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Entanglement swapping using nondegenerate optical parametric amplifier

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Abstract

We introduce a simple, experimentally realizable, entanglement swapping protocol for continuous variables by exploring nondegenerate optical parametric amplifier. Due to adopting the bright EPR beams and the simple direct measurement for Bell-state, the entanglement swapping and the verification of entanglement swapping is within the reach of current technology and significantly simplify the implementation. © 2002 Elsevier Science B.V. All rights reserved.

Entanglement is central to all branches of the emerging field of quantum information and quantum computation. Entanglement swapping [1] may entangle two quantum systems that have never directly interacted with each other, which may be useful in establishing nonlocal correlations over very large distances and other applications [2–4]. The entanglement swapping of single-photon polarization states has been realized experimentally by type II parametric down conversion. All these investigations have only referred to discrete-variable systems in finite-dimensional Hilbert spaces. The schemes for continuous variable entanglement swapping were proposed that polarization-entangled states of single photons are teleported using squeezed-state entanglement [5], and both entangled states are produced with squeezed light [6]. Entanglement swapping is really a special case of the quantum teleportation. The scheme for continuous variable entanglement swapping [6] is based on the preliminary experimental demonstration of continuous variable teleportation of coherent state [7]. In the experiment of Ref. [7], an entanglement source was built from two single-mode phase squeezed vacuum states combined at a beamsplitter. The Bell-state measurement at Alice needs two sets of balanced homodyne detectors and local oscillators (LO's). Recently quantum dense coding for continues variables has been experimentally accomplished by means of exploiting bright EPR beam with anticorrelation of amplitude quadratures and correlation of phase quadratures, which is generated from a nondegenerate optical parametric amplifier (NOPA) operating in the state of deamplification [8]. In this Letter we will propose a continuous variable entanglement swapping scheme based on the experiment [8] in which the sources of

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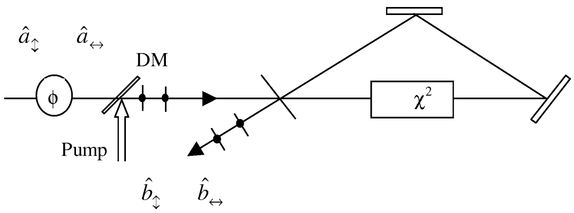

Fig. 1. The schematic diagram for phase sensitive NOPA. DM, dichroic mirror.

entanglement are the bright EPR from the deamplification of NOPA and the Bell-state measurement is obtained from the direct detection implemented with two photodetectors and two RF splitters [9,10]. We consider more general condition than the Ref. [6] in which the squeezed vacuum states are regarded as the minimum-uncertainty state. In the practical systems, the generated squeezed state is always not the minimum-uncertainty state ($\langle\delta X^2\rangle\langle\delta Y^2\rangle = 1$, δX and δY are the quantum fluctuation of the quadrature components), and there is large excess noise in the unsqueezed quadrature phase component ($\langle\delta Y^2_{\mathrm{unsquee}}\rangle > 1/\langle\delta X^2_{\mathrm{squee}}\rangle$) [8,11,12]. We will demonstrate the level of entanglement produced from entanglement swapping depend not only on the squeezing and also the large excess noise in the unsqueezed quadrature component. Applying the direct joint measurement of Bell-state gives us simultaneously a protocol for the experimental verification of entanglement swapping.

The schematic diagram for phase sensitive NOPA is shown in Fig. 1. Two coherent input signals $a_{\updownarrow}$ and $a_{\leftrightarrow}$ with same frequency ω_0 and orthogonal polarizations are injected into the NOPA. For simplification and without losing generality, we assume that the polarizations of the injected signal and idler field are orientated along the vertical and horizontal directions, and their intensities and original phases before NOPA are considered to be identical. The amplifier is pumped with the second harmonic wave of $\omega_p = 2\omega_0$ and the amplitude $a_p \gg a_{\updownarrow}, a_{\leftrightarrow}$, in this case the pump field can be considered as a classical field without depletion during the amplification process. The output signal and idler fields polarized along the vertical and horizontal directions are denoted with $b_{\updownarrow}$ and $b_{\leftrightarrow}$. We define the operators of the light fields at the center frequency ω_0 in the rotating frame,

