Tunable double optomechanically induced transparency in an optomechanical system

Peng-Cheng Ma1,2,3, Jian-Qi Zhang2†, Yin Xiao1, Mang Feng2‡ and Zhi-Ming Zhang1‡
1Laboratory of Nanophotonic Functional Materials and Devices (SIPSE), and Laboratory of Quantum Engineering and Quantum Materials, South China Normal University, Guangzhou 510006, China
2State Key Laboratory of Magnetic Resonance and Atomic and Molecular Physics, Wuhan Institute of Physics and Mathematics, Chinese Academy of Sciences, Wuhan 430071, China
3School of Physics and Electronic Electrical Engineering, Huaiyin Normal University, Huaian 223300, China

We study the dynamics of a driven optomechanical cavity coupled to a charged nanomechanical resonator via Coulomb interaction, in which the tunable double optomechanically induced transparency (OMIT) can be observed from the output field at the probe frequency by controlling the strength of the Coulomb interaction. We calculate the splitting of the two transparency windows, which varies near linearly with the Coulomb coupling strength in a robust way against the cavity decay. Our double-OMIT is much different from the previously mentioned double-EIT or double-OMIT, and might be applied to measure the Coulomb coupling strength.

PACS numbers: 42.50.Wk, 46.80.+j, 41.20.Cv

I. INTRODUCTION

Recently, significant theoretical and experimental efforts have been paid on studying the characteristic and application of nanomechanical resonators (NRs) 1–4. NRs own some important properties, such as phonon induced transparency 5, phonon blockade 6, and high harmonic generation 7, and can be employed in many applications as, for example, single photon source 8, single phonon source 9, biological sensor 10, quantum information processing 11, and quantum metrology 12,13.

In combination with an optical cavity, an NR turns to be an optomechanical system 14–17, in which the NR interacts with the cavity mode via the radiation pressure force and enables observation of the NR-induced quantum mechanical behaviors from the output light of the cavity. Until now, there have been a lot of theoretical predictions in such optomechanical systems, for example, photon blockade 18, Kerr effect 19, optomechanically induced transparency (OMIT) 20, quantum information transfer 21, normal-mode splitting 22, and some of them have been demonstrated experimentally, such as, OMIT 23–26, slow light 27, frequency transfer 27, and normal-mode splitting 28.

The present work is focused on the OMIT effect in the optomechanical cavity. The OMIT is a kind of induced transparency caused by the radiation pressure in an optomechanical system 23–26, which stands at the center of current studies for optomechanics. We have noticed recent OMIT-relevant work on four-wave mixing 29, superluminal and ultraslow light propagation 30,31, quantum router 32, and precision measurement of electrical charge 33. On the other hand, double electromagnetic induced transparency (EIT) 34–36 is a hot topic over recent years, which extends conventional EIT to the one with double transparency windows, and discovers some new physics and applications. This arises a question: what would happen in an OMIT with two transparency windows (i.e., double-OMIT)? To the best of our knowledge, there have been a few theoretical schemes 37–39 for the double-OMIT with different models, using a nonlinear crystal or a qubit in an optomechanical cavity 37,38, and using a ring cavity with two movable mirrors 39. However, in all the schemes mentioned above, the frequency of the transparency light for the double-EIT/OMIT cannot be changed due to the fixed coupling for splitting the transparency windows.

In this work, we demonstrate a tunable double-OMIT observable in an optomechanical system, in which the two NRs are charged and the two transparency windows are split due to the Coulomb interaction. Specifically, our optomechanical system consists of an optomechanical cavity and a NR outside, as sketched in Fig. 1, where the NR of the optomechanical cavity (i.e., NR1) not only couples to the cavity field by the radiation pressure, but also interacts with the NR outside the cavity (i.e., NR2) through the tunable Coulomb interaction, which can be controlled by the bias voltages on the NRs.

Compared with the conventional OMIT with a single transparency window 23–26, our scheme owns some favorable features: (i) The two output lights with different frequencies are controlled by a single driving light; (ii) Our scheme is robust to the cavity decay, and the transparency windows are with narrow profiles; (iii) We find that the two windows of the double-OMIT are apart near linearly with respect to the Coulomb coupling strength. The feature reminds us of a practical application of the double-OMIT for precisely detecting the Coulomb coupling strength. In this context, we have to emphasize...
that our proposal is essentially different from the previous ideas [38, 39], where the double-OMIT is caused by the frequency difference between the two NRs and the frequencies of the transparency lights are fixed. In contrast, our studied double-OMIT can be observed even for two identical NRs, and the frequency of the transparency light can be selected by tuning the Coulomb coupling under a constant driving light.

