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Generating EPR beams in a cavity optomechanical system

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We propose a scheme to produce continuous variable entanglement between phase-quadrature amplitudes of two light modes in an optomechanical system. For proper driving power and detuning, the entanglement is insensitive with bath temperature and Q of mechanical oscillator. Under realistic experimental conditions, we find that the entanglement could be very large even at room temperature.

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Entanglement is the key resource of the field of quantum information. Light is the perfect medium to distribute entanglement among distant parties. Entangled light with continuous variable (CV) entanglement between phase-quadrature amplitudes of two light modes is widely used in teleportation, entanglement swap, dense coding, etc. [1]. This type of entangled state is also called Einstein-Podolsky-Rosen (EPR) state. The EPR beams have been generated experimentally by a nondegenerate optical parameter amplifier [2], or Kerr nonlinearity in an optical fiber [3]. The later one is simpler and more reliable. The Kerr nonlinearity is used to generate two independent squeezed beams. With interference at a beam splitter, the EPR entanglement is obtained between output beams. However, Kerr nonlinearity in fiber is very weak, which limits entanglement between output beams.

It was found that strong Kerr nonlinearity appeared in an optomechanical system consisting of a cavity with a movable boundary [4, 5, 6]. Besides, the single-mode squeezing could be made insensitive with thermal noise [5], which makes the scheme very attractive. However, the frequency of output squeezed beams cannot be made identical, which makes interference difficult. Then it was generalized to two-mode schemes in order to generate EPR beams without interference [7, 8, 9, 10, 11]. However, they are either very sensitive to thermal noises [7, 8, 10, 11], or requiring ultrahigh mechanical oscillator Q $\sim 10^8$ to suppress thermal noise effects [9], which is 2 to 3 orders higher than the present available parameters [12]. The practical scheme to generate EPR beams in an optomechanical system needs to overcome these problems.

In this paper we propose a practical scheme to produce EPR beams in an optomechanical system, which consists of a whispering-gallery mode(WGM) cavity with a movable boundary. We find that, similarly as the single-mode scheme [5], the thermal noise in the two-mode scheme can be greatly suppressed by adiabatically eliminating an oscillator mode. By precisely tuning the laser power and detuning, the oscillation mode is adiabatically eliminated and two output sideband modes are entangled. Unlike the cavity-free scheme [9], our scheme requires modest oscillator Q. Besides, the output light is continuous in our scheme, other than pulse in Ref. [9]. The most attractive feature of our scheme is that the entanglement between output beams is nearly not changed under different bath temperature and Q of the mechanical oscillator. Within the experimentally available parameters [13, 14], we find the maximum two-mode squeezing could be higher than 16 dB under room temperature. The entanglement of formation (EOF) between two modes is larger than 5 [15]. Since the coupling efficiency between cavity and fiber could be larger than 99% in the WGM cavity system [16], we neglect the coupling induced noises in this paper.

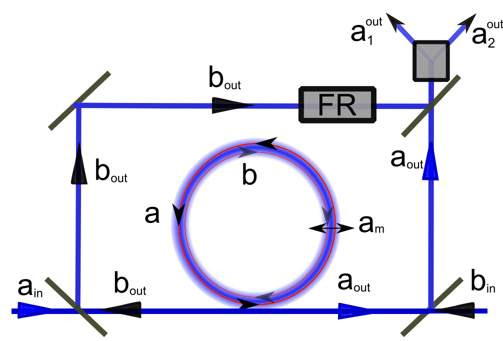

FIG. 1: (Color online) Experimental setup. Cavity modes a and b, which are driven by four lasers, couple to the mechanical mode a_m.

As shown in Fig. 1, we consider an optomechanical system consisting of a WGM cavity with a movable boundary. There are two cavity modes a and b with the same frequency but the opposite momentum. They are coupling with the same mechanical oscillation mode a_m and driven by four lasers, two from the right-hand side with frequencies ω_L and $\omega_{L'}$, the other two from the left-hand side with frequencies ω_L and $\omega_{L'}$. $a_{\rm in}$ and $b_{\rm in}$ are input lights. $a_1^{\rm out}$ and $a_2^{\rm out}$ are output lights. Two lower mirrors have very high probability ($> 99\%$) to reflect the driving lasers. So we neglect the reflecting induced noise for $a_{\rm out}$ and $b_{\rm out}$. The system Hamiltonian is $H = H_0 + H_d + H_I$, where [17, 18, 19, 20]

$$
\begin{aligned}
H_0 =&\hbar\omega_p(a^\dagger a + b^\dagger b) + \hbar\omega_m a_m^\dagger a_m \\
H_d =&\hbar(\frac{\Omega_a}{2}a + \frac{\Omega_b}{2}b)e^{-i\omega_L t} + \hbar(\frac{\Omega_a'}{2}a + \frac{\Omega_b'}{2}b)e^{-i\omega_L' t} + \text{H.c.} \\
H_I =&\hbar\nu(a^\dagger b + ab^\dagger) + \hbar\eta\omega_m(a^\dagger a + b^\dagger b)(a_m + a_m^\dagger).
\end{aligned} \tag{1}
$$

