Multiphoton Interference as a Tool to Observe Families of Multiphoton Entangled States

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***Abstract*—Spontaneous parametric downconversion in combination with linear optics was successfully used to observe a variety of multiphoton entangled states. Yet, experiments performed so far lacked flexibility, as each of the various setups was useful for only a particular multiphoton entangled state. In this paper, we describe how, by using multiphoton interference, one can observe entire families of multiphoton entangled states in the very same linear optical setup. Our method thus goes beyond the commonly used two-photon interference and turns out to be a very useful tool for state observation. We will discuss the interference of four and six photons at different types of beam splitters and show which families of entangled states are observable. The benefits of this approach are demonstrated in a four-photon interference experiment by observing a variety of highly entangled multiphoton states.**

***Index Terms*—Frequency conversion, interference, nonlinear optics, parametric devices, ultrafast optics.**

I. Introduction

MULTIPARTITE entanglement is an important nonclassical resource for applications of quantum information. In principle, many physical systems are well suited for experimental realizations of multipartite entangled quantum states. So far, photonic qubits allowed observations of the biggest variety of multipartite entangled states. In order to describe and categorize all these quantum states, among others, the criterion of equivalence under stochastic local operations and classical communication (SLOCC) was introduced [1]– [4]. It has already been shown [2] that for four qubits infinitely many SLOCC-inequivalent four-qubit entangled pure states exist. This classification is quite useful for multiparty quantum communication applications, where each SLOCC-inequivalent state has the potential to lead to a particular nonclassical application. Hence, a flexible method to observe—and finally to apply—many SLOCC-inequivalent states is surely desirable.

To observe multiphoton entangled states, usually, a combination of spontaneous parametric downconversion (SPDC) with linear optical elements is used. To this end, indistinguishability of photons originating from different SPDC sources or emissions is required in order to achieve multiphoton interference [5], [6] enabling the observation of entangled states. However, most earlier experiments relied on enforced indistigushability of just two photons [7]–[9]. In this paper, we will demonstrate how four- and six-photon interference is of additional benefit as it allows one to observe whole families of entangled quantum states in a single setup. This breaks with the common approach to design a particular linear optical experiment for each quantum state. Besides possible applications in quantum communication, multiphoton interference was also proposed to be a useful tool to entangle distant atoms [10]–[16], or to improve precision measurements [17]– [21]. Previously, it was studied with respect to photon bunching and multipath interference at a beam splitter (BS) [22]–[25] as a generalization of the Hong–Ou–Mandel effect [26].

We will discuss, in Section II, the interference of four and six photons on different types of BSs and analyze the potential of these cases with respect to the observation of SLOCC-inequivalent entangled states. In Section III, we describe a particular experimental implementation using four-photon interference at a polarizing BS, which was recently performed by us [27], with special emphasis placed on the analysis of the entanglement of the various states. Finally, in Section IV, we summarize our main findings.

II. Multiphoton Entanglement via Multiphoton Interference

In the following, we will discuss how multiphoton interference can be used to observe various multipartite entangled states. We will study four- and six-photon interference at different kinds

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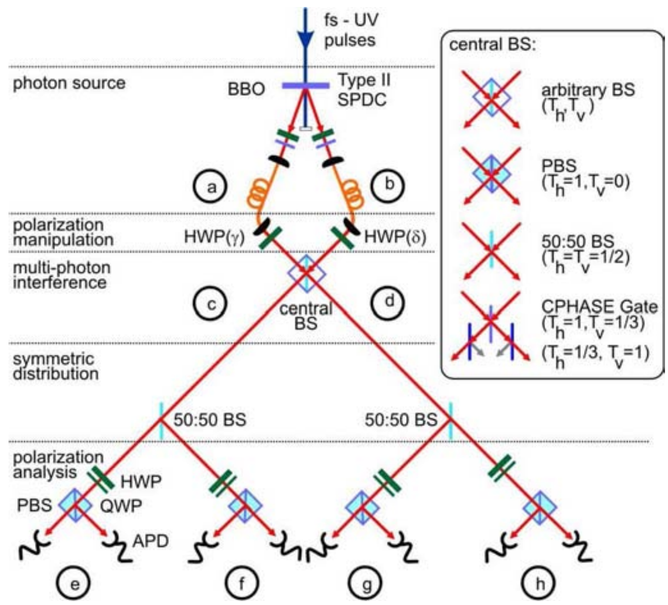

Fig. 1. Schematic experimental setup for using four-photon interference at a central BS to observe families of entangled four-photon states. BSs having different transmittances for horizontal (T_h) and vertical (T_v) polarization are used as the central BS (shown on the right). The polarization analysis of the photon states in modes e, f, g, h is performed with an HWP and a QWP in front of a PBS.

of BSs. This approach will be shown to be superior to conventional state observation schemes relying also on linear optics, most of which served to observe a single state only. With our method, it is possible to observe a multitude of different and relevant multipartite entangled states in a single linear optical setup.

The general experimental scheme to achieve this is as follows (Fig. 1). We start with a photon source that delivers $2n$ photons to two spatially distinct modes (labeled a and b), such that n photons occupy each spatial mode. Such a source is given by a type-II noncollinear SPDC that generates in its nth order emission $2n$ photons. Next, we will change the polarization state of the photons via a half-wave plate (HWP) in mode a (we will, additionally, also consider an HWP in mode b). A similar approach was pursued in a recent experiment, where an HWP at the beginning of the optical setup was used to continuously vary between three photon states out of the same entanglement class [28]. Subsequently, the photons interfere at a BS with a certain transmittance T_h (T_v) for horizontally (vertically) polarized photons. Its output modes (labeled c and d) are split by polarization independent BSs into $2n$ modes that have an equal output probability. Under the condition of having a single photon in each of the $2n$ modes, we observe the desired states. In the following, we will discuss the cases $n = 2$ and $n = 3$, and thus, the interference of four and six photons, respectively.

