Purification of Single and Entangled Photons by Wavepacket Shaping

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Single photons and entangled photons lie at the heart of photonic quantum technologies, whose optimal performances are normally reached when the purity of the single or entangled photons is high. However, the multiphoton emission, dissipation, and decoherence in practical realizations always lead to the degradation of the single- and entangled-photon quality. The purification of single or entangled photons is thus valuable to restore the quantum states and enhance the performance of quantum technologies. The applications of wavepacket shaping, an emerging quantum optics tool to manipulate the single- and entangled-photon wavefunctions, to purify single and entangled photons are reviewed. In particular, by modulating the single photons emitted from optically excited room-temperature quantum dots, it is shown that the fast-decaying multiphoton emission can be eliminated to obtain a low value of $g^{(2)}(0)$ that is independent of the excitation power. It is also shown that the two-photon interference and polarization entanglement of the non-degenerate biphotons from spontaneous parametric down-conversion can be restored after modulating the biphotons with a periodic function. The works have potential applications in long-distance quantum communication and linear optical quantum computation.

1. Introduction

Photonic quantum technologies, such as quantum communication,[1–6] quantum computation,[7] and quantum random number generation,[8,9] benefit from single and entangled photons with high quality. For example, highly pure single photons can provide significantly improved security and communication rate in quantum key distribution.[10–12] Highly entangled photons are also demanded for implementing the long-distance quantum communication[13] and large cluster states.[14] However, the single and entangled photons in the practical realizations generally suffer the degraded quality owing to the multiphoton emission, dissipation, and decoherence. The purification of single and entangled photons is thus useful to restore their quantum states as well as the performance of quantum technologies. In this article, we review the purification of single and entangled photons by shaping their quantum wavepackets.

The wavepacket shaping of single and entangled photons has enriched photonic quantum technologies in various aspects. For example, in a quantum network comprising strongly coupled cavities and atoms, the quantum state mapping between distant nodes can be nearly perfect by sending single photons with a time-symmetric waveform.[15] The storage or absorption of single photons by atomic ensembles can also be optimized or suppressed if employing a single-photon waveform determined by the time-reversal technique.[16] To manipulate the quantum wavepackets, the modulation of single and entangled photons by modulators has emerged as a powerful and straightforward quantum optics tool. Like many other methods that shape the laser pulses pumping the single atom, ion, or quantum dot in the single- or entangled-photon emission,[17–23] the temporally long wavepacket of the single or entangled photons plays an indispensable role. The electro-optic modulation of a single-photon wavepacket was first demonstrated by the Harris group at Stanford University.[24] In their experiment, single anti-Stokes photons were heralded by detecting the Stokes photons from the spontaneous four-wave mixing in cold atoms.[25] The detection of the Stokes photons in turn establishes the time origin for an electro-optic intensity modulator, which modulated the single photons into the waveforms of rectangular pulses, a Gaussian, and an exponential decay. Although this experiment only demonstrated single photons of controllable waveforms, the modulation of wavepackets is not limited to the waveform shaping. The phase shaping of single photons, which finds potential applications in quantum key distribution,[26,27] where the phase changes of single photons are to be measured, and linear optical quantum computation,[28,29] where arbitrary phases may be created or compensated, was later demonstrated by the Rempe...
group at Max-Planck-Institute for Quantum Optics using a coupled single-atom cavity quantum electrodynamics system.[30] The single photons were generated by pumping an atom trapped in a high-finesse cavity with laser pulses, which provided the necessary time origin for the electro-optic phase modulation.

The modulation of wavepacket can also be applied to entangled photons. Interestingly, the time origin is not necessary for modulating the entangled photons although it can be established by modulating the pump fields[31,32] or heralding the entangled photons by detecting auxiliary photons.[33,34] Harris[35] showed that, if each photon of the entangled pair is modulated by the same sinusoidal functions at the same frequency, the Glauber correlation function (or the temporal wavepacket) of the entangled photons will be modulated by the intensity correlation of the modulation functions. Namely, the absolute common phase of the sinusoidal modulation on each photon is insignificant. For example, if each photon is modulated by \(\cos(\omega t)\), the entangled-photon wavepacket is modulated by \(\frac{2 \cos(2\omega t)}{\omega}\) where \(\omega\) is the modulation frequency and \(t\) is the time delay between the detected photon pair.