$$\hat{O}(t) = \hat{o}(t)e^{i\omega_0 t}, \tag{1}$$

there $O = [\hat{a}_{\updownarrow}, \hat{a}_{\leftrightarrow}, \hat{b}_{\updownarrow}, \hat{b}_{\leftrightarrow}]$ are the field envelope operators and $o = [\hat{A}_{\updownarrow}, \hat{A}_{\leftrightarrow}, \hat{B}_{\updownarrow}, \hat{B}_{\leftrightarrow}]$ are the field operators corresponding to input and output signal and idler fields. By the Fourier transformation we have

$$\hat{O}(\Omega) = \frac{1}{\sqrt{2\pi}} \int dt\, \hat{O}(t)e^{-i\Omega t}. \tag{2}$$

Here, the fields are described as functions of the modulation frequency Ω with commutation relation $[\hat{O}(\Omega), \hat{O}^+(\Omega')] = 2\pi\delta(\Omega - \Omega')$. A practical light field can be decomposed to a carrier $\hat{O}(0)$ oscillating at the centre frequency ω_0 with an average amplitude (O_{ss}) which equals to the amplitude of its steady state field, and surrounded by "noise side-bands" $\hat{O}(\Omega)$ oscillating at frequency $\omega_0 \pm \Omega$ with zero average amplitude [13]

$$\langle\hat{O}(\Omega = 0)\rangle = O_{ss}, \qquad \langle\hat{O}(\Omega \neq 0)\rangle = 0. \tag{3}$$

The noise spectral component at frequency Ω is the hereodyne mixing of the carrier and the noise side-bands. The amplitude and phase quadratures are expressed by

$$\begin{aligned} \hat{X}_O(\Omega) &= \hat{O}(\Omega) + \hat{O}^+(-\Omega), \\ \hat{Y}_O(\Omega) &= \frac{1}{i}\big[\hat{O}(\Omega) - \hat{O}^+(-\Omega)\big], \end{aligned} \tag{4}$$

with

$$\big[\hat{X}_O(\Omega), \hat{Y}_O\big(\Omega'\big)\big] = i\delta\big(\Omega + \Omega'\big). \tag{5}$$

The input–output Heisenberg evolutions of the field modes of the NOPA are given by [14]

$$\begin{aligned} \hat{b}_{0\updownarrow} &= \mu\hat{a}_{0\updownarrow} + \nu\hat{a}^+_{0\leftrightarrow}, & \hat{b}_{0\leftrightarrow} &= \mu\hat{a}_{0\leftrightarrow} + \nu\hat{a}^+_{0\updownarrow}, \\ \hat{b}_{+\updownarrow} &= \mu\hat{a}_{+\updownarrow} + \nu\hat{a}^+_{+\leftrightarrow}, & \hat{b}_{+\leftrightarrow} &= \mu\hat{a}_{+\leftrightarrow} + \nu\hat{a}^+_{+\updownarrow}, \\ \hat{b}_{-\updownarrow} &= \mu\hat{a}_{-\updownarrow} + \nu\hat{a}^+_{-\leftrightarrow}, & \hat{b}_{-\leftrightarrow} &= \mu\hat{a}_{-\leftrightarrow} + \nu\hat{a}^+_{-\updownarrow}, \end{aligned} \tag{6}$$

where $\hat{a}$, $\hat{a}^+$ and $\hat{b}$, $\hat{b}^+$ denote the annihilation and creation operators of the input and the output modes. The subindex 0 and $\pm$ stand for the central mode at frequency ω_0 and the side-bands at frequency $\omega_0 \pm \Omega$, respectively. The parameters $\mu = \cosh r$ and $\nu = e^{i\theta_p}\sinh r$ are the function of the squeezing factor r ($r \propto L\chi^2|a_p|$, L is the nonlinear crystal length, χ^2 is the effective second-order susceptibility of the nonlinear crystal in NOPA, a_p is the amplitude of pump field) and the phase θ_p of pump field. In the following calculation the phase θ_p is set to zero as the reference of relative phases of all other light fields. For bright optical field, the quadratures of the output orthogonal polarization modes at a certain rotated phase θ are expressed by