A. Four-Photon Interference

Let us start with the interference of four photons at a BS. To this end, we consider four photons of the second-order SPDC emission, which are in the state [29] (to assure indistinguishability of photons coming from different SPDC pairs, one must use filters of spectral width narrower than that of the pulsed pump [5], [6])

$$\begin{aligned}&\propto ({a_h}^\dagger {b_v}^\dagger + {a_v}^\dagger {b_h}^\dagger)^2\,|\mathrm{vac}\rangle\\ &= [({a_h}^\dagger {b_v}^\dagger)^2 + ({a_v}^\dagger {b_h}^\dagger)^2 + 2{a_h}^\dagger {a_v}^\dagger {b_h}^\dagger {b_v}^\dagger]\,|\mathrm{vac}\rangle \qquad (1)\end{aligned}$$

where ${m_i}^\dagger$ denotes the creation operator of a photon in mode m having polarization i and $|\mathrm{vac}\rangle$ is the vacuum state. Here and in the following, we neglect all higher order emissions, and thus, implicitly assume low conversion efficiency, e.g., due to a weak pump beam. The influence of high conversion efficiency on the state quality, in particular, for low detection efficiency, is known [30], [31] and strongly depends on the particular parameters, which will be subjected to further investigation. The HWP in mode a transforms the polarization state of the photons according to ${a_h}^\dagger \to \cos(2\gamma){a_h}^\dagger + \sin(2\gamma){a_v}^\dagger$ and ${a_v}^\dagger \to \sin(2\gamma){a_h}^\dagger - \cos(2\gamma){a_v}^\dagger$, where γ is the orientation of the optical axis with respect to the polarization of the impinging photons. Subsequently, the photons interfere on a BS with the transmittances T_h and T_v, where we assume a lossless BS, i.e., $T_h + R_h = 1$ and $T_v + R_v = 1$ hold, with R_i being reflectance of the BS. The BS transforms the photon state from input mode a to the superposition ${a_i}^\dagger \to \sqrt{T_i}{c_i}^\dagger + \mathrm{i}\sqrt{1-T_i}{d_i}^\dagger$ and from input mode b to ${b_i}^\dagger \to \sqrt{T_i}{d_i}^\dagger + \mathrm{i}\sqrt{1-T_i}{c_i}^\dagger$, where c and d are the output modes of the BS and $\mathrm{i} = \sqrt{-1}$.

1) Arbitrary BS: We first use a central BS with T_h and T_v arbitrary, before we focus on three particular parameter sets. Splitting its two output modes into four final modes by two polarization independent BSs ($T_h = T_v = 1/2$) yields the states in modes e, f, g, h (up to normalization, for the notation of states, see Table I)

$$\begin{aligned}&a_{bs}\,|GHZ_4'\rangle + b_{bs}\,|\psi^+\rangle \otimes |\psi^+\rangle\\ &\quad + c_{bs}\,|HHHH\rangle + d_{bs}\,|VVVV\rangle\\ &\quad - e_{bs}(|VHHH\rangle + |HVHH\rangle\\ &\quad - |HHVH\rangle - |HHHV\rangle)\\ &\quad - f_{bs}(|VVVH\rangle + |VVHV\rangle\\ &\quad - |VHVV\rangle - |HVVV\rangle) \qquad (2)\end{aligned}$$

where each amplitude depends on the three parameters γ, T_h, and T_v in the following way:

$$\begin{aligned}a_{bs} &= \frac{1}{\sqrt{2}}(T_v - R_h - 2T_hT_v\\ &\quad + (T_v - R_h - 2T_hT_v - 4\sqrt{R_hT_hR_vT_v})\cos 4\gamma)\\ b_{bs} &= 2\Bigg(\sqrt{R_hT_hR_vT_v} + \left(\frac{1}{2} - T_v + \sqrt{R_hT_hR_vT_v}\right.\\ &\qquad \left. + T_h(-1 + 2T_v)\right)\cos 4\gamma\Bigg)\\ c_{bs} &= (6R_hT_h - 1)(\sin 2\gamma)^2\\ d_{bs} &= (6R_vT_v - 1)(\sin 2\gamma)^2\end{aligned}$$

TABLE I
VARIOUS MULTIPARTITE ENTANGLED STATES THAT ARE CONTAINED IN DIFFERENT FAMILIES OF STATES, WHERE EACH FAMILY CAN BE OBSERVED WITH A SINGLE SETUP

