All-optical transistor based on a cavity optomechanical system with a Bose-Einstein condensate

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We propose a scheme of an all-optical transistor based on a coupled Bose-Einstein condensate cavity system. The calculated results show that, in such an optomechanical system, the transmission of the probe beam is strongly dependent on the optical pump power. Therefore, the optical pump field can serve as a "gate" field of the transistor, effectively controlling the propagation of the probe field (the "signal" field). The scheme proposed here may have potential applications in optical communication and quantum information processing.

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I. INTRODUCTION

An all-optical transistor [1–3] is a device by which the propagation of the optical "signal" field can be effectively controlled by another optical "gate" field. Due to the fact that the photonic signals are faster than their electronic counterparts and that they do not perturb each other, the all-optical transistor has many interesting applications in optical communication [4], optical quantum computation [5], quantum information processing [6], and so on. However, because of the weakness of the photon-photon interaction, the realization of an all-optical transistor is challenging. Recently, Chang *et al*. [1] have put forward an approach to realizing strong nonlinear interactions between single photons by exploiting the strong coupling between an individual optical emitter and surface plasmons confined to a conducting nanowire, and by using such a system they have realized a single-photon transistor. With a single dye molecule, Hwang *et al*. [2] have demonstrated that a tightly focused laser beam (the signal beam) can coherently be attenuated or amplified by effectively regulating the second gate beam, which controls the degree of the population inversion. Therefore, such a molecule can also operate as an all-optical transistor. Most recently, Li and Zhu [3] also proposed a quantum optical transistor consisting of a single quantum dot embedded in a photonic crystal nanocavity.

Today, a cavity optomechanical system, generally composed of an optical cavity with one movable end mirror [7–10] or with a micromechanical membrane inside it [11,12], has become an attractive subject. In this system, the interaction between a single mode of the optical cavity and the mechanical resonator induced by the radiation pressure can effectively be controlled by an external optical pump laser. Therefore, a strong-coupling regime can easily be reached just by regulating the pump power [7]. With such a strong coupling, cooling a micromechanical resonator to its ground state also becomes possible using dynamical backaction [8,9]. Most importantly, in the strong-coupling optomechanical system, the motion of the oscillator can easily modulate the optical path length and thus the frequency of the cavity, creating sidebands above (the anti-Stokes sideband) and below (the Stokes sideband) the drive frequency. For a near-resonant probe laser, due to its interference with the induced sidebands, the propagation of the probe beam can be strongly altered by the optomechanical system, which has been demonstrated theoretically by Agarwal and Huang [10] and experimentally by Weis *et al*. [13], Safavi-Naeini *et al*. [14], and Teufel *et al*. [15].

In the present paper, we will propose a novel scheme to realize an all-optical transistor based on a new cavity optomechanical system, i.e., a coupled Bose-Einstein condensate (BEC)-cavity system [16–28] consisting of an optical cavity with a BEC trapped in it, which has been proposed and has attracted much attention in recent years. For such a coupled system, the collective motion of the BEC can serve as a mechanical oscillator coupled to the cavity field. On one hand, the cavity field will affect the motion degrees of the freedom of the BEC through the exchange of momentum between the cavity field and the BEC. On the other hand, the collective density excitation of the BEC offers feedback on the cavity field by the dependence of the optical path length on the atomic density distribution within the spatially periodic cavity mode structure [17]. Brennecke *et al*. [16,17] have experimentally realized the strong coupling of the collective oscillation of the BEC to the quantized cavity field with all atoms in the same motional quantum state. In this strong-coupling optomechanical system, a near-resonant probe beam can effectively be modulated by the optical pump beam. Therefore, an all-optical transistor can be expected to be realized based on this strongly coupled BEC-cavity system. The paper is organized as follows. Section II gives the theoretical model and method. Results and discussion are shown in Sec. III. A summary is given by Sec. IV.