Since the pioneer works by the Harris and Rempe groups, a number of applications have been demonstrated.[30,36–43] The fermionic spatial behavior of two photons was observed using the phase-shaped single photons.[40] A new technique of measuring the wavepacket of short entangled photons by “slow” detectors that would otherwise be incapable of resolving the temporal shape due to the poor timing resolution was developed.[36] The recovery of single photons, which were initially hidden in a massive amount of noise photons with the same frequency, was demonstrated by shaping the single photons with a set of complementary pseudo-random phase sequences.[37] By modulating the single photons into an exponential growth waveform, the efficient single-photon absorption by an atomic ensemble[38] and the efficient single-photon loading into an optical cavity[39] were also demonstrated. The optical precursor of a single photon was observed by using step- and square-modulated single-photon waveforms,[40] where the upper bound of the information flow at the single-photon level is also confirmed to be the speed of light at vacuum. The time-resolved detection of the photon-surface-plasmon coupling was demonstrated by exploiting the waveform-controlled single photons.[41] Recently, by modulating the single and entangled photons, we restored the degraded single-photon purity of the room-temperature colloidal quantum dots[42] and the degraded polarization entanglement from the two-photon interference,[43] which will be the focuses of this review with more details described as follows.

2. Purification of Single Photons

The optical properties and scalability of semiconductor quantum dots are striking among the available single-photon sources today.[44] While enormous effort has been devoted to the single-photon sources based on the self-assembled quantum dots at cryogenic temperatures, single-photon emission was also demonstrated with the colloidal core/shell quantum dots.[45] At the room temperature, the colloidal quantum dots exhibit high quantum yield and high photostability. However, due to the broad biexciton spectrum that spectrally overlaps with the single-exciton spectrum, the single-photon purity of the room-temperature colloidal quantum dots was never comparable to their cryogenic-temperature counterparts. We describe below how the modulation of single photons can be useful to increase the single-photon purity of the colloidal quantum dots at the room temperature.[42]

2.1. Experimental Setup

Figure 1 illustrates our experimental setup, where CdSeTe/ZnS quantum dots (7.5-nm radius, 705-nm peak emission) are spin-coated on a cover glass (CS) and optically excited by focused laser pulses using a 100x microscopic objective (MO). The emitted single photons are collected by the same MO, separated from the laser pulses at the dichroic mirror (DM), and detected by the single photon counting modules (SPCMs). To generate waveform-controlled single photons, an acousto-optic modulator (AOM) driven by a function generator is used to modulate the waveform of the single photons. Other optical elements shown are mirror (M), filter (F), pinhole (PH), lens (L), and beam splitter (BS).

2.2. Measurement and Modulation of Single-Photon Wavepacket

The time-resolved photoluminescence of the single quantum dot is then confirmed by the antibunching statistics (an unique feature of single photons) in the Hanbury-Brown–Twiss measurement (Figure 2b). The typical brightness and fluorescence spectrum of such a single quantum dot at different excitation powers are shown in Figure 2c,d, respectively.
constant \( \tau_1 \approx 2 \) ns is contributed by the multiphoton emission mainly from the biexciton transition.\(^{[46]}\) The slow decay of time constant \( \tau_2 \approx 138 \) ns, on the other hand, results from the radiative single-exciton transition,\(^{[47]}\) which determines the temporal width of the single photons. This long temporal length allows us to manipulate the wavepacket by the acousto-optic modulators. As the polarization of the fluorescence from a single colloidal quantum dot depends on the quantum dot’s orientation, an acousto-optic modulator (which is insensitive to the incident polarization compared to the electro-optic modulators) is more appropriate than an electro-optic modulator to modulate the photons.

To establish the time origin for the modulation, the excitation laser is triggered by the function generator (after a controlled time delay) which drives the acousto-optic modulator. By doing so, the single photons and the electric signal driving the acousto-optic modulator arrive at the modulator simultaneously. Figure 3b shows the waveform of the single photons modulated...
Figure 4. The calculated a) biexciton percentage and b) relative total counts at different time offsets when the single photons is modulated by a Heaviside step function. c) Upper panel shows the waveforms of the unmodulated (red) and modulated (black, \( t_0 = 45 \text{ ns} \)) single photons. Lower panel shows the Hanbury–Brown and Twiss measurements of the modulated (red) and modulated (black, \( t_0 = 45 \text{ ns} \)) single photons. d) The measured biexciton percentage at different time offsets.