$\lvert W_3\rangle$	$=\frac{1}{\sqrt{3}}(\lvert HHV\rangle+\lvert HVH\rangle+\lvert VHH\rangle)$
$\lvert\overline{W}_3\rangle$	$=\frac{1}{\sqrt{3}}(\lvert VVH\rangle+\lvert VHV\rangle+\lvert HVV\rangle)$
$\lvert GHZ_4\rangle$	$=\frac{1}{\sqrt{2}}(\lvert HHHH\rangle+\lvert VVVV\rangle)$
$\lvert GHZ'_4\rangle$	$=\frac{1}{\sqrt{2}}(\lvert HHVV\rangle+\lvert VVHH\rangle)$
$\lvert W_4\rangle$	$=\frac{1}{2}(\lvert HHHV\rangle+\lvert HHVH\rangle+\lvert HVHH\rangle+\lvert VHHH\rangle)$
$\lvert\overline{W}_4\rangle$	$=\frac{1}{2}(\lvert VVVH\rangle+\lvert VVHV\rangle+\lvert VHVV\rangle+\lvert HVVV\rangle)$
$\lvert D_4^{(2)}\rangle$	$=\frac{1}{\sqrt{6}}(\lvert HHVV\rangle+\lvert HVHV\rangle+\lvert VHHV\rangle+\lvert HVVH\rangle+\lvert VHVH\rangle+\lvert VVHH\rangle)$
$\lvert D_4^{(2)'}\rangle$	$=\frac{1}{\sqrt{6}}(\lvert HHHH\rangle+\lvert HVHV\rangle+\lvert VHHV\rangle+\lvert HVVH\rangle+\lvert VHVH\rangle+\lvert VVVV\rangle)$
$\lvert\psi^+\rangle\otimes\lvert\psi^+\rangle$	$=\frac{1}{2}(\lvert HVHV\rangle+\lvert VHVH\rangle+\lvert HVVH\rangle+\lvert VHHV\rangle)$
$\lvert\psi^-\rangle\otimes\lvert\psi^-\rangle$	$=\frac{1}{2}(\lvert HVHV\rangle+\lvert VHVH\rangle-\lvert HVVH\rangle-\lvert VHHV\rangle)$
$\lvert\phi^+\rangle\otimes\lvert\phi^+\rangle$	$=\frac{1}{2}(\lvert HHHH\rangle+\lvert HHVV\rangle+\lvert VVHH\rangle+\lvert VVVV\rangle)$
$\lvert\phi^-\rangle\otimes\lvert\phi^-\rangle$	$=\frac{1}{2}(\lvert HHHH\rangle-\lvert HHVV\rangle-\lvert VVHH\rangle+\lvert VVVV\rangle)$
$\lvert\Psi_4^-\rangle$	$=\sqrt{2/3}\lvert GHZ'_4\rangle-\sqrt{1/3}\lvert\psi^+\rangle\otimes\lvert\psi^+\rangle$
$\lvert\Psi_4^{-'}\rangle$	$=\sqrt{2/3}\lvert GHZ_4\rangle-\sqrt{1/3}\lvert\psi^+\rangle\otimes\lvert\psi^+\rangle$
$\lvert\Psi_4^+\rangle$	$=\sqrt{2/3}\lvert GHZ_4\rangle+\sqrt{1/3}\lvert\psi^+\rangle\otimes\lvert\psi^+\rangle$
$\lvert GHZ_6\rangle$	$=\frac{1}{\sqrt{2}}(\lvert HHHHHH\rangle+\lvert VVVVVV\rangle)$
$\lvert GHZ_6^-\rangle$	$=\frac{1}{\sqrt{2}}(\lvert HHHHHH\rangle-\lvert VVVVVV\rangle)$
$\lvert\Psi_6^+\rangle$	$=\frac{1}{\sqrt{2}}\lvert GHZ_6^-\rangle+\frac{1}{2}(\lvert\overline{W}_3\rangle\otimes\lvert\overline{W}_3\rangle-\lvert W_3\rangle\otimes\lvert W_3\rangle)$

We use the notation for polarization encoded qubits in different spatial modes, where $H(V)$ stands for horizontal (vertical) polarization and encodes a logical 0(1). The notation of, e.g., $\lvert HHHH\rangle$ is an abbreviated form of $\lvert HHHH\rangle=\lvert H\rangle_m\otimes\lvert H\rangle_n\otimes\lvert H\rangle_o\otimes\lvert H\rangle_p$, where the subscripts denote the spatial mode of each photon.

$$e_{bs}=\frac{1}{2}(\sqrt{R_hR_v}-3T_h\sqrt{R_hR_v}+2\sqrt{T_hT_v}-3T_h^{3/2}\sqrt{T_v})\sin 4\gamma$$
$$f_{bs}=\frac{1}{2}(\sqrt{R_hR_v}(1-3T_v)+\sqrt{T_hT_v}(2-3T_v))\sin 4\gamma. \quad (3)$$

These states appear in several entanglement families of the four-qubit SLOCC classification introduced recently [2], [4]. To obtain a clearer insight into these states, we will discuss the following three particular BSs in more detail.

2) Polarizing Beam Splitter: By using a polarizing beam splitter (PBS) with $T_h=1$ and $T_v=0$, the family of states

$$\lvert\Psi_4(\gamma)\rangle=a_4(\gamma)\lvert GHZ_4\rangle+b_4(\gamma)\lvert\psi^+\rangle\otimes\lvert\psi^+\rangle \quad (4)$$

with

$$a_4(\gamma)=\frac{\sqrt{2}(1-\cos 4\gamma)}{\sqrt{5-4\cos(4\gamma)+3\cos 8\gamma}}$$
$$b_4(\gamma)=\frac{(2\cos 4\gamma)}{\sqrt{5-4\cos(4\gamma)+3\cos 8\gamma}} \quad (5)$$

and $a_4(\gamma)^2+b_4(\gamma)^2=1$ is obtained [27] [Fig. 2(a)]. The states $\lvert\Psi_4(\gamma)\rangle$ form a superposition of the well-known $\lvert GHZ_4\rangle$ state, a highly entangled four-qubit state, and a product of two Bell states, a biseparable state.By using the SLOCC classification of [2], we can attribute the family $\lvert\Psi_4(\gamma)\rangle$ to the generic entanglement class G_{abcd} of four-qubit entangled states. States of this class form a continuous set of SLOCC-inequivalent states, i.e., for each particular value of $\gamma\in[0,\pi/8]$, we obtain an SLOCC-inequivalent state [27] with a probability $p_4(\gamma)=(5-4\cos(4\gamma)+3\cos(8\gamma))/48$ [Fig. 2(a)]. Recently, we accomplished the experimental realization of $\lvert\Psi_4(\gamma)\rangle$ [27], which will be discussed in Section III in more detail.

