Generating multiphoton Greenberger-Horne-Zeilinger states with weak cross-Kerr nonlinearity

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We propose a scheme to generate polarization-entangled multiphoton Greenberger-Horne-Zeilinger states with weak cross-Kerr nonlinearity based on controlled bus rotation and subsequent homodyne measurement. Our method is simple in operation and has high success probabilities with near perfect fidelities in an ideal case.

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Greenberger-Horne-Zeilinger (GHZ) states [1], one kind of multipartite entangled states, play an important role in nonlocality test of quantum mechanics [2] and in quantum-information processing (QIP) [3]. Much effort has been devoted to generating GHZ states in various physical systems over the past few years [4]. In particular, matter-qubit GHZ states have been produced experimentally in microwave cavity QED [5] and ion trap [6]. Compared with most massive qubits, flying photon qubits are more suitable for long-distance quantum communication and more robust against decoherence because of their fast speed and weak interaction with environment [7]. Photon GHZ states have been successfully demonstrated by the multifold coincidence-measurement method [8]. In these experiments, however, one does not know whether or not photon qubits are in the desired states unless these photons are detected and destroyed. So the output entangled states cannot be used as the input state for further experimental operations. This drawback would limit their efficient applications in large-scale QIP [9].

One possible approach for generating photon GHZ states that overcomes this drawback utilizes optical cross-Kerr effect. This type of approach involves two optical field modes, one is called signal mode and the other probe mode. A weak interaction between photons in these two modes is induced by passing them through nonlinear Kerr media. The Hamiltonian describing this interaction is given by [10] $H_{\mathrm{cK}}=-\hbar\chi a^{\dagger}ab^{\dagger}b$, where χ is proportional to the third-order nonlinear susceptibility $\chi^{(3)}$ and $a^{\dagger}$ and a ($b^{\dagger}$ and b) represent the creation and annihilation operators of the signal mode (probe mode), respectively. After passing through the Kerr medium, a photon in signal mode accumulates a phase shift $\theta=\chi t$ (t is the effective interaction time) that is proportional to the photon number in the probe mode. If the nonlinearity is large enough, such as $\theta=\pi$, then it can be used to implement quantum gates [11,12] and generate entangled states [13]. Unfortunately, natural Kerr media has extremely small nonlinearities, with a typical magnitude of $\theta\approx 10^{-18}$ rad [14,15]. However, it was suggested recently that the nonlinearities might be largely improved, say $\theta\sim 10^{-2}$, with electromagnetically induced transparency (EIT) [16,17]. In addition, the weak nonlinearity can be compensated by using a probe coherent state field $|\alpha\rangle$ with a very large amplitude. So this kind of small-but-not-tiny Kerr nonlinearities could still be effectively used for QIP [17–21].

In this Brief Report we present a scheme to generate GHZ states of multiple polarization-encoded photons using this weak-nonlinearity-based method. As shown in Fig. 1 strong probe mode interacts successively with multiple signal-mode photons, each causing a conditional phase rotation in the probe mode. Subsequent P $[P=-i(b-b^{\dagger})/\sqrt{2}]$ homodyne measurement [22] of the probe mode will project the photons in the signal mode into the desired entangled states. Such a scheme has nearly perfect fidelities and higher success probabilities than the one in Ref. [23]. In addition, our proposal has several other advantages. First, our scheme uses only one probe mode instead of multiple as in some parity-gate-based methods [24,25]. This makes our scheme simpler to implement. Second, our scheme uses P homodyne measurement, which requires considerably smaller strength of the coherent state in the probe mode than the X $[X=(b+b^{\dagger})/\sqrt{2}]$ measurement in other schemes [20]. Therefore our scheme is more feasible and more robust against decoherence [26]. Moreover, our method is not limited to the all-optical implication and can be also used in other bus-mediated QIP architectures with a cross-Kerr-like interaction Hamiltonian [27,28].