$$\hat{X}_{\hat{b}_{\updownarrow}}(\theta) = \frac{b^*_{0\updownarrow}\hat{b}_{+\updownarrow}e^{-i\theta} + b_{0\updownarrow}\hat{b}^+_{-\updownarrow}e^{i\theta}}{|b_{0\updownarrow}|}$$
$$= \hat{b}_{+\updownarrow}e^{-i(\theta+\varphi)} + \hat{b}^+_{-\updownarrow}e^{i(\theta+\varphi)},$$
$$\hat{X}_{\hat{b}_{\leftrightarrow}}(\theta) = \hat{b}_{+\leftrightarrow}e^{-i(\theta+\varphi)} + \hat{b}^+_{-\leftrightarrow}e^{i(\theta+\varphi)}, \tag{7}$$

where

$$\varphi = \arg(b_{0\updownarrow}) = \arg(b_{0\leftrightarrow}) = \arg\bigl(e^{i\Phi} + e^{-i\Phi}\tanh r\bigr)$$

is the phase of the modes $\hat{b}_{0\updownarrow}$, $\hat{b}_{0\leftrightarrow}$ relative to θ_p and Φ is the phase of the modes $\hat{a}_{0\updownarrow}$, $\hat{a}_{0\leftrightarrow}$ relative to θ_p. Taking $\theta = 0$ and $\theta = \pi/2$ in Eq. (7), the amplitude and phase quadrature of the output field are obtained

$$\hat{X}_{\hat{b}_{\updownarrow}} = \hat{X}_{\hat{b}_{\updownarrow}}(0) = \hat{b}_{+\updownarrow}e^{-i\varphi} + \hat{b}^+_{-\updownarrow}e^{i\varphi},$$
$$\hat{X}_{\hat{b}_{\leftrightarrow}} = \hat{X}_{\hat{b}_{\leftrightarrow}}(0) = \hat{b}_{+\leftrightarrow}e^{-i\varphi} + \hat{b}^+_{-\leftrightarrow}e^{i\varphi},$$
$$\hat{Y}_{\hat{b}_{\updownarrow}} = \hat{X}_{\hat{b}_{\updownarrow}}\left(\frac{\pi}{2}\right) = -i\bigl(\hat{b}_{+\updownarrow}e^{-i\varphi} - \hat{b}^+_{-\updownarrow}e^{i\varphi}\bigr),$$
$$\hat{Y}_{\hat{b}_{\leftrightarrow}} = \hat{X}_{\hat{b}_{\leftrightarrow}}\left(\frac{\pi}{2}\right) = -i\bigl(\hat{b}_{+\leftrightarrow}e^{-i\varphi} - \hat{b}^+_{-\leftrightarrow}e^{i\varphi}\bigr). \tag{8}$$

When the injected subharmonic signal and harmonic pump field are in phase ($\Phi = \varphi = 0$), the maximum parametric amplification is achieved [15]. The difference of the amplitude quadratures and the sum of the phase quadratures between two orthogonal polarization modes are

$$\hat{X}_{\hat{b}_{\updownarrow}} - \hat{X}_{\hat{b}_{\leftrightarrow}} = e^{-r}\hat{X}_{\hat{a}_{\updownarrow}} - e^{-r}\hat{X}_{\hat{a}_{\leftrightarrow}},$$
$$\hat{Y}_{\hat{b}_{\updownarrow}} + \hat{Y}_{\hat{b}_{\leftrightarrow}} = e^{-r}\hat{Y}_{\hat{a}_{\updownarrow}} - e^{-r}\hat{Y}_{\hat{a}_{\leftrightarrow}}. \tag{9}$$