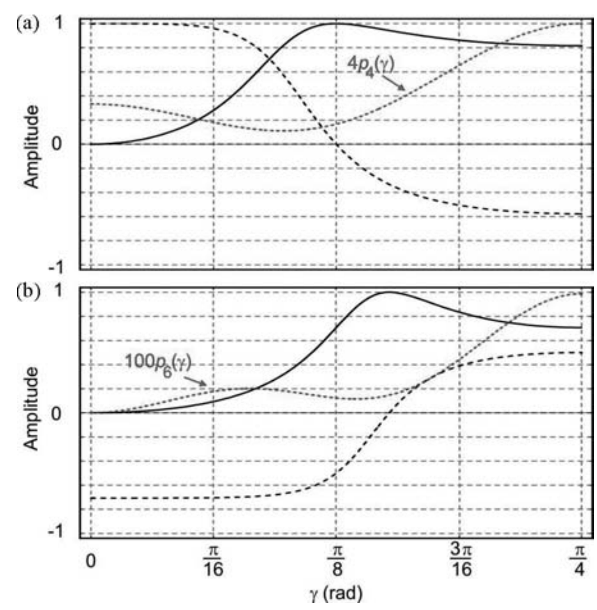

Fig. 2. Amplitudes. (a) $a_4(\gamma)$ (solid) and $b_4(\gamma)$ (dashed) for the family of states $\lvert\Psi_4(\gamma)\rangle$. (b) $a_6(\gamma)$ (solid) and $b_6(\gamma)$ (dashed) for $\lvert\Psi_6(\gamma)\rangle$. Further, (a) $p_4(\gamma)$ and (b) $p_6(\gamma)$ denote the probability of each state to be observed in the corresponding linear optical setup (dotted).

At this point, let us highlight the well-known states of the family

$$\begin{aligned} \gamma = 0 &\rightarrow |\psi^+\rangle \otimes |\psi^+\rangle \\ \gamma = \frac{\pi}{12} &\rightarrow |D_4^{(2)\prime}\rangle \\ \gamma = \frac{1}{2}\arctan\frac{1}{\sqrt{2}} &\rightarrow |\Psi_4^+\rangle \\ \gamma = \frac{\pi}{8} &\rightarrow |GHZ_4\rangle \\ \gamma = \frac{\pi}{4} &\rightarrow |\Psi_4^{-\prime}\rangle. \end{aligned} \tag{6}$$

The states $|\psi^+\rangle \otimes |\psi^+\rangle$ and $|\Psi_4^+\rangle, |\Psi_4^{-\prime}\rangle$ are local unitary (LU) equivalent to $|\psi^-\rangle \otimes |\psi^-\rangle$ and $|\Psi_4^-\rangle$, respectively, which are the two basis states for decoherence-free communication of a qubit [32]. The state $|D_4^{(2)\prime}\rangle$ is LU equivalent to $|D_4^{(2)}\rangle$, which belongs to the family of Dicke states [33]. A remarkable property of $|D_4^{(2)}\rangle$ is that it allows to obtain by a single projective measurement states out of the two inequivalent three-qubit SLOCC entanglement classes [1], [34]. The state $|GHZ_4\rangle$ is a graph state [35] and can be used for numerous applications, e.g., for multiparty quantum secret sharing [36], dense coding [37], and simulating anyonic statistics [38]. While all of these states have been previously realized in dedicated linear optical setups [7], [34], [39]–[41], now, it is possible to observe all of them in a single setup only.

Finally, the usage of an additional HWP(δ) in mode b in front of the PBS adds another tuning parameter δ. However, it turns out that the angle dependence changes simply into $\gamma \rightarrow \gamma + \delta$, resulting in the same states as before.

3) 50:50 BS: Another commonly used BS is given by $T_h = T_v = 1/2$. There, we obtain the states (up to normalization)

$$\begin{aligned} &4\sqrt{2}(\cos\gamma)^2(\sin\gamma)^2\,|GHZ_4\rangle \\ &\quad - (1 + 3\cos 4\gamma)/\sqrt{2}\,|GHZ_4'\rangle \\ &\quad + 2(\cos 2\gamma)^2\,|\psi^+\rangle \otimes |\psi^+\rangle. \end{aligned} \tag{7}$$

Let us mention particularly interesting states of this family

$$\begin{aligned} \gamma = 0 &\rightarrow |\Psi_4^-\rangle \\ \gamma = \frac{\pi}{8} &\rightarrow \frac{1}{\sqrt{2}}(|\phi^-\rangle \otimes |\phi^-\rangle + |\psi^+\rangle \otimes |\psi^+\rangle) \\ \gamma = \frac{1}{4}\arccos\left(-\frac{1}{3}\right) &\rightarrow |\Psi_4^+\rangle \\ \gamma = \frac{\pi}{6} &\rightarrow \sqrt{\frac{3}{4}}\,|GHZ_4\rangle + \sqrt{\frac{1}{4}}\,|D_4^{(2)}\rangle \\ \gamma = \frac{\pi}{4} &\rightarrow |\phi^+\rangle \otimes |\phi^+\rangle. \end{aligned} \tag{8}$$

The states given for $\gamma = 0$, $\gamma = (1/4)\arccos(-1/3)$, and $\gamma = \pi/4$ are LU equivalent to states of the family $|\Psi_4(\gamma)\rangle$. However, for other values of γ, we find different states, e.g., the state $1/\sqrt{2}(|\phi^-\rangle \otimes |\phi^-\rangle + |\psi^+\rangle \otimes |\psi^+\rangle)$ $(\gamma = \pi/8)$ is a superposition of two biseparable states and $\sqrt{3/4}\,|GHZ_4\rangle + \sqrt{1/4}\,|D_4^{(2)}\rangle$ $(\gamma = \pi/6)$ is a superposition of all distinct permutations of an even number of vertically polarized photons.

