the power of pump field on the degree of squeezing in Fig. 4. Comparing Fig. 2(c) and Fig. 4, we can find that the degree of squeezing is enhanced with increasing power of pump field, while there is no effect on the separation of the two optimal squeezing points. This is due to the fact that a larger power will input more photons which will effectively overcome the effects of the thermal noise.

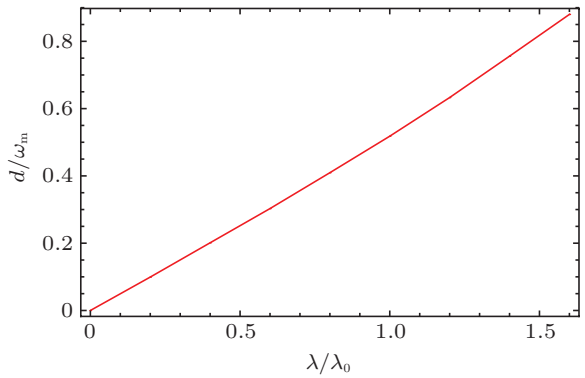


Fig. 3. (color online) Separation between two optimal squeezing points as a function of the Coulomb coupling strength λ (units of $\lambda_0 = \pi \times 10^7$ Hz/m^2) when $\wp_{\rm p} = 10$ mW, $T_1 = 2$ mK, and $T_2 = 3$ mK.

Finally, we consider the ponderomotive squeezing with non-identical NMOs since the two identical NMOs considered above are theoretically simplified but rarely existing experimentally. For two charged NMOs with different frequencies, as plotted in Fig. 5, where we have set the frequencies of NMOs $\omega_1 = \omega_{\rm m} \pm 0.1 \times 2\pi \times 10^7$ Hz and $\omega_2 = \omega_{\rm m} \mp 0.1 \times 2\pi \times 10^7$ Hz, the degree of squeezing at the optimal squeezing points changes with respect to the case of identical NMOs. That is to say, the degree of squeezing at frequency $\omega_{\rm L}$ decreases while the degree of squeezing at frequency $\omega_{\rm R}$ is en-

hanced in case of $\omega_1 > \omega_2$, as seen in Fig. 5(a), but the opposite changes in degree of squeezing at these two points take place in case of $\omega_1 < \omega_2$, as shown in Fig. 5(b). Thus, this provides us a good method for controlling the degree of squeezing and finding the optimal squeezing.

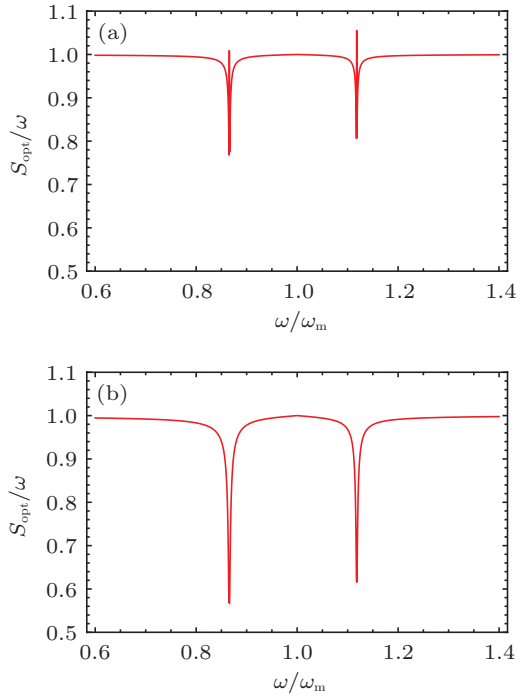


Fig. 4. (color online) Squeezing spectrum with different power of the pump field when $\lambda = 0.5\lambda_0$, $T_1 = 2$ mK, $T_2 = 3$ mK. (a) $\wp_{\rm p} = 5$ mW; (b) $\wp_{\rm p} = 15$ mW.