This paper is structured as follows. In Sec. II we present the model and the analytical expressions of the optomechanical system and obtain the steady-state mean values. Sec. III includes numerical calculations for the double-OMIT based on recent experimental parameters. The feasibility of precision measurement of the Coulomb coupling strength between the two NRs is discussed in Sec. IV and we also justify the robustness of our approach against the cavity decay. The last section is a brief conclusion.

II. THE MODEL AND THE SOLUTIONS

For the system in Fig. 1, the Hamiltonian is given by,

$$H_{\text{whole}} = \hbar \omega_c c^\dagger c + \left( \frac{p_1^2}{2m_1} + \frac{1}{2} m_1 \omega_1^2 q_1^2 \right) + \left( \frac{p_2^2}{2m_2} + \frac{1}{2} m_2 \omega_2^2 q_2^2 \right) - \hbar g c^\dagger q_1 + H_C$$

$$+ i \hbar \varepsilon_1 (c^\dagger e^{-i \omega_1 t} - h. c.) + i \hbar (c^\dagger \varepsilon_p e^{-i \omega_p t} - h. c.),$$

where the first term is for the single-mode cavity field with frequency $\omega_c$ and annihilation (creation) operator $c$ ($c^\dagger$). The second (third) term describes the vibration of the charged NR$_1$ (NR$_2$) with frequency $\omega_1$ ($\omega_2$), effective mass $m_1$ ($m_2$), position $q_1$ ($q_2$) and momentum operator $p_1$ ($p_2$). NR$_1$ couples to the cavity field due to the radiation pressure with the coupling strength $g = \frac{\hbar}{L}$ with $L$ being the cavity length.

The fifth term $H_C$ in Eq. (1) presents the Coulomb coupling between the charged NR$_1$ and NR$_2$ [40], where the NR$_1$ and NR$_2$ take the charges $C_1 V_1$ and $-C_2 V_2$, with $C_1(C_2)$ and $V_1(-V_2)$ being the capacitance and the voltage of the bias gate, respectively. So the Coulomb coupling between NR$_1$ and NR$_2$ is given by

$$H_C = -\frac{C_1 V_1 C_2 V_2}{4 \pi \varepsilon_0 r_0} (q_1 - q_2),$$

where $r_0$ is the equilibrium distance between NR$_1$ and NR$_2$, $q_1$ and $q_2$ represent the small displacements of NR$_1$ and NR$_2$ from their equilibrium positions, respectively. In the case of $r_0 \gg q_1, q_2$, with the second order expansion, the Hamiltonian above is rewritten as

$$H_C = -\frac{C_1 V_1 C_2 V_2}{4 \pi \varepsilon_0 r_0} \left[ 1 - \frac{q_1 - q_2}{r_0} + \left( \frac{q_1 - q_2}{r_0} \right)^2 \right],$$

where the linear term may be absorbed into the definition of the equilibrium positions, and the quadratic term includes a renormalization of the oscillation frequency for both NR$_1$ and NR$_2$. This implies a reduced form

$$H_C = \hbar \lambda q_1 q_2,$$

where $\lambda = \frac{C_1 V_1 C_2 V_2}{2 \pi \varepsilon_0 r_0 \hbar \omega_c} [41, 42].$

The last two terms in Eq. (1) describe the interactions between the cavity field and the two input fields, respectively. The strong (week) pump (probe) field owns the frequency $\omega_p$ ($\omega_p$) and the amplitude $\varepsilon_1 = \sqrt{2 \kappa \delta \rho / \hbar \omega_c}$ ($\varepsilon_p = \sqrt{2 \kappa \rho_p / \hbar \omega_c}$), where $\rho$ ($\rho_p$) is the power of the pump (probe) field and $\kappa$ is the cavity decay rate.