When we additionally use an HWP(δ) in mode b in front of the 50:50 BS, the states

$$\begin{aligned} &\propto \sqrt{2}(\sin 2(\gamma+\delta))^2(|GHZ_4\rangle - |GHZ_4'\rangle) \\ &\quad + 2(\cos 2(\gamma+\delta))^2\,|\psi^+\rangle \otimes |\psi^+\rangle \\ &\quad + (\sin 4(\gamma+\delta))(|\overline{W}_4\rangle - |W_4\rangle) \end{aligned} \tag{9}$$

are obtained. New states can be observed compared to using only one HWP. For example, terms with an odd number of vertically polarized photons $(|\overline{W}_4\rangle - |W_4\rangle)$ also appear now.

4) CPHASE BS: Another well-known BS is given by $T_h = 1$ and $T_v = 1/3$. It was used in combination with two attenuation BSs of reversed splitting ratio ($T_h = 1/3$ and $T_v = 1$) to construct an all-optical controlled phase gate (CPHASE) [42]–[44]. When the CPHASE is used as the central overlap BS, one obtains the states

$$\begin{aligned} &a_{cp}\,|GHZ_4'\rangle + b_{cp}\,|\psi^+\rangle \otimes |\psi^+\rangle \\ &\quad + c_{cp}(-|HHHH\rangle + 3\,|VVVV\rangle) \\ &\quad + d_{cp}(|VHVV\rangle + |HVVV\rangle \\ &\qquad - |VVVH\rangle - |VVHV\rangle) \\ &\quad + \frac{d_{cp}}{3}(|VHHH\rangle + |HVHH\rangle \\ &\qquad - |HHVH\rangle - |HHHV\rangle) \end{aligned} \tag{10}$$

with $a_{cp} = -\sqrt{2}(\cos 2\gamma)^2$, $b_{cp} = -(\cos 4\gamma)$, $c_{cp} = (\sin 2\gamma)^2$, and $d_{cp} = 3/2\sin 4\gamma$. These states have a similar complexity to the states observed with an arbitrary BS.

An additional HWP(δ) in mode b leads to the states

$$\begin{aligned} &a_{cp2}\,|GHZ_4'\rangle + b_{cp2}\,|\psi^+\rangle \otimes |\psi^+\rangle \\ &\quad + c_{cp2}(-|HHHH\rangle + 3\,|VVVV\rangle) \\ &\quad + d_{cp2}(|W_4\rangle + 3\,|\overline{W}_4\rangle) \end{aligned} \tag{11}$$

with $a_{cp2} = -\sqrt{2}(\cos 2(\gamma+\delta))^2$, $b_{cp2} = \cos 4(\gamma+\delta)$, $c_{cp2} = (\sin 2(\gamma+\delta))^2$, and $d_{cp2} = \sin 4(\gamma+\delta)$.

We can also directly use the four output modes of the attenuation BSs instead of distributing two of the four output modes via two 50:50 BSs. Then, we obtain

$$\begin{aligned} &\cos 4\gamma\,|VHVH\rangle + 2(\sin 2\gamma)^2\,|HHHH\rangle \\ &\quad + \sin 4\gamma(-|VHHH\rangle + |HHVH\rangle). \end{aligned} \tag{12}$$

Note that qubits in modes f and h can be factored from the aforementioned state

$$\begin{aligned} &|HH\rangle \otimes (\cos 4\gamma\,|VV\rangle + 2(\sin 2\gamma)^2\,|HH\rangle \\ &\quad + \sqrt{2}\sin 4\gamma\,|\psi^-\rangle) \end{aligned} \tag{13}$$

(we changed the qubit order into f, h, e, g). This comes from the fact that the attenuation BSs reflect only H polarized photons.

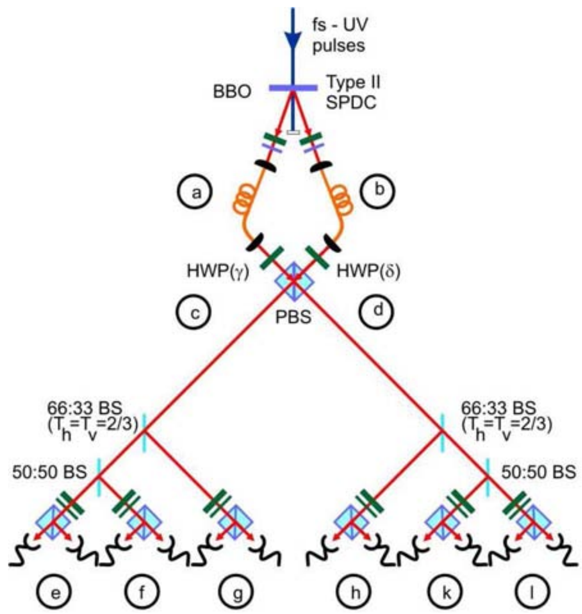

Fig. 3. Schematic experimental setup for using six-photon interference at a PBS to observe a family of entangled six-photon states in modes e, f, g, h, k, l.

B. Six-Photon Interference

Extending the level of interference further, let us consider six-photon interference. In Section II-A, the most successful approach to observe a family of states was to use interference at a *polarizing* BS [27]. Therefore, we will examine in the following the interference of the third-order SPDC with the photon state $(a_h{}^\dagger b_v{}^\dagger + a_v{}^\dagger b_h{}^\dagger)^3\,|\,\mathrm{vac}\rangle$ at a PBS too. Again, before the photons interfere, their polarization state is changed with an HWP(γ) in mode a (additionally, we also consider an HWP(δ) in mode b). The output modes of the PBS are split into six spatial modes by four polarization independent BSs (see Fig. 3).