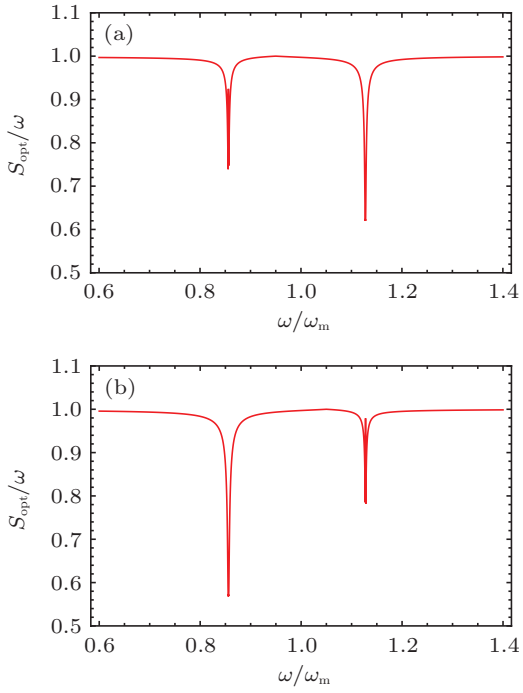


Fig. 5. (color online) The squeezing spectrum with different frequencies of NMOs when $\lambda = 0.5\lambda_0$, $\wp_{\rm p} = 10$ mW, $T_1 = 2$ mK, $T_2 = 3$ mK. (a) $\omega_1 = \omega_{\rm m} + 0.1 \times 2\pi \times 10^7$ Hz, $\omega_2 = \omega_{\rm m} - 0.1 \times 2\pi \times 10^7$ Hz; (b) $\omega_1 = \omega_{\rm m} - 0.1 \times 2\pi \times 10^7$ Hz, $\omega_2 = \omega_{\rm m} + 0.1 \times 2\pi \times 10^7$ Hz.

4. Conclusion

In conclusion, we analyzed the properties of the ponderomotive squeezing in an optomechanical system with two charged NMOs. We show that a squeezed light can be achievable in this system due to the introduction of Coulomb interaction between two NMOs. We also find that the degree of squeezing can be flexibly controlled by adjusting the Coulomb coupling strength, the laser power, and the frequencies of the two NMOs. In addition, we demonstrate the splitting of the two optimal squeezing points varies near linearly with the Coulomb interaction, which might be applied to measure the Coulomb coupling strength.