In a frame rotating with the frequency $\omega_p$ of the pump field, the Hamiltonian of the total system Eq.(1) can be rewritten as,

$$H = \hbar \Delta_c c^\dagger c + \left( \frac{p_1^2}{2m_1} + \frac{1}{2} m_1 \omega_1^2 q_1^2 \right)$$

$$+ \left( \frac{p_2^2}{2m_2} + \frac{1}{2} m_2 \omega_2^2 q_2^2 \right) - \hbar g c^\dagger q_1 + \hbar \lambda q_1 q_2$$

$$+ i \hbar \varepsilon_1 (c^\dagger e^{-i \omega_2 t} - h. c.) + i \hbar (c^\dagger \varepsilon_p e^{-i \omega_p t} - h. c.),$$

(2)

where $\Delta_c = \omega_c - \omega_1$ is the detuning of the pump field from the bare cavity, and $\delta = \omega_p - \omega_1$ is the detuning of the probe field from the pump field. Considering photon losses from the cavity and the Brownian noise from the environment, we may describe the dynamics of the system governed by Eq. (2) using following nonlinear quantum Langevin equations [32],

$$\dot{q}_1 = \frac{p_1}{m_1},$$

$$\dot{p}_1 = -m_1 \omega_1^2 q_1 - \hbar \lambda q_2 + \hbar g c^\dagger c - \gamma_1 p_1 + \sqrt{2 \gamma_1} \xi_1(t),$$

$$\dot{q}_2 = \frac{p_2}{m_2},$$

$$\dot{p}_2 = -m_2 \omega_2^2 q_2 - \hbar \lambda q_1 - \gamma_2 p_2 + \sqrt{2 \gamma_2} \xi_2(t),$$

$$\dot{c} = -[\kappa + i(\Delta_c - \gamma_1)] c + \varepsilon_1 + \varepsilon_p e^{-i \omega_p t} + \sqrt{2 \kappa c_{in}},$$

(3)
where $\gamma_1$ and $\gamma_2$ are the decay rates for NR$_1$ and NR$_2$, respectively. The quantum Brownian noise $\xi_1$ ($\xi_2$) comes from the coupling between NR$_1$ (NR$_2$) and its own environment with zero mean value. $\epsilon_o$ is the input vacuum noise operator with zero mean value. Under the mean field approximation $\langle Qc \rangle = \langle Q \rangle \langle c \rangle$, the mean value equations are given by

$$
\langle q_1 \rangle = \frac{\langle p_1 \rangle}{m_1}, \quad \langle p_1 \rangle = -m_1 \omega_1^2 \langle q_1 \rangle - \hbar \lambda \langle q_2 \rangle + \hbar g \langle \phi \rangle \langle c \rangle - \gamma_1 \langle p_1 \rangle, \quad \langle q_2 \rangle = \frac{\langle p_2 \rangle}{m_2}, \quad \langle p_2 \rangle = -m_2 \omega_2^2 \langle q_2 \rangle - \hbar \lambda \langle q_1 \rangle - \gamma_2 \langle p_2 \rangle, \quad \langle \phi \rangle = -[\hbar \kappa + i(\Delta_c - g \langle q_1 \rangle)] \langle c \rangle + \epsilon_l + \epsilon_p e^{-i\delta t}, \quad (4)
$$

which is a set of nonlinear equations and the steady-state response in the frequency domain is composed of many frequency components. We suppose the solution with the following form

$$
\langle q_1 \rangle = q_{1s} + q_1 e^{-i\delta t} + q_1 e^{i\delta t}, \quad \langle p_1 \rangle = p_{1s} + p_1 e^{-i\delta t} + p_1 e^{i\delta t}, \quad \langle q_2 \rangle = q_{2s} + q_2 e^{-i\delta t} + q_2 e^{i\delta t}, \quad \langle p_2 \rangle = p_{2s} + p_2 e^{-i\delta t} + p_2 e^{i\delta t}, \quad \langle c \rangle = c_s + c_+ e^{-i\delta t} + c_- e^{i\delta t}, \quad (5)
$$

where each quantity contains three items $O_s$, $O_-$, and $O_+$ (with $O \in \{q_1, q_2, p_1, p_2, c\}$), corresponding to the responses at the frequencies $\omega_1$, $\omega_p$, and $2\omega_c - \omega_p$, respectively. In the case of $O_s \gg O_\pm$, Eq. (4) can be solved by treating $O_\pm$ as perturbations. After substituting Eq. (5) into Eq. (4), and ignoring the second-order terms, we obtain the steady-state mean values of the system as