II. SYSTEM AND METHOD

The system considered in this paper is shown in Fig. 1, where a BEC of N ^{87}Rb atoms in ground state is trapped within an optical ultrahigh-finesse Fabry-Perot cavity (the finesse is $\mathcal{F} = 3.5 \times 10^5$ [28]) by a crossed-beam dipole trap. The cavity field is driven by a pump-laser field accompanying a weak probe-laser field along the cavity axis. The probe photons transmitted from the cavity are monitored by a detector. In a rotating frame at a driving field frequency $\omega_{\rm pu}$, the Hamiltonian of the BEC-cavity system in the case of weak atom-atom interactions and a shallow external trapping potential can be expressed as follows [17]:

$$H = H_M + H_C + H_{\rm CM} + H_{\rm IN}. \tag{1}$$

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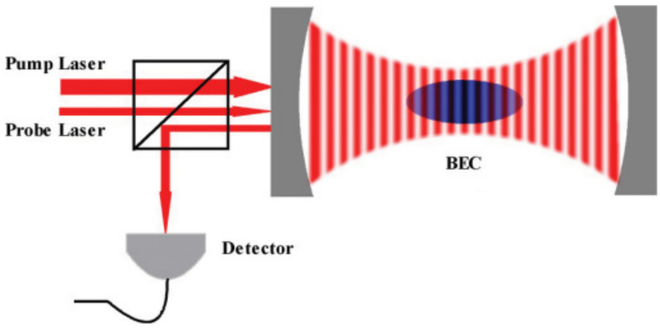

FIG. 1. (Color online) Schematic diagram for a coupled BEC-cavity system in the simultaneous presence of a strong pump laser and a weak probe laser.

The first term presents the energy of the Bogoliubov mode of the collective oscillation of the BEC with $H_M = \hbar\omega_m a^\dagger a$, where ω_m and a ($a^\dagger$) denote the oscillation frequency and the annihilation (creation) operator of the Bogoliubov mode, respectively. As demonstrated experimentally by Brennecke *et al.* [17], the BEC-cavity dynamics can be understood as follows. The macroscopically occupied zero-momentum state is coupled to the symmetric superposition of the $\pm 2\hbar k$ momentum state via absorption and stimulated emission of cavity photons, where k is the wave vector of the cavity photons. This physical process can be well described by the Bogoliubov mode oscillating at the frequency $\omega_m = \frac{1}{\hbar}\frac{(2\hbar k)^2}{2m} = 4 \times \frac{\hbar k^2}{2m} = 4\omega_{\rm rec}$ (m is the mass of the atom and $\omega_{\rm rec}$ is the recoil frequency), with the atom-atom interactions being neglected at this stage [17]. The second term gives the energy of the cavity mode with $H_C = \hbar\Delta_c c^\dagger c$, where c ($c^\dagger$) denotes the annihilation (creation) operator of the cavity mode and $\Delta_c = \omega_c - \omega_{\rm pu} + \frac{1}{2}NU_0$ is a shifted cavity-pump detuning. Here ω_c is the cavity oscillation frequency and $\frac{1}{2}NU_0$ is the frequency shift of the empty cavity resonance induced by the BEC. N is the number of the condensed atoms. $U_0 = \frac{g_0^2}{\Delta_a}$ is the maximum light shift which an atom experiences in the cavity mode, where g_0 is the maximum coupling strength between a single atom and a single polarized intracavity photon, and Δ_a is the detuning between the pump-laser frequency and atomic D_2 line transition frequency. For a large detuning Δ_a, the excited atomic level can be eliminated adiabatically. The third term describes the coupling between the BEC and the cavity with $H_{\rm CM} = \hbar g(a^\dagger + a)c^\dagger c$, where $g = \frac{U_0}{2}\sqrt{\frac{N}{2}}$ is a collectively enhanced coupling strength. The last term shows the classical light inputs including the pump light and probe light with frequencies $\omega_{\rm pu}$ and $\omega_{\rm pr}$, respectively, and $H_{\rm IN} = i\hbar[(\varepsilon_{\rm pu} + \varepsilon_{\rm pr}e^{-i\delta t})c^\dagger - (\varepsilon_{\rm pu} + \varepsilon_{\rm pr}e^{i\delta t})c]$, where $\varepsilon_{\rm pu}$ and $\varepsilon_{\rm pr}$ are related to the laser power P, respectively, by $|\varepsilon_{\rm pu}| = \sqrt{2P_{\rm pu}\kappa/\hbar\omega_{\rm pu}}$ and $|\varepsilon_{\rm pr}| = \sqrt{2P_{\rm pr}\kappa/\hbar\omega_{\rm pr}}$ (here κ is the decay rate of the cavity amplitude) [8,10], and $\delta = \omega_{\rm pr} - \omega_{\rm pu}$ is the probe-pump detuning.