2.3. Elimination of Multiphotons

The emission of single photons from colloidal quantum dots is contaminated by the multiphoton emission. This is evident by the non-vanishing counts at the time delay \( \tau \approx 0 \mu s \) in Figure 2b, where the antibunching is supposed to happen. Nevertheless, the multiphoton emission decays in a few nanoseconds as can be seen in Figure 3a. This fast decay rate makes the elimination of the multiphoton emission possible. For example, if one can modulate the single photons with a Heaviside-step function \( H(t - t_0) \) at a proper time offset \( t_0 \), the multiphoton component can be mostly removed at a cost of reducing the single-photon rate. This is illustrated in Figure 4a, where we calculate the percentage of the photons \( \beta \) from the multiphoton (biexciton) emission as a function of the time offsets. The calculation assumes an emission consisting of 4% fast-decaying biexcitons and 96% slow-decaying single excitons with time constants of 2 and 138 ns, respectively. A significant drop of the biexciton percentage can be seen after the time offset increases to a few nanoseconds. When the time delay is 6 ns, the biphotons are reduced to 0.01% while the single-photon rate reduces by 10% (Figure 4b).

In practice, the modulation waveform is smoothed at the rising edge due to the finite rise time of the modulator. Figure 4c (upper panel) shows the wavepackets of the unmodulated (red) and modulated (black) single photons, with the corresponding Hanbury–Brown and Twiss measurement given in the lower panel. When the chosen time offset is larger than 2 ns, the area of the peak centered at \( \tau = 0 \mu s \) reduces noticeably. At large \( t_0 \), the biexciton emission is completely eliminated and the peak centered at \( \tau = 0 \mu s \) vanishes. These observations are consistent with the fact that the lifetime of the biexciton state is \( \tau_1 \approx 2 \) ns. In Figure 4d, we also plot the measured biphoton percentage versus the time offsets. To estimate the biphotons percentages, we reconstruct the time trace of the single-exciton emission by fitting the photoluminescence at large time offsets (>100 ns) with a single exponential decay. The time trace of the biexciton emission is then obtained by subtracting that of the single-exciton emission from the total emission. One can see that the biexciton emission (percentage) is nearly eliminated at large time offsets.

The biphoton percentage is approximately related to the normalized photon correlation function at zero time delay by \( g^{(2)}(0) \approx 2\beta^{4\beta} \) if \( \beta \) is small and only the exciton and biexciton emission are considered. The \( g^{(2)}(0) \) function gives an upper
Figure 5. Normalized photon correlation functions $g^{(2)}(0)$ at zero time delay. a) The time offset of the modulation function is varied from 16 to 45 ns for five quantum dots. b) The excitation power is varied from 1.4 to 7 $P_{\text{sat}}$ for four quantum dots. $g^{(2)}(0)$ of the unmodulated and modulated single photons are shown by the open and solid symbols, respectively.

bound on the probability of the multiphoton emission$^{[49]}$ and can be used to characterize the purity of the single-photon sources. In practice, the dark counts of the detectors, the leakage light of the excitation laser, and the higher-order multiphoton emission can also contribute to $g^{(2)}(0)$. For pulsed sources, $g^{(2)}(0)$ is equal to the ratio of the peak area at zero time delay ($\tau = 0 \mu s$) to the average of the peak areas at nonzero time delays ($\tau \neq 0 \mu s$).

**Figure 5a** shows the $g^{(2)}(0)$ from five quantum dots at the excitation power of $6.7P_{\text{sat}}$ before ($t_0 = -50$ ns) and after ($t_0 > 15$ ns) the modulation. The $g^{(2)}(0)$ of the single photons initially varies from 0.04 to 0.08. After the modulation, both the magnitude and variation of $g^{(2)}(0)$ decrease. Although the strength and lifetime of the biexciton emission from different quantum dots vary, we are able to achieve $g^{(2)}(0) = 0.02$ for all quantum dots at large time offset by controlling the waveform of the single photons. **Figure 5b** shows the $g^{(2)}(0)$ of the single photons from four quantum dots at different excitation powers. Without the modulation, the probability of the biexciton emission increases with the excitation power, resulting in a degradation of the single-photon purity. The measured $g^{(2)}(0)$ of the unmodulated single photons (open symbols) increases up to 0.07 at an excitation power of $6.7P_{\text{sat}}$. However, with the biexciton emission eliminated by modulating the wavepacket, the emitted single photons preserve the high purity ($g^{(2)}(0) \approx 0.01$) even at high excitation power.