To illustrate the essential features of our scheme, it is sufficient to first consider how to produce a *three*-photon GHZ state. The schematic setup of our proposal is shown in Fig. 1. The photon in the signal mode is prepared in a superposition state of horizontal and vertical polarizations while the probe beam is initially set in a coherent state superposition (CSS) $c(|\alpha\rangle+|-\alpha\rangle)$ with a normalization factor $c=1/\sqrt{2+2e^{-2\alpha^2}}$. Without loss of generality, we have assumed the amplitude α is positive. Then the initial state of the system that consists of three photons in the signal mode and a coherent probe beam is

$$
\begin{aligned}
|\psi\rangle &= \bigotimes_{k=1}^{3}\frac{1}{\sqrt{2}}(|H\rangle_k+|V\rangle_k)\otimes c(|\alpha\rangle+|-\alpha\rangle)\\
&=\frac{c}{2\sqrt{2}}(|HHH\rangle+|HHV\rangle+|HVH\rangle+|VHH\rangle+|VVH\rangle\\
&\quad+|VHV\rangle+|HVV\rangle+|VVV\rangle)\otimes(|\alpha\rangle+|-\alpha\rangle),
\end{aligned}
\tag{1}
$$

where $|H\rangle_k$ ($|V\rangle_k$) is the horizontal (vertical) polarization state of the kth photon in the signal mode. Several schemes to produce a strong CSS with weak cross-Kerr nonlinearity have been proposed [29,30]. Very recently the generation of a traveling-wave optical CSS with a small α have been experimentally demonstrated by subtracting one photon from a squeezed vacuum state [31,32]. In principle, arbitrary large CSS can be produced out of arbitrary small CSSs by the

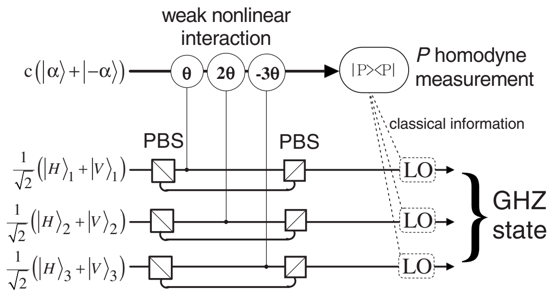

FIG. 1. Experimental scheme for generating the three-photon GHZ state. PBS represents polarization beam splitters and LO means local operation. For details, see text.

process described by Lund *et al.* [33] if quantum optical memory is available.

As shown in Fig. 1, the three photon qubits in the signal mode are split individually on the polarization beam splitters (PBSs) into two spatial modes, one of which then interacts the probe beam in the Kerr nonlinear medium. After recombined by the last PBS array, the whole system evolves into the following state:

$$
\begin{aligned}
|\psi'\rangle = \frac{c}{2\sqrt{2}}[&(|HHH\rangle + |VVV\rangle)(|\alpha\rangle + |-\alpha\rangle) + |HVV\rangle(|\alpha e^{i\theta}\rangle \\
&+ |-\alpha e^{i\theta}\rangle) + |VHH\rangle(|\alpha e^{-i\theta}\rangle + |-\alpha e^{-i\theta}\rangle) + |VHV\rangle \\
&\times(|\alpha e^{2i\theta}\rangle + |-\alpha e^{2i\theta}\rangle) + |HVH\rangle(|\alpha e^{-2i\theta}\rangle + |-\alpha e^{-2i\theta}\rangle) \\
&+ |HHV\rangle(|\alpha e^{3i\theta}\rangle + |-\alpha e^{3i\theta}\rangle) + |VVH\rangle(|\alpha e^{-3i\theta}\rangle \\
&+ |-\alpha e^{-3i\theta}\rangle)]. \qquad (2)
\end{aligned}
$$

In the next step, we implement the P-quadrature homodyne measurement on the probe beam [17]. The resulting three-photon state in the signal mode is then

$$
\begin{aligned}
|\psi''\rangle = \frac{1}{N}(&2g_0 \cos\tau_0|\mathrm{GHZ}_0\rangle + g_{1+}|\mathrm{GHZ}_{1+}\rangle + g_{1-}|\mathrm{GHZ}_{1-}\rangle \\
&+ g_{2+}|\mathrm{GHZ}_{2+}\rangle + g_{2-}|\mathrm{GHZ}_{2-}\rangle + g_{3+}|\mathrm{GHZ}_{3+}\rangle \\
&+ g_{3-}|\mathrm{GHZ}_{3-}\rangle), \qquad (3)
\end{aligned}
$$