Using the experimental layout described before, one obtains the family of states

$$|\,\Psi_6(\gamma)\rangle = a_6(\gamma)\,|\,GHZ_6^-\rangle + b_6(\gamma)(\,|\,\overline{W}_3\rangle \otimes |\,\overline{W}_3\rangle - |\,W_3\rangle \otimes |\,W_3\rangle) \quad (14)$$

where

$$a_6(\gamma) = \frac{2(\sin 2\gamma)^2}{\sqrt{7 + 4\cos 4\gamma + 5\cos 8\gamma}}$$
$$b_6(\gamma) = -\frac{1 + 3\cos 4\gamma}{\sqrt{2}\sqrt{7 + 4\cos 4\gamma + 5\cos 8\gamma}} \quad (15)$$

and $a(\gamma)^2 + 2b(\gamma)^2 = 1$ [see Fig. 2(b)]. We observe these states with a probability of $p_6(\gamma) = (2/81)(\cos\gamma\sin\gamma)^2(7 + 4\cos 4\gamma + 5\cos 8\gamma)$ [Fig. 2(b)]. Let us highlight two states that are well known to be useful for quantum information

$$\gamma = \arccos\sqrt{\frac{(3+\sqrt{3})}{6}} \rightarrow |\,GHZ_6^-\rangle$$
$$\gamma = \frac{\pi}{4} \rightarrow |\,\Psi_6^+\rangle. \quad (16)$$

The state $|\,GHZ_6^-\rangle$ is a graph state and could be already observed experimentally in a dedicated linear optical setup [9]. With the described method, it is not only possible to observe this state but also the entire family $|\,\Psi_6(\gamma)\rangle$. For example, the state $|\,\Psi_6^+\rangle$ can be used for telecloning. It is LU equivalent (the necessary local transformation is $\sigma_z \otimes \sigma_z \otimes \sigma_z \otimes \mathbb{1} \otimes \mathbb{1} \otimes \mathbb{1}$) to the telecloning state described in [45] for $M = 3$ recipients, where $2M = 6$ qubits are necessary. An LU-equivalent state of $|\,\Psi_6^+\rangle$ was recently observed using a different configuration, without the central BS [46].

Finally, we note that the usage of an additional HWP(δ) in mode b in front of the PBS leads to the family of states

$$|\,\Psi_6'(\gamma)\rangle = a_6(\gamma)\,|\,GHZ_6\rangle + b_6(\gamma)(\,|\,\overline{W}_3\rangle \otimes |\,\overline{W}_3\rangle + |\,W_3\rangle \otimes |\,W_3\rangle) \quad (17)$$

which is LU equivalent to the family $|\,\Psi_6(\gamma)\rangle$.

III. Experimental Realization of a Family of Four-Photon Entangled States

Let us now move to the experimental realization of one of the presented schemes. We implemented the interference of four photons at a PBS using as a photon source the second-order emission of a noncollinear type-II SPDC process. The general layout of the experiment was described in Section II-A and Fig. 1, which leads to the observation of the family of states [see Section II-A2 and Fig. 2(a)]

$$|\,\Psi_4(\gamma)\rangle = a_4(\gamma)\,|\,GHZ_4\rangle + b_4(\gamma)\,|\,\psi^+\rangle \otimes |\,\psi^+\rangle. \quad (18)$$

A. Experimental Setup

A frequency-doubled Ti:sapphire laser emits femtosecond UV pulses with a repetition rate of 81 MHz and a power of 600 mW at 390 nm. The UV pulses pump a 2-mm-thick β-barium borate (BBO) crystal, which is cut for type-II noncollinear SPDC (see Fig. 1). Its second-order emission yields the desired four photons necessary for the interference at the PBS. Walk-off effects in the BBO crystal due to birefringence are compensated by an HWP flipping the polarization state of each photon and a 1-mm-thick BBO crystal [47]. The spatial modes a and b are defined by coupling the SPDC emission into single-mode fibers. In mode a, an HWP(γ) is placed before the photons in modes a and b interfere at the PBS. Interference filters centered around the degenerate wavelength of 780 nm with a full-width at half-maximum of 3 nm are placed in output modes c and d (not shown in Fig. 1) to define the spectral modes of the SPDC photons. Further, in mode c, an additional HWP is placed (not shown in Fig. 1), which flips the polarization of the photons. Subsequently, each mode is split by a polarization independent BS, whose birefringence is compensated by a pair of birefringent, perpendicularly oriented yttrium vanadate crystals in each output mode (not shown in Fig. 1). Finally, the polarization state of each photon is analyzed with an HWP and quarter-wave plate (QWP) in front of a PBS. The outputs of the PBS are coupled into multimode fibers, which guide the photons to silicon avalanche photodiodes (APDs). The detection signals are fed into a coincidence unit capable

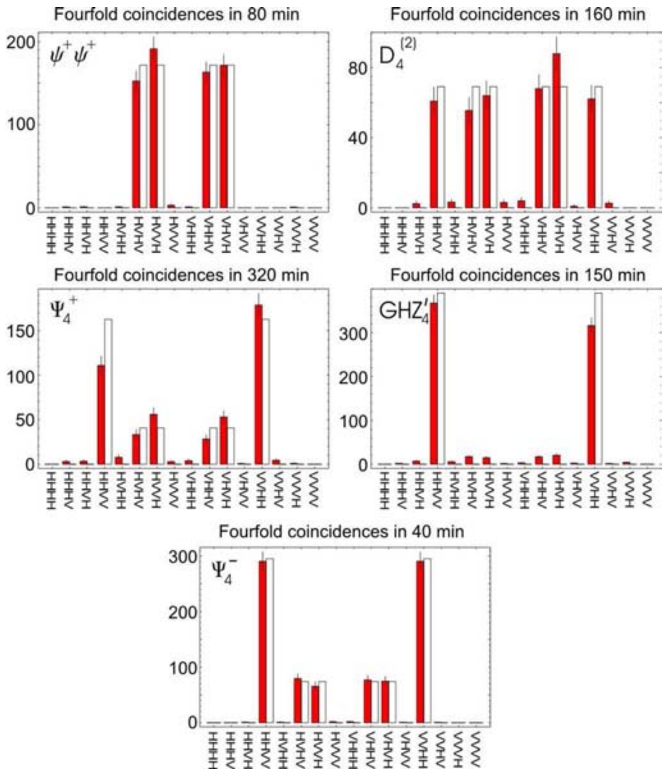

Fig. 4. Recorded counts in the computational basis for the states $|\psi^+\rangle \otimes |\psi^+\rangle$, $|D_4^{(2)}\rangle$, $|\Psi_4^+\rangle$, $|GHZ_4'\rangle$, and $|\Psi_4^-\rangle$. Clearly, the different contributions of the $|GHZ_4'\rangle$ and $|\psi^+\rangle \otimes |\psi^+\rangle$ terms are observable. Open bars without error show expected counts.