Here a, b and a_m are the annihilation operators for the optical and mechanical modes, ω_p and ω_m are their angular frequency. Ω_j with $j = a, b$ is the driving amplitude and defined as $\Omega_j = 2\sqrt{P_j/\hbar\omega_L\tau}$, where P_j is the input laser power and $\tau = 1/\gamma$ is the photon loss rate into the output modes. ν is the coupling strength between cavity modes a and b. For the WGM cavity system, it ranges from 100 MHz to 10 GHz [19, 20]. The dimensionless parameter $\eta = (\omega_p/\omega_m)(x_m/R)$ is used to characterize optomechanical coupling, with $x_m = \sqrt{\hbar/m\omega_m}$ the zero-point motion of the mechanical resonator mode [21], m its effective mass, and R a cavity radius. In typical WGM cavity systems we find $\eta \sim 10^{-4}$.

We define the normal modes $a_1 = (a+b)/\sqrt{2}$ and $a_2 = (a-b)/\sqrt{2}$. We suppose the conditions that $\Omega_a - \Omega_b = 0$ and $\Omega_a' + \Omega_b' = 0$ are satisfied. The Hamiltonian can be written as

$$
\begin{aligned}
H =&\hbar(\omega_p + \nu)a_1^\dagger a_1 + \hbar(\omega_p - \nu)a_2^\dagger a_2 + \hbar\omega_m a_m^\dagger a_m \\
&+ \hbar(\frac{\Omega_1}{2}a_1 e^{-i\omega_L t} + \frac{\Omega_2}{2}a_2 e^{-i\omega_L' t} + \text{H.c.}) \\
&+ \hbar\eta\omega_m(a_1^\dagger a_1 + a_2^\dagger a_2)(a_m^\dagger + a_m),
\end{aligned} \tag{2}
$$

where $\Omega_1 = \Omega_a + \Omega_b$ and $\Omega_2 = \Omega_a' - \Omega_b'$. We define the detuning $\Delta_1 = \omega_L - \omega_p - \nu$ and $\Delta_2 = \omega_L' - \omega_p + \nu$. As shown in Fig. 1, with beam splitters and Faraday rotator, we can get the output mode of a_1 and a_2. We assume both cavity and oscillator modes are weakly dissipating at rates γ and γ_m, respectively, where $\gamma_m \ll \omega_m$. We can get quantum Langevin equations [22]

$$
\begin{aligned}
\dot{a}_j &= i\Delta_j a_j - i\eta\omega_m a_j(a_m + a_m^\dagger) - i\frac{\Omega_j}{2}a_j - \frac{\gamma}{2}a_j(t) \\
&\quad +\sqrt{\gamma}a_j^{\text{in}} \quad \text{for } j = 1, 2, \qquad (3)\\
\dot{a}_m &= -i\eta\omega_m\sum_{j=1}^{2}a_j^\dagger a_j - (i\omega_m + \frac{\gamma_m}{2})a_m + \sqrt{\gamma_m}a_m^{\text{in}}, \qquad (4)
\end{aligned}
$$

where thermal noise inputs are defined as correlation functions $\langle a_m^{\text{in}\dagger}(t), a_m^{\text{in}}(t')\rangle = n_m\delta(t-t')$, $\langle a_m^{\text{in}\dagger}(t), a_m^{\text{in}\dagger}(t')\rangle = \langle a_m^{\text{in}}(t), a_m^{\text{in}}(t')\rangle = 0$, $\langle a_j^{\text{in}\dagger}(t), a_j^{\text{in}}(t')\rangle = \langle a_j^{\text{in}\dagger}(t), a_j^{\text{in}\dagger}(t')\rangle = \langle a_j^{\text{in}}(t), a_j^{\text{in}}(t')\rangle = 0$, with n_m the thermal occupancy number of thermal bath for oscillator mode. We suppose cavity modes couple with vacuum bath.