References

[1] Kippenberg T J and Vahala K J 2008 *Science* **321** 1172
[2] Marquardt F and Girvin S M 2009 *Physics* **2** 40
[3] Verlot P, Tavernarakis A, Briant T, Cohadon P F and Heidmann A 2010 *Phys. Rev. Lett.* **104** 133602
[4] Mahajan S, Kumar T, Bhattacherjee A B and ManMohan 2013 *Phys. Rev. A* **87** 013621
[5] Huang S and Agarwal G S 2011 *Phys. Rev. A* **83** 043826
[6] Han Y, Cheng J and Zhou L 2011 *J. Phys. B: At. Mol. Opt. Phys.* **44** 165505
[7] Ma Y H and Zhou L 2013 *Chin. Phys. B* **22** 024204
[8] Karuza M, Biancofiore C, Bawaj M, Molinelli C, Galassi M, Natali R, Tombesi P, Diuseppe G and Vitali D 2013 *Phys. Rev. A* **88** 013804
[9] Shu J 2011 *Chin. Phys. Lett.* **28** 104203
[10] Yan X B, Yang L, Tian X D, et al. 2014 *Acta Phys. Sin.* **63** 204201 (in Chinese)
[11] Yan X B, Gu K H, Fu C B, et al. 2014 *Chin. Phys. B* **23** 114201
[12] Rocheleau T, Ndukum T, Machlin C, Hertzberg J, Clerk A and Schwab K 2010 *Nature* **463** 72
[13] Miao H, Danilishin S, Müller-Ebhardt H and Chen Y 2010 *New J. Phys.* **12** 083032
[14] Liu Y C, Xiao Y F, Luan X S and Wong C W 2015 *Sci. China-Phys. Mech. Astro* **58** 050305
[15] Yan Y, Gu W J and Li G X 2015 *Sci. China-Phys. Mech. Astro.* **58** 050306
[16] Vitali D, Gigan S, Ferreira A, Böhm H R, Tombesi P, Guerreiro A, Vedral V, Zeilinger A and Aspelmeyer M 2007 *Phys. Rev. Lett.* **98** 030405
[17] Palomaki T A, Teufel J D, Simmonds R W and Lehnert K W 2013 *Science* **342** 710
[18] Joshi C, Akram U and Milburn G J 2014 *New J. Phys.* **16** 023009
[19] Liao J Q, Wu Q Q and Nori F 2014 *Phys. Rev. A* **89** 014302
[20] Ghobadi R, Kumar S, Pepper B, Bouwmeester D, Lvovsky A I and Simon C 2014 *Phys. Rev. Lett.* **112** 080503
[21] Zheng Q and Zhang D 2013 *Chin. Phys. Lett.* **30** 024213
[22] Wu Q, Xiao Y and Zhang Z M 2015 *Chin. Phys. B* **24** 104208
[23] Woolley M J, Doherty A C, Milburn G J and Schwab K C 2008 *Phys. Rev. A* **78** 062303
[24] Sete E A and Eleuch H 2012 *Phys. Rev. A* **85** 043824
[25] Asjad M, Agarwal G S, Kim M S, Tombesi P, Di Giuseppe G and Vitali D 2014 *Phys. Rev. A* **89** 023849
[26] Abadie J, Abbott B P, Abott T D, *et al.* 2011 *Nat. Phys.* **7** 962
[27] Taylor M A, Janousek J, Daria V, Knittel J, Hage B, Bachor H A and Bowen W P 2013 *Nat. Photon.* **7** 229
[28] Braunstein S L and Van L P 2005 *Rev. Mod. Phys.* **77** 513
[29] Eberle T, Steinlechner S, Bauchrowitz J, Händchen V, Vahlbruch H, Mehmet M, Müller-Ebhardt H and Schnabel R 2010 *Phys. Rev. Lett.* **104** 251102
[30] Braginsky V and Manukin A 1967 *Sov. Phys. JETP* **25** 653
[31] Fabre C, Pinard M, Bourzeix S, Heidmann A, Giacobino E and Reynaud S 1994 *Phys. Rev. A* **49** 1337
[32] Mancini E and Tombesi P 1994 *Phys. Rev. A* **49** 4055
[33] Mancini S, Mankó V I and Tombesi P 1997 *Phys. Rev. A* **55** 3042

[34] Brooks D W C, Botter T, Schreppler S, Purdy T P, Brahms N and Stamper-Kum D M 2012 *Nature* **448** 476
[35] Safavi-Naeini A H, Gröblacher S, Hill J T, Chan J, Aspelmeyer M and Painter O 2013 *Nature* **500** 185
[36] Purdy T P, Yu P L, Peterson R W, Kampel N S and Regal C A 2013 *Phys. Rev. X* **3** 031012
[37] Xiao Y, Yu Y F and Zhang Z M 2014 *Opt. Express* **22** 017979
[38] Min F X, Xiao X Y, Yu Y F and Zhang Z M 2015 *Chin. Phys. B* **24** 050301
[39] Ma P C, Zhang J Q, Xiao Y, Feng M and Zhang Z M 2014 *Phys. Rev. A* **90** 043825
[40] Chen R X, Shen L T and Zheng S B 2015 *Phys. Rev. A* **91** 022326
[41] Hensinger W K, Utami D W, Goan H S, Schwab K, Monroe C and Milburn G J 2005 *Phys. Rev. A* **72** 041405(R)
[42] Agarwal G S and Huang M 2012 *Phys. Rev. A* **85** 021801(R)
[43] Genes C, Vitali D, Tombei D, Gigan S and Aspelmeyer M 2008 *Phys. Rev. A* **77** 033804
[44] Kleckner D, Marshall W, de Dood Michiel J A, Dniyari K N, Pors B J, Irvine W T M and Bouwmeester D 2006 *Phys. Rev. Lett.* **96** 173901
[45] Teufel J D, Li D, Allman M S, Cicak K, Sirois A J, Whittaker J D and Simmonds K R W 2011 *Nature* **471** 204