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Quantum switch for continuous variable teleportation

Junxiang Zhang, Changde Xie and Kunchi Peng

The Key Laboratory of Quantum Optics, Ministry of Education, People's Republic of China and Institute of Opto-electronics, Shanxi University, Taiyuan, Shanxi 030006, People's Republic of China

E-mail: kcpeng@sxu.edu.cn

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Abstract
We propose a quantum teleportation scheme in which a quantum state is teleported from the sending station (Alice) to either of two receiving stations (Bob1, Bob2). In this scheme, two pairs of EPR beams with identical frequency and constant phase relation are used to produce two pairs of conditional entangled beams by composing their modes on two beamsplitters. One output of a beamsplitter is sent to Alice and the two outputs of the other beamsplitter are sent to Bob1 and Bob2. Which receiving station actually receives the teleported state can be decided by correlating the in-phase or out-of-phase quadrature components of two two-mode squeezed vacuum states. The switch system manipulated by squeezed state light might be developed as a practical quantum switch device for the communication and teleportation of quantum information.

Keywords: Teleportation, EPR pairs, squeezing

Quantum information technology aims to achieve performance in communication and computation systems superior to those based on classical physics by utilizing the nonlocal quantum correlations of entangled states. It has also significantly improved understanding of the quantum systems involved in the factual realization of a quantum computer, and raised many interesting problems such as the encoding of information [1], entanglement of states [2], quantum cryptography [3], quantum information manipulation [4] and quantum communication of teleportation [5]. One of the most striking features of quantum information is quantum teleportation. In the quantum teleportation scheme, quantum information of an unknown state is faithfully transmitted from a sender to a remote receiver via initially shared EPR pairs which function as a quantum information channel for the faithful transmission. Quantum teleportation was originally proposed for a discrete variable in finite-dimensional Hilbert space, later it was successfully developed from a discrete quantum system into a quantum system for continuous variables [6–9]. Teleportation of optical fields holds great promise due to the power of the required optical tools and the maturity of relevant optical communications technology. Quantum teleportation represents the basic building block of future quantum communication networks between distant parties [10].

Some attempts have been made to enhance the performance of a quantum teleportation system. One of them is to teleport a quantum state from the sender to either of two receivers using three-particle entanglement at a certain condition of measurement [11]. Because of the experimental difficulties in generating the multiparticle-entanglement state [12], this teleportation scheme has not been implemented in experiments. In this paper, we propose a novel scheme to teleport a quantum state from Alice to two different receivers in turn using a two-mode squeezed state as the quantum switch to manipulate the transmission route. In this scheme, the EPR entangled beams shared by Alice and the two Bobs are produced by mixing a pair of two-mode squeezed state lights with identical frequency and constant phase relation on two 50% beamsplitters. As in the usual teleportation scheme [9], one performs a certain joint measurement on the unknown input quantum state and one output of a beamsplitter at Alice, and Alice's measurement results are split into two identical parts by radio frequency (RF) power splitters and transmitted to two Bob receivers, then the receivers perform

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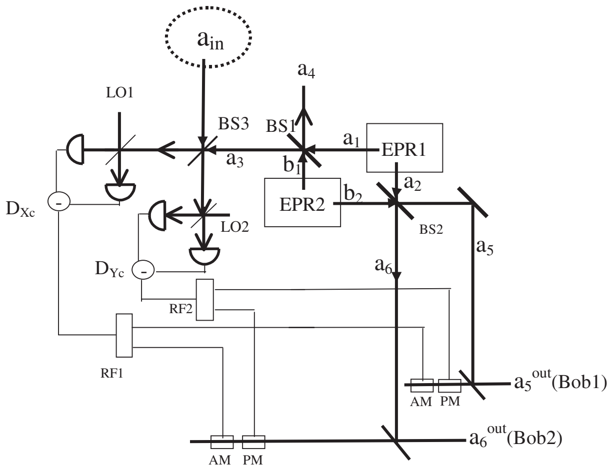