To simplify Eqs. (3) and (4), we apply a shift to normal coordinate, $a_j \to a_j + \alpha_j$, $a_m \to a_m + \beta$. α_j and β are c numbers, which are chosen to cancel all c number terms in the transformed equations. We find they should fulfill the following requirements: $\beta \simeq -\eta(|\alpha_1|^2 + |\alpha_2|^2)$, and $i\Delta_j\alpha_j + 2i\eta^2\omega_m\alpha_j(|\alpha_1|^2 + |\alpha_2|^2) - \frac{\gamma}{2}\alpha_j - i\frac{\Omega_j}{2} = 0$. Because $\gamma_m \ll \omega_m$, the imaginary part of β can be neglected. In the limit $\Delta_j \gg 2\eta^2 w_m(|\alpha_1|^2 + |\alpha_2|^2)$, we find $\alpha_j \simeq \Omega_j/\sqrt{\gamma^2 + 4\Delta_j^2}$. In the limit $|\alpha| \gg |\langle a_p\rangle|$, the Langevin equations are linearized as

$$
\begin{aligned}
\dot{a}_j &= -i\eta\omega_m\alpha_j(a_m + a_m^\dagger) + (i\Delta_j' - \frac{\gamma}{2})a_j + \sqrt{\gamma}a_j^{\text{in}}, \qquad (5)\\
\dot{a}_m &= -i\eta\omega_m\sum_{p=1}^{2}(\alpha_p^* a_p + \alpha_p a_p^\dagger) - (i\omega_m + \frac{\gamma_m}{2})a_m \\
&\quad +\sqrt{\gamma_m}a_m^{\text{in}}, \qquad (6)
\end{aligned}
$$

where $j = 1, 2$ and $\Delta_j' = \Delta_j + 2\eta^2\omega_m(|\alpha_1|^2 + |\alpha_2|^2)$. We suppose $\Delta_1' < 0$ and $\Delta_2' > 0$. We define $\delta = (\Delta_2' - \Delta_1')/2 - \omega_m$ and $d = -(\Delta_1' + \Delta_2')/2$.

In the limit $\omega_m \gg \delta, d, \gamma, \gamma_m$, the Langevin equations (5) and (6) can be simplified as

$$
\begin{aligned}
\dot{a}_1 =& -ida_1 - i\eta\omega_m\alpha_1 a_m - \frac{\gamma}{2}a_1 + \sqrt{\gamma}a_1^{\text{in}}, \\
\dot{a}_2 =& -ida_2 - i\eta\omega_m\alpha_2 a_m^\dagger - \frac{\gamma}{2}a_2 + \sqrt{\gamma}a_2^{\text{in}}, \\
\dot{a}_m =& i\delta a_m - i\eta\omega_m(\alpha_1^* a_1 + \alpha_2 a_2^\dagger) - \frac{\gamma_m}{2}a_m + \sqrt{\gamma_m}a_m^{\text{in}}.
\end{aligned} \tag{7}
$$

With proper detuning and input power, we can always tune the cavity mode amplitude $\alpha_1 = \alpha_2 = \alpha$. Define the Fourier components of the intracavity field by $a(t) = \frac{1}{\sqrt{2\pi}}\int_{-\infty}^{\infty} e^{-i\omega(t-t_0)}a(\omega)d\omega$. In the limit $\delta \gg \omega, \gamma_m$, we can adiabatically eliminate the a_m mode. We get $a_m(\omega) \simeq \eta\frac{\omega_m}{\delta}(\alpha^* a_1 + \alpha a_2^\dagger) - \frac{\sqrt{\gamma_m}}{i\delta}a_m^{\text{in}}$. Then we have quantum Langevin equations for $a_1(\omega)$ and $a_2^\dagger(-\omega)$

$$
\begin{aligned}
-i\omega a_1(\omega) =& -ig' a_1(\omega) - iga_2^\dagger(-\omega) - \frac{\gamma}{2}a_1(\omega) \\
&+ \sqrt{\gamma}a_1^{\text{in}}(\omega) + \sqrt{\tilde{a}_m}a_m^{\text{in}}(\omega), \\
-i\omega a_2^\dagger(-\omega) =& ig' a_2^\dagger(-\omega) + iga_1(\omega) - \frac{\gamma}{2}a_2^\dagger(-\omega) \\
&+ \sqrt{\gamma}a_2^{\dagger\text{in}}(-\omega) - \sqrt{\tilde{\gamma}_m}a_m^{\text{in}}(\omega),
\end{aligned} \tag{8}
$$

where $g = \eta^2|\alpha|^2\omega_m^2/\delta$, $\tilde{\gamma}_m = (\eta|\alpha|\omega_m/\delta)^2\gamma_m$, and $g' = g + d$. In Eq. (8), we neglect the phase of α because it is not important.