In what follows, we deal with the mean response of the system to the probe field in the presence of the coupling field, and $\langle c\rangle$, $\langle c^\dagger\rangle$, and $\langle Q\rangle$ are the expectation values of operators $c, c^\dagger$, and Q [which is defined as $Q = (a + a^\dagger)/\sqrt{2}$], respectively. Taking the damping terms into consideration and according to the Heisenberg equation of motion and the communication relations $[c,c^\dagger] = 1$ and $[a,a^\dagger] = 1$, the temporal evolutions of a and Q can be obtained and the corresponding equations are given by

$$\frac{d\langle c\rangle}{dt} = -(i\Delta_c + \kappa)\langle c\rangle - i\sqrt{2}g\langle Q\rangle\langle c\rangle + \varepsilon_{\rm pu} + \varepsilon_{\rm pr}e^{-i\delta t}, \tag{2}$$

$$\frac{d^2\langle Q\rangle}{dt^2} + \gamma_m\frac{d\langle Q\rangle}{dt} + \omega_m^2\langle Q\rangle = -\omega_m g\sqrt{2}\langle c^\dagger\rangle\langle c\rangle, \tag{3}$$

where γ_m is the damping rate of the Bogoliubov mode of the collective oscillation of the BEC. To solve these equations, we make the ansatz as follows:

$$\begin{aligned}\langle c(t)\rangle &= c_0 + c_+e^{-i\delta t} + c_-e^{i\delta t},\\ \langle Q(t)\rangle &= Q_0 + Q_+e^{-i\delta t} + Q_-e^{i\delta t}.\end{aligned} \tag{4}$$

Substituting Eq. (4) into Eqs. (2) and (3), respectively, equating terms with the same time dependence, and working to the lowest order in $\varepsilon_{\rm pr}$ but to all orders in $\varepsilon_{\rm pu}$, we can get

$$\begin{aligned}c_0 &= \frac{\varepsilon_{\rm pu}}{i\Delta_c + i\sqrt{2}gQ_0 + \kappa},\\ c_+ &= \frac{\varepsilon_{\rm pr} - i\sqrt{2}gQ_+c_0}{i\sqrt{2}gQ_0 + i\Delta_c + \kappa - i\delta},\\ c_- &= \frac{-i\sqrt{2}gQ_-c_0}{i\sqrt{2}gQ_0 + i\Delta_c + \kappa + i\delta},\end{aligned} \tag{5}$$

and

$$\begin{aligned}Q_0 &= -\frac{g\sqrt{2}|c_0|^2}{\omega_m},\\ Q_+ &= \frac{-\omega_m g\sqrt{2}(c_0^*c_+ + c_-^*c_0)}{\omega_m^2 - i\delta\gamma_m - \delta^2},\\ Q_- &= \frac{-\omega_m g\sqrt{2}(c_0^*c_- + c_+^*c_0)}{\omega_m^2 + i\delta\gamma_m - \delta^2}.\end{aligned} \tag{6}$$