of registering all $2^8 = 256$ possible detection events between all eight detectors. The errors on following data are deduced from Poissonian counting statistics and errors on independently determined relative detection efficiencies.

Under the condition of detecting a single photon in each mode e, f, g, h, the family of states

$$|\Psi_4(\gamma)'\rangle = a_4(\gamma)\,|GHZ_4'\rangle + b_4(\gamma)\,|\psi^+\rangle \otimes |\psi^+\rangle \qquad (19)$$

is observed. Note that the family $|\Psi_4(\gamma)'\rangle$ differs (by an LU operation) from $|\Psi_4(\gamma)\rangle$, i.e., by a polarization flip in modes e, f, which is performed by the additional HWP in mode c.

To show the power of the experimental setup, we choose five known states of the family $|\Psi_4(\gamma)'\rangle$, namely $|\psi^+\rangle \otimes |\psi^+\rangle$, $|D_4^{(2)}\rangle$, $|\Psi_4^+\rangle$, $|GHZ_4'\rangle$, and $|\Psi_4^-\rangle$ (with $\gamma = 0$, $\gamma = \pi/12$, $\gamma = (1/2)\arctan(1/\sqrt{2})$, $\gamma = \pi/8$, and $\gamma = \pi/4$, respectively), and record for each state the counts in the computational basis. This demonstrates that we are able to observe different states in a single experimental setup simply by changing the angle setting of HWP(γ). Fig. 4 shows the 16 possible measurement outcomes for these five states. Open bars show the theoretically expected coincidences, with the scaling chosen such to give the same sum of counts. Clearly, a good agreement between experiment and theory is found. Deviations originate from higher order emissions of the SPDC that give undesired contributions. Additionally, an imperfect interference at the PBS further adds noise. Nevertheless, a clear transition between the states $|\psi^+\rangle \otimes |\psi^+\rangle$ and $|GHZ_4'\rangle$ can be observed.

B. Detecting the Entanglement

Let us now discuss the detection of different degrees of entanglement for the observed states. First, we want to exclude that any of the states is separable. Further, we want to show that the expected four-partite entanglement is also found in the observed states.

A simple criterion to exclude separability was recently introduced in [48]. It is based on the correlations of a state. In our case, we have to consider the correlation tensor $\hat{T}$ of a four-qubit state ρ with

$$\rho = \frac{1}{16} \sum_{\mu_1,\ldots,\mu_4=0}^{3} T_{\mu_1,\ldots,\mu_4} (\sigma_{\mu_1} \otimes \sigma_{\mu_2} \otimes \sigma_{\mu_3} \otimes \sigma_{\mu_4}) \qquad (20)$$

where $\sigma_{\mu_n} \in \{\mathbb{1}, \sigma_x, \sigma_y, \sigma_z\}$ is the σ_{μ_n} th Pauli matrix of the nth qubit (with $\sigma_0 = \mathbb{1}$) and $T_{\mu_1,\ldots,\mu_4} \in [-1, 1]$ are the components of the correlation tensor $\hat{T}$. The values $T_{\mu_1,\ldots,\mu_4}$ are given by the expectation value $T_{\mu_1,\ldots,\mu_4} = Tr[\rho(\sigma_{\mu_1} \otimes \sigma_{\mu_2} \otimes \sigma_{\mu_3} \otimes \sigma_{\mu_4})]$. For fully separable states, it holds that

$$T_4^{\max} \geq \sum_{j_1,\ldots,j_4} T^2_{j_1,\ldots,j_4} \qquad (21)$$

where $T_4^{\max}$ is the maximal value of the four-qubit correlation function and is given by $T_4^{\max} = \max_{\vec{o}_1\otimes\vec{o}_2\otimes\vec{o}_3\otimes\vec{o}_4} (\hat{T}, \vec{o}_1 \otimes \vec{o}_2 \otimes \vec{o}_3 \otimes \vec{o}_4)$, with $\vec{o}_n = (T_x^{(n)}, T_y^{(n)}, T_z^{(n)})$ being a 3-D unit vector describing a pure state of the nth qubit [48].

The detection of entanglement can be very simple. As soon as the sum of squared correlations

$$\sum_{j_1,\ldots,j_4} T^2_{j_1,\ldots,j_4} \qquad (22)$$

exceeds unity, our experimental states are entangled [48]. Fig. 5 shows the correlations $T_{1,1,1,1} \equiv T_{x^{\otimes 4}}$, $T_{2,2,2,2} \equiv T_{y^{\otimes 4}}$, and $T_{3,3,3,3} \equiv T_{z^{\otimes 4}}$. When we sum the squares of, e.g., $T_{x^{\otimes 4}}$ and $T_{z^{\otimes 4}}$, we find that for all states, $(T_{x^{\otimes 4}})^2 + (T_{z^{\otimes 4}})^2 > 1$, and thus, all states are entangled. The same is found for $(T_{x^{\otimes 4}})^2 + (T_{y^{\otimes 4}})^2$. Hence, we conclude that the experimental states contain at least some entanglement.