Under the limit $r \to \infty$, the output orthogonal polarization modes are the perfect EPR beams with quadrature amplitude correlation and quadrature phase anticorrelation [7]. When the injected subharmonic signal and harmonic pump field are out of phase, i.e., $\Phi = \varphi = \pi/2$, NOPA operates at parametric deamplification [16]. Therefore the sum of the amplitude quadratures and the difference of the phase quadratures of the orthogonal polarization modes are as follows

$$\hat{X}_{\hat{b}_{\updownarrow}} + \hat{X}_{\hat{b}_{\leftrightarrow}} = e^{-r}\hat{Y}_{\hat{a}_{\updownarrow}} - e^{-r}\hat{Y}_{\hat{a}_{\leftrightarrow}},$$
$$\hat{Y}_{\hat{b}_{\updownarrow}} - \hat{Y}_{\hat{b}_{\leftrightarrow}} = -e^{-r}\hat{X}_{\hat{a}_{\updownarrow}} + e^{-r}\hat{X}_{\hat{a}_{\leftrightarrow}}. \tag{10}$$

Obviously, the EPR beams with the quadrature amplitude anticorrelation and quadrature phase correlation are obtained for $r > 0$ [9,10].

Now we consider the experimental NOPA system operating at parametric deamplification. The parameters μ and ν in NOPA process not only are the function of the squeezing factor r but also the function of the amplifying factor r' of unsqueezed quadrature component as

$$\mu = \frac{e^{r'} + e^{-r}}{2}, \qquad \nu = \frac{e^{r'} - e^{-r}}{2},$$

here, $r' \geqslant r$, which satisfy the uncertainty relations for the quadratures. The amplitude and phase quadratures of two orthogonal polarization output modes are obtained

$$\hat{X}_{\hat{b}_{\updownarrow}} = \mu\hat{Y}_{\hat{a}_{\updownarrow}} - \nu\hat{Y}_{\hat{a}_{\leftrightarrow}},$$
$$\hat{X}_{\hat{b}_{\leftrightarrow}} = \mu\hat{Y}_{\hat{a}_{\leftrightarrow}} - \nu\hat{Y}_{\hat{a}_{\updownarrow}},$$
$$\hat{Y}_{\hat{b}_{\updownarrow}} = -\mu\hat{X}_{\hat{a}_{\updownarrow}} - \nu\hat{X}_{\hat{a}_{\leftrightarrow}},$$
$$\hat{Y}_{\hat{b}_{\leftrightarrow}} = -\mu\hat{X}_{\hat{a}_{\leftrightarrow}} - \nu\hat{X}_{\hat{a}_{\updownarrow}}. \tag{11}$$

In Ref. [10] bright entangled EPR beam can be produced by combining two bright amplitude squeezed beams on a 50% beamsplitter. Conversely two bright amplitude squeezed beams may be obtained by combining bright entangled EPR beam on a 50% beamsplitter. The amplitude and phase quadrature of two bright amplitude squeezed beams produced from bright EPR beam are obtained from Eq. (11)

$$\hat{X}_{s_1} = \frac{1}{2}e^{-r}\bigl[\hat{Y}_{\hat{a}_{\updownarrow}} + \hat{Y}_{\hat{a}_{\leftrightarrow}} - \hat{X}_{\hat{a}_{\updownarrow}} + \hat{X}_{\hat{a}_{\leftrightarrow}}\bigr],$$
$$\hat{Y}_{s_1} = \frac{1}{2}e^{r'}\bigl[\hat{Y}_{\hat{a}_{\updownarrow}} - \hat{Y}_{\hat{a}_{\leftrightarrow}} - \hat{X}_{\hat{a}_{\updownarrow}} - \hat{X}_{\hat{a}_{\leftrightarrow}}\bigr],$$
$$\hat{X}_{s_2} = \frac{1}{2}e^{-r}\bigl[\hat{Y}_{\hat{a}_{\updownarrow}} + \hat{Y}_{\hat{a}_{\leftrightarrow}} + \hat{X}_{\hat{a}_{\updownarrow}} - \hat{X}_{\hat{a}_{\leftrightarrow}}\bigr],$$
$$\hat{Y}_{s_2} = \frac{1}{2}e^{r'}\bigl[-\hat{Y}_{\hat{a}_{\updownarrow}} + \hat{Y}_{\hat{a}_{\leftrightarrow}} - \hat{X}_{\hat{a}_{\updownarrow}} - \hat{X}_{\hat{a}_{\leftrightarrow}}\bigr]. \tag{12}$$