Solving Eqs. (5) and (6), we obtain

$$\omega_0\left[\kappa^2 + \left(\Delta_c - \frac{2g^2}{\omega_m}\omega_0\right)^2\right] = \varepsilon_{\rm pu}^2 \tag{7}$$

and

$$c_+ = \varepsilon_{\rm pr}\left[\frac{(\kappa - i\delta) - i(\Delta_c - C)}{(\kappa - i\delta)^2 + (\Delta_c - C)^2 - D}\right], \tag{8}$$

with $C = A\omega_m\omega_0(1 + B)$, $D = A^2B^2\omega_m^2\omega_0^2$, $A = \frac{2g^2}{\omega_m^2}$, $B = \frac{\omega_m^2}{\omega_m^2 - i\delta\gamma_m - \delta^2}$, and $\omega_0 = |c_0|^2$.

To investigate the optical property of the output field for this coupling system, using the input-output relation which is valid for a one-sided open cavity [13,29,30], $c_{\rm out}(t) = c_{in}(t) - \sqrt{2\kappa}c(t)$, where $c_{\rm in}$ and $c_{\rm out}$ are the input and output operators, respectively, we can obtain the expectation value of the output field as

$$\begin{aligned}\langle c_{out}(t)\rangle = {}&(\varepsilon_{\rm pu}/\sqrt{2\kappa} - \sqrt{2\kappa}c_0)e^{-i\omega_{\rm pu}t} + (\varepsilon_{\rm pr}/\sqrt{2\kappa} - \sqrt{2\kappa}c_+)\\ &\times e^{-i(\omega_{\rm pu}+\delta)t} - \sqrt{2\kappa}c_-e^{-i(\omega_{\rm pu}-\delta)t}.\end{aligned} \tag{9}$$

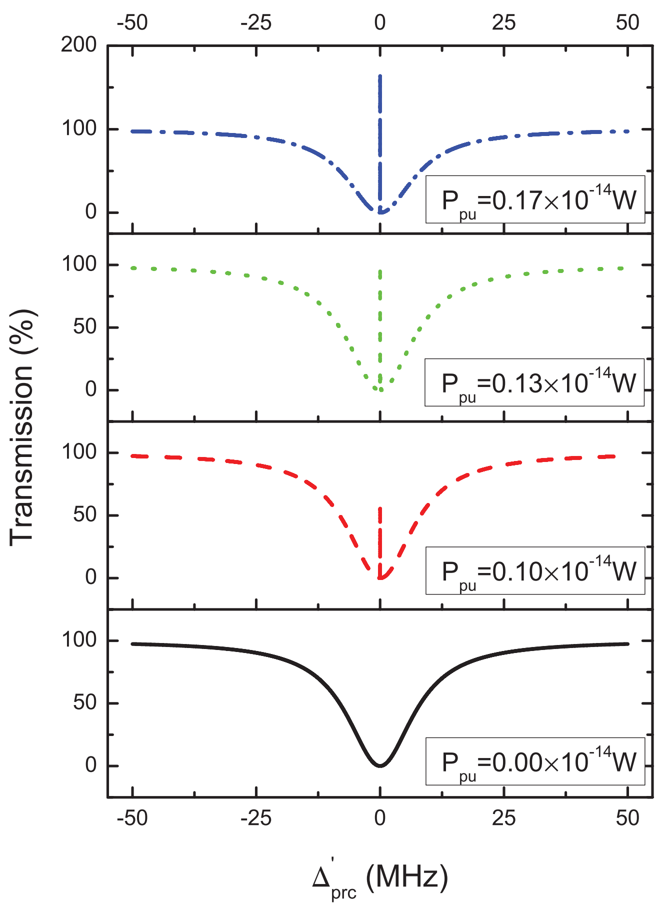

FIG. 2. (Color online) The transmission $|T_{\rm pr}|^2$ of the probe beam as a function of the shifted probe-cavity detuning $\Delta'_{\rm prc}$ for $P_{\rm pu} = 0$, 0.10, 0.13, and 0.17 × 10^{-14} W, respectively. Other parameters are $N = 1.2 \times 10^5$, $g_0 = 2\pi \times 10.9$ MHz, $\kappa = 2\pi \times 1.3$ MHz, $\Delta_a = 2\pi \times 32$ GHz, $\gamma_m = 2\pi \times 0.4$ kHz and $\omega_{\rm rec} = 2\pi \times 3.8$ kHz.