Denote $\vec{a}(\omega) = \binom{a_1(\omega)}{a_2^\dagger(-\omega)}$, $\vec{a}^{\text{in}}(\omega) = \binom{a_1^{\text{in}}(\omega)}{a_2^{\text{in}\dagger}(-\omega)}$ and $\vec{a}_m^{\text{in}}(\omega) = \binom{a_m^{\text{in}}(\omega)}{-a_m^{\text{in}}(\omega)}$. We get the following matrix equation

$$
\mathbf{A}\vec{a}(\omega) = \sqrt{\gamma}\vec{a}^{\text{in}}(\omega) + \sqrt{\tilde{\gamma}_m}\vec{a}_m^{\text{in}}, \tag{9}
$$

where

$$
\mathbf{A} = \begin{pmatrix} -i\omega + \frac{\gamma}{2} + ig' & ig \\ -ig & -i\omega + \frac{\gamma}{2} - ig' \end{pmatrix}.
$$

Using boundary conditions $a_j^{\text{out}}(\omega) = -a_j^{\text{in}}(\omega) + \sqrt{\gamma}a_j(\omega)$ for $j = 1, 2$, we can calculate output field as

$$
\begin{aligned}
a_1^{\text{out}}(\omega) &= G(\omega)a_1^{\text{in}}(\omega) - H(\omega)a_2^{\text{in}\dagger}(-\omega) + I(\omega)a_m^{\text{in}}(\omega), \\
a_2^{\text{out}\dagger}(\omega) &= G(\omega)^* a_2^{\text{in}\dagger}(\omega) - H(\omega)^* a_1^{\text{in}}(-\omega) - I(\omega)^* a_m^{\text{in}}(-\omega),
\end{aligned} \tag{10}
$$

where $G(\omega) = (\omega^2 + \frac{\gamma^2}{4} + g^2 - g'^2 - ig'\gamma)/\Delta(\omega)$, $H(\omega) = ig\gamma/\Delta(\omega)$, $I(\omega) = (-i\omega + \frac{\gamma}{2} - ig' + ig)\sqrt{\gamma\tilde{\gamma}_m}/\Delta(\omega)$, and $\Delta(\omega) = (-i\omega + \frac{\gamma}{2})^2 + g'^2 - g^2$.

Let us define the dimensionless position and momentum operators of fields $X_j^{\rm out}(\omega) = \left[a_1^{\rm out}(\omega) + a_1^{\rm out\dagger}(-\omega)\right]$ and $P_j^{\rm out}(\omega) = \left[a_j^{\rm out}(\omega) - a_j^{\rm out\dagger}(-\omega)\right]/i$, for $j = 1, 2$. We define the correlation matrix of the output field as $V_{ij} = \langle(\xi_i\xi_j + \xi_j\xi_i)/2\rangle$, where $\xi = (X_1^{\rm out}, P_1^{\rm out}, X_2^{\rm out}, P_2^{\rm out})$. We calculate the correlation matrix with Eq. (10). Up to local unitary transformation, the standard form of it is

$$V_S = \begin{pmatrix} n & 0 & k_x & 0 \\ 0 & n & 0 & -k_x \\ k_x & 0 & n & 0 \\ 0 & -k_x & 0 & n \end{pmatrix}, \qquad (11)$$

where $n = \{(\omega^2 + \frac{\gamma^2}{4} + g^2 - g'^2)^2 + (g'^2 + g^2)\gamma^2 + [(\omega + g' - g)^2 + \frac{\gamma^2}{4}]\gamma\tilde{\gamma}_m(2n_m + 1)\}/|\Delta(\omega)|^2$, $k_x = \sqrt{V_{14}^2 + V_{24}^2}$, where $V_{14} = -2g\gamma(w^2 + \frac{\gamma}{4} + g^2 - g'^2)/|\Delta(\omega)|^2$, $V_{24} = \{2g'g\gamma^2 + [(\omega + g' - g)^2 + \frac{\gamma^2}{4}]\gamma\tilde{\gamma}_m(2n_m + 1)\}/|\Delta(\omega)|^2$. This is the symmetric Gaussian state. The EOF for the symmetric Gaussian states is defined as [15]

$$\begin{aligned} E_F = C_+(n - k_x)\log_2[C_+(n - k_x)] \\ -C_-(n - k_x)\log_2[C_-(n - k_x)] \end{aligned} \qquad (12)$$

where $C_\pm(x) = (x^{-1/2} \pm x^{1/2})^2/4$. V describes an entanglement state if and only if $n - k_x < 1$. Based on the standard form of matrix (11), we also find that $\langle\delta^2(X_1 + X_2)\rangle = \langle\delta^2(P_1 - P_2)\rangle = n - k_x$. We define the two-mode squeezing as $S = -10\log_{10}(n - k_x)$.