$$
P_{1s} = P_{2s} = 0, \quad \epsilon_1 = \frac{\hbar g |c_s|^2}{m_1 \omega_1^2 - \hbar^2 \lambda^2 / m_2 \omega_2^2}, \quad \epsilon_2 = \frac{\hbar \lambda g_{1s}}{m_2 \omega_2^2}, \quad \epsilon_l = \frac{\hbar \lambda g_{1s}}{m_2 \omega_2^2}, \quad |c_s|^2 = \frac{|\epsilon_l|}{\Delta^2 + \kappa^2},
$$

with $\Delta = \Delta_c - g \langle q_{1s} \rangle$, and

$$
c_+ = \frac{\kappa - i(\Delta + \delta)|\delta^2 - \omega_1^2 + i\delta \gamma_1| - G - 2i\omega_1 \beta}{\Delta^2 - (\delta + i\kappa)^2}|\delta^2 - \omega_1^2 + i\delta \gamma_1| - G + 4\Delta \omega_1 \beta, \quad (7)
$$

where $\beta = \frac{\epsilon_l^2 \lambda^2}{2 \omega_2}$ and $G = \frac{\hbar^2 \lambda^2}{m_2 (\delta^2 - \omega_1^2 + i\delta \gamma_1)}$. When there is no Coulomb coupling $\lambda$ (i.e., $G = 0$) between the two NRs, Eq. (7) is reduced to Eq. (5) in Ref. 20. However, different from the output field in Ref. 20 involving a single center frequency for the single-mode OMIT, there are two centers with different frequencies in our scheme due to the Coulomb interaction. As a result, under the actions of the radiation pressure and the probe light, two OMITs with different centers are reconstructed, yielding the double-OMIT.

Making use of the input-output relation of the cavity

$$
\epsilon_{out}(t) + \epsilon_p e^{-i\delta t} + \epsilon_l = 2\kappa |c|, \quad (6)
$$

and

$$
\epsilon_{out}(t) = \epsilon_{out} + \epsilon_{out} \epsilon_p e^{-i\delta t} + \epsilon_{out} \epsilon_p e^{i\delta t}, \quad (8)
$$

we obtain

$$
\epsilon_{out} = 2\kappa |c| + 1 = 2\kappa |c|, \quad (9)
$$

which can be measured by homodyne technique. This output light $\epsilon_{out}$ is of the same frequency $\omega_p$ as the probe field. Defining

$$
\epsilon_T = \epsilon_{out} + 1 = 2\kappa |c|, \quad (10)
$$

yields the real and imaginary parts, with $Re[\epsilon_T]$ and $Im[\epsilon_T]$, representing the absorption and dispersion of the optomechanical system, respectively.

### III. DOUBLE-OMIT IN THE OUTPUT FIELD

We present below the feasibility of the tunable double-OMIT in the optomechanical system, and the relationship between the double-OMIT and the Coulomb interaction between the two NRs. As an estimate for Eq. (8), we employ the parameters from the recent experiment
in the observation of the normal-mode splitting. For simplicity, we first consider two identical NRs in our numerics, which is not essentially different in physics from the case of two different NRs. We will also treat the different NRs later.

As shown in Fig. 2, the absorption \( Re[\varepsilon_T] \) and dispersion \( Im[\varepsilon_T] \) of the output field are plotted as functions of \( \delta_c/\omega_1 = (\delta - \omega_1)/\omega_1 \) for different Coulomb couplings. We may find that the output lights for the probe field behave from the double-OMIT to the single-mode OMIT with diminishing Coulomb coupling. The physics behind the double-OMIT phenomenon can be understood from the interference \[23, 46\] and the level configuration in Fig. 3.

The OMIT originates from the radiation pressure coupling an optical mode to a mechanical mode. The simultaneous presence of the pump and probe fields generates a radiation-pressure force oscillating at the frequency difference \( \delta = \omega_p - \omega_1 \). If this frequency difference is close to the resonance frequency \( \omega_2 \) of NR\(_1\), the mechanical mode starts to oscillate coherently. This in turn gives rise to the Stokes- and anti-Stokes scattering of light from the strong pump field. If the system is operated within the resolved-sideband regime \( \kappa \ll \omega_1 \), the Stokes scattering is strongly suppressed since it is highly off-resonant with the optical cavity. We can therefore assume that only an anti-Stokes field with frequency \( \omega_p = \omega_1 + \omega_1 \) builds up inside the cavity. However, since this field is degenerate with the probe field sent into the cavity, the two fields interfere destructively, suppressing the case of a single transparency window for the output beam. Thus the OMIT occurs. As it depends on quantum interference, the OMIT is sensitive to phase disturbances. The coupling between NR\(_1\) and NR\(_2\) not only adds a fourth level, as shown in Fig. 3, but also breaks down the symmetry of the OMIT interference, and thereby produces a spectrally sharp bright resonance within the OMIT line shape. Then the single OMIT transparency window is split into two transparency windows, which yields the double-OMIT.