The transmission coefficient of the probe beam, defined as the ratio of the output and input field amplitudes at the probe frequency, is then given by [13,14]

$$T_{\rm pr}(\omega_{\rm pr}) = \frac{\varepsilon_{\rm pr}/\sqrt{2\kappa} - \sqrt{2\kappa}c_+}{\varepsilon_{\rm pr}/\sqrt{2\kappa}} = 1 - 2\kappa\left[\frac{(\kappa - i\delta) - i(\Delta_c - C)}{(\kappa - i\delta)^2 + (\Delta_c - C)^2 - D}\right]. \quad (10)$$

III. RESULTS AND DISCUSSION

To illustrate the numerical results, we choose the realistic parameters of the BEC-cavity system as follows [17,20]: $N = 1.2 \times 10^5$, $g_0 = 2\pi \times 10.9$ MHz, $\kappa = 2\pi \times 1.3$ MHz, $\Delta_a = 2\pi \times 32$ GHz, $\gamma_m = 2\pi \times 0.4$ kHz, $\lambda_{\rm pu} = 780$ nm, and $\omega_{\rm rec} = 2\pi \times 3.8$ kHz. In Fig. 2, the transmission $|T_{\rm pr}|^2$ of the probe beam is plotted as a function of the shifted probe-cavity detuning $\Delta'_{\rm prc} = \omega_{\rm pr} - \omega_c - \frac{1}{2}NU_0$ with a blue cavity-pump detuning, i.e., $\Delta_c = -\omega_m$, for $P_{\rm pu} = 0$, 0.10, 0.13, and 0.17 × 10^{-14} W, respectively. This figure shows that the transmitted spectrum of the probe laser can effectively be modulated by the pump laser. In the absence of the pump laser, the probe laser cannot transmit the cavity. However, when applying a pump laser driving the cavity with a blue cavity-pump detuning, significant transmissions of the probe laser can be obtained at the resonant region, i.e., $\Delta'_{\rm prc} = 0$. When the pump power $P_{\rm pu} = 0.13 \times 10^{-14}$ W, the probe laser can completely transmit the cavity as shown in the figure. More importantly, as the power of the pump laser increases further, the transmission of the probe laser can even be amplified significantly. For example, when $P_{\rm pu} = 0.17 \times 10^{-14}$ W, the probe transmission is amplified by about 75%. In the present experimental conditions, the light loss in the cavity due to scattering can be controlled at a very low level. For example, in the work by Colombe *et al.* [18] the light loss is about 56 ppm. Thus, from the numerical simulations, we can obtain that when the pump power $P_{\rm pu}$ is larger than 0.13 × 10^{-14} W, the transmission of the probe beam can effectively be switched on. These phenomena can be understood as follows. The collective motion of the BEC can be made an analogy to a mechanical resonator with resonance frequency ω_m. The radiation pressure force at the beat frequency δ between the probe and pump photons drives the collective motion of the BEC near its resonant frequency under resonant cavity-pump detuning, i.e., $\Delta_c = \pm\omega_m$. When the beam frequency δ is close to the resonance frequency ω_m of the BEC, the mechanical mode starts to oscillate coherently, which will induce Stokes ($\omega_s = \omega_{\rm pu} - \omega_m$) and anti-Stokes ($\omega_{as} = \omega_{\rm pu} + \omega_m$) scattering of light from the strong intracavity field. For the red detuning, i.e., $\Delta_c = \omega_m$, the anti-Stokes field interferes with the probe field for the near-resonant probe ($\Delta'_{\rm prc} \approx 0$) and the effective interaction Hamiltonian between the cavity modes and mechanical phonon modes has the form of the well-known beam-splitter one in quantum optics [31]. However, for the blue detuning, i.e., $\Delta_c = -\omega_m$, it is the Stokes field that interferes with the near-resonant probe field and thus modifies the probe spectrum. Furthermore, the effective interaction Hamiltonian has the form of the parametric amplification