Now, let us demonstrate that the experimental states exhibit the expected genuine four-partite entanglement. Note that the state $|\psi^+\rangle \otimes |\psi^+\rangle$ is a biseparable state, i.e., a product of two entangled pairs, and thus, is the only state that is not genuine four-partite entangled. To show genuine n-partite entanglement, we use the method of entanglement witnesses [49]. Generally, an entanglement witness that detects a genuine four-partite entangled state $|\xi\rangle$ is given by the operator $\mathcal{W}_{\xi,\alpha} = \alpha\mathbb{1}^{\otimes 4} - |\xi\rangle\langle\xi|$. Thereby, the constant α is the maximal overlap of $|\xi\rangle$ with all biseparable states (B-S), i.e., $\alpha = \max_{|\phi\rangle \in \text{B-S}} |\langle\phi|\xi\rangle|^2$. This construction guarantees that $Tr[(\mathcal{W}_{\xi,\alpha})\rho_{\text{B-S}}]$ is positive for all biseparable states $\rho_{\text{B-S}}$, but negative for $|\xi\rangle$. The power of entanglement witnesses stems from the fact that their expectation value is also negative for states close to $|\xi\rangle$. Hence, a negative expectation value $Tr[(\mathcal{W}_{\xi,\alpha})\rho_{\text{exp}}] = Tr[(\alpha\mathbb{1}^{\otimes 4} - |\xi\rangle\langle\xi|)\rho_{\text{exp}}] =$

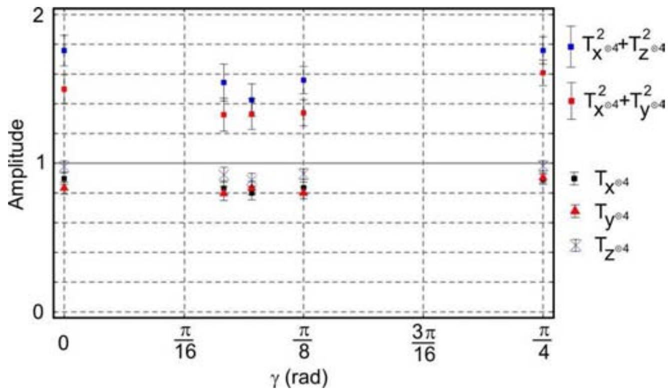

Fig. 5. Experimental correlations $T_{x^{\otimes 4}}$, $T_{y^{\otimes 4}}$, and $T_{z^{\otimes 4}}$ for the states $|\psi^+\rangle \otimes |\psi^+\rangle$, $|D_4^{(2)}\rangle$, $|\Psi_4^+\rangle$, $|GHZ_4'\rangle$, and $|\Psi_4^-\rangle$ (expected value is 1). From these, the values of $(T_{x^{\otimes 4}})^2 + (T_{y^{\otimes 4}})^2$ and $(T_{x^{\otimes 4}})^2 + (T_{z^{\otimes 4}})^2$ are deduced, which allow to demonstrate entanglement, if either of these values exceeds unity. All states fulfill this condition, and thus are entangled.

TABLE II
DETECTION OF GENUINE FOUR-PARTITE ENTANGLEMENT VIA ENTANGLEMENT WITNESSES

State	Entanglement Witness	Expectation value
$\lvert D_4^{(2)}\rangle$	$\mathcal{W}_{D_4^{(2)},2/3}$	-0.142 ± 0.014
$\lvert \Psi_4^+\rangle$	$\mathcal{W}_{\Psi_4^+,3/4}$	-0.030 ± 0.013
$\lvert GHZ_4'\rangle$	$\mathcal{W}_{GHZ_4',1/2}$	-0.330 ± 0.008
$\lvert \Psi_4^-\rangle$	$\mathcal{W}_{\Psi_4^-,3/4}$	-0.182 ± 0.008

All experimental states yield negative expectation values, and thus, are genuine four-partite entangled.

$\alpha - \langle \xi | \rho_{\exp} | \xi \rangle = \alpha - F_\xi(\rho_{\exp}) < 0$ will signal genuine four-partite entanglement of the experimental state $\rho_{\exp}$, where $F_\xi(\rho_{\exp})$ is the fidelity of state $\rho_{\exp}$ with respect to $|\xi\rangle$. Hence, by measuring the fidelity, one can directly compute the expectation value of the corresponding entanglement witness.

Entanglement witnesses that detect the states $|D_4^{(2)}\rangle$, $|\Psi_4^+\rangle$, $|GHZ_4'\rangle$, and $|\Psi_4^-\rangle$ have already been constructed. The corresponding witnesses are given by $\mathcal{W}_{D_4^{(2)},2/3}$ in [34] and [50], $\mathcal{W}_{\Psi_4^+,3/4}$ in [51], $\mathcal{W}_{GHZ_4',1/2}$ in [52], and $\mathcal{W}_{\Psi_4^-,3/4}$ in [51], respectively. To determine their expectation values, we measured the fidelity of the corresponding state (see [27] for details). The result is shown in Table II. We see that all experimental states are four-partite entangled as the expectation values of the corresponding witnesses are below zero.

Hence, we could indeed show that selected states of the family $|\Psi_4(\gamma)'\rangle$ are not only entangled, but also genuine four-partite entangled. We again stress that all these states are SLOCC-inequivalent and that, so far, different experimental setups were necessary to observe each state.

IV. Conclusion

We have shown how multiphoton interference at different types of BSs can be used to observe different families of multiphoton entangled states. The photons were generated by higher order emissions of an SPDC source. The combination of polarization rotations in the BS input modes with multiphoton interference at polarization-dependent BSs provides the necessary ingredients for the powerful scheme we presented. We implemented one particular case experimentally that allowed us to observe an entire family of four-photon entangled states. Our method opens the way for flexible linear optical experiments in the future and surely can also be applied in other areas of quantum information, e.g., for linear optical quantum computing.

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