**IV. MEASUREMENT OF THE COUPLING STRENGTH BETWEEN NR\(_1\) AND NR\(_2\)**

To further explore the characteristic of the tunable double-OMIT, we plot the absorption \( Re[\varepsilon_T] \) as functions of \( \delta_c/\omega_1 \) and \( \lambda \). One can find from Fig. 4 that only a single transparency window appears at \( \delta_c = 0 \) for \( \delta = \omega_1 \) in the absence of the Coulomb coupling, and the single transparency window is split into two transparency windows once the Coulomb coupling is present. The two transparency windows are more and more apart with the increase of \( \lambda \). The two minima of the absorption in Fig. 4 can be evaluated by

\[
\frac{d Re[\varepsilon_T]}{d \delta_c}|_{\delta_c = \delta_{c+}} = 0, \quad \frac{d Re[\varepsilon_T]}{d \delta_c}|_{\delta_c = \delta_{c-}} = 0, \quad \delta_{c+} \quad \text{and} \quad \delta_{c-} \quad \text{are the points with absorption minima.}
\]

So the separation of the minima is \( d = |\delta_{c+} - \delta_{c-}| \), as plotted in Fig. 5, where the almost linear increase of \( d \) with \( \lambda \) within the regime \( \lambda = \{0, 15\lambda_0\} \) reminds us of the possibility to detect the Coulomb coupling strength between NR\(_1\) and NR\(_2\) by measuring the separation \( d \) in the absorption spectrum \( Re[\varepsilon_T] \) of the output field. From Fig. 5, one can calculate the measuring sensitivity by \( \delta_{c+}/\lambda \) on the order of \( 10^{-31} \) m\(^2\). Considering a Coulomb coupling change \( \Delta F \) due to a slight deviation \( q_1 \), we have \( \Delta F = \hbar q_1 \). Provided \( q_1 \approx 0.1 \) nm, we may assess \( \partial \Delta F/\partial d \) to be of the order of \( 10^{-13} \) N/Hz, implying the possible precision of measuring \( \Delta F \) decided by the resolution of \( d \) in the absorption spectrum.

Fig. 6 presents the variation of the absorption \( Re(\varepsilon_T) \) with respect to \( \delta_c/\omega_1 \) for different cavity decay rates, where the maxima (i.e., the points A, B and C) and the minima (i.e., the points D and E) of the curves remain unchanged in the parameter changes, but the profiles of the transparency window become narrower and sharper with the cavity decay rate \( \kappa \) increasing. Provided a fixed
become wider and wider with strength of the radiation pressure, which is reflected in the radiation pressure and makes it less precise in detecting the driving light, the bigger cavity decay will disturb the radiation pressure. The spectrum of the output becomes narrower for the larger displacement of the NR. In this way, our scheme can be robust against the cavity decay rate.

Moreover, for two different NRs, the results will be slightly different from those above for identical NRs. Considering \( \omega_1 \neq \omega_2 \) in the calculation, we have plotted in Fig. 7 the absorption of the double-OMIT with larger separations of the minima in comparison with the identical NR case. It implies a more sensitivity to the coupling strength \( \lambda \) in the case of two different NRs. With respect to the \( \omega_1 = \omega_2 \) case, the absorption curves move rightward (leftward) in the case of \( \omega_1 > \omega_2 \) (\( \omega_1 < \omega_2 \)). The enhancement of the sensitivity to the Coulomb force can be calculated by Eq. (9) and \( d = |\delta_{\omega_+} - \delta_{\omega_-}| \), and is exemplified in Fig. 7 as 1.139 (1.529) times using \( \omega_2 = 1.1\omega_1 \) (\( \omega_2 = 0.8\omega_1 \)).