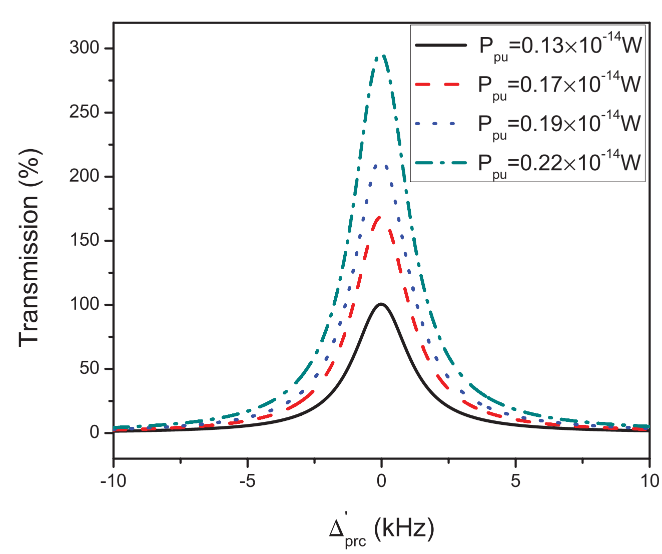

FIG. 3. (Color online) The magnified transparency window as a function of $\Delta'_{\rm prc}$ for $P_{\rm pu} = 0.13$, 0.17, 0.19, and 0.22 × 10^{-14} W, respectively. Other parameters are $N = 1.2 \times 10^5$, $g_0 = 2\pi \times 10.9$ MHz, $\kappa = 2\pi \times 1.3$ MHz, $\Delta_a = 2\pi \times 32$ GHz, $\gamma_m = 2\pi \times 0.4$ kHz, and $\omega_{\rm rec} = 2\pi \times 3.8$ kHz.

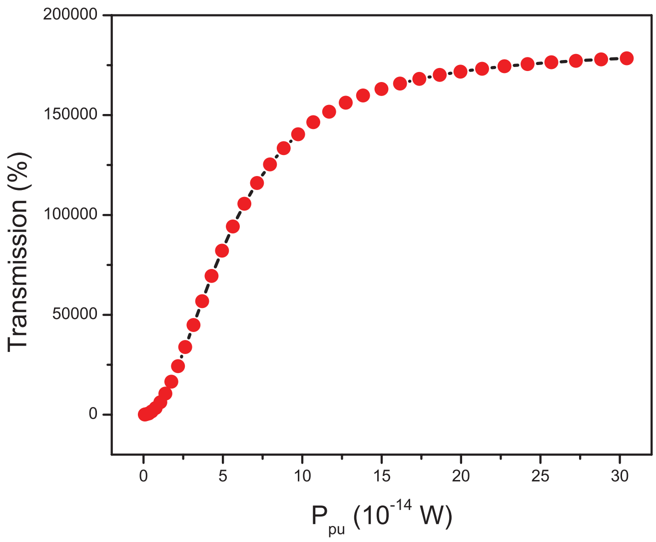

FIG. 4. (Color online) The probe transmission $|T_{\text{pr}}|^2$ on resonance as a function of the pump power P_{pu} with $\Delta_c = -\omega_m$. Other parameters are $N = 1.2 \times 10^5$, $g_0 = 2\pi \times 10.9$ MHz, $\kappa = 2\pi \times 1.3$ MHz, $\Delta_a = 2\pi \times 32$ GHz, $\gamma_m = 2\pi \times 0.4$ kHz, and $\omega_{\text{rec}} = 2\pi \times 3.8$ kHz.

one, $H_{\text{CM}} = \hbar G(c^\dagger a^\dagger + ac)$ with $G = g\sqrt{\omega_0}$, and, with strong pump powers, the probe optical signal can significantly be amplified [14].