When the input modes $\hat{a}_{\updownarrow}$, $\hat{a}_{\leftrightarrow}$ of NOPA are the coherent state, $\langle\delta\hat{X}^2_{\hat{a}_{\updownarrow}}\rangle = \langle\delta\hat{X}^2_{\hat{a}_{\leftrightarrow}}\rangle = \langle\delta\hat{Y}^2_{\hat{a}_{\updownarrow}}\rangle = \langle\delta\hat{Y}^2_{\hat{a}_{\leftrightarrow}}\rangle = 1$, we can readily write up the variances of two bright amplitude squeezed beams

$$\bigl\langle\delta\hat{X}^2_{s_1}\bigr\rangle = \bigl\langle\delta\hat{X}^2_{s_2}\bigr\rangle = e^{-2r},$$
$$\bigl\langle\delta\hat{Y}^2_{s_1}\bigr\rangle = \bigl\langle\delta\hat{Y}^2_{s_2}\bigr\rangle = e^{2r'}.$$

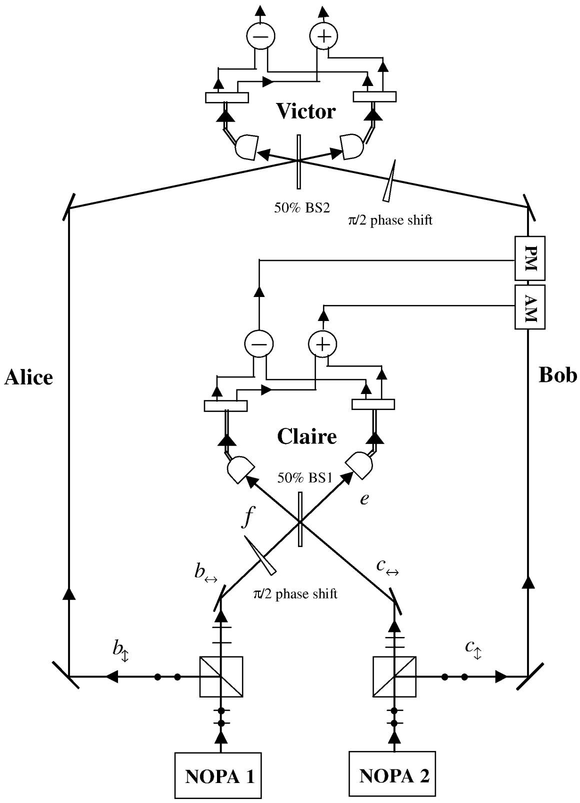

Fig. 2. Entanglement swapping using two NOPAs.

The Heisenberg inequality of two bright amplitude squeezed beams are given

$$\langle\delta\hat{X}^2_{s_{1(2)}}\rangle\langle\delta\hat{Y}^2_{s_{1(2)}}\rangle = e^{2(r'-r)} \geqslant 1. \tag{13}$$

When the r' is equal to r, two bright amplitude squeezed beams are the minimum-uncertainty state. In the practical NOPA systems, the r' is always larger than the r. Thus the two bright amplitude squeezed beams produced from bright EPR beam are not the minimum-uncertainty state.