where

$$
\begin{aligned}
|\mathrm{GHZ}_0\rangle &= (1/\sqrt{2})(|HHH\rangle + |VVV\rangle), \\
|\mathrm{GHZ}_{1\pm}\rangle &= (1/\sqrt{2})(e^{\mp i\tau_{1\pm}}|HVV\rangle + e^{\pm i\tau_{1\pm}}|VHH\rangle), \\
|\mathrm{GHZ}_{2\pm}\rangle &= (1/\sqrt{2})(e^{\mp i\tau_{2\pm}}|VHV\rangle + e^{\pm i\tau_{2\pm}}|HVH\rangle), \\
|\mathrm{GHZ}_{3\pm}\rangle &= (1/\sqrt{2})(e^{\mp i\tau_{3\pm}}|HHV\rangle + e^{\pm i\tau_{3\pm}}|VVH\rangle), \qquad (4)
\end{aligned}
$$

and the coefficients are

$$
g_0 = \pi^{-1/4}e^{-1/2p^2}, \quad \tau_0 = \sqrt{2}\alpha p,
$$

$$
g_{n\pm} = \pi^{-1/4}e^{-(p \mp \sqrt{2}\alpha \sin n\theta)^2/2}\ (n = 1,2,3),
$$

$$
\tau_{n\pm} = \sqrt{2}\alpha\cos(n\theta)\big[p \mp (\sqrt{2}/2)\sin(n\theta)\big],
$$

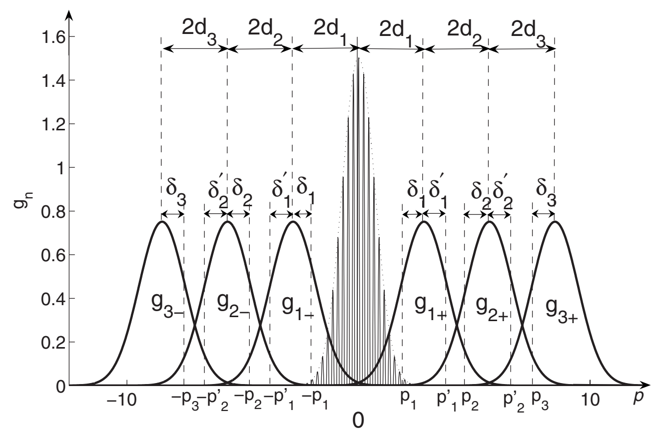

FIG. 2. Peak functions $2g_0|\cos\tau_0|$ and $g_{n\pm}$ $(n=1,2,3)$ of the homodyne measurement result p with $\theta=0.01$ and $\alpha=200$.

$$
\begin{aligned}
|N|^2 = {}&4g_0^2\cos^2\tau_0 + g_{1+}^2 + g_{1-}^2 + g_{2+}^2 + g_{2-}^2 + g_{3+}^2 + g_{3-}^2 \\
&+ 2g_{1+}g_{1-}\cos(\tau_{1+}+\tau_{1-}) + 2g_{2+}g_{2-}\cos(\tau_{2+}+\tau_{2-}) \\
&+ 2g_{3+}g_{3-}\cos(\tau_{3+}+\tau_{3-}). \qquad (5)
\end{aligned}
$$

In Fig. 2, we plot $2g_0|\cos\tau_0|$ and $g_{n\pm}$ $(n=1,2,3)$ as a function of the homodyne measurement result p. We observe that $2g_0|\cos\tau_0|$ oscillates fast and $g_{n\pm}$ are six Gaussian curves with the peaks located at $\sqrt{2}\alpha\sin(n\theta)$, respectively. The distances between the neighboring peaks are $d_n=\sqrt{2}\alpha\{\sin(n\theta) - \sin[(n-1)\theta]\}$, which are usually referred to as the *distinguishabilities* of the measurement [28]. Because the distinguishabilities are approximately in proportion to $\alpha\theta$ when $\theta \ll 1$, we can use a strong coherent state probe light to ensure good distinguishability, that is, the overlaps between the peaks are very small. The following discussion will be made on the assumption that good distinguishability is already achieved.