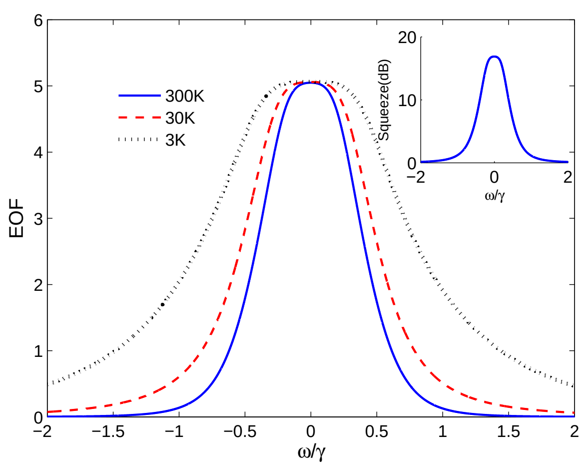

FIG. 2: (Color online) EOF for different temperature. $|\alpha|$ is 1000, $\delta/2\pi = 10$ MHz, $d = 0.07\gamma$. The maxima squeezing is larger than 16 dB when $T = 300$ K.

We now estimate the bath noise influence in experimentally accessible conditions [13]. The cavity resonant frequency $\omega_p = 2\pi \times 300$ THz. The oscillator frequency $\omega_m = 2\pi \times 73.5$ MHz. The mechanical Q factor is about 30 000. The cavity and oscillator modes decay rates are $\gamma = 2\pi \times 3.2$ MHz and $\gamma_m = \omega_m/Q$, respectively. The cavity radius is $R = 38$ μm. The dimensionless coupling parameter is $\eta \simeq 10^{-4}$. We find if $0 < d \ll \gamma$, the thermal noise does not decrease the maximum entanglement with strong enough input power. This is because we adiabatically eliminate the mechanical mode and suppress the effects of thermal noise. As show in Fig. 2, the temperature change does not change the maximum entanglement with proper driving and detuning. But the higher the temperature, the less the entanglement spectrum width. The embeded figure shows that the maximum squeezing could be larger than 16 dB at room temperature.

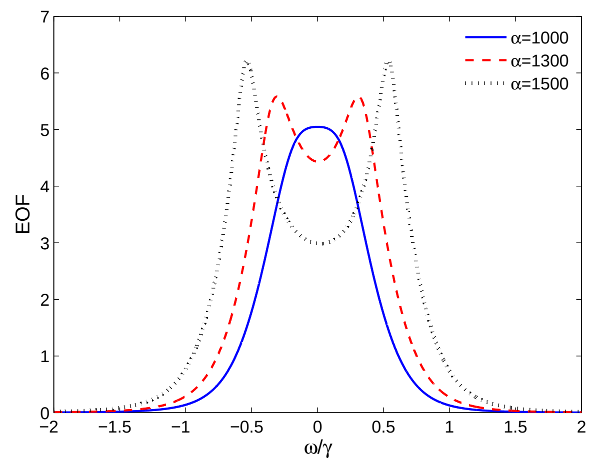

FIG. 3: (Color online) EOF for different cavity mode amplitude α. Here we adopt $\delta/2\pi = 10$ MHz, $T = 300$ K, $\gamma_p/2\pi = 3.2$MHz, and $d = 0.07\gamma$.

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As shown in Fig. 3, the bigger the cavity mode amplitude α, the larger the output entanglement. Because $\alpha_j \simeq \Omega_j/\sqrt{\gamma^2 + 4\Delta_j^2}$, the output entanglement is proportional to driving amplitude. But the peak of entanglement is splitted into two symmetric peaks when driving is very strong. The splitting distance is proportional to driving power. Increasing driving power can decrease the entanglement too. This is because adiabatical elimination condition $\omega \ll \delta$ are not valid around peaks for very strong driving. So the driving power should be neither too big nor too small. For the specific α and δ, we find there is an optimum d which makes entanglement maximum and the entanglement peaks appear near $\omega = 0$. The optimum d is $d_o = \sqrt{(\eta^2\omega^2\alpha^2/\delta)^2 + \gamma^2/4} - (\eta\omega\alpha)^2/\delta$, corresponding to squeezing $S_o = -10\log_{10}(4d_o^2/\gamma^2)$ and entanglement which is obtained from Eq. (12) with $n - kx = 4(d_o/\gamma)^2$. It is obvious that the higher the input power, the smaller the optimum d. In the mean time, we find that decreasing the mechanical Q factor nearly does not change the entanglement spectrum if d is around its optimum value and the condition $\omega_m/Q \ll \delta$ is fulfilled. Leaving other parameters unchanged, Q could be as low as 300. Considering the difficulty of increasing the mechanical oscillator Q, the above finding makes our scheme more practical.

We also test the stability of our scheme. As shown in Fig. 4, the optimum d is around 0.07γ if $\alpha = 1000$, $\delta/2\pi = 10$ MHz. To maintain such high entanglement, we need to precisely control the d down to $0.02\gamma \sim 2\pi\times 60$ kHz. d is defined as $d = -(\Delta_1' + \Delta_2')/2 = -(\Delta_1 + \Delta_2)/2 - 4\eta^2\omega_m|\alpha|^2$. The higher entanglement is needed, the more precise detuning and driving power is required at the same time. To maintain the entanglement as high as Fig. 4, the laser spectrum width should be less than 60 kHz and the driving power fluctuation should be less than 1%. The lower entanglement between two beams is needed to maintain, the larger the optimum d is. Therefore higher fluctuations of detuning and driving power are allowed.

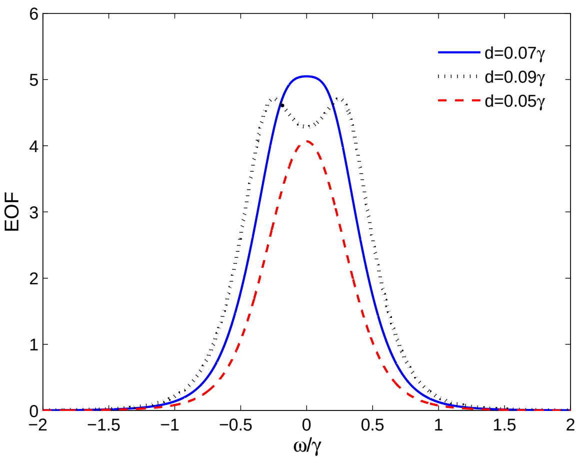

FIG. 4: (Color online) EOF for different d. Here we adopt $\omega_m/2\pi = 73.5$ MHz, $T = 300$K, $\gamma_m = \omega_m/30\,000$, $\gamma/2\pi = 3.2$ MHz, and $|\alpha| = 1000$.

.

Before conclusion, we briefly discuss the approximations we used. Our scheme needs the steady states existing, which requires $\langle a_j^\dagger a_j\rangle \ll |\alpha|^2$. During numerical calculation, $\langle a^\dagger a\rangle$ is in the order of 10^3, which is much less than $|\alpha|^2 \sim 10^6$. The other two approximations are rotating wave approximation $\omega_m \gg \delta, d, \gamma, \gamma_m$ and adiabatical elimination $\delta \gg \omega, \gamma_m$, which can be fulfilled independently. For $\alpha \sim 10^3$, the driving amplitude Ω is in the order of 10^{11} Hz, which is much lower than the distance between adjacent cavity modes $\Delta\omega = c/(Rn_0) \sim 5\times 10^{12}$ Hz, where c is the light speed in a vacuum, n_0 the refractive index of silica. Therefore the approximation that one laser only drives one cavity mode is valid. Laser power is needed in the order of 10 mW, which is available in the laboratory.

In conclusion, we have proposed a scheme to generate EPR lights in an optomechanical system. Two sideband modes, which couple with the mechanical mode, are driven by lasers. After adiabatically eliminating the the mechanical mode, we find that the output sideband modes are highly entangled. The higher power of the driving laser, the larger entanglement of the output light. To maintain the entanglement, we need to precisely control the driving power and laser frequency at the same time. With proper parameters, the entanglement is insensitive to the thermal noise and mechanical Q factor. We test the scheme by experimental available parameters. Though in this paper we fucus on WGM cavity systems, our scheme can be realized in other optomechanical systems, as long as the mechanical mode frequency is much larger than the cavity decay rate.

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