For our entanglement swapping scheme as shown in Fig. 2, we need two NOPA: a bright entangled state of two orthogonal polarization output modes $\hat{b}_{\updownarrow}$, $\hat{b}_{\leftrightarrow}$ and a bright entangled state of two orthogonal polarization output modes $\hat{c}_{\updownarrow}$, $\hat{c}_{\leftrightarrow}$. Let us introduce Alice, Bob, Claire and Victor to illustrate the whole protocol with entanglement swapping and subsequent experimental verification of entanglement swapping. Alice and Claire shall share the entangled state of modes $\hat{b}_{\updownarrow}$ and $\hat{b}_{\leftrightarrow}$ while Claire and Bob are sharing the other entangled state of modes $\hat{c}_{\updownarrow}$ and $\hat{c}_{\leftrightarrow}$. Thus, initially Alice and Bob do not share an entangled state. However, we will see that Alice and Bob can establish the entanglement of mode $\hat{b}_{\updownarrow}$ and $\hat{c}_{\updownarrow}$ with information about Claire's measurement results. Let us assume Bob obtains the classical results from Claire. Claire perform a joint measurement for mode $\hat{b}_{\leftrightarrow}$ and $\hat{c}_{\leftrightarrow}$ by the direct measurement of Bell-state. $\hat{b}_{\leftrightarrow}$ is phase-shifted $\pi/2$ then is mixed with $\hat{c}_{\leftrightarrow}$ on a 50% beamsplitter (BS1). The bright output beams, $\hat{e}$ and $\hat{f}$, are directly detected by D_1 and D_2. The e and f are given by

$$\hat{e} = \frac{\sqrt{2}}{2}\big(\hat{b}_{\leftrightarrow} + i\hat{c}_{\leftrightarrow}\big),$$
$$\hat{f} = \frac{\sqrt{2}}{2}\big(\hat{b}_{\leftrightarrow} - i\hat{c}_{\leftrightarrow}\big). \tag{14}$$

Each of the detected photocurrents is divided into two parts by the RF power splitters. The sum and difference of the divided photocurrents are expressed by

$$i_+ = \frac{1}{\sqrt{2}}\big(\hat{X}_{\hat{b}_{\leftrightarrow}} + \hat{X}_{\hat{c}_{\leftrightarrow}}\big),$$
$$i_- = \frac{1}{\sqrt{2}}\big(\hat{Y}_{\hat{b}_{\leftrightarrow}} - \hat{Y}_{\hat{c}_{\leftrightarrow}}\big). \tag{15}$$

Then the photocurrents are sent to amplitude and phase modulators in the receiver (Bob), respectively. The amplitude and phase modulators transform the photocurrents into mode $\hat{c}_{\updownarrow}$. The output beam from modulators is found to be

$$\hat{c}'_{\updownarrow} = \hat{c}_{\updownarrow} + \sqrt{2}\,g_{\mathrm{swap}} i_+ + i\sqrt{2}\,g_{\mathrm{swap}} i_-, \tag{16}$$

where g_{swap} describe Bob's (suitably normalized) amplitude and phase gain for the transformation from photocurrent to output beam. Now Alice and Bob send mode $\hat{b}_{\updownarrow}$ and $\hat{c}'_{\updownarrow}$, respectively, to Victor, then Victor perform another joint measurement by the direct measurement of Bell-state to verify the entanglement swapping. The sum and difference of the divided photocurrents at Victor are expressed by

$$i^v_+ = \frac{1}{\sqrt{2}}\big(\hat{X}_{\hat{b}_{\updownarrow}} + \hat{X}_{\hat{c}'_{\updownarrow}}\big)$$
$$= \frac{1}{\sqrt{2}}\big[\big(\hat{X}_{\hat{b}_{\updownarrow}} + g_{\mathrm{swap}}\hat{X}_{\hat{b}_{\leftrightarrow}}\big) + \big(\hat{X}_{\hat{c}_{\updownarrow}} + g_{\mathrm{swap}}\hat{X}_{\hat{c}_{\leftrightarrow}}\big)\big],$$

$$i_-^v = \frac{1}{\sqrt{2}}\left(\hat{Y}_{\hat{b}_\updownarrow} - \hat{Y}_{\hat{c}'_\updownarrow}\right)$$
$$= \frac{1}{\sqrt{2}}\left[\left(\hat{Y}_{\hat{b}_\updownarrow} - g_{\text{swap}}\hat{Y}_{\hat{b}_\leftrightarrow}\right) + \left(\hat{Y}_{\hat{c}_\updownarrow} - g_{\text{swap}}\hat{Y}_{\hat{c}_\leftrightarrow}\right)\right]. \tag{17}$$