Let us consider the case where the result of the P homodyne measurement is near the center of one of the six side peaks. Suppose, for example, the result is $p_2=\sqrt{2}\sin(2\theta) - \delta_2$, which is near the peak of g_{2+} as seen from Fig. 2. In this case, the polarization photon state $|\psi''\rangle$ in Eq. (3) is very close to the entangled state $|\mathrm{GHZ}_{2+}\rangle$. The fidelity of the resulting state with respect to $|\mathrm{GHZ}_{2+}\rangle$ is

$$
F_{|\mathrm{GHZ}_{2+}\rangle}(p_2) = |\langle \mathrm{GHZ}_{2+}|\psi''\rangle|^2 \approx \frac{1}{1+e^{-4d_2(d_2-\delta_2)}}. \qquad (6)
$$

For given α and θ, we can decide δ_2 from the needed high fidelity. If $\theta=0.01$ and $\alpha=400$, $F_{|\mathrm{GHZ}_{2+}\rangle}(p_2)$ will be 0.9999 with $\delta_2 \approx 0.712d_2$. Similarly, if the measurement result is p_2', as shown in Fig. 2, the resulting fidelity with respect to $|\mathrm{GHZ}_{2+}\rangle$ will be

$$
F_{|\mathrm{GHZ}_{2+}\rangle}(p_2') \approx \frac{1}{1+e^{-4d_3(d_3-\delta_2')}}. \qquad (7)
$$

Assume the minimal acceptable fidelity of the resulting state is $F_{\min}$. By appropriately choosing δ_2 and δ_2', we can have $F_{|\mathrm{GHZ}_{2+}\rangle}(p_2)=F_{|\mathrm{GHZ}_{2+}\rangle}(p_2')=F_{\min}$. Then, as long as the measurement result p is in the regime of $p_2<p<p_2'$, we can get

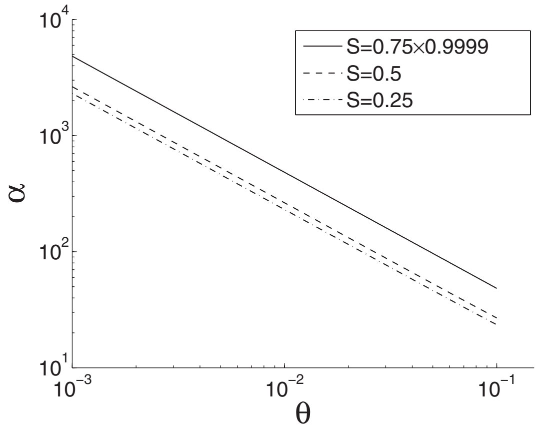

FIG. 3. Logarithmic plots of the α as a function of the θ for the minimal fidelity $F_{\min}=0.9999$ and several total success probabilities S of obtaining the three-photon GHZ state.

the polarization photon state $|\mathrm{GHZ}_{2+}\rangle$ with the fidelity higher than $F_{\min}$. The probability of such an event is

$$P_2 = \int_{p_2}^{p_2'} f(p)dp \approx \frac{1}{16}[\mathrm{erf}(\delta_2) + \mathrm{erf}(\delta_2')], \tag{8}$$

where $\mathrm{erf}(x)$ is the error function and $f(p)$ is the probability density distribution of the projection with the form

$$f(p) = \langle p|\rho_{\mathrm{probe}}|p\rangle \approx \tfrac{1}{8}(4g_0^2\cos^2\tau_0 + g_{1+}^2 + g_{1-}^2 + g_{2+}^2 + g_{2-}^2 + g_{3+}^2 + g_{3-}^2). \tag{9}$$

In the above equation, $\rho_{\mathrm{probe}}=\mathrm{Tr}_{\mathrm{signal}}[|\psi'\rangle\langle\psi'|]$ is the reduced density operator for the probe field. Finally, according to the measurement result p, simple local rotations using phase shifters can be performed via a feed-forward process to remove the relative phase τ_{2+}, that is, to transform $|\mathrm{GHZ}_{2+}\rangle$ into the standard GHZ state $|\mathrm{GHZ}_0\rangle$, which is independent of p.

Similar analysis can be made when the measurement result p is near the centers of the other five side peaks in Fig. 2. Because of the rapid oscillation near $p=0$, we simply discard the measurement results in the regime between $-p_1$ and p_1. Therefore the total success probability S of obtaining $|\mathrm{GHZ}_0\rangle$ with the fidelity higher than $F_{\min}$ is

$$S \approx \tfrac{1}{8}[\mathrm{erf}(\delta_1) + \mathrm{erf}(\delta_1') + \mathrm{erf}(\delta_2) + \mathrm{erf}(\delta_2') + \mathrm{erf}(\delta_3) + 1], \tag{10}$$

where δ_i and δ_i' are marked out in Fig. 2. As the distinguishabilities increase, this probability will approach 3/4, which is much higher than 1/4, the highest success probability achieved in Ref. [23].