### 3. Purification of Entangled Photons

Two-photon interference$^{[50,51]}$ has been the workhorse in many quantum optics experiments and applications.$^{[52-54]}$ By interfering non-entangled and non-orthogonally polarized photon pair at a beam splitter, the photons exiting through different ports can be entangled in the polarization degree of freedom. However, if the frequencies of the photon pair are different, the interference visibility reduces and the entanglement degrades or disappears. We describe below how the lost two-photon interference as well as the resulting entanglement and non-locality can be restored by modulating the wavepacket of the entangled photons.$^{[43]}$ The degree of the restored entanglement can achieve full recovery if proper modulation functions are applied to the entangled photons.

#### 3.1. Experimental Setup

Our experimental setup is illustrated in **Figure 6**. Collinear type-II quasi-phase-matched biphotons at 795 nm are generated by the doubly resonant parametric downconversion in a 1-cm-long periodically poled KTP crystal, pumped by a frequency-doubled and stabilized continuous wave laser at 397.5 nm (linewidth $<2$ MHz). To achieve single-mode operation and high spectral brightness without the external filtering,$^{[55-57]}$ the crystal’s end faces are spherically polished and high-reflection coated at the signal and idler wavelengths to realize a monolithic cavity. In addition, one of the end faces is also high-reflection coated at the pump wavelength to implement the double-pass pumping. The biphotons have a linewidth of $\Delta \omega = \left[ (\sqrt{\Gamma_1^s + 6\Gamma_1^i} + \Gamma_1^i - \Gamma_2^i - \Gamma_3^i) / 2 \right]^{1/2} \approx 5$ MHz$^{[58]}$, that is narrower than the natural linewidth of the $^{87}$Rb D1 and D2 lines, where $\Gamma_1 = 1/(20.7 \text{ ns})$ and $\Gamma_2 = 1/(24.4 \text{ ns})$ are the decay constants of the double-exponential temporal wavepacket. The spatial modes of the signal and idler photons are very close to a Gaussian, with a slight divergence angle corrected by an additional lens.

#### 3.2. Preparation and Measurement of Polarization Entanglement

The temporally long biphotons are entangled in the polarization degree of freedom by selecting pairs exiting through different ports of the beam splitter. The frequency difference of the signal and idler photons, which determines the distinguishability in the two-photon interference and the degree of entanglement, is tuned by adjusting the phase-matching condition and measured by the beat frequency of the time-resolved two-photon interference. When the biphotons are degenerate, the beat vanishes and the two-photon interference results in the polarization-entangled state $(|H_1V_1]_a + |V_1H_1]_b)/\sqrt{2}$ with $a$ and $b$ denoting the ports each photon exits. The entanglement can be characterized by the concurrence$^{[48]}$. To calculate the concurrence, we first compute the square root of the eigenvalues of $\rho(\sigma_1 \otimes \sigma_1)\rho^*(\sigma_2 \otimes \sigma_2)$ in descending order. {\sqrt{\lambda_1, \ \sqrt{\lambda_2, \ \sqrt{\lambda_3, \ \sqrt{\lambda_4}}}$. The concurrence is then obtained by $C = \max(0, \sqrt{\lambda_1 - \sqrt{\lambda_2} - \sqrt{\lambda_3} - \sqrt{\lambda_4})$, whose value ranges from 0 (for non-entangled states) to 1 (for maximally entangled states). Using this entanglement measure
Figure 6. Schematic of the experimental setup. ECDL: external cavity diode laser, SHG: second harmonic generation, BS: beam splitter, Pol.: polarizers, $\lambda/2$: half-wave plates, $\lambda/4$: quarter-wave plates, EOM: electro-optic intensity modulators, SPCM: single-photon counting modules.

Figure 7. Density matrices of the biphotons with a frequency difference of a) 50 MHz and b) 100 MHz. c) The CHSH inequality $|S| \leq 2$ of the biphotons requires shorter coincidence windows to be violated when the frequency differences from the top to bottom are 0, 10, 20, 40, 60, 80, and 100 MHz.

and the density matrix $\rho$ reconstructed by the quantum state tomography, we obtain the concurrence $C = 0.71$ and the purity $\text{Tr}[\rho^2] = 0.81$ for degenerate biphotons.

The concurrence declines rapidly as the frequency difference increases. Figure 7a,b show the tomographic reconstruction of the density matrices for nondegenerate biphotons with a frequency difference of 50 and 100 MHz, respectively. For both frequency differences, the concurrence and the purity are 0 and 0.45, respectively. Similar decline of entanglement can also be observed in the violation of the Clauser–Horne–Shimony–Holt (CHSH) inequality, which places constraints on the value $S$ of a combination of four polarization correlation probabilities analyzed by two possible settings for each photon. If the CHSH inequality is violated, quantum entanglement is necessary to explain the correlations or nonlocality. To verify the CHSH inequality, we arrange each set of the quarter-waveplate, $\phi$ value of the CHSH inequality and observe the non-locality.

3.3. Revival of Quantum Entanglement

To recover the quality of entanglement and the ability to observe nonlocality, the indistinguishability in two-photon interference must be improved. This can be accomplished by modulating the wavepacket. The modulation of the temporal wavepacket at the frequency difference of the signal and idler photons ensures a constant phase between the two-photon states in the two-photon interference. Such a modulation equivalently generates signal (or idler) photons at the carrier frequency or sidebands of the idler (or signal) photons. With the two-photon wavefunction in our experiment denoted by $\phi_{ij}(t) = \phi_{ij}(t) = e^{-\pi t/2\sigma_0}$, the quantum state exiting through different ports of the beam splitter is given by

$$|\nu\rangle = \frac{1}{\sqrt{2}} \int dtdr |\phi_{ij}(t)\rangle \hat{a}^+_j(t)|0\rangle + e^{i\Delta\omega t}\phi_{ij}(t)\hat{a}^+_j(t)|0\rangle e^{i\omega_0 t}|0\rangle,$$

where $\hat{a}^+_j(t) = m_j(t)\hat{a}_{j+0}(t)\hat{a}^+_{j+0}(t)$ is the creation operator of the $j$-polarized photon at ith port after (before) the amplitude modulation $m_j(t)$, and $\Delta\omega/2\pi = (\omega_1 - \omega_0)/2\pi$ is the frequency difference of the biphoton. Without the modulation $(m_j(t) = 1)$, the coherence $\zeta(\Delta\omega) = (1/4\sigma_0)\int_0^{\infty} e^{-\pi t^2/\sigma_0} e^{-i\Delta\omega t} dt = 1/2(1 + \theta^2)$ with $\theta = \Delta\sigma \sigma_{0}$ in the density matrix, $\rho = 1/2(|H_1 V_2\rangle|V_1 H_2\rangle + |V_1 H_2\rangle|H_1 V_2\rangle + |V_1 H_2\rangle|V_1 H_2\rangle + |H_1 V_2\rangle|V_1 H_2\rangle)$. is gradually
lost as the frequency distinguishability increases. When \( \Delta \omega \gg \tau_0^{-1} \), the coherence as well as the concurrence \( C = 2 \zeta (\Delta \omega) \) (solid curve in Figure 8a) vanishes; the entanglement does not exist anymore. However, if we modulate the signal and idler photons by synchronized square-wave functions \( |m_i(t)|^2 = |S(t)|^2 \) (the biphoton wavepacket is accordingly modulated by triangular function \( M(\tau) = (1/T) \int_{-\tau/2}^{\tau/2} |S(t)|^2 |S(t + \tau)|^2 dt \) twice the frequency with the modulation period denoted by \( T \), the coherence \( \zeta (\Delta \omega) = \int_{-\infty}^{\infty} M(\tau) e^{-i\Delta \omega \tau} d\tau = \int_{-\infty}^{\infty} M(\tau) e^{-i|\tau|/\tau_s} d\tau = (\pi \theta + \pi s \theta^2 + \theta^3)/[2\pi(\pi - \theta)(1 + \theta^2)] \), which peaks at the modulating frequency of \( \Delta \omega/2\tau_s \), remains finite even if \( \Delta \omega \gg \tau_0^{-1} \). The corresponding concurrence (dot curve in Figure 8a) and purity (Figure 8b) are thus restored at large frequency differences.

The non-classical time correlation of the biphotons also allows to restore the entanglement by modulating one photon of the pair conditionally (on the detection of another photon) with the convolution of the modulation functions. If one photon is modulated by a cosinusoidal function \( (1 + e^{i \Delta \omega \tau})/2 \), both the coherence \( \zeta (\Delta \omega) = (2 + 7\theta^2 + 2\theta^3)/(2 + \theta^2)(1 + 4\theta^2) \) and concurrence (dash-dot curve in Figure 8a) survive at large frequency differences as well; the entanglement is restored again. We note that the concurrence at large frequency differences of \( \approx 100 \text{ MHz} \) does not place a limitation on the entanglement restored by shaping the biphotons. The amount of entanglement that can be restored depends on the modulation function used. For example, if the biphotons are modulated by a periodic sinc \(^2 \) function \( \int_{-\infty}^{\infty} |m_i(t)|^2 |m_i(t, \tau)|^2 dt = [(1/s) \sum_{\omega=\Delta \omega} \exp(i \omega \Delta \omega \tau_s)]^2 \) with \( s = 100 \) (which reduces to the cosinusoidal function if \( s = 2 \)), the entanglement can be restored nearly to its quality at degeneracy (dash curve in Figure 8a).

To modulate the entangled-photon wavepacket, we use a function generator to drive the electro-optic intensity modulators in both signal and idler channels synchronously. Figure 9a shows the wavepacket of the biphotons, which has a frequency difference of 50 MHz and is modulated by a triangular function at the same frequency. By tomographically reconstructing the density matrix in Figure 9b, we observe increased concurrence \( C = 0.28 \) and purity \( \text{Tr}[\rho^2] = 0.5 \) compared to 0 and 0.45, respectively, in the unmodulated case. In Figure 9c, we modulate only the signal photons of the biphotons with a frequency difference of 100 MHz conditionally by a cosinusoidal function at 100 MHz. For this purpose, the detection of the idler photons by single-photon counting module is used to trigger the function generator which drives the electro-optic modulator. In addition, the signal photons are optically delayed by a 50-m-long optical fiber to ensure that the arrival of the signal photons at the modulator is synchronized with the start of the amplitude modulation. The density matrix is shown in Figure 9d. The concurrence \( C = 0.32 \) and purity \( \text{Tr}[\rho^2] = 0.5 \) both increase again. The CHSH inequality (Figure 9e) can also be violated at larger frequency differences and coincidence windows. For example, the inequality with a frequency difference of 20 MHz (light blue curve), which would not be violated previously...
for coincidence windows larger than 19 ns (Figure 7c), is violated all the way up to 100 ns.

4. Conclusion

We have reviewed the purification of single and entangled photons by the modulation of the single-photon and entangled-photon wavepackets. By temporally shaping the single photons emitted from the colloidal quantum dots at room temperature, we eliminate the biexciton emission and obtain $g^{(2)}(0) = 0.01$, which has only been achieved previously at the cryogenic temperature among the solid-state single-photon sources. Our work thus provides a novel way of preparing high-purity single photons from room-temperature quantum dots and allows the generation rate to increase with the excitation power without compromising the single-photon quality. We note that the generation rate will eventually be limited by the inverse of the wavepacket's temporal length, which can be optimized by exploiting the Purcell effect or selecting a different kind of quantum dot. We also note that, in contrast to the wavepacket engineering achieved with single ions, atomic ensembles, or nonlinear crystals, the wavepacket shaping utilized in our work manipulates the temporal envelope but not the waveform of the single photons.

By modulating the biphoton's temporal wavepacket, we have also observed the revival of quantum interference, entanglement and nonlocality that would otherwise be degraded or lost due to the frequency distinguishability of biphotons. Our study shows that the amount of the restored entanglement is only limited by the forms of modulation and can achieve full recovery if the modulation function is properly designed. Quantum entanglement is at the heart of many photonic quantum technologies. The storage of entanglement of narrowband biphotons in quantum memories is important for realizing quantum repeaters or large cluster states. Previously, various schemes of the entanglement purification of arbitrary unknown states, such as using the controlled NOT (CNOT) gates, linear optics with two ideal sources of polarization-entangled photons, (phase-stabilized) spatial entanglement with or without the quantum non-demolition measurements, or polarization-entangled photon pairs as part of the parity check have been proposed or demonstrated. The wavepacket shaping demonstrated in our work, which does not exploit the CNOT gates, additional (ideal) entangled-photon sources, or (phase-stabilized) spatial entanglement, is particularly useful for purifying the polarization entanglement degraded by nondegenerate entangled photons.

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Conflict of Interest

The authors declare no conflict of interest.
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