We have to mention that the robustness discussed above is limited within the resolved regime (\( \kappa < \omega_1 \)) where the double-OMIT works. In contrast, the unresolved regime (\( \kappa > \omega_1 \)) blurs the sideband transitions, which makes the quantum interference unavailable.

V. CONCLUSION

In conclusion, we have demonstrated the feasibility of the tunable double-OMIT in the optomechanical system under the Coulomb interaction between two charged NRs. To our knowledge, this is the first proposal for the tunable double-OMIT in the optomechanical system. Although our proposal is in principle extendable to other interactions, such as the dipole-dipole coupling, the Coulomb coupling, as a long-range interaction, is easier to control, and thereby more practical. We have to emphasize that our double-OMIT is neither a simple extension of the conventional OMIT nor a simple transformation of the previously discussed double-EIT. Due to narrow profiles of the transparency windows and robustness against dissipation, the double OMIT might be employed...
for precisely detecting the Coulomb coupling strength. Therefore, we argue that our scheme have paved a new avenue towards the study of the OMIT with more transparency windows as well as the relevant application.

ACKNOWLEDGMENTS

PCM thanks Lei-Lei Yan and Wan-Lu Song for their helps in the numerical simulation. JQZ thanks Yong Li for the helpful discussion. This work was supported by the "973" Program (Grant No. 2011CBA00200, No. 2012CB922102 and No. 2013CB921804), the Major Research Plan of the NSFC (Grant No. 91121023), the NSFC (Grants No. 61378012, No. 60978009, No. 11274352 and No. 11304366), the SRFDPHEC (Grant No. 20124407110009), the PCSIRT (Grant No. IRT1243), China Postdoctoral Science Foundation (Grant No. 2013M531771 and No. 2014T70760), Natural Science Funding for Colleges and Universities in Jiangsu Province (Grant No. 12KJD140002), and Program for Excellent Talents of Huaiyin Normal University(No. 11HSQNZ07).