In Fig. 3, the magnified transparency window of the transmitted probe beam is presented as a function of Δ'_{prc} with a blue cavity-pump detuning, i.e., $\Delta_c = -\omega_m$ for $P_{\text{pu}} = 0.13$, 0.17, 0.19, and 0.22×10^{-14} W. As shown in Fig. 3, a transparency window of several kilohertz is achieved around the resonance ($\Delta'_{\text{prc}} \approx 0$). On resonance, the transmission coefficient is largely magnified while increasing the power of the pump laser. In Fig. 4, the transmission $|T_{\text{pr}}|^2$ of the probe beam on resonance is shown as a function of the pump-laser power P_{pu} with $\Delta_c = -\omega_m$. This figure shows that with the enhancement of the pump power the probe transmission increases greatly at first and finally reaches saturation at a critical pump power. Moreover, it can be obviously seen that in the interval of 1.9–5.2 $\times 10^{-14}$ W for the pump power P_{pu} the system presents a greater switching efficiency. For this case, the probe transmission can be amplified by about 200% while increasing every 0.01×10^{-14} W of the pump power. When $P_{\text{pu}} > 5.2 \times 10^{-14}$ W, the switching efficiency starts to reduce and then approaches saturation at $P_{\text{pu}} \approx 20 \times 10^{-14}$ W.

In view of the fact that the transmission of the probe beam can be effectively modulated by the pump beam, we propose a scheme for realizing an all-optical transistor based on the BEC-cavity system. In such a transistor, the pump laser serves as a gate beam. For the blue detuning of the pump laser, i.e., $\Delta_c = -\omega_m$, the resonant probe beam (i.e., $\Delta'_{\text{prc}} = 0$) can be attenuated or amplified depending on the power of the gate beam. In the absence of the pump beam, the probe beam can fully be attenuated. However, as the pump beam is present with an ideal power ($P_{\text{pu}} = 0.13 \times 10^{-14}$ W), the cavity becomes completely transparent to the probe beam. With the further increase of the pump power, the probe beam can significantly be amplified.

IV. CONCLUSION

In conclusion, we have theoretically investigated the light propagation in the BEC-cavity system and proposed a scheme for an all-optical transistor based on this coupled system. The calculated results show that with a blue detuning between the pump and effective cavity frequency the transmission of the probe beam from the cavity can be controlled effectively by the pump beam. In the absence of the pump beam, the probe beam cannot transmit the cavity. However, while driving the cavity by a pump beam with an ideal power, the cavity becomes completely transparent to the probe beam. As the pump power increases further, the transmission of the probe beam can be amplified greatly. Therefore, the BEC-cavity system can work as an optical transistor, where the pump beam serves as a gate beam which can effectively control the transmission of the probe beam just by regulating its power. Finally, we hope that this all-optical transistor proposed here will be implemented by current experiments in the near future.

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[1] D. E. Chang, A. S. Sørensen, E. A. Demler, and M. D. Lukin, Nat. Phys. **3**, 807 (2007).
[2] J. Hwang, M. Pototschnig, R. Lettow, G. Zumofen, A. Renn, S. Götzinger, and V. Sandoghdar, Nature (London) **460**, 76 (2009).
[3] J. J. Li and K. D. Zhu, Nanotechnology **22**, 055202 (2011).
[4] H. M. Gibbs, Optical Bistability: *Controlling Light with Light* (Academic, Orlando, 1985).
[5] J. L. O'Brien, Science **318**, 1567 (2007).
[6] D. Bouwmeester, A. Ekert, and A. Zeilinger, *The Physics of Quantum Information* (Springer, Berlin, 2000).
[7] S. Gröblacher, K. Hammerer, M. R. Vanner, and M. Aspelmeyer, Nature (London) **460**, 724 (2009).
[8] C. Genes, D. Vitali, P. Tombesi, S. Gigan, and M. Aspelmeyer, Phys. Rev. A **77**, 033804 (2008).
[9] I. Wilson-Rae, N. Nooshi, W. Zwerger, and T. J. Kippenberg, Phys. Rev. Lett. **99**, 093901 (2007).
[10] G. S. Agarwal and S. Huang, Phys. Rev. A **81**, 041803 (2010).
[11] J. D. Thompson, B. M. Zwickl, A. M. Jayich, F. Marquardt, S. M. Girvin, and J. G. E. Harris, Nature (London) **452**, 72 (2008).
[12] J. C. Sankey, C. Yang, B. M. Zwickl, A. M. Jayich, and J. G. E. Harris, Nature Phys. **6**, 707 (2010).
[13] S. Weis, R. Rivière, S. Deléglise, E. Gavartin, O. Arcizet, A. Schliesser, and T. J. Kippenberg, Science **330**, 1520 (2010).

[14] A. H. Safavi-Naeini, T. P. Mayer Alegre, J. Chan, M. Eichenfield, M. Winger, Q. Lin, J. T. Hill, D. E. Chang, and O. Painter, Nature (London) **472**, 69 (2011).
[15] 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).
[16] F. Brennecke, T. Donner, S. Ritter, T. Bourdel, M. Köhl, and T. Esslinger, Nature (London) **450**, 268 (2007).
[17] F. Brennecke, S. Ritter, T. Donner, and T. Esslinger, Science **322**, 235 (2008).
[18] Y. Colombe, T. Steinmetz, G. Dubois, F. Linke, D. Hunger, and J. Reichel, Nature (London) **450**, 272 (2007).
[19] T. P. Purdy, D. W. C. Brooks, T. Botter, N. Brahms, Z. Y. Ma, and D. M. Stamper-Kurn, Phys. Rev. Lett. **105**, 133602 (2010).
[20] S. Ritter, F. Brennecke, K. Baumann, T. Donner, C. Guerlin, and T. Esslinger, Appl. Phys. B **95**, 213 (2009).
[21] D. Nagy, G. Kónya, G. Szirmai, and P. Domokos, Phys. Rev. Lett. **104**, 130401 (2010).
[22] D. Nagy, P. Domokos, A. Vukics, and H. Ritsch, Eur. Phys. J. D **55**, 659 (2009).
[23] J. M. Zhang, F. C. Cui, D. L. Zhou, and W. M. Liu, Phys. Rev. A **79**, 033401 (2009).
[24] A. B. Bhattacherjee, Phys. Rev. A **80**, 043607 (2009).
[25] K. Zhang, W. Chen, M. Bhattacharya, and P. Meystre, Phys. Rev. A **81**, 013802 (2010).
[26] M. Paternostro, G. De Chiara, and G.-M. Palma, Phys. Rev. Lett. **104**, 243602 (2010).
[27] G. Szirmai, D. Nagy, and P. Domokos, Phys. Rev. A **81**, 043639 (2010).
[28] A. Öttl, S. Ritter, M. Köhl, and T. Esslinger, Rev. Sci. Instrum. **77**, 063118 (2006).
[29] C. W. Gardiner and P. Zoller, *Quantum Noise* (Springer, New York, 2004).
[30] C. W. Gardiner and M. J. Collett, Phys. Rev. A **31**, 3761 (1985).
[31] V. Fiore, Y. Yang, M. C. Kuzyk, R. Barbour, L. Tian, and H. Wang, Phys. Rev. Lett. **107**, 133601 (2011).