Figure 3 shows the amplitude α of the probe mode as a function of the nonlinearity θ for the minimal fidelity $F_{\min}=0.9999$ and several different total success probabilities S of obtaining GHZ states. Approximately, α decreases with the increase of θ logarithmically as can be seen. This is because the distinguishabilities $d_n \approx \alpha\theta$ as $\theta \ll 1$. When $\theta=0.01$, α is about 483 with $S\approx 0.75$. In Ref. [20], X-quadrature homodyne detection is performed instead of P detection as in our scheme. Although the success probability of obtaining the GHZ state can be near 1, the distinguishabilities in their scheme are approximately proportional to $\alpha\theta^2$. Therefore for the same $\theta=0.01$, α would be the order of 10^5, which is much larger than our result and puts a constraint on the practical implementation. This shows that our method can be used more efficiently in the regime of weak cross-Kerr nonlinearity.

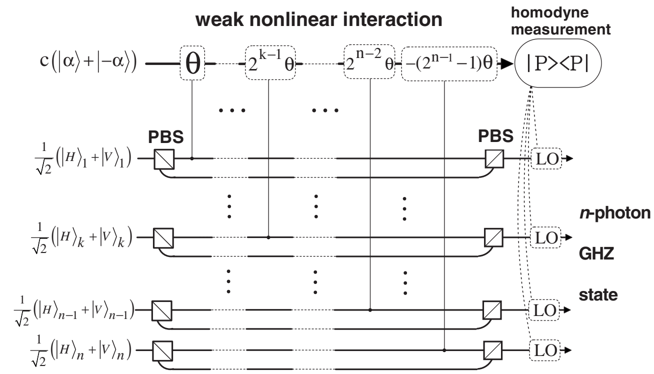

FIG. 4. Experimental scheme for generating the n-photon GHZ states.

The above scheme can be easily extended to generate n-photon GHZ states. The general initial state is

$$\frac{1}{\sqrt{2^n}}(|H\rangle_1 + |V\rangle_1)\cdots(|H\rangle_n + |V\rangle_n)c(|\alpha\rangle + |-\alpha\rangle). \tag{11}$$

After n successive cross-Kerr nonlinearity interactions, as shown in Fig. 4, the whole system becomes

$$\frac{1}{\sqrt{2^n}}[(|HH\cdots H\rangle + |VV\cdots V\rangle)(|\alpha\rangle + |-\alpha\rangle) + \cdots + |H\cdots HV\rangle \times (|\alpha e^{2^{n-1}i\theta}\rangle + |-\alpha e^{2^{n-1}i\theta}\rangle) + |V\cdots VH\rangle(|\alpha e^{-2^{n-1}i\theta}\rangle + |-\alpha e^{-2^{n-1}i\theta}\rangle)]. \tag{12}$$

Then we perform P quadrature measurement on the probe beam and convert the photons to standard GHZ states by a classical feed-forward control conditioned on the P measurement results. By choosing proper α, the success probability of $1-(1/2^{n-1})$, which approaches 1 as n increases, can be achieved with near perfect fidelity. In Ref. [23], the success probability is only $1/2^{n-1}$.

Our scheme generally depends on the precision of the homodyne measurement. If the results are not accurate enough, the relative phases $\tau_{n\pm}$ could not be exactly eliminated to get the standard GHZ state. So the uncertainty in the homodyne measurement must be very small, and this can be experimentally achieved by using a much more intense local oscillator than the probe mode. Another important point with regard to the practical implementation of our scheme is that the decoherence effects during the entanglement generation process in nonlinear media must be low [34]. Also, it is a tremendous experimental challenge to generate a strong probe field in a CSS [35,36].

In conclusion, we have presented an effective scheme to generate multiphoton GHZ states with high success probabilities and near perfect fidelities. This scheme takes advantage of the nonlinear Kerr effect. It uses only the basic tools in quantum optical laboratories and can be implemented in the regime of the weak cross-Kerr nonlinearity. We expect it be realized in the near future with the fast development of the relevant techniques.

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