[1] K. C. Schwab and M. L. Roukes, Phys. Today 58, 36 (2005).
[2] X. Zhou, F. Hocke, A. Schliesser, A. Marx, H. Huebl, R. Gross, and T. J. Kippenberg, Nat. Phys. 9, 179 (2013).
[3] P. C. Ma, Y. Xiao, Y. F. Yu, and Z. M. Zhang, Opt. Express 22, 3621 (2014).
[4] J. Q. Zhang, S. Zhang, J. H. Zou, L. Chen, W. Yang, Y. Li, and M. Feng, Opt. Express 21, 29695 (2013).
[5] H. Okamoto, A. Gourgout, C. Y. Chang, K. Onomitsu, I. Mahboob, E. Y. Chang, and H. Yamaguchi, Nat. Phys. 9, 480 (2013).
[6] Y. X. Liu, A. Miranowicz, Y. B. Gao, J. Bajer, C. P. Sun, and F. Nori, Phys. Rev. A 82, 032101 (2010).
[7] H. Xiong, L. G. Si, X. Y. Lü, X. X. Yang, and Y. Wu, Opt. Lett. 38, 353 (2013).
[8] L. Qiu, L. Gan, W. Ding, and Z. Y. Li, J. Opt. Soc. Am. B 30, 1683 (2013).
[9] K. V. Kepesidis, S. D. Bennett, S. Portolan, M. D. Lukin, and P. Rabl, Phys. Rev. A 88, 064105 (2013).
[10] K. Eom, H. S. Park, D. S. Yoon, and T. Kwon, Phys. Rep. 503, 115 (2011).
[11] P. Rabl, S. J. Kolkowitz, F. H. L. Kopplens, J. G. E. Harris, P. Zoller, and M. D. Lukin, Nat. Phys. 6, 602 (2010).
[12] W. He, J. J. Li and K. D. Zhu, Opt. Lett. 35, 339 (2010).
[13] J. J. Li and K. D. Zhu, Appl. Phys. Lett. 94, 063116 (2009).
[14] T. J. Kippenberg and K. J. Vahala, Opt. Express 15, 17172 (2007).
[15] T. J. Kippenberg and K. J. Vahala, Science 321, 1172 (2008).
[16] F. Marquardt and S. M. Girvin, Physics 2, 40 (2009).
[17] M. Aspelmeyer, P. Meystre, and K. Schwab, Phys. Today 65, 29 (2012).
[18] P. Rabl Phys. Rev. Lett. 107, 063601 (2011).
[19] Z. R. Gong, H. Ian, Y. X. Liu, C. P. Sun, and F. Nori, Phys. Rev. A 80, 065801 (2009).
[20] G. S. Agarwal and S. Huang, Phys. Rev. A 81, 041803 (2010).
[21] C. H. Dong, V. Fiore, M. C. Kuzyk, and H. L. Wang, Science 338, 1609 (2012).
[22] J. M. Dobrinint, I. Wilson-Rae, and T. J. Kippenberg, Phys. Rev. Lett. 101, 263602 (2008).
[23] S. Weiss, R. Rivière, S. Deléglise, E. Gavartin, O. Arcizet, A. Schliesser, and T. J. Kippenberg, Science 330, 1520 (2010).
[24] A. H. Safavi-Naeini, T. P. M. Alegre, J. Chan, M. Eichenfield, M. Winger, Q. Lin, J. T. Hill, D. Chang, and O. Painter, Nature (London) 472, 69 (2011).
[25] M. Karuza, C. Biancofiore, M. Bawaj, C. Molinelli, M. Galassi, R. Natali, P. Tombesi, G. Di Giuseppe, and D. Vitali, Phys. Rev. A 88, 013804 (2013).
[26] H. Xiong, L. G. Si, A. S. Zheng, X. Yang and Y. Wu, Phys. Rev. A 86, 013815 (2012).
[27] J. T. Hill, A. H. Safavi-Naeini, J. Chan, and O. Painter, Nat. Commun. 3, 1196 (2012).
[28] S. K. Gröblacher Hammerer, M. R. Vanner, and M. Aspelmeyer, Nature (London) 460, 724 (2009).
[29] S. Huang and G. S. Agarwal, Phys. Rev. A 81, 033830 (2010).
[30] J. D. Teufel, D. Li, M. S. Allman, K. Cicak, A. J. Sirois, J. D. Whittaker, and R. W. Simmonds, Nature (London) 471, 204 (2011).
[31] D. Tarhan, S. Huang, and Ö. E. Müstecaplioglu, Phys. Rev. A 87, 013824 (2013).
[32] G. S. Agarwal and S. Huang, Phys. Rev. A 85, 021801(R) (2012).
[33] J. Q. Zhang, Y. Li, M. Feng, and Y. Xu, Phys. Rev. A 86, 053806 (2012).
[34] Z. B. Wang, K. P. Marzlin, and B. C. Sanders, Phys. Rev. Lett. 97, 063901 (2006).
[35] B. W. Shiau, M. C. Wu, C. C. Lin, and Y. C. Chen, Phys. Rev. Lett. 106, 193006 (2011).
[36] S. J. Li, X. D. Yang, X. M. Cao, C. H. Zhang, C. D. Xie, and H. Wang, Phys. Rev. Lett. 101, 073602 (2008).
[37] S. Shahidani, M. H. Naderi, and M. Soltanolkotabi, Phys. Rev. A 86, 053813 (2013).
[38] H. Wang, X. Gu, Y. X. Liu, A. Miranowicz and F. Nori, arXiv:1402.2764 (2014).
[39] S. Huang, J. Phys. B: At. Mol. Opt. Phys. 47, 055504 (2014).
[40] W. K. Hensinger, D. W. Utami, H. S. Goan, K. Schwab, C. Monroe, and G. J. Milburn, Phys. Rev. A 72, 041305(R) (2005).
[41] C. N. Ren, J. Q. Zhang, L. B. Chen and Y. J. Gu, arXiv:1402.6434 (2014).
[42] L. Tian and P. Zoller, Phys. Rev. Lett. 93, 266403 (2004).
[43] C. Genes, D. Vitali, P. Tombesi, S. Gigan, and M. Aspelmeyer, Phys. Rev. A 77, 033804 (2008).
[44] S. Huang and G. S. Agarwal, Phys. Rev. A 83, 032823 (2011).
[45] D. F. Walls and G. J. Milburn, Quantum Optics (Springer-Verlag, Berlin) (1994).
[46] J. A. Sedlacek, A. Schwettmann, H. Kübler, R. Löw, T.
Pfau, and J. P. Shaffer, Nature Phys. 8 819 (2012).