Figure 1. Schematic of the quantum teleportation switch, BS: beamsplitter, LO: local oscillator for homodyne detection, $D_{\rm Xc}$: homodyne detector of in-phase quadrature component, $D_{\rm Yc}$: homodyne detector of out-of-phase quadrature component, RF: RF power splitters, AM: amplitude modulator, PM: phase modulator.

linear displacement of the two outputs of the other beamsplitter in order to produce the desired input state. Converting the squeezing parameters or correlations between the in-phase and out-of-phase quadrature components, which can be easily realized by converting the relative phase of the pump field and injected field of an OPA between 0 and $\pi/2$ [13], the original unknown input quantum state can be conditionally mimicked at either of the two Bobs in turn, so that the squeezing direction of the two-mode squeezed state light plays the role of a quantum switch between two spatially separated receivers.

The scheme of the proposed system is shown in figure 1. A pair of two-mode squeezed state lights are used as the two EPR sources (EPR1 and EPR2). In the Heisenberg representation, the in-phase quadrature and out-of-phase quadrature phases $\hat{X}$ and $\hat{Y}$ for two modes of each two-mode squeezed state are expressed as follows [12, 14]:

$$\begin{aligned}
\hat{X}_{a1} &= (\mathrm{e}^{r_a}\hat{X}_{a1}^{(0)} + \mathrm{e}^{-r_a}\hat{X}_{a2}^{(0)})/\sqrt{2},\\
\hat{Y}_{a1} &= (\mathrm{e}^{-r_a}\hat{Y}_{a1}^{(0)} + \mathrm{e}^{r_a}\hat{Y}_{a2}^{(0)})/\sqrt{2}\\
\hat{X}_{a2} &= (\mathrm{e}^{r_a}\hat{X}_{a1}^{(0)} - \mathrm{e}^{-r_a}\hat{X}_{a2}^{(0)})/\sqrt{2},\\
\hat{Y}_{a2} &= (\mathrm{e}^{-r_a}\hat{Y}_{a1}^{(0)} - \mathrm{e}^{r_a}\hat{Y}_{a2}^{(0)})/\sqrt{2}\\
\hat{X}_{b1} &= (\mathrm{e}^{r_b}\hat{X}_{b1}^{(0)} + \mathrm{e}^{-r_b}\hat{X}_{b2}^{(0)})/\sqrt{2},\\
\hat{Y}_{b1} &= (\mathrm{e}^{-r_b}\hat{Y}_{b1}^{(0)} + \mathrm{e}^{r_b}\hat{Y}_{b2}^{(0)})/\sqrt{2}\\
\hat{X}_{b2} &= (\mathrm{e}^{r_b}\hat{X}_{b1}^{(0)} - \mathrm{e}^{-r_b}\hat{X}_{b2}^{(0)})/\sqrt{2},\\
\hat{Y}_{b2} &= (\mathrm{e}^{-r_b}\hat{Y}_{b1}^{(0)} - \mathrm{e}^{r_b}\hat{Y}_{b2}^{(0)})/\sqrt{2},
\end{aligned} \tag{1}$$

where the subscripts a_1, a_2 and b_1, b_2 are designated for the two modes in the two-mode squeezed state from EPR1 and EPR2, respectively. The superscript '(0)' denotes initial coherent modes, r_a and r_b are the correlation parameters between the a_1 and a_2 as well as the b_1 and b_2 modes.

For the case of finite squeezing, when the uncertainty product for the variances of the two inferences $\langle\Delta^2\hat{X}_{\rm inf}\rangle\langle\Delta^2\hat{Y}_{\rm inf}\rangle = \langle\Delta^2(\hat{X}_1 - g_x\hat{X}_2)\rangle\langle\Delta^2(\hat{Y}_1 + g_y\hat{Y}_2)\rangle$ is less than the limit of unity associated with the Heisenberg uncertainty relation, the EPR paradox for continuous variables is demonstrated [15, 16], here $g_{x,y}$ is the scaling factor for minimizing the variances of $\langle\Delta^2\hat{X}_{\rm inf}\rangle$ and $\langle\Delta^2\hat{Y}_{\rm inf}\rangle$. From equation (1) we can calculate $\langle\Delta^2\hat{X}_{\rm inf}\rangle\langle\Delta^2\hat{Y}_{\rm inf}\rangle = (1/\cosh 2r_{a(b)})^2$, once squeezing exists, i.e. $r_{a(b)} > 0$, the uncertainty product $\langle\Delta^2\hat{X}_{\rm inf}\rangle\langle\Delta^2\hat{Y}_{\rm inf}\rangle < 1$, thus the imperfect squeezed state corresponds to the non-ideal EPR pairs. Under the limit of perfect correlation $r_{a(b)} \to \infty$, the uncertainty product $\langle\Delta^2\hat{X}_{\rm inf}\rangle\langle\Delta^2\hat{Y}_{\rm inf}\rangle = 0$ and $g_{x,y} = 1$, the two-mode squeezed state approaches the eigenstates of $\hat{X}_{a(b)1} - \hat{X}_{a(b)2}$ and $\hat{Y}_{a(b)1} + \hat{Y}_{a(b)2}$, which corresponds to perfect EPR pairs having the in-phase quadrature phase correlation and out-of-phase quadrature phase anticorrelation. Otherwise, if $-r_{a(b)} > 0$, the uncertainty product $\langle\Delta^2(\hat{X}_{a(b)1}+g_x\hat{X}_{a(b)2})\rangle\langle\Delta^2(\hat{Y}_{a(b)1}-g_y\hat{Y}_{a(b)2})\rangle < 1$. Especially for $r_{a(b)} \to -\infty$, it approaches the eigenstates of $\hat{X}_{a(b)1}+\hat{X}_{a(b)2}$ and $\hat{Y}_{a(b)1} - \hat{Y}_{a(b)2}$, which are also the perfect EPR state with the anticorrelated in-phase quadrature components and correlated out-of-phase quadratures components. Therefore we say that the two-mode squeezed vacuum state produces EPR entanglement of quadrature phase components. When $r = 0$ we have $\langle\Delta^2\hat{X}_{\rm inf}\rangle\langle\Delta^2\hat{Y}_{\rm inf}\rangle = 1$, that is the classical limit of EPR entanglement.

In order to perform a teleportation switch, the sender has to share the entanglement with two receivers for the different cases. Initially, mode a_1 shares the entanglement with mode a_2 and mode b_1 shares the entanglement with mode b_2, then we mix mode a_1 and a_2 with the mode b_1 and b_2 at the beamsplitters BS1 and BS2, respectively. The output modes of the two beamsplitters are

$$\begin{aligned}
\hat{X}_3 &= (\hat{X}_{a1} - \hat{X}_{b1})/\sqrt{2}, &\quad \hat{Y}_3 &= (\hat{Y}_{a1} - \hat{Y}_{b1})/\sqrt{2}\\
\hat{X}_4 &= (\hat{X}_{a1} + \hat{X}_{b1})/\sqrt{2}, &\quad \hat{Y}_4 &= (\hat{Y}_{a1} + \hat{Y}_{b1})/\sqrt{2}\\
\hat{X}_5 &= (\hat{X}_{a2} + \hat{X}_{b2})/\sqrt{2}, &\quad \hat{Y}_5 &= (\hat{Y}_{a2} + \hat{Y}_{b2})/\sqrt{2}\\
\hat{X}_6 &= (\hat{X}_{a2} - \hat{X}_{b2})/\sqrt{2}, &\quad \hat{Y}_6 &= (\hat{Y}_{a2} - \hat{Y}_{b2})/\sqrt{2}.
\end{aligned} \tag{2}$$

From equation (2), we get

$$\begin{aligned}
&\langle\Delta^2(\hat{X}_3 - g_x\hat{X}_5)\rangle\langle\Delta^2(\hat{Y}_3 + g_y\hat{Y}_5)\rangle\\
&= \frac{2 + \exp[2(r_a + r_b)] + \exp[-2(r_a + r_b)]}{\exp[2r_a] + \exp[-2r_a] + \exp[2r_b] + \exp[-2r_b]},\\
&\langle\Delta^2(\hat{X}_3 - g_x\hat{X}_6)\rangle\langle\Delta^2(\hat{Y}_3 + g_y\hat{Y}_6)\rangle\\
&= \frac{2 + \exp[2(r_a - r_b)] + \exp[-2(r_a - r_b)]}{\exp[2r_a] + \exp[-2r_a] + \exp[2r_b] + \exp[-2r_b]}.
\end{aligned} \tag{3}$$

The dependence of the uncertainty product for the variances of the two inferences of $\langle\Delta^2(\hat{X}_3 - g_x\hat{X}_5)\rangle\langle\Delta^2(\hat{Y}_3 + g_y\hat{Y}_5)\rangle$ and $\langle\Delta^2(\hat{X}_3 - g_x\hat{X}_6)\rangle\langle\Delta^2(\hat{Y}_3 + g_y\hat{Y}_6)\rangle$ upon the squeezing of two two-mode squeezed states is shown in figures 2 and 3. Figure 3 shows that for both cases of $r_a > 0$ and $r_b > 0$, $\langle\Delta^2(\hat{X}_3 - g_x\hat{X}_6)\rangle\langle\Delta^2(\hat{Y}_3 + g_y\hat{Y}_6)\rangle < 1$ which demonstrates the EPR entanglement for modes 3 and 6. In figure 2 it is shown that only for the case of $r_a > 0$ and $r_b < 0$, modes 3 and 5 will present the EPR entanglement.

In the case of perfect squeezing, the relations between modes 3 and 5, 6 can be written as

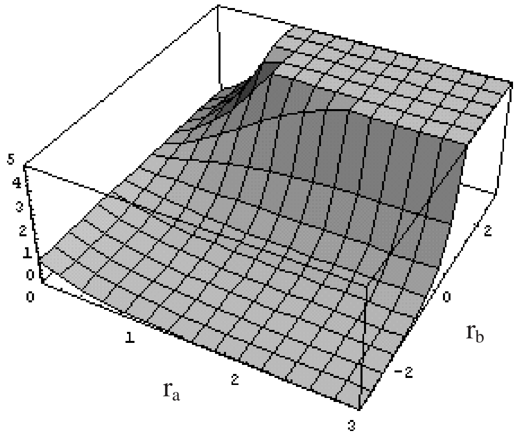

Figure 2. The uncertainty product for the variances of the two inferences of $\langle\Delta^2(\hat{X}_3 - g_x\hat{X}_5)\rangle\langle\Delta^2(\hat{Y}_3 + g_y\hat{Y}_5)\rangle$ is plotted versus the squeezing parameters r_a and r_b of two two-mode squeezed states. Demonstration of the EPR paradox requires $\langle\Delta^2(\hat{X}_3 - g_x\hat{X}_5)\rangle\langle\Delta^2(\hat{Y}_3 + g_y\hat{Y}_5)\rangle < 1$. It is obvious that when $r_a > 0$ and $-r_b > 0$, the product will go below 1 which demonstrates the EPR entanglement between modes 3 and 5.

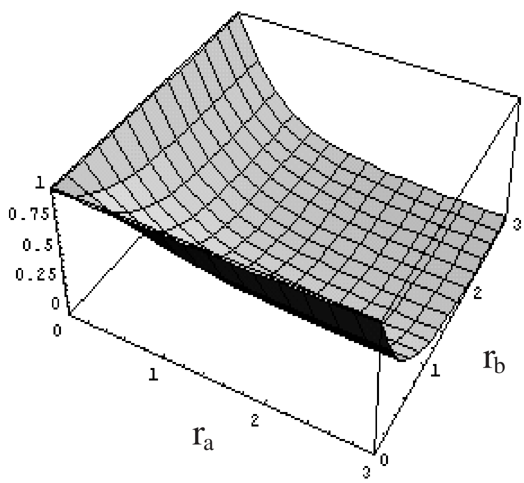

Figure 3. The uncertainty product for the variances of the two inferences of $\langle\Delta^2(\hat{X}_3 - g_x\hat{X}_6)\rangle\langle\Delta^2(\hat{Y}_3 + g_y\hat{Y}_6)\rangle$ is plotted versus the squeezing parameters r_a and r_b of two two-mode squeezed states. The EPR entanglement between modes 3 and 6 requires that both two two-mode squeezed state are squeezed, i.e. $r_a > 0$ and $r_b > 0$.

$$
\begin{aligned}
\hat{X}_3 - \hat{X}_5 &= [(\hat{X}_{a1} - \hat{X}_{a2}) - (\hat{X}_{b1} + \hat{X}_{b2})]/\sqrt{2} \\
\hat{Y}_3 + \hat{Y}_5 &= [(\hat{Y}_{a1} + \hat{Y}_{a2}) - (\hat{Y}_{b1} - \hat{Y}_{b2})]/\sqrt{2} \\
\hat{X}_3 - \hat{X}_6 &= [(\hat{X}_{a1} - \hat{X}_{a2}) - (\hat{X}_{b1} - \hat{X}_{b2})]/\sqrt{2} \\
\hat{Y}_3 + \hat{Y}_6 &= [(\hat{Y}_{a1} + \hat{Y}_{a2}) - (\hat{Y}_{b1} + \hat{Y}_{b2})]/\sqrt{2}.
\end{aligned} \tag{4}
$$

It is obvious that the modes 3 and 5 become the perfect EPR pairs if one of the two-mode squeezed states is the eigenstate of $\hat{X}_{a1} - \hat{X}_{a2}$ and $\hat{Y}_{a1} + \hat{Y}_{a2}$ (i.e. $r_a \to \infty$), and the other one is the eigenstate of $\hat{X}_{b1} + \hat{X}_{b2}$ and $\hat{Y}_{b1} - \hat{Y}_{b2}$ (i.e. $r_b \to -\infty$). Meanwhile, mode 3 will be perfectly entangled with mode 6 at the conditions of $r_a \to \infty$, $r_b \to \infty$. According to the different entanglement condition the quantum information from a sender can be controllably transmitted to Bob1's mode 5 or Bob2's mode 6.

In Alice's station an unknown input quantum state represented by the quadrature operators $\hat{X}_{\text{in}}$ and $\hat{Y}_{\text{in}}$ is superposed with mode 3 at the beamsplitter BS3. Two balanced homodyne detectors D_{Xc} and D_{Yc} are used to measure the observables of the in-phase quadrature phase $\hat{X}_{\text{c}} = (\hat{X}_{\text{in}} - \hat{X}_3)/\sqrt{2}$ of one output of BS3 and the out-of-phase quadrature phase $\hat{Y}_{\text{c}} = (\hat{Y}_{\text{in}} + \hat{Y}_3)/\sqrt{2}$ of the other output of BS3. The resulting classical outcomes are scaled by the operators

$$
\begin{aligned}
\hat{X}_{\text{c}} &= [\hat{X}_{\text{in}} - (\hat{X}_{a1} - \hat{X}_{b1})/\sqrt{2}]/\sqrt{2} \\
\hat{Y}_{\text{c}} &= [\hat{Y}_{\text{in}} + (\hat{Y}_{a1} - \hat{Y}_{b1})/\sqrt{2}]/\sqrt{2}.
\end{aligned} \tag{5}
$$

Now we rewrite modes 5 and 6 in the following forms:

$$
\begin{aligned}
\hat{X}_5 &= \hat{X}_{\text{in}} + [(\hat{X}_{a2} - \hat{X}_{a1}) + (\hat{X}_{b2} + \hat{X}_{b1})]/\sqrt{2} - \sqrt{2}\hat{X}_{\text{c}} \\
\hat{Y}_5 &= \hat{Y}_{\text{in}} + [(\hat{Y}_{a2} + \hat{Y}_{a1}) + (\hat{Y}_{b2} - \hat{Y}_{b1})]/\sqrt{2} - \sqrt{2}\hat{Y}_{\text{c}} \\
\hat{X}_6 &= \hat{X}_{\text{in}} + [(\hat{X}_{a2} - \hat{X}_{a1}) - (\hat{X}_{b2} - \hat{X}_{b1})]/\sqrt{2} - \sqrt{2}\hat{X}_{\text{c}} \\
\hat{Y}_6 &= \hat{Y}_{\text{in}} + [(\hat{Y}_{a2} + \hat{Y}_{a1}) - (\hat{Y}_{b2} + \hat{Y}_{b1})]/\sqrt{2} - \sqrt{2}\hat{Y}_{\text{c}}.
\end{aligned} \tag{6}
$$

After measurement which yields the classical results X_{c} and Y_{c}, Bob1's mode 5 and Bob2's mode 6 in equation (6) collapse into [17]

$$
\begin{aligned}
\hat{X}_5 &= \hat{X}_{\text{in}} + [(\hat{X}_{a2} - \hat{X}_{a1}) + (\hat{X}_{b2} + \hat{X}_{b1})]/\sqrt{2} - \sqrt{2}X_{\text{c}} \\
\hat{Y}_5 &= \hat{Y}_{\text{in}} + [(\hat{Y}_{a2} + \hat{Y}_{a1}) + (\hat{Y}_{b2} - \hat{Y}_{b1})]/\sqrt{2} - \sqrt{2}Y_{\text{c}} \\
\hat{X}_6 &= \hat{X}_{\text{in}} + [(\hat{X}_{a2} - \hat{X}_{a1}) - (\hat{X}_{b2} - \hat{X}_{b1})]/\sqrt{2} - \sqrt{2}X_{\text{c}} \\
\hat{Y}_6 &= \hat{Y}_{\text{in}} + [(\hat{Y}_{a2} + \hat{Y}_{a1}) - (\hat{Y}_{b2} + \hat{Y}_{b1})]/\sqrt{2} - \sqrt{2}Y_{\text{c}}.
\end{aligned} \tag{7}
$$

Due to the entanglement between modes a_3 and a_5, or a_3 and a_6 for the cases of $r_a \to \infty$, $r_b \to -\infty$ or $r_a \to \infty$, $r_b \to \infty$, the measurements lead to the collapse of the modes a_5 or a_6 into a state which differs from the unknown input state in the classical phase-space displacement. Thus for the possibility to recover the input state at the two locations, each of the classical outcomes is divided into two identical parts with RF power splitters (RF1 and RF2) [18], and the classical information is sent separately to the two remote locations for performing appropriate displacements on modes 5 and 6:

$$
\begin{aligned}
\hat{X}_5 &\to \hat{X}_5^{\text{out}} = \hat{X}_5 + \sqrt{2}g_1X_{\text{c}} \\
\hat{Y}_5 &\to \hat{Y}_5^{\text{out}} = \hat{Y}_5 + \sqrt{2}g_1Y_{\text{c}} \\
\hat{X}_6 &\to \hat{X}_6^{\text{out}} = \hat{X}_6 + \sqrt{2}g_2X_{\text{c}} \\
\hat{Y}_6 &\to \hat{Y}_6^{\text{out}} = \hat{Y}_6 + \sqrt{2}g_2Y_{\text{c}}.
\end{aligned} \tag{8}
$$

According to equations (1), (5), (7) and (8), the outgoings of two modes become

$$
\begin{aligned}
\hat{a}_5^{\text{out}} &= g_1\hat{a}_{\text{in}} + \frac{1+g_1}{2}[\mathrm{e}^{-r_a}(-\hat{X}_{a2}^{(0)} + \mathrm{i}\hat{Y}_{a1}^{(0)}) + \mathrm{e}^{r_b}(\hat{X}_{b1}^{(0)} - \mathrm{i}\hat{Y}_{b2}^{(0)})] \\
&\quad + \frac{1-g_1}{2}[\mathrm{e}^{r_a}(\hat{X}_{a1}^{(0)} - \mathrm{i}\hat{Y}_{a2}^{(0)}) + \mathrm{e}^{-r_b}(-\hat{X}_{b2}^{(0)} + \mathrm{i}\hat{Y}_{b1}^{(0)})] \\
\hat{a}_6^{\text{out}} &= g_2\hat{a}_{\text{in}} + \frac{1+g_2}{2}[\mathrm{e}^{-r_a}(-\hat{X}_{a2}^{(0)} + \mathrm{i}\hat{Y}_{a1}^{(0)})
\end{aligned}
$$

$$+\mathrm{e}^{-r_b}(\hat{X}_{b2}^{(0)} - \mathrm{i}\hat{Y}_{b1}^{(0)})] + \frac{1-g_2}{2}[\mathrm{e}^{r_a}(\hat{X}_{a1}^{(0)} - \mathrm{i}\hat{Y}_{a2}^{(0)})$$
$$+\mathrm{e}^{r_b}(-\hat{X}_{b1}^{(0)} + \mathrm{i}\hat{Y}_{b2}^{(0)})], \quad (9)$$

where the parameters g_1 and g_2 describe the normalized gain of two teleportation processes from a sender to Bob1 and Bob2.

Equation (9) shows that both the output modes $\hat{a}_5^{\text{out}}$ and $\hat{a}_6^{\text{out}}$ contain some information about the teleported state but it is not exactly the input state due to some additional noise from the quantum channels.

For the case of ideal measurement process $g_1 = 1$ and $r_a \to \infty$, $r_b \to -\infty$. Equation (9) becomes $\hat{a}_5^{\text{out}} = \hat{a}_{\text{in}}$. So perfect quantum teleportation is accomplished at the Bob1 receiver.

If $g_2 = 1$ and $r_a \to \infty$, $r_b \to \infty$, we have $\hat{a}_6^{\text{out}} = \hat{a}_{\text{in}}$, then the unknown quantum state is perfectly mimicked at the Bob2 receiver.

The fidelity quantifying the quality of teleportation is defined for a coherent input state ($|\alpha\rangle$) by $F = \langle\alpha|\hat{\rho}_{\text{out}}|\alpha\rangle$ [19,20], it describes the match between the input and the teleported state. Up to a factor π, this fidelity is the Q function of the teleported field evaluated for α:

$$F = \pi Q_{\text{tel}}(\alpha) = \frac{2}{\sqrt{(\langle\delta^2\hat{X}_{\text{out}}\rangle + 1)(\langle\delta^2\hat{Y}_{\text{out}}\rangle + 1)}} \times \exp\left[-2\frac{(1-g)^2|\alpha_{\text{in}}|^2}{\sqrt{(\langle\delta^2\hat{X}_{\text{out}}\rangle + 1)(\langle\delta^2\hat{Y}_{\text{out}}\rangle + 1)}}\right], \quad (10)$$

where $\delta^2\hat{X}_{\text{out}}$ and $\delta^2\hat{Y}_{\text{out}}$ are the variance of in-phase quadrature and out-of-phase quadrature phases of the output mode, g describes a normalized gain for the transformation from classical values to complex field amplitude performed by the Bob receivers, it will be g_1 for Bob1 and g_2 for Bob2. Using equation (9), they are given by

$$\langle\delta^2\hat{X}_{\text{out}}\rangle = g_{1(2)}^2\langle\delta^2\hat{X}_{\text{in}}\rangle + \left(\frac{1+g_{1(2)}}{2}\right)^2[\mathrm{e}^{-2r_a} + \mathrm{e}^{\pm 2r_b}] + \left(\frac{1-g_{1(2)}}{2}\right)^2[\mathrm{e}^{2r_a} + \mathrm{e}^{\mp 2r_b}]$$
$$\langle\delta^2\hat{Y}_{\text{out}}\rangle = g_{1(2)}^2\langle\delta^2\hat{Y}_{\text{in}}\rangle + \left(\frac{1+g_{1(2)}}{2}\right)^2[\mathrm{e}^{-2r_a} + \mathrm{e}^{\pm 2r_b}] + \left(\frac{1-g_{1(2)}}{2}\right)^2[\mathrm{e}^{2r_a} + \mathrm{e}^{\mp 2r_b}]. \quad (11)$$

In equation (11), the symbol '$\pm$' and '1(2)' represent that the teleportation is accomplished at the output mode of $\hat{a}_5$ or $\hat{a}_6$ respectively. $\langle\delta^2\hat{X}_{\text{in}}\rangle$, $\langle\delta^2\hat{Y}_{\text{in}}\rangle$ are the variances of the input coherent state, then we have $\langle\delta^2\hat{X}_{\text{in}}\rangle = \langle\delta^2\hat{Y}_{\text{in}}\rangle = 1$.

In the classical system without quantum correlation $r_a = 0$ and $r_b = 0$, we obtain $F = 1/2$ for the normalized gain $g_{1(2)} = 1$, so the classical limit of teleportation in this system remains the same as the usual teleportation system for continuous variables [20]. According to equations (10) and (11), the best optimal fidelity for quantum teleportation occurs around $g = 1$. In this case the fidelity becomes $F = 2/([\mathrm{e}^{-2r_a} + \mathrm{e}^{\pm 2r_b}] + 2)$, thus to meet the requirement of the quantum teleportation $F > 1/2$, it requires that $r_a > 0$ and $r_b < 0$ for Bob1, and it requires that $r_a > 0$ and $r_b > 0$ for Bob2.

For experiments the most important work is to establish two EPR beam sources with identical frequency and constant phase relation. Two identical degenerate [9] or nondegenerate optical parametric amplifiers [16, 21] (DOPA or NOPA) pumped by an identical laser can be used to produce the required two two-mode squeezed states. The mature parametric technique is beneficial to complete the proposed prototype. The correlation relation between two modes of EPR pairs can be manipulated by converting the relative phase between the pump field and injected signal field of the OPA between 0 and π. For the parametric deamplification (the pump field and the injected field are in phase of 0) the two-mode amplitude squeezing is completed which corresponds to the EPR beams with the quadrature amplitude correlation and quadrature phase anticorrelation between two modes [22]. For a polarization nondegenerate parametric amplification (the pump field and the injected field are out of phase, i.e. the relative phase is π) the two-mode phase squeezing is obtained which corresponds to the quadrature amplitude anticorrelation and quadrature phase correlation EPR beams [16, 21, 23, 24].

In conclusion, we propose a quantum switching system for sending controllably an unknown quantum state to either of two remote receivers. The control condition is only to convert the squeezed component of one of two two-mode squeezed states between its quadrature amplitude and phase. The conditional teleportation system might be developed as a practical quantum switching in future quantum communication. The well known optical parametric technique provides great convenience for its experimental demonstration.

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