From Eqs. (11) and (17) we can readily write out the variances of the sum and difference of the divided photocurrents

$$V_{i_+^v} = V_{i_-^v} = \left(\frac{e^{r'} + e^{-r}}{2} - g_{\text{swap}}\frac{e^{r'} - e^{-r}}{2}\right)^2 + \left(\frac{e^{r'} - e^{-r}}{2} - g_{\text{swap}}\frac{e^{r'} + e^{-r}}{2}\right)^2, \tag{18}$$

where the input modes of NOPA all are the coherent state and two NOPAs have the same squeezing factor r and the amplifying factor r' of unsqueezed quadrature component. The mode $\hat{b}_\updownarrow$ and $\hat{c}'_\updownarrow$ possess the quantum entanglement when $V_{i_+^v} < 1$ and $V_{i_-^v} < 1$. The smaller is the variances of the sum and difference of the divided photocurrents, the large is the entanglement of mode $\hat{b}_\updownarrow$ and $\hat{c}'_\updownarrow$. There is an optimum gain for the maximum entanglement, which one can easily find by minimizing $V_{i_+^v}$ and $V_{i_-^v}$ to be

$$g_{\text{swap}}^{\text{opt}} = \frac{e^{2r'} - e^{-2r}}{e^{2r'} + e^{-2r}},$$
$$V_{i_+^v}^{\text{opt}} = V_{i_-^v}^{\text{opt}} = \frac{2e^{2(r'-r)}}{e^{2r'} + e^{-2r}}. \tag{19}$$

From Eq. (19) it is obvious that the quality of the entanglement from entanglement swapping depends on not only squeezing factor r but also the amplifying factor r' of unsqueezed quadrature component as shown in Fig. 3. The variances of joint measurement of NOPA 1 or 2 only depend on the squeezing factor r, and the quality of entanglement of NOPA 1 or 2 is larger than that of the entanglement from entanglement swapping (curve (e) of Fig. 3). Under the same r, the larger is the amplifying factor r' of unsqueezed quadrature component, the smaller is the quality of the entanglement from entanglement swapping. When r' has infinite large noise, the creation of entanglement between mode $\hat{b}_\updownarrow$ and $\hat{c}_\updownarrow$ is possible only with more than 3 dB squeezing ($e^{-2r} < 0.5$) and in this case $g_{\text{swap}}^{\text{opt}}$ is unity gain.

We proposed easily realized scheme of entanglement swapping for continuous variables using two

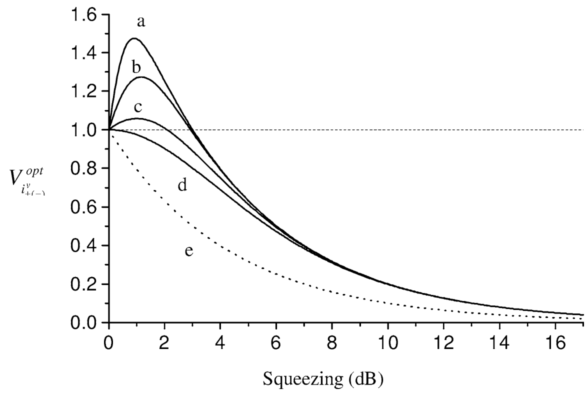

Fig. 3. Variance of Victor's joint measurement, (a) $r' = 10r$, (b) $r' = 5r$, (c) $r' = 2r$, (d) $r' = r$, (e) variance of joint measurement of NOPA 1 or 2.

NOPA operating in the state of deamplification. We pointed out the quality of the entanglement from entanglement swapping depend on not only squeezing factor r but also the amplifying factor r' of unsqueezed quadrature component. Due to adopting the bright EPR beams and the simple direct measurement for Bell-state, the entanglement swapping and the verification of entanglement swapping is within the reach of current technology and significantly simplify the implementation. On the one hand, the joint measurement of the entangled beam is an important operation in quantum information, such as for dense coding, on the other hand, the local measurement with classical communication is another important method which has been used to demonstrate the EPR-type entanglement [11]. It may reduces the loss of the transmission of quantum state from Alice and Bob to Victor.

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