

Optimal filtering and generation of entangled photons for quantum applications in the presence of noise

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Filtering is commonly used in quantum optics to reject noise photons, and also to enable interference between independent photons. However, filtering the joint spectrum of photon pairs can reduce the inherent coincidence probability or loss-independent heralding efficiency. Here we investigate filtering for multiphoton applications based on entanglement and interference (e.g., quantum teleportation). We multiplex C-band entangled photons and C-band classical communications into the same long-distance fibers, which enables scalable low-loss quantum networking but requires filtering of spontaneous Raman scattering noise from classical light. Using tunable-bandwidth filters, low-jitter detectors, and polarization filters, we distribute wavelength-degenerate time-bin entangled photons in the C band (1536.5 nm) over 25 km/25 km standard fibers (50 km) each carrying 10 Gbit/s C-band classical data signals. Narrow filtering enables milliwatt-level C-band power, which exceeds the power achieved by comparable studies by roughly an order of magnitude and could feasibly support terabit per second classical rates. We evaluate how performance depends on pump and filter bandwidths, multipair emission, filter shapes, loss, phase matching, and how quantum information is measured. We find a trade-off between reducing noise impact and increasing single-mode purity, and discuss mitigation methods toward optimal multiphoton applications. Importantly, these results apply to noise in free space and in quantum devices (sources, frequency converters, switches, detectors, etc.) and provide insight into filter-induced degradation of single-photon purity and rates even in noise-free environments.

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I. INTRODUCTION

The deployment of advanced quantum optical technology is a widespread initiative of quantum research, which would enable quantum-enhanced encryption, sensing, and distributed quantum computation [1–3]. However,

real-world environments can have high background noise that obscures the detection of single photons, limiting the fidelity or distance over which quantum information can be communicated. Noise photons can originate in the transmission medium, including from spontaneous Raman scattering (SpRS) when quantum and classical networks coexist in the same deployed fibers [4–27] and ambient light in free-space quantum links [28–33], or in quantum devices including some photon sources [34,35], quantum frequency converters [36,37], or optical switches and shutters [38–40]. Noise photons are most commonly rejected by frequency, time, polarization, and spatial filtering around the known modes of the photons generated at the quantum light source [Fig. 1(a)]. The design of filters for individual detectors is a well-studied subject (see, e.g., Refs. [30,31,41–43]), including time-bandwidth limits and incoherent filtering versus coherent filtering.

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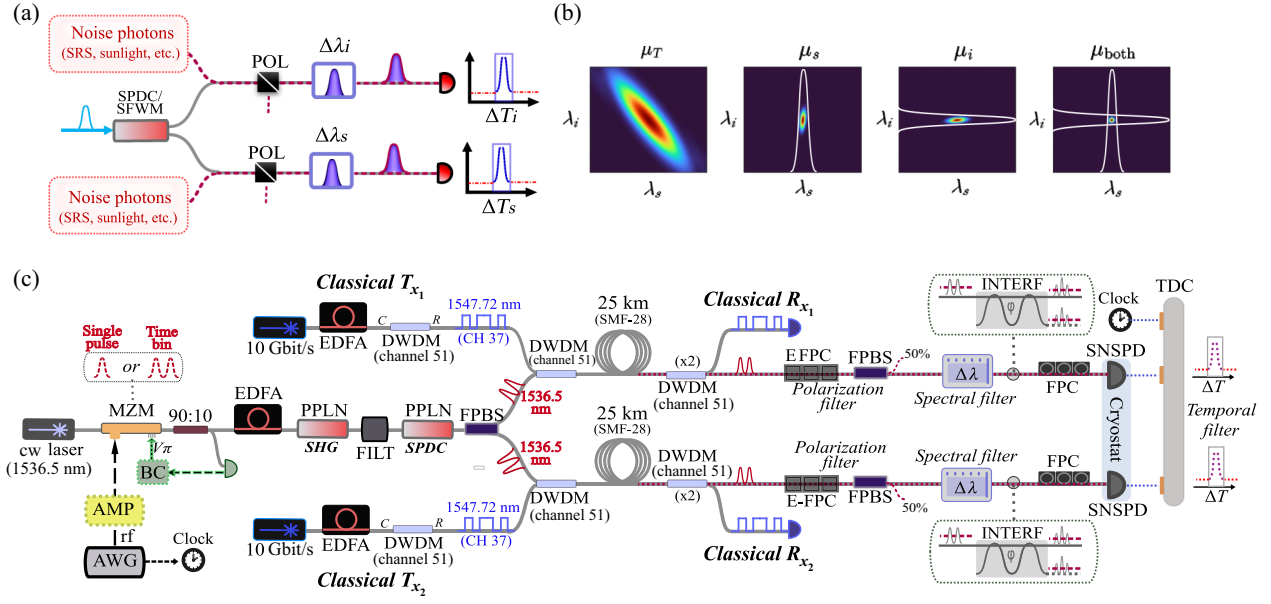


FIG. 1. (a) Conceptual diagram for polarization, frequency, and time filtering around photon pairs in the presence of background noise photons. (b) Effect of filtering a source's joint spectral correlations on the number of photons within the signal (μ_s) and idler (μ_i) filter passbands, which lead to single-photon counts, and their overlap (μ_{both}), which leads to true coincidence counts. When $\mu_s, \mu_i > \mu_{\text{both}}$, not all detected photons will yield true coincidences. (c) Experimental diagram. Single-pulse or time-bin entangled photon pairs generated by type-II SPDC at degenerate wavelengths (1536.5 nm) are distributed over two 25 km single-mode fibers to copropagate with 10 Gbit/s C-band classical signals (1547.72 nm), which generates strong SpRS noise. We filter SpRS and SPDC photons using tunable-bandwidth filters (32–600 pm). Polarization is filtered with a FPBS and an EFPC to actively correct time-dependent polarization rotations. Temporal filtering is applied in postprocessing of the counts from low-jitter SNSPDs and a TDC. Michelson interferometers (INTERF) are inserted for time-bin experiments to perform measurements along the time-bin Bloch sphere [71]. AMP, amplifier; BC, bias controller; POL, polarizer; rf, radio frequency.

However, multiphoton quantum applications can involve more complex physics, as the design of each individual filter contributes to the overall fidelity of nonlocal multiphoton correlations. Furthermore, the physics of quantum sources must also be considered. Entangled photons are most commonly created by generating correlated photon pair states via spontaneous parametric down-conversion (SPDC) or spontaneous four-wave mixing (SFWM), where filtering is also often used to increase their single-mode purity to achieve Hong-Ou-Mandel (HOM) interference [44] between independent sources [45]. This enables operations such as Bell state measurements [46], and underlies applications such as quantum state teleportation [46], entanglement swapping [47], Greenberger-Horne-Zeilinger state generation [48], photonic quantum computing [49,50], and operations [51,52].

Unfortunately, filtering a photon pair source's joint spectral amplitude (JSA) can degrade the “filter heralding efficiency” (FHE) [53–60], i.e., the probability that each detected signal/idler photon will be detected in coincidence with a photon at the idler/signal detector, even in a lossless environment. This issue is not resolved by simply increasing the pump power, as systems with an imperfect FHE are more susceptible to degradation of fidelity

from multiphoton pair emission at frequencies that never yield “true” coincidences between single-photon pairs [55,58,60]. Although many studies have explored engineering the source to obtain high single-mode purity without filtering [50,54,58,61–70], noisy environments still require narrow external filtering of one or more photons.

If the reduction in coincidence rates exceeds the rejection of noise, increases in fidelity may be hindered compared with increases in systems with ideal FHE. This further influences the design of multiphoton applications when filtering also affects the single-mode purity.

In this paper, we explore these concepts in the context of filtering SpRS noise photons generated when classical communication signals propagate in the same long-distance fibers as entangled photons, which would enable large-scale quantum networking within the existing fiber infrastructure and the ability to use classical signals to perform quantum node synchronization, routing, and error mitigation. Although “quantum-classical coexistence” in fiber has been well explored with the use of weak coherent state sources and a single detector [4,6–10,21], recent studies have focused on identifying the challenges facing more complex applications using entangled photons [5,12,13,16–19,23–25] and quantum teleportation [20,22,27].

Since the SpRS spectrum has a broad bandwidth spanning the full 1260–1625 nm range of telecom wavelengths (more than 53 THz) [12,14,15], filtering is critical to maximize quantum and classical capabilities. However, allocation of quantum and classical signals to different telecom bands suppresses SpRS due to the multimode Raman gain spectrum and the phonon population, both of which are highly dependent on the frequency separation of the signals [6–8,11,12,14,15]. For example, the use of O-band (1260–1360 nm) quantum sources and narrow filtering can enable applications based on weak coherent states [10], entangled photons [12], and quantum teleportation [20,22] to operate alongside C-band (1530–1565 nm) classical power on the order of 20 dBm. However, use of the O band sacrifices transmission loss for noise suppression (approximately 0.33 dB/km in standard fiber), whereas quantum-classical coexistence in the C band using dense wavelength-division multiplexing (DWDM) enables both network paradigms to leverage the lowest-loss telecom wavelengths (approximately 0.18 dB/km). The main challenge lies in the significantly higher Raman gain and phonon population, which makes it difficult to engineer discrete-variable systems that can tolerate more than -10 dBm of C-band power [5,7–9,11,13,16,19,21–23]. Relatively high classical rates have been achieved by attenuating classical power, but much higher power is common in the current fiber infrastructure. C-band quantum networking alongside milliwatt-level C-band classical communications has thus been a long-standing goal.

To experimentally investigate this, we distribute either single-pulse or time-bin entangled photon pairs at degenerate wavelengths (1536.5 nm) over 50 km of standard single-mode fiber, where each photon copropagates over 25 km fibers with a 10 Gbit/s C-band (1547.72 nm) classical data signal [Fig. 1(c)]. We first study the filtering of correlated photon pairs and show that imperfect FHE has a negative effect on performance. Then we apply narrow spectrotemporal and polarization filtering to distribute pulsed time-bin entangled photons in the C band alongside milliwatt-level C-band power, which is an order of magnitude higher than in comparable C-band/C-band experiments using entangled photons [5,13,16,18,19,23]. We then quantify how optimizing FHE, loss, and pump power could increase tolerance to classical power or noise.

By simulating various system designs, we find that fidelity depends on the pump pulse width, filter bandwidth and shape, pump power, and phase matching. For example, flat-top filters are shown to achieve higher noise tolerance and purity compared with Gaussian filters. Furthermore, we find that the duration of the pump pulse can notably impact performance since the FHE is lessened at wider filter bandwidths, where optimizing the entanglement fidelity can come at the expense of single-mode purity. We consider the impact of multipair emission for each case, showing that the trade-off between the

single-detector signal-to-noise ratio (SNR) and degradation due to multipairs is more nuanced for systems with an unideal FHE.

This work highlights further considerations for the design of sources and receivers beyond single-detector filtering and applies to many scenarios beyond quantum-classical integration, such as noise in free-space links and quantum devices. Our results and methods have implications for many experiments using filtering to either reduce noise or increase single-mode purity and for high-repetition-rate sources requiring short pulses and wider filters. Furthermore, because of the 1536.5 nm wavelength of both entangled photons, this study provides insight into the challenges of deploying quantum memories based on erbium ions [72,73] into noisy fiber networks carrying classical signals. The ability to copropagate high classical power in nearby C-band DWDM channels shows promise for scalable low-loss entanglement-based networks in existing infrastructure. More broadly, our thorough investigation of the role of multipair emission given imperfect FHE in noise-free environments provides useful methods for designing general multiphoton applications.

Multiple experimental and simulated scenarios are presented in this work to explore each of the concepts relevant for high-performance quantum communication in noisy environments. In Sec. II, we demonstrate the impact of tight spectral filtering on two entangled photons traveling through 50 km of fiber. The trade-off of decreased filter heralding efficiency and increased accidental coincidence count suppression is discussed in detail to set the framework for optimal filtering. In Sec. III, optimal filtering techniques established in Sec. II are showcased in a high-noise scenario where a C-band/C-band entanglement distribution coexistence experiment is demonstrated. While these filtering parameters work for our specific experiment, in Sec. IV, we simulate filters with different spectral shapes and sources with various pump widths to illustrate the complex relationship between filtering bandwidth, pulse width, coincidence-to-accidental ratio (CAR), and spectral purity in both dark and noisy environments. We conclude with an in-depth discussion in Sec. V aimed toward how the lessons learned in experiments and simulations can be applied to future entangled photon source design.

II. FILTERING NOISE AND CORRELATED PHOTON PAIRS

Our experimental design is shown in Fig. 1(c). We generate C-band photon pairs (1536.5 nm) using type-II SPDC and multiplex them into fibers using DWDM systems [International Telecommunication Union (ITU) channel 37] to co-propagate with 10 Gbit/s 1547.72 nm (ITU channel 51) classical data signals over 25 km of standard single-mode fiber (Corning SMF-28).

Frequency-domain filtering is studied with the use of tunable-bandwidth optical filters (EXFO XTA-50/U) with a full width at half maximum (FWHM) bandwidth range of 32–600 pm. Temporal filtering is applied by our defining a detection time window around the pulse arrival times during postprocessing of the count time stamps from our superconducting-nanowire single-photon detectors (SNSPDs; Quantum Opus) and time-to-digital converter (TDC; Swabian Timetagger Ultra). They have a root-mean-square timing jitter of 45 and 8 ps, respectively, enabling tight filtering. Each SNSPD has a dead time of less than 25 ns, an efficiency of more than 90%, and dark count rate of less than 100 counts/s, which means that the vast majority of noise comes from SpRS.

Polarization filtering is applied by our placing fiber polarizing beam splitters (FPBSs; Thorlabs) at the end of each fiber. Although Raman gain has some polarization dependence [74], wavelength-dependent birefringence and polarization mode dispersion over long-distance fiber cause SpRS photons to be approximately unpolarized [12,20,75], where the FPBSs reject approximately 50% of the noise exiting the fiber. To correct for time-dependent rotations of polarization in the fibers due to environmental disturbances, we implement active alignment for each polarization filter by maximizing single-photon counts using electronic fiber polarization controllers (Luna Inc.).

Photon pairs are generated by type-II SPDC in a periodically poled lithium niobate waveguide (PPLNWG; HC Photonics). Our source is pumped at a repetition rate of 200 MHz by our intensity-modulating a 1536.5 nm continuous wave (cw) laser (Pure Photonics PPCL550) using a 20 GHz bandwidth lithium niobate Mach-Zehnder modulator (MZM; Exail MXER-LN-20). The MZM is driven by amplified radio-frequency pulses originating from an arbitrary waveform generator (AWG; Tektronix AWG70002), which produces pulses with a temporal width of approximately 50 ps (FWHM). A 90:10 splitter is used to tap the pump light as input to a MZM bias controller (Exail) to stabilize the extinction ratio of the pulses. The pump pulses are then amplified with the use of an erbium-doped fiber amplifier (EDFA; Pritel PMFA-20) and injected into a PPLNWG to undergo second-harmonic generation (SHG), producing pulses centered on 768.25 nm. After filtering out of the 1536.5 nm pump, the 768.25 nm pulses drive type-II SPDC in a second PPLNWG (HC Photonics). SPDC generates a two-mode squeezed vacuum state [76], which is a probabilistic source of correlated photon number states in the Fock state basis with an imperfect but high single-photon fidelity in the low-gain regime. The photons in each mode (signal and idler) are orthogonally polarized and separated with the use of a FPBS. Each photon is subsequently multiplexed with the 10 Gbit/s C-band classical light with use of the pass and reflect ports of a 100 GHz DWDM system (ITU channel 37) to copropagate over 25 km SMF-28 fiber spools.

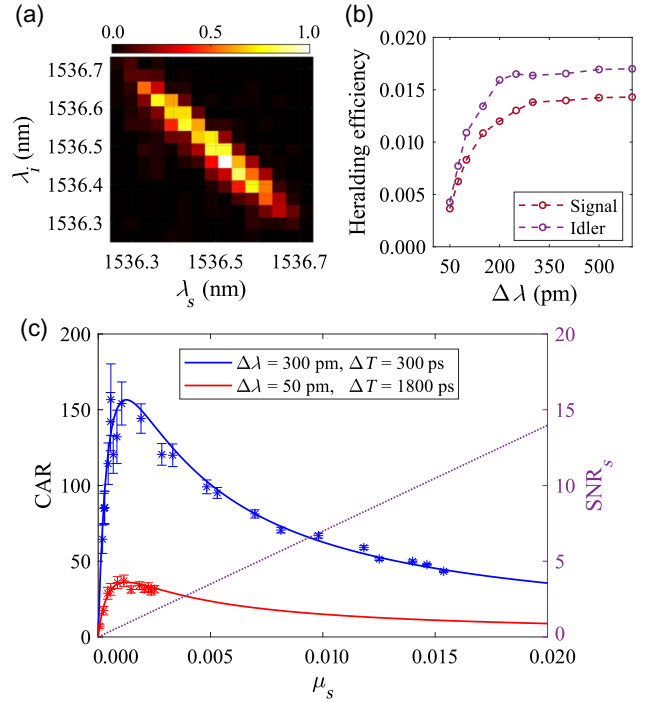


FIG. 2. (a) Measured JSI of the type-II SPDC photon-pair source. (b) Heralding efficiency as a function of filter bandwidth $\Delta\lambda$. (c) CAR as a function of μ_s for $\Delta\lambda\Delta T = 300$ pm $\times$ 300 ps or $\Delta\lambda\Delta T = 50$ pm $\times$ 1800 ps, which have different FHEs but the same background count probability since $\Delta\lambda\Delta T$ is fixed. The classical power is adjusted to receive $P_R = -27$ dBm after each 25 km fiber. The solid lines show the prediction based on our model. Even when each case has the same single-detector SNR (purple), the twofold CAR for $\Delta\lambda = 50$ pm is notably lower than that for $\Delta\lambda = 300$ pm.

The measured joint spectral intensity (JSI) of our source is shown in Fig. 2(a), and is prefiltered by a 100-GHz DWDM system and measured to a resolution of 32 pm. Figure 2(b) shows the measured heralding efficiency as we decrease the bandwidth of the filters for each photon with the classical sources OFF (negligible noise), which is defined as $\eta'_s = C/S_i$ ($\eta'_i = C/S_s$), where C and S_j are the coincidence and single-detector count rates, respectively. It is clear in Fig. 2(b) that η'_j is notably impacted as we decrease the filter bandwidth.

Each of the classical data signals is generated by small-form-factor pluggable transceivers (FS SFP-10G-40), which are controlled by field-programmable gate arrays programmed by QUANTUM INSTRUMENTATION CONTROL KIT [77]. The field-programmable gate arrays run an integrated bit error ratio test to measure error rates in each classical channel, and require a minimum received power P_R of -18 dBm to communicate with a sufficiently low bit error rate. The quantum and classical light experienced approximately 5 dB transmission loss over the 25 km fibers. The minimum launch power P_0 to operate each

classical link was -12.5 dBm when we account for a 0.5 dB loss in the demultiplexing DWDM systems. To test the limits of classical power and background noise, the classical signals are amplified by an EDFA before being multiplexed with the quantum signal. We filter the classical signals before transmission using DWDM systems to prevent 1536.5 nm amplified spontaneous emission from the laser and the EDFA from entering the 25 km fibers. At the quantum receivers, we cascaded multiple DWDM systems until negligible noise counts were observed when we directly injected 1547.72 nm light within the power range tested. The process described above ensures that the only appreciable noise comes from SpRS along the fibers and is within the passband of the 1536.5 nm filters.

To model how filtering affects single-fold and twofold detection rates, we note that an imperfect FHE implies that the mean numbers of photons that are generated at frequencies within the signal (μ_s), idler (μ_i), and overlap (μ_{both}) of the filter passbands are not equal. This is represented pictorially in Fig. 1(b). When filtering the JSA leads to $\mu_j > \mu_{\text{both}}$, the FHE is given by $\delta_s = \mu_{\text{both}}/\mu_i < 1$, $\delta_i = \mu_{\text{both}}/\mu_s < 1$. This means that even when the true signal/idler channel efficiencies are ideal ($\eta_s = \eta_i = 1$), not all detected single photons result in twofold coincidence events between single-photon pairs. However, SPDC and SFWM sources also emit multiphoton pairs, which leads to accidental coincidences or imperfect fidelity for entangled pair sources [78,79].

For an unfiltered JSA represented by $f(\omega_s, \omega_i)$ with $\iint d\omega_s d\omega_i |f(\omega_s, \omega_i)|^2 \equiv 1$, we define μ_T as the mean number of photon pairs per pulse generated across the unfiltered JSI. When filters with a normalized filter transmission function $g_j(\omega_j)$ are applied to each photon, we define the mean number of photons per pulse within the passbands of the signal/idler filters individually as (see Appendix B)

$$\mu_j = \mu_T \iint d\omega_s d\omega_i |f(\omega_s, \omega_i) g_j(\omega_j)|^2, \quad (1)$$

which quantifies the photons that can lead to signal (μ_s) or idler (μ_i) single-detector counts, and the number of photons within both filters

$$\mu_{\text{both}} = \mu_T \iint d\omega_s d\omega_i |f(\omega_s, \omega_i) g_s(\omega_s) g_i(\omega_i)|^2, \quad (2)$$

which quantifies the number of photons that can lead to “true” coincidence counts. However, multipairs populate the entire JSI, which becomes more problematic when $\mu_j > \mu_{\text{both}}$ because multipairs are detected at frequencies that do not contribute to single-pair coincidences, and will have a greater impact on fidelity compared with a system with $\mu_s = \mu_i = \mu_{\text{both}}$.

To model our experiment in the low-gain regime, we modify the models in Refs. [78–80] by substituting μ for

μ_s and μ_i [Eq. (1)] and including a filtered noise term to obtain the single-detector count probability per gate (see Appendix B)

$$S_j = \mu_j \eta_j + (\eta_{r_j} \alpha_{\text{pol}} \mathcal{R}_j \Delta\lambda_j + d_j) \Delta T_j \quad (3)$$

and replace μ by μ_{both} [Eq. (2)] to obtain the twofold coincidence probability

$$C \approx \mu_{\text{both}} \eta_s \eta_i + S_s S_i. \quad (4)$$

In the above, we represent an arbitrary noise photon source to be filtered by the parameter $\mathcal{R}_j$, which has the unit of $\text{nm}^{-1} \text{s}^{-1}$. Since noise photons are typically filtered to a narrow bandwidth ($\Delta\lambda < 1$ nm), we assume that the noise is approximately constant across the passbands of the frequency-domain and time-domain filters, which is an accurate approximation for broadband SpRS generated by an uncorrelated pump in long-distance fiber. $\Delta\lambda_j$ is the FWHM bandwidth of the spectral filter and d_j is the detector’s dark count rate per second. ΔT_j is the detection time window, which temporally filters both background noise and dark counts. We assume at present that ΔT_j is greater than the temporal width of the photon pairs (see Sec. V). The channel efficiency for each photon can be broken down as $\eta_j = \eta_{c_j} \eta_{\text{ch}_j} \eta_{r_j} \cos^2(\theta_{\text{pol}_j})$ where η_{c_j} , η_{ch_j} , and η_{r_j} account for the loss at the source, transmission channel, and receiver, respectively. Maximizing η_{c_j} is critical for our experiment, since minimizing the loss acting solely on the source photons will maximize the SNR in each detector. α_{pol} is a polarization filtering parameter that reduces the noise count rate, and is equal to $1/2$ for unpolarized noise. The term $\cos^2(\theta_{\text{pol}})$ represents the transmission of source photons through the polarization filter.

The accidental coincidence probability is modeled as $A = S_s S_i$, which is the same as in Refs. [78–80] but with the new definitions for μ_s and μ_i . We can then define the CAR to be C/A , which is the two-photon analogue of the single-photon SNR. When the background noise is negligible, the heralding efficiency including channel loss is defined as $\eta'_s = C/S_i \approx \delta_s \eta_s$, $\eta'_i = C/S_s \approx \delta_i \eta_i$, and $\delta_s = \mu_{\text{both}}/\mu_i$ are $\delta_i = \mu_{\text{both}}/\mu_s$ are the FHEs. When one is considering imperfect FHE of both photons, it is useful to use the “pair-symmetric heralding efficiency” (PSHE), which is defined as $\delta_{\text{PS}} \equiv \sqrt{\delta_s \delta_i}$ [57].

To confirm that imperfect FHE affects performance and how this differs from single-detector filtering, we first compare two scenarios with unequal FHEs but the same single-detector SNR. To do this, we evaluate two cases of spectrotemporal filtering, in which one has a wider spectral bandwidth but narrower temporal filtering ($\Delta\lambda_1 = 300$ pm, $\Delta T_1 = 300$ ps) and the other has a narrower bandwidth but wider temporal filtering ($\Delta\lambda_2 = 50$ pm, $\Delta T_2 = 1800$ ps), such that $\Delta\lambda_1 \Delta T_1 = \Delta\lambda_2 \Delta T_2$. For a single detector, the ratio

of the signal and noise counts is $S_j/\mathcal{N}_j = \mu_j \eta_j / (\eta_r \alpha_{\text{pol}} \mathcal{R}_j \Delta \lambda_j T_j)$, assuming dark counts are negligible. Since $S_j/\mathcal{N}_j \propto 1/\Delta \lambda \Delta T$, two different receivers with the same value for the product $\Delta \lambda \Delta T$ have an equivalent SNR whenever the μ_j values are equal. This also holds for the CAR when $\delta_{\text{PS}} = 1$. However, we predict that even if the SNR is the same, the CAR of a system with the lowest δ_{PS} will be more susceptible to noise. For our system, Fig. 2(b) shows that the FHE when $\Delta \lambda = 50$ pm is significantly lower than the FHE when $\Delta \lambda = 300$ pm, which implies that 50 pm filtering should have a lower CAR.

We first measure the CAR as a function of μ_s for 300 pm/300 ps filtering versus 50 pm/1800 ps filtering. For this test, the classical sources are attenuated to obtain a received power P_R of -27 dBm at the output of each 25 km fiber spool. This gives a noise count rate within $\Delta \lambda \Delta T$ of 5180.4 and 5050.5 counts/s for signal and idler detectors, respectively.

The results are shown in Fig. 2(c). For 300 pm/300 ps filtering we measured a maximum CAR of approximately 155, while 50 pm/1800 ps filtering results in a much lower maximum of approximately 40. This clear difference between the two cases unequivocally confirms that filtering the JSA has a notable impact. Since the CAR differed even when the SNR was the same, this also confirms a clear difference between filtering photon pairs compared with filtering single photons.

We also observe some consequences of balancing the SNR and degradation due to multipair emission during SPDC. Interestingly, the CAR can reach a high value even when the SNR is insufficiently low. We attribute this to the temporal and photon-number correlations for SPDC/SFWM, which are only benefited from in coincidence experiments. However, the CAR begins to decrease as the SNR increases. The maximum CAR occurs near $\mu_s \sim 10^{-3}$, but decreases as $\mu_s \rightarrow 0$ and $\mu_s \rightarrow \infty$. When the source's pump power is low, multipair emission is optimized at the expense of a lower SNR. When the SNR is increased, multiphoton emission begins to degrade the CAR and eventually becomes a dominant source of "noise." The maximum μ_s for 50 pm filtering in Fig. 2(c) was more limited than the maximum for 300 pm filtering due to the narrower bandwidth, which requires higher pump power to produce the same number of photons. However, our results show that simply increasing the source's pump power does not compensate for this when $\delta_{\text{PS}} < 1$, which is partially because each case is not affected equally by accidental coincidences due to multipair emission (see Sec. IV).

Instead of fixing the noise count probability, we now fix $\Delta T = 300$ ps and tune $\Delta \lambda$. We also increase the launch power P_0 of the classical signals to -12.5 dBm so that P_R is greater than the minimum receiver sensitivity of the 10 Gbit/s transceivers. This results in SpRS noise counts of 1477.4 and 1040.1 (counts/pm)/s for the signal and idler

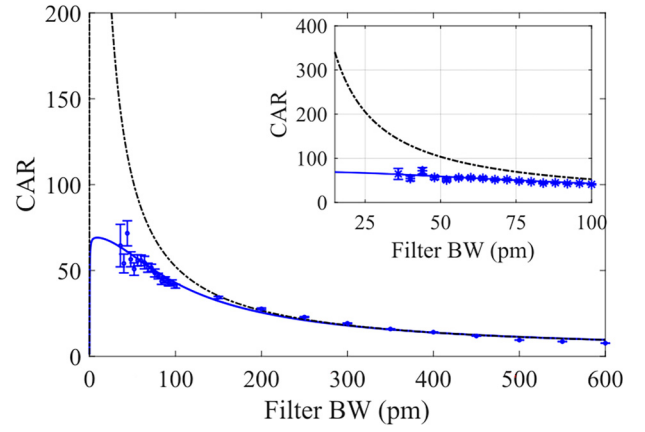


FIG. 3. CAR as a function of the spectral filter bandwidth (BW) at the signal/idler receivers ($\Delta \lambda_s, \Delta \lambda_i$) for a fixed SPDC pump power while each C-band classical source transmits $P_0 = -12.5$ dBm into each 25 km/25 km fiber. The black line shows the simulation for an ideal system, and the blue line shows the results for $\delta_j(\Delta \lambda_j) < 1$. The inset shows the results within the range of $\Delta \lambda_j = 20$ –100 pm.

detectors, respectively, where the difference is attributed mainly to unequal detection efficiencies. For this measurement, we fix the SPDC pump power to its maximum setting. We note that the SPDC pump power does not increase as $\Delta \lambda$ is varied, which does not create a perfect comparison, since the number of SPDC photons and the number of SpRS photons change with $\Delta \lambda$. Narrowing $\Delta \lambda$ therefore decreases multipair emission and noise count rates, but also does not increase the SNR. However, it allows us to observe some of the consequences of narrow filtering when ΔT is fixed, in contrast to the previous measurement. In Sec. IV, we simulate the case where we increase the SPDC pump power as $\Delta \lambda$ is decreased to maintain a fixed source intensity.

The results are shown in Fig. 3. We observe that the measured CAR (blue) begins to diverge significantly from the ideal model assuming a perfect FHE (dashed black line) for $\Delta \lambda < 200$ pm, which coincides with the trend in the change in heralding efficiency in Fig. 2(b). For example, a $\chi_{\text{ideal}}(32 \text{ pm}) = 155$ is predicted when $\delta_{\text{PS}} = 1$, while the measured result is $\chi_{\text{meas}}(32 \text{ pm}) = 71.3 \pm 5.0$. Since a decrease in FHE usually coincides with an increase in single-mode purity [57], this shows a potential trade-off between minimizing the impact of noise photons and maximizing the raw fidelity of applications requiring high-visibility HOM interference. We explore the consequences of this in more detail in Sec. IV.

III. C-BAND TIME-BIN ENTANGLEMENT COEXISTING WITH C-BAND CLASSICAL COMMUNICATIONS

We now investigate filtering for a concrete application—namely, entanglement-based networking in

long-distance fibers carrying classical communications. Although polarization entanglement [5,12,16,20,23] and time-energy entanglement [13] have been studied, to our knowledge the impact of SpRS on pulsed time-bin entanglement has not yet been investigated. Pulsed time-bin [81–83] encoding fundamentally requires a spectrally broadband pump, especially at higher repetition rates, meaning that narrow filtering can reasonably impact the FHE of many systems.

To generate entangled photons, we use the AWG to drive the MZM with two equal-amplitude radio-frequency pulses with a time separation of 347 ps at the same 200 MHz repetition rate. Since our 1536.5 nm cw laser has a coherence time of approximately 30 μ s, the pump is in a coherent superposition of being in “early” and “late” time bins: $|\psi\rangle_p = \frac{1}{\sqrt{2}}(|E_p\rangle + |L_p\rangle)$. After amplification by the EDFA and SHG at the first PPLNWG, the time-bin encoded 768.25 nm pump probabilistically generates single-photon pairs in the $|\Phi^+\rangle = \frac{1}{\sqrt{2}}(|E_s, E_i\rangle + |L_s, L_i\rangle)$ Bell state [84]. Michelson 1-bit delay interferometers are then inserted at each receiver to perform projective measurements on the time-bin Bloch sphere [71].

In this experiment, we fix the filtering to $\Delta\lambda_j = 50$ pm (6.5 GHz) and $\Delta T_j = 200$ ps for both photons. This means that each filter is close to being time-bandwidth limited and also gives a ratio between the temporal width of the filtered photon and the pump pulse of $\tau_{ph}/\tau_p \approx 1.5$ (see Appendix A), which implies that both photons are forward compatible with HOM interference-based applications.

However, the FHE is notably impacted, and we measured it to be 0.23 and 0.21 for signal and idler photons, respectively. The heralding efficiencies including loss were $\eta'_s = 0.0012$ and $\eta'_i = 0.0014$, which are lower than previous measurements due to the interferometers. At the expense of the tightly filtered background noise and increased spectral purity, this reduction is expected to hinder the ability to maintain fidelity compared to a source providing higher filter heralding efficiency at the same filter bandwidth.

Figure 4 shows the results of distributing the 1536.5 nm time-bin entangled pairs over 25 km spools to copropagate with the 10 Gbit/s C-band classical signals. Figure 4(a) shows a histogram of the signal detector’s single-photon counts per 10 ps bin width relative to the arrival of our clock at the TDC when $P_0 = -4.8$ dBm. We plot the source and noise counts individually to distinguish how the interferometer affects the arrival time of time-bin versus SpRS photons. For the time-bin photons, we observe pulses within three distinct time slots, with the middle bin (X and Y bases) having double the counts of the two side bins (Z basis) due to photons combining after taking the short and long paths through the interferometer. However, the noise has equal intensity across all time, indicating that SpRS from the classical source is generated randomly in time relative to the source. The noise intensity at the output of each interferometer port is 1/2 of the input, resulting from their randomly exiting the second beam splitter of the interferometer. The SpRS noise count

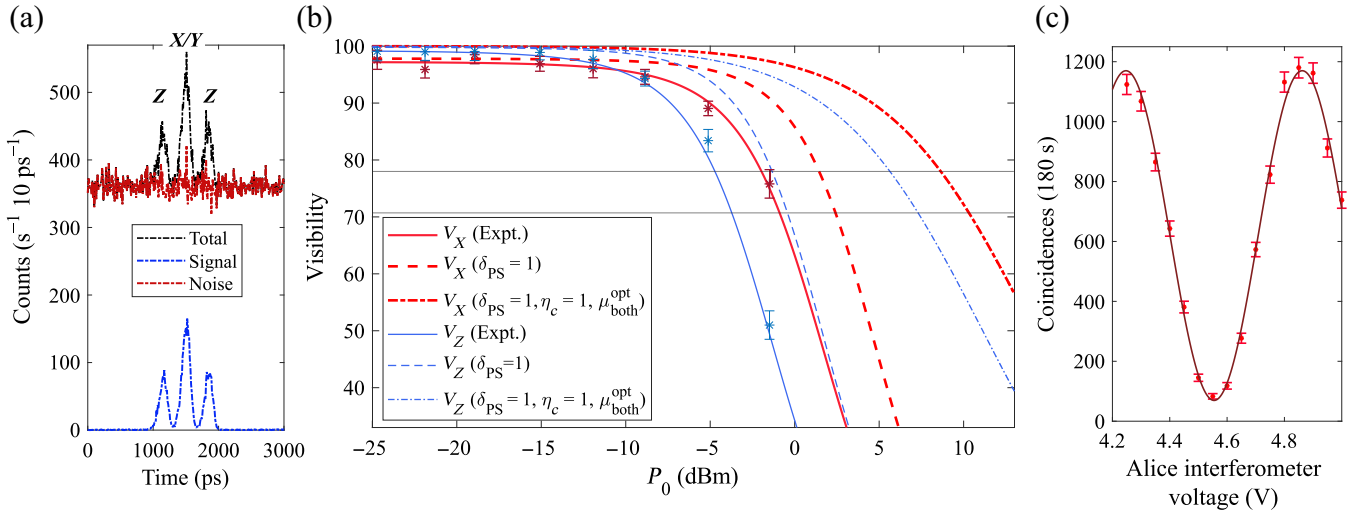


FIG. 4. Results for distributing 1536.5 nm time-bin entangled photons over 25 km/25 km fibers with a copropagating 10 Gbit/s C-band (1547.72 nm) classical data signal in each fiber. The filtering is set to $\Delta\lambda = 50$ pm and $\Delta T = 200$ ps for each photon. (a) Histogram of single counts relative to the clock for signal time-bin photons, SpRS noise, and total counts at the output of Alice’s interferometer when the classical launch power $P_0 = -4.8$ dBm. (b) Entanglement visibility in the X and Z bases vs P_0 . The color-coded solid lines show the predictions given our experimental parameters. The dashed lines are simulations when parameters such as ideal PSHE ($\delta_{PS} = 1$), source loss ($\eta_c = 1$), or the entangled pair source’s pump power at each P_0 value are optimized. The horizontal black lines show the limits for quantum key distribution and verifying nonlocality. (c) Two-photon interference fringe when $P_0 = -4.8$ dBm, where $V_X = (89.6 \pm 1.6)\%$.

rates increased linearly as a function of classical power, with 145 793.8 and 158 694.0 (counts/s)/mW in the signal and idler detectors, respectively.

Figure 4(b) shows how the quality of the time-bin entangled photons changes as we increase P_0 in each fiber. Non-local two-photon interference (TPI) fringes are observed in coincidence counts between the middle bins when we vary Alice's interferometer phase (ϕ_A) while fixing Bob's (ϕ_B) [Fig. 4(c)], which cycles through projective measurements along the equator of the Bloch sphere [71]. To quantify the degree of entanglement, we calculate $V = (R_{ll} - R_{lk}) / (R_{ll} + R_{lk})$ for each measurement, where R_{ll} and R_{lk} are the maximum and minimum coincidences, respectively. When the classical sources are OFF ($P_0 = 0$ mW), we measure $V_Z = (99.1 \pm 0.9)\%$ and $V_X = (97.1 \pm 1.6)\%$ in the Z and X bases, respectively, where the X basis is more degraded due to imperfections in the interferometers. The visibility without additional noise is imperfect mainly due to multipair emission during SPDC. For these measurements we kept a fixed SPDC pump power, which we estimated yielded $\mu_s = 1.0 \times 10^{-3}$, $\mu_i = 1.1 \times 10^{-3}$, and $\mu_{\text{both}} = 2.5 \times 10^{-4}$.

As we increase P_0 , Fig. 4(b) shows that our system is able to tolerate relatively high C-band classical power for a C-band/C-band quantum/classical wavelength allocation. In the X basis, we predict that $V > 78\%$ could be maintained at $P_0 = -1.7$ dBm and $V > 70.7\%$ could be maintained up to $P_0 = -0.4$ dBm, which are the limits for quantum key distribution [85] and verifying nonlocality [86], respectively. However, V_Z could maintain these values up to $P_0 = -4.6$ dBm and $P_0 = -3.5$ dBm, respectively. We attribute this to the unequal SNR of each bin in the output of the interferometers, as can be seen in the single-photon measurement in Fig. 4(a). This could be mitigated by measuring the Z basis in an alternative way, such as taking out the interferometer (at the expense of measuring X or Y) or using a switch or splitter to bypass the interferometer for Z -basis measurements (see Appendix C).

To investigate how imperfect FHE limited our results and how closer-to-optimal systems could be designed, we model time-bin entanglement generation, filtering, and measurement. We assume that the mean number of photons generated per bin is given by $\mu_j \equiv \mu_j^E = \mu_j^L$, which is required for the state to be maximally entangled. Building off the models in Refs. [78,87], we model the coincidence and singles count probabilities for the X and Y bases as (see Appendix C)

$$C^{xy}(\phi) = \frac{1}{4} \mu_{\text{both}} \eta_s \eta_i \left(\frac{1}{2} + \frac{1}{2} V_{\text{int}} \cos(\phi) \right) + S_s^{xy} S_i^{xy}, \quad (5)$$

$$S_j^{xy} = \frac{1}{2} \mu_j \eta_j + \left(\frac{1}{2} \eta_{r_j} \alpha_{\text{pol}} \mathcal{R}_j \Delta \lambda_j + d_j \right) \Delta T_j, \quad (6)$$

where $\phi \equiv \phi_A + \phi_B$ and V_{int} accounts for the visibility of the interferometer. The background noise term is $1/2$ the input due to splitting and recombining at the two interferometer beam splitters. The $1/4$ reduction in the coincidence probability is due to how the time-bin photons are distributed across the three bins of the two interferometer outputs [78]. The Z basis results in a $1/16$ reduction in coincidence rates, leading to a higher susceptibility to incident noise compared with the X and Y bases when only the output ports of the interferometers are used. In Appendix C, we investigate different measurement schemes and compare time-bin with polarization entanglement distribution in noisy environments.

In Fig. 4(b), we plot the predictions according to our experimental parameters (solid lines) and investigate the impact of optimizing different parameters (dashed lines). We find that our source was not generating enough photons to obtain the highest possible visibility by overcoming the high noise rates. We predict that the maximum classical power P_0 where we could still maintain $V_X > 70.7\%$ is 2.11 dBm. We also simulate our system assuming the same pump power but with perfect FHE to quantify how much the FHE impacted our measurements, which we predict could have tolerated P_0 up to approximately 3 dBm. Another parameter that can be optimized is the loss of photons at the source, since the SpRS photons do not experience this loss. In our system, photons can be lost due to coupling from the PPLNWG into fiber and insertion loss in the FPBS and DWDM systems, which totaled approximately 2.1 dB. Assuming ideal PSHE, source loss, and optimization of the pair generation probability for each P_0 , the simulation in Fig. 4(b) predicts that $V_{X/Y} > 70.7\%$ can be maintained with $P_0 \approx 9.8$ dBm in each fiber.

Even without further optimization, filtering in multiple degrees of freedom has enabled us to achieve notably higher C-band power compared with previous C-band entangled photon pair experiments. However, the simulations described above can guide the design of future quantum networks to have an even higher noise tolerance.

IV. CONSIDERATIONS IN THE DESIGN OF OPTIMAL SOURCES AND FILTERS

Having explored the effects of filtering for our particular experiment, we perform simulations to investigate other considerations in the design of sources and filters that can apply more generally.

We simulate the impact of the pump pulse width for photon pair generation, which will change the unfiltered JSA and the behavior of the FHE versus filter bandwidth. Flat-top spectral filters were also shown to provide a higher FHE compared with Gaussian-shaped filters [57], which indicates that the filter shape could impact performance in the context of this study. Furthermore, we investigate the role of multiphoton pair emission in each case to gain more

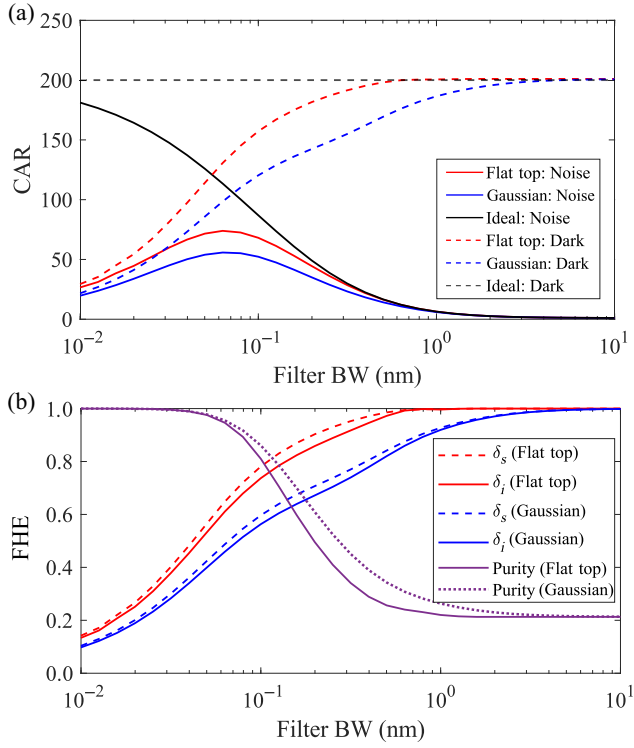


FIG. 5. Impact of spectral filter shape: (a) CAR as a function of filter bandwidth (BW) for either Gaussian or flat-top filter shapes when $P_0 = -12.5$ dBm. The dashed lines show the results excluding background noise. The SPDC source is adjusted to keep a constant $\mu_s(\Delta\lambda) = 0.005$. (b) Signal and idler filter heralding efficiencies for flat-top (red) and Gaussian (blue) filters and spectral purity for flat-top (solid line) and Gaussian (dashed line) filters.

insight into its impact in both noisy and noiseless environments. We estimate the single-mode purity for each case to compare the noise impact on twofold coincidence fidelity with the raw HOM visibility that could be achieved with the filtered photons.

In Fig. 5, we compare the effects of changing the bandwidth of Gaussian or flat-top filters in noisy and noise-free environments. We simulate the CAR of a photon-pair source transmitting 1536.5 nm photons over the same 25 km/25 km fiber spools with the 1547.72 nm classical signals set to $P_0 = -12.5$ dBm. Temporal filtering is fixed at $\Delta T = 300$ ps. In this simulation, we adjust the SPDC pump power so that $\mu_s(\Delta\lambda) = 0.005$ for each bandwidth setting such that the SNR increases as the filter is restricted, which is in contrast to our experimental measurement in Fig. 3. However, coincidence experiments have a more nuanced behavior. Figure 5(a) shows the CAR as we vary $\Delta\lambda$ for noisy (solid lines) and noiseless (dashed lines) scenarios. The black lines are the results for an ideal source that has a constant $\delta_j(\Delta\lambda) = 1$. Figure 5(b) shows a simulation of $\delta_j(\Delta\lambda)$ for each filter type and the estimated spectral purity.

Figure 5(a) demonstrates that flat-top filters achieve a higher CAR when background noise is or is not present when $\delta_{PS} < 1$. In the noisy case, we see that the filter bandwidth needs to be less than 1 nm to filter out enough noise for the CAR to increase to practical values; however, reduction of the bandwidth to less than this begins to decrease the FHE [Fig. 5(b)]. Interestingly, we see that the CAR begins to decrease at $\Delta\lambda < 0.05$ nm, even though the single-detector SNR continues to increase with the rejection of more SpRS noise. This is attributed to the increased sensitivity to multipair emission when the FHE is imperfect, which can be seen by the simulation that excludes background noise in Fig. 5(a) (dashed lines), which decreases according to the FHE in Fig. 5(b). To correct for this, one would need to reduce the SPDC pump power, which means that even in ideal environments, $\delta_{PS} < 1$ affects the coincidence rate and the CAR. In a noisy environment, attenuating the pump would mean reducing the SNR at each detector, showing a more nuanced trade-off between maximizing the SNR and minimizing degradation due to multipair emission (see Appendix B). The flat-top filters outperform the Gaussian filters mainly due to the higher FHE versus noise rejection. Importantly, flat-top filters achieve the highest CAR for a given spectral purity, which implies that they are closer to optimal for $N > 2$ -photon applications.

In Fig. 6 we show the CAR as a function of the source's pump pulse width τ_p , which changes how the FHE varies with the filter bandwidth [$\delta_j(\tau_p, \Delta\lambda_j)$] and the pump power, which varies the impact of multiphoton pair emission. This is similar to our measurements in Fig. 2(c), where the FHE was varied without the background noise being changed, and has implications for the design of high-rate pulsed sources that require short pump pulses and filtered photons. Figures 6(a)–6(c) show simulations of the CAR for a source pumped by transform-limited Gaussian pump pulses between $\tau_p = 2$ ps and $\tau_p = 200$ ps (FWHM) for both the noisy case (top) and the noise-free case (bottom) for $\Delta\lambda_j = 32$ pm, $\Delta\lambda_j = 100$ pm, or $\Delta\lambda_j = 500$ pm. We calculate the relevant parameters for each generated JSA using a Gaussian approximation developed in Refs. [57,88] for the pump, phase matching, and filter functions. We ignore any effects of dispersion to maintain a more general analysis.

Figure 6 clearly shows that performance is highly dependent on how the source is pumped both in intensity and in pump pulse width relative to $\Delta\lambda$. This is because $\delta_j(\tau_p, \Delta\lambda_j)$ decreases more rapidly as a function of $\Delta\lambda_j$ for short pulses compared with longer pulses [Fig. 6(d)], which is due to the broadening of the pump spectral bandwidth and thus the width of the unfiltered JSA. This means that as the pump is shortened, the FHE becomes problematic at wider filter bandwidths, wherein the SpRS noise photons are not as tightly filtered.

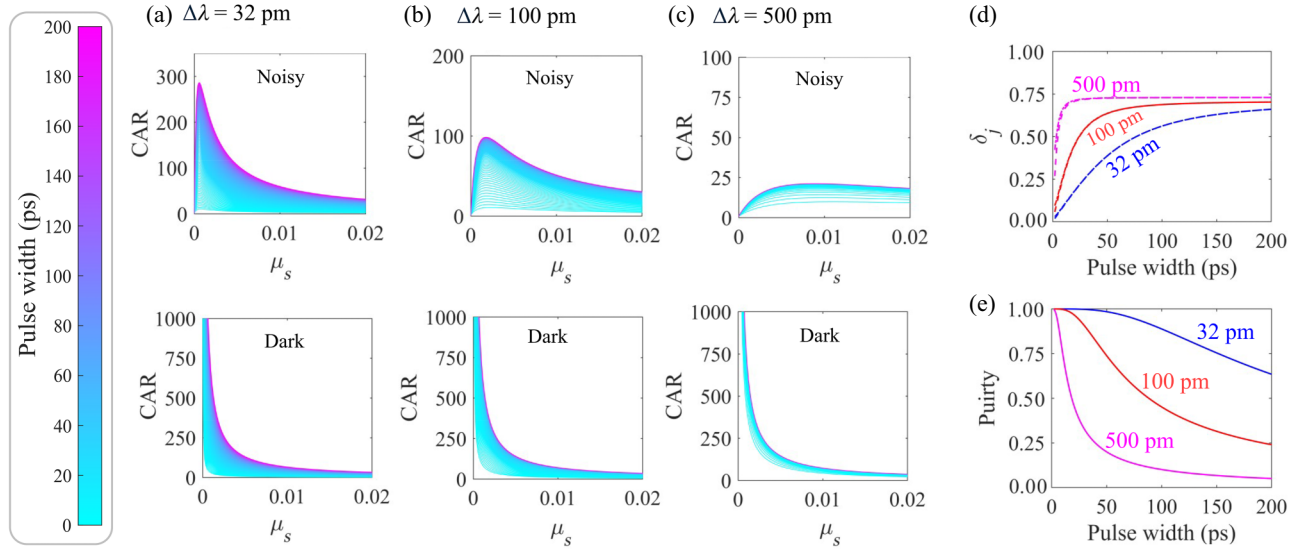


FIG. 6. Impact of the choice of pump pulse width τ_p and pair generation rate μ_s for differing filter bandwidths $\Delta\lambda$. We simulate both noisy (top) and background-noise-free (bottom) scenarios, where the noisy case has $P_0 = -12.5$ dBm in each 25 km fiber. (a)–(c) CAR vs μ_s for (a) $\Delta\lambda_s = \Delta\lambda_i = 32$ pm, (b) $\Delta\lambda_s = \Delta\lambda_i = 100$ pm, and (c) $\Delta\lambda_s = \Delta\lambda_i = 500$ pm, where $\Delta T_s = \Delta T_i = 300$ ps. Each simulation is color-coded according to the pump pulse width (left), ranging from $\tau_p = 2$ ps (teal) to $\tau_p = 200$ ps (purple). The CAR significantly depends on τ_p and $\Delta\lambda$, which is due to how the FHE $[\delta_j(\tau_p, \Delta\lambda_j)]$ changes on the basis of the relationship between these parameters [(d)]. (e) Estimated spectral purity vs τ_p for each $\Delta\lambda$ value to illustrate the compatibility of each $\Delta\lambda$ value with HOM interference-based systems.

Although this implies for this case study that optimal performance occurs when one increases τ_p while minimizing $\Delta\lambda$, this will notably reduce the spectral purity of the photons [Fig. 6(d)], in which case optimization of parameters depends on the application. This again shows a trade-off between the twofold fidelity and the optimality for applications such as state teleportation and entanglement swapping. We note that narrow pump pulses could allow tighter temporal filtering without cutting into the joint temporal correlations, which could change this comparison if we simulated a variable ΔT according to τ_p and the filter coherence time. However, for the narrower pulses, this is conditioned on having ultrahigh-resolution detectors.

In the noise-free simulations at the bottom of Figs. 6(a)–6(c), we see that the CAR as a function of the source intensity also depends on the relationship between τ_p and $\Delta\lambda$, with short pulses having a lower CAR for the same number of single photons. This again indicates the importance of mitigating accidental coincidences between higher-order Fock states occupying incompatible frequencies (see Appendix B).

An important conclusion from these results is that an imperfect FHE will cause a system to be more susceptible to noise that is independent of the filter bandwidth, such as intrinsic detector dark counts or pump leakage. For each plot in Figs. 6(a)–6(c), the background noise rate was kept constant, but the fidelity still varied according to the FHE. For example, this implies that an imperfect FHE

limits the maximum distance or attenuation a quantum signal can experience when limited by intrinsic detector noise.

To get a quantitative sense of exactly how much the CAR is affected by the FHE, we can derive a simple equation for the maximum CAR that can be achieved after optimization of the source's pump power to balance multipair emission and the SNR for a given δ_j , η_j , and $\mathcal{D}_j$, where $\mathcal{D}_j$ is a noise count probability representing an arbitrary background noise source. We find that maximum fidelity and $\mu_{\text{both}}^{\text{opt}}$ are both directly related to the PSHE, $\delta_{\text{PS}} \equiv \sqrt{\delta_s \delta_i}$. Maximization with respect to μ_{both} yields the following expression for the maximum CAR (see Appendix B):

$$\chi_{\text{max}} = \frac{\delta_{\text{PS}} \sqrt{\eta_s \eta_i \mathcal{D}_s \mathcal{D}_i}}{2 \mathcal{D}_s \mathcal{D}_i + \delta_{\text{PS}} \sqrt{\frac{\mathcal{D}_s \mathcal{D}_i}{\eta_s \eta_i} \left(\frac{\eta_s \mathcal{D}_i}{\delta_i} + \frac{\eta_i \mathcal{D}_s}{\delta_s} \right)}} + 1, \quad (7)$$

which is maximized when $\mu_{\text{both}}^{\text{opt}} = \delta_{\text{PS}} \sqrt{\mathcal{D}_s \mathcal{D}_i / \eta_s \eta_i}$. As an example, in the special case of $\delta_s \approx \delta_i \equiv \delta_{\text{PS}}$, both CAR_{max} and $\mu_{\text{both}}^{\text{opt}}$ simply decrease linearly with δ_{PS} . If the FHE of each photon is measured as $\delta_{\text{PS}} \approx 0.2$, for example, then both the maximum CAR and the coincidence rate are 20% of that which would be obtained with an ideal PSHE. For an entangled source with a maximum TPI visibility of 96% when $\delta_{\text{PS}} = 1$, $\delta_{\text{PS}} = 0.2$ would limit visibility to only 81.5%. This clearly shows that careful attention to the

design of sources and filters can have a significant impact on the achievable rates and fidelity in noisy environments.

V. DISCUSSION AND OUTLOOK

In summary, we have studied various methods of filtering photon pairs generated from SPDC and SFWM sources in both noisy and ideal conditions, which has broad implications for optimal source and receiver designs for many multiphoton quantum applications. We further conducted the a study of pulsed time-bin entanglement distribution coexisting with classical communication in fiber, discussing some special considerations for optimal sources, filters, and qubit measurement systems. Using narrow filtering, we demonstrated the copropagation of C-band time-bin entangled photons and 10 Gbit/s C-band classical data over a total of 50 km of standard fiber and showed that our system could tolerate milliwatt-level C-band power. Although immediately applicable to entanglement-based applications, the study of twofold fidelity alongside single-mode purity provides valuable insight into the consequences of filtering for more complex teleportation-based applications.

Regarding the design of optimal sources for $N > 2$ -photon applications, it is important to maximize high single-mode purity, PSHE, and noise rejection simultaneously. On the basis of the results in this paper, the optimal pair source for mitigating source or background noise (e.g., SpRS or sunlight) would produce spectrally pure photons without requiring external filtering of the JSA, enabling high-visibility interference between independent sources, but both signal *and* idler photons are narrow band such that noise photons can be narrowly filtered in each arm without affecting those produced at the source. For bandwidth-independent noise (detector dark counts, pump leakage, etc.), the main concerns are maximizing the PSHE and purity. Most studies attempting to resolve the trade-off between FHE and purity focus on engineering sources to avoid external filtering [50,54,58,61–70]. Some approaches can reasonably satisfy the above requirements for high-noise operation; however, some sources may not be as optimal when considering both photons operating in noisy conditions. For example, one approach is to engineer a nonlinear medium such that the group velocity of the signal or idler modes matches the pump during photon pair generation. Closer-to-ideal noise rejection could be achieved if the JSA were circular and narrow, but some of these source designs can yield an unfiltered JSA such that one photon is spectrally narrow and the other photon is more broadband. Because of this asymmetry, the broader photon's filter would either have to pass more noise photons or filter its spectrum, which would decrease the PSHE. Generally, narrowly filtering both photons while maintaining high PSHE and purity is a challenge [57]. Another approach is to have the nonlinear medium placed within

a properly designed cavity such that photon pairs can be generated within ultranarrow bandwidths [64,67,73], in which case isolating one of the cavity modes could, in principle, enable the maintenance of high PSHE while filtering noise on the order of approximately 100 MHz. We note that significant progress has been made recently with appropriately designed integrated photonic sources [50,68,69], which can, in principle, also have narrow bandwidths and PSHE. Beyond pair generation using SPDC or SFWM, atomlike entangled-photon sources could be an alternative means to satisfy the requirements of optimized source engineering, where quantum dot sources have recently been investigated to enable tight filtering in noisy free-space quantum communications [32].

In this study, our focus was mainly on optimizing system designs for discrete-variable multiphoton applications employing incoherent filters. Because of its simple implementation and relatively low cost, incoherent filtering has to date been the most common noise-mitigation method in quantum-classical coexistence, free-space communication, and noisy devices. However, coherent filtering can offer much greater noise rejection and can potentially outperform optimal source engineering for multiphoton experiments. Coherent filtering includes detectors that selectively pass only a particular set of spectrotemporal modes [41,43,89]. For example, quantum frequency conversion with a tailored pump pulse can selectively convert the frequency of only a single temporal mode, leaving all others unchanged [41,89]. Single-mode filtering at all detectors could simultaneously maximize PSHE, single-mode purity, and noise rejection, assuming ideal mode-selection efficiency. We note that similar concepts apply to the measurement of continuous-variable quantum states, which has been shown to significantly mitigate the impact of SpRS noise from classical communications compared with discrete-variable systems [90].

For time-domain filtering, we note that generating or filtering photons to ultranarrow spectral bandwidths limits the pump repetition rate and, importantly, governs how narrow a time-domain filter can be before cutting into the joint temporal amplitude or intensity. Instead of or in combination with short-duration pump pulses, independent photons can also interfere by use of a detection time resolution that is less than the coherence time of the photons [48,91]. Using Fourier transformations, our model could be generalized to account for narrow frequency-time filtering, which will be explored in future studies. Beyond this, the model could be expanded to more accurately predict the role of multipair emission in the high-gain regime, a multimode analysis, and experiments involving more than two photons. Our results imply that for multiple pair sources that each have the same PSHE, N -fold coincidences would scale as $C_N \propto \delta_{\text{PS}}^{N/2}$ with visibility negatively affected by multipair emission according to $\mu_s, \mu_i > \mu_{\text{both}}$ for each source.

From the perspective of advancing the field of coexisting quantum and classical fiber networks, the power levels achieved in this study indicate that aggregate C-band classical data rates exceeding terabits per second could be communicated alongside each entangled photon by spreading the power across multiple attenuated DWDM channels and using transceivers with a higher receiver sensitivity [21]. SpRS could be more suppressed with the use of different C-band classical wavelengths, since our 1536.5 nm/1547.72 nm channel selection is within a Raman gain peak in the C-band SpRS spectrum [8,21]. However, this would require attenuating classical signals to levels lower than that typically used in such high-rate classical communications. We note that other noise sources beyond SpRS (e.g., four-wave mixing) may also need to be considered when more DWDM channels are included. Importantly, our results apply to any wavelength scheme for wavelength-division multiplexing quantum-classical integration, such as systems using O-band quantum or classical sources. For example, the use of wavelength engineering to obtain weaker SpRS could enable wider band filtering, which could increase the PSHE or performance of high-rate pulsed sources. The experimental design also provides valuable information on how networks of erbium-ion quantum memories could potentially be impacted by classical signals due to the approximately 1536 nm center wavelength of both entangled photons [72,73]. On the basis of our results, it is promising that these networks could feasibly operate alongside high-rate classical communications, but further work is needed to evaluate these challenges due to the strong SpRS observed here and the unique physics.

Although we used quantum-classical networking as an experimental case study, we believe that our results will be of general interest for other noisy environments such as ambient light in free-space quantum applications [28–33] or noise due to SpRS from the pump in SFWM-based photon-pair sources [34,35], quantum frequency conversion devices [36,37], and all-optical switches or shutters [38–40]. Furthermore, the analysis of multipair emission when the PSHE is imperfect could be useful for general experiments regardless of additional background noise.

ACKNOWLEDGMENTS

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

APPENDIX A: CHARACTERIZATION OF FREQUENCY-DOMAIN AND TIME-DOMAIN FILTERS

Here we characterize the properties of our tunable filter passbands used in the experiment for spectral filtering. Figure 7 shows the normalized filter transmission of various bandwidth settings used in the experiment. We see that the filters are approximately flat-top with a slight slope at wider bandwidths. At narrow bandwidths they approach a more Gaussian-like shape due to the slope of the edge of the filter shape. We also characterize the filtered photons in the time domain. We plot the broadening of the photon's coherence time relative to the source's pump width as a function of reducing $\Delta\lambda$. We first measure time-domain traces of the signal photon's temporal shape for filter settings ranging from $\Delta\lambda = 32$ pm to $\Delta\lambda = 600$ pm [Fig. 8(a)]. We then attenuate the pump pulse that drives SPDC to the single-photon level and directly send this pulse to the same SNSPD. It is clear from Fig. 8(a) that as the filter is restricted, the SPDC photons begin to broaden more widely than the pump, and approach the pump width as the filter is widened. Figure 8(b) shows the ratio between the FWHM of the filtered photons relative to the measured pump pulse. This is the characteristic property needed to cause photons from independent SPDC or SFWM sources to interfere [45]. Figure 8(a) shows that for $\Delta\lambda = 50$ pm, which is what we used for the time-bin entanglement study, a 200 ps coincidence window is just narrow enough to tightly filter noise without cutting into the joint temporal correlations or source counts.

APPENDIX B: IMPACT OF MULTIPHOTON PAIR EMISSION AFTER SPECTRAL FILTERING

To model these experiments, it is critical to include the effects of higher-order Fock state emission during SPDC or SFWM. The spontaneous parametric processes of SPDC or SFWM generate a two-mode squeezed vacuum

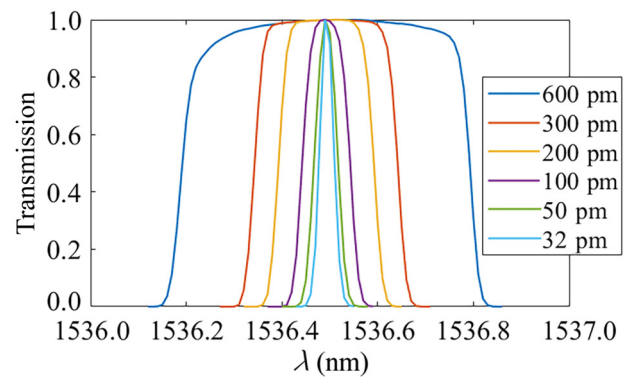


FIG. 7. Shape of the tunable-bandwidth filters as a function of bandwidth.

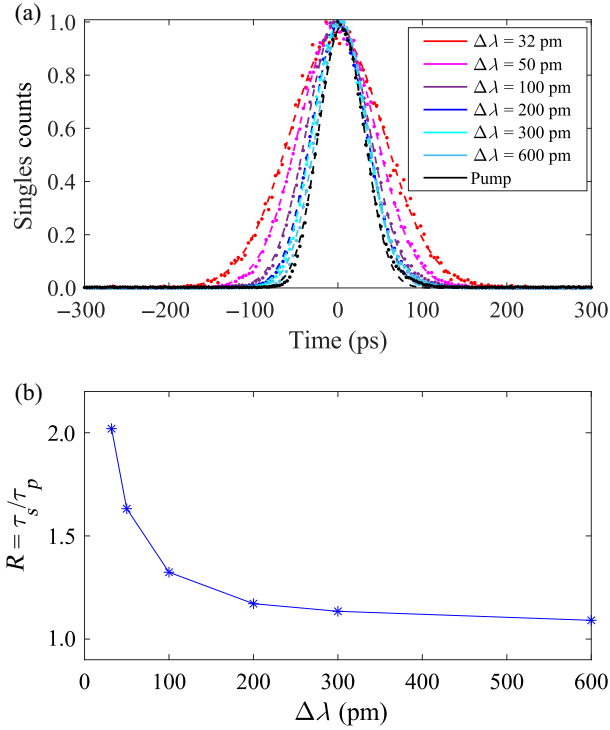


FIG. 8. (a) Normalized singles counts relative to the clock arrival time for different settings of $\Delta\lambda$. The black curve shows the 768.25 nm SPDC pump pulse when attenuated to the single-photon level and injected into the same SNSPD. (b) Ratio of the signal photon's temporal width relative to the pump pulse's temporal width ($R = \tau_s/\tau_p$) vs $\Delta\lambda$.

(TMSV) state [92]. In the Fock state basis, a single-mode TMSV source has the quantum state (before loss) of $|\xi\rangle_q = \sum_{m=0}^{\infty} (-1)^m \sqrt{\mu^m / (1 + \mu)^{m+1}} |m_s, m_i\rangle$, where m is the number of photons in signal and idler modes, which follows a thermal photon number distribution as a function of the mean number of photon pairs per pulse μ produced in the medium: $P_{th}(\mu, m) = \mu^m / (1 + \mu)^{m+1}$ [92]. For our purposes, only the $m = 1$ term is desired, and $m > 1$ leads to accidental coincidence counts or nonzero coincidences in an entangled photon pair's TPI fringe [78,79]. Many sources, especially when not filtered, have many modes emitting μ_k mean number of photons per mode. A multimode source with M modes has a photon number distribution ranging from thermal ($M = 1$) to Poisson ($M \rightarrow \infty$) as a function of $\mu = \sum \mu_k$, since it is a convolution of M thermal sources. However, a Poisson model is a good approximation to predict the impact of multiphoton emission when $\mu \ll 1$ for any number of modes [78], which significantly simplifies modeling. Since we operate at $\mu \ll 1$ in our experiment, we build on this approximation and leave a full multimode analysis in the high-gain regime to future work.

When μ , the efficiency in signal/idler channels η_j , and the total background noise count probability $\mathcal{D}_j$ are all

much less than 1, the single-detector (S_j), coincidence (C), and accidental coincidence (A) count probabilities per detection window for a SPDC or SFWM experiment are often modeled as [78–80]

$$S_j = \mu \eta_j + \mathcal{D}_j, \quad (\text{B1})$$

$$C = \mu \eta_s \eta_i + S_s S_i, \quad (\text{B2})$$

$$A = S_s S_i. \quad (\text{B3})$$

However, these equations ignore the properties of the source's unfiltered JSA, $f(\omega_s, \omega_i)$, and the impact of external filtering. The approach in this paper is to note that imperfect FHE means that the numbers of photons that can lead to signal, idler, and coincidence detections are unequal.

To include this, we build on Refs. [57,60] and propose the definitions for the mean number of photons that pass signal ($j = s$) or idler ($j = i$) filters as

$$\mu_j = \mu_T \iint d\omega_s d\omega_i |f(\omega_s, \omega_i) g_j(\omega_j)|^2 \quad (\text{B4})$$

and the number of photons that pass both filters as

$$\mu_{\text{both}} = \mu_T \iint d\omega_s d\omega_i |f(\omega_s, \omega_i) g_s(\omega_s) g_i(\omega_i)|^2, \quad (\text{B5})$$

where μ_T is the mean number of photons generated when the source is not filtered and $g_j(\omega_j)$ are functions representing the signal/idler normalized filter transmission functions. Using these definitions, we modify the commonly used equations to include the effects of filtering a source's joint spectrum.

The twofold coincidence probability per pulse is modeled by our substituting μ_{both} for μ in Eq. (B2) to obtain

$$C \approx \mu_{\text{both}} \eta_s \eta_i + S_s S_i, \quad (\text{B6})$$

where the single-detector count probability is obtained by our substituting μ_s and μ_i for μ in Eq. (B1) to obtain

$$S_s = \mu_s \eta_s + \mathcal{D}_s, \quad (\text{B7})$$

$$S_i = \mu_i \eta_i + \mathcal{D}_i, \quad (\text{B8})$$

where $\mathcal{D}_j$ is the total background noise count probability, which we modeled in more detail in the main text to account for the physics of our specific experiment.

Accidental coincidences are modeled similarly to Eq. (B3) but with the new definitions for S_s and S_i , which is given by the product of signal and idler single-photon probabilities:

$$A = S_s S_i. \quad (\text{B9})$$

When $\mu_j = \mu_{\text{both}}$, the FHE is ideal, which gives the typical equations for the probability of coincidence and

singles counts $C = \mu\eta_s\eta_i + S_sS_i$ and $S_j = \mu\eta_j + \mathcal{D}_j$ in Refs. [78–80].

In a noisy environment, the CAR is given by

$$\chi = \frac{\mu_{\text{both}}\eta_s\eta_i}{(\mu_s\eta_s + \mathcal{D}_s)(\mu_i\eta_i + \mathcal{D}_i)} + 1. \quad (\text{B10})$$

It can be useful to arrange this using only one parameter to account for the source intensity. Since $\mu_i = \mu_{\text{both}}/\delta_s$ and $\mu_s = \mu_{\text{both}}/\delta_i$, we substitute these into Eq. (B10) to obtain

$$\chi = \frac{\mu_{\text{both}}\eta_s\eta_i}{(\delta_i^{-1}\mu_{\text{both}}\eta_s + \mathcal{D}_s)(\delta_s^{-1}\mu_{\text{both}}\eta_i + \mathcal{D}_i)} + 1. \quad (\text{B11})$$

Taking the derivative of this function with respect to μ_{both} , we can find the mean photon number that maximizes the CAR by balancing increases to the SNR when increasing μ_T with the inevitable degradation of the CAR caused by multipair emission.

Meyer-Scott *et al.* [57] proposed a parameter called the “pair-symmetric filter heralding efficiency”, $\delta_{\text{PS}} \equiv \sqrt{\delta_s\delta_i}$, which quantifies the impact of the FHE of both photons. Using this definition, we can express the optimal μ_{both} that maximizes the CAR as

$$\mu_{\text{both}}^{\text{opt}} = \delta_{\text{PS}} \sqrt{\frac{\mathcal{D}_s\mathcal{D}_i}{\eta_s\eta_i}}. \quad (\text{B12})$$

Substituting this into Eq. (B11), we obtain the maximum achievable CAR given a particular channel efficiency, background rate, and FHE:

$$\chi_{\text{max}} = \frac{\delta_{\text{PS}}\sqrt{\eta_s\eta_i\mathcal{D}_s\mathcal{D}_i}}{2\mathcal{D}_s\mathcal{D}_i + \delta_{\text{PS}}\sqrt{\frac{\mathcal{D}_s\mathcal{D}_i}{\eta_s\eta_i}}\left(\frac{\eta_s\mathcal{D}_i}{\delta_i} + \frac{\eta_i\mathcal{D}_s}{\delta_s}\right)} + 1. \quad (\text{B13})$$

Thus, the maximum CAR that can be achieved clearly depends on δ_{PS} , and the coincidence rate at which the CAR is maximized also decreases linearly with the PSHE.

A common method for using experimental data to obtain or predict key parameters for photon pair experiments is to use the experimental result for the CAR in noiseless environments, CAR_{dark} , to obtain the mean photon number generated in the source. This is often done with the equation $\chi_{\text{dark}} \approx 1 + 1/\mu$, which when rearranged gives $\mu \approx 1/(\chi_{\text{dark}} - 1)$ [78]. One can also predict the CAR as a function of μ using this simple relationship.

However, when $\delta_{\text{PS}} < 1$, this is no longer valid. In the case of imperfect FHE in the low-gain regime, we find that the CAR in noise-free environments is related to the pair

generation probability by the equation

$$\chi_{\text{dark}} = \frac{\mu_{\text{both}}}{\mu_s\mu_i} + 1. \quad (\text{B14})$$

Using the relationships $\mu_i = \mu_{\text{both}}/\delta_s$ and $\mu_s = \mu_{\text{both}}/\delta_i$, we can rewrite this equation as

$$\chi_{\text{dark}} = \frac{\delta_s\delta_i}{\mu_{\text{both}}} + 1, \quad (\text{B15})$$

which can be used to estimate μ_{both} with the use of experimentally measured parameters for the CAR and FHEs. Equation (B15) can be useful to predict the CAR based on the pair generation probability. This shows that for a given coincidence rate (proportional to μ_{both}) the CAR is reduced by a factor of δ_{PS}^2 compared with the ideal case.

1. Equal heralding efficiency approximation

If we assume that the FHEs for each photon are approximately equal ($\delta_s = \delta_i \equiv \delta_{\text{PS}}$), which is not universally true and depends on the source’s phase-matching angle and bandwidth [57] or the similarity between signal and idler filters, this implies that $\mu_s \approx \mu_i \equiv \mu_{\text{si}}$. This leads to the relationship $\mu_{\text{both}} \approx \delta_{\text{PS}}\mu_{\text{si}}$. Substituting this into the model, we obtain the following equations:

$$C \approx \delta_{\text{PS}}\mu_{\text{si}}\eta_s\eta_i + S_sS_i, \quad (\text{B16})$$

$$S_j \approx \mu_{\text{si}}\eta_j + \mathcal{D}_j. \quad (\text{B17})$$

In this approximation, the CAR excluding any background noise simplifies to

$$\begin{aligned} \chi_{\text{dark}} &\approx \frac{\delta_{\text{PS}}}{\mu_{\text{si}}} + 1 \\ &\approx \frac{\delta_{\text{PS}}^2}{\mu_{\text{both}}} + 1, \end{aligned} \quad (\text{B18})$$

which means that one can use the measured CAR to estimate the experimental parameters:

$$\mu_{\text{si}} \approx \frac{\delta_{\text{PS}}}{\chi_{\text{dark}} - 1}, \quad (\text{B19})$$

$$\mu_{\text{both}} \approx \frac{\delta_{\text{PS}}^2}{\chi_{\text{dark}} - 1}. \quad (\text{B20})$$

Equation (B19) resembles the equation derived in Ref. [60]; however, our work expands this approach to more accurately quantify the number of photons that lead to singles and coincidence counts.

With use of the equal-FHE approximation, the equation for the CAR accounting for background noise is given by

$$\chi = \frac{\delta_{\text{PS}} \mu_{si} \eta_s \eta_i}{(\mu_{si} \eta_s + \mathcal{D}_s)(\mu_{si} \eta_i + \mathcal{D}_i)} + 1, \quad (\text{B21})$$

and the value of μ_{si} that maximizes this function is given by

$$\mu_{si}^{\text{opt}} = \sqrt{\frac{\mathcal{D}_s \mathcal{D}_i}{\eta_s \eta_i}}. \quad (\text{B22})$$

Interestingly, μ_{si}^{opt} does not depend on δ_{PS} , which means that the optimal number of signal/idler photons is roughly the same regardless of the value of the FHE.

However, the maximum CAR that can be achieved for a given η_j , $\mathcal{D}_j$, and δ_{PS} is given by

$$\chi_{\text{max}}(\delta_{\text{PS}}) = \frac{\delta_{\text{PS}} \sqrt{\eta_s \eta_i \mathcal{D}_s \mathcal{D}_i}}{2\mathcal{D}_s \mathcal{D}_i + \sqrt{\frac{\mathcal{D}_s \mathcal{D}_i}{\eta_s \eta_i}} (\mathcal{D}_s \eta_i + \mathcal{D}_i \eta_s)} + 1. \quad (\text{B23})$$

The optimal $\mu_{\text{both}}^{\text{opt}}$ to maximize the CAR is still given by $\mu_{\text{both}}^{\text{opt}} = \delta_{\text{PS}} \sqrt{\frac{\mathcal{D}_s \mathcal{D}_i}{\eta_s \eta_i}}$. Thus, the single count rate that optimizes the CAR is independent of the PSHE, but the maximum CAR and the single-pair coincidence rate at which it is obtained are reduced by a factor of δ_{PS} .

In Fig. 9, we use the approximation $\delta_{\text{PS}} = \delta_s = \delta_i$ and plot the CAR as a function of μ_{both} for different values

of δ_{si} for dark [Fig. 9(a)] and noisy [Fig. 9(b)] environments. We analyze an example case with true channel efficiencies $\eta_s = \eta_i = 0.1$ and arbitrary background noise count probabilities $\mathcal{D}_s = \mathcal{D}_i = 10^{-4}$. We also plot in Fig. 9(c) the single-pair coincidence probability as a function of the number of signal/idler photons $C_{\text{pair}} = \delta_{\text{PS}} \mu_{si} \eta_s \eta_i$. In Fig. 9(d), we plot the purity P of the photons as a function of δ_{PS} using a Gaussian approximation for the filters, pump, and phase-matching function [57] (see Appendix D).

As δ_{PS} decreases, the source becomes generally much more susceptible to accidental coincidences caused by multiphoton emission at incompatible frequencies for single-pair coincidence detection [Fig. 9(a)]. Compared with the ideal $\delta_{\text{PS}} = 1$, the CAR is reduced approximately by a factor of δ_{PS} for the same single-photon rate, and by δ_{PS}^2 for the same coincidence rate. Thus, not only does $\delta_{\text{PS}} < 1$ reduce the coincidence rate by δ_{PS} for the same number of detected single photons, but it also reduces the CAR. To achieve the same CAR as for $\delta_{\text{PS}} = 1$, one needs to reduce the pump power of the source until the true single-pair coincidence is decreased by a factor of δ_{PS}^2 compared with the ideal case, resulting in an even greater reduction in coincidence rates when a desired target fidelity is factored in.

Figure 9(b) highlights the consequences of this effect in noisy environments. As δ_{PS} decreases, the CAR decreases for all values of μ_{both} , even though noise and loss are unchanged. We can also see that as δ_{PS} is reduced, μ_{both} must be further reduced to obtain the maximum CAR. This is similar to the experiment in Fig. 2(c) and the simulations

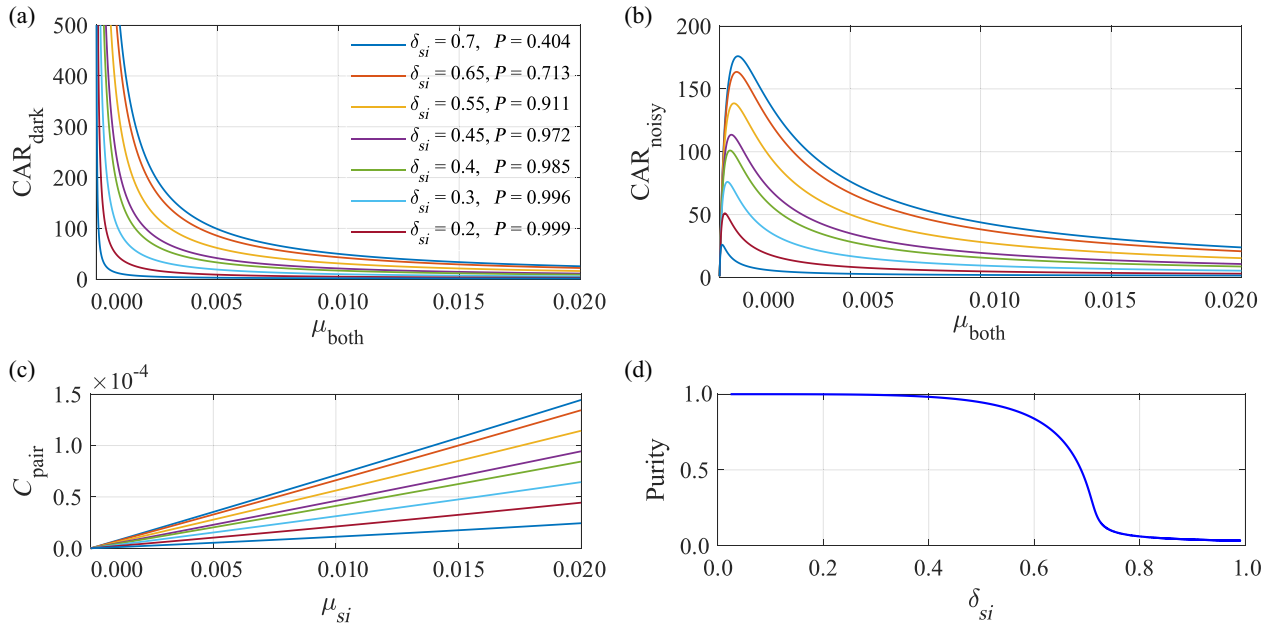


FIG. 9. (a) CAR vs μ_{both} in a noise-free environment for different values of $\delta_{si} \equiv \delta_s = \delta_i$. (b) CAR when the signal/idler background count probability $\mathcal{D}_j = 10^{-4}$ and $\eta_j = 0.1$. (c) Single-pair coincidence probability vs the number of generated signal/idler photons, $C_{\text{pair}} = \delta_{\text{PS}} \mu_{si} \eta_s \eta_i$, for different values of δ_{si} . (d) Simulated single-mode purity as a function of δ_{si} .

in Fig. 6, since $\mathcal{D}_j$ is the same but δ_{PS} is varied. Importantly, this is also similar to the impact of the intrinsic detector dark count noise, which does not change with either the pump or filter properties, but its impact depends on δ_{PS} . This shows, for example, that reduced heralding efficiency can limit long-distance quantum communications excluding additional noise from coexisting classical signals.

Figure 9(d) shows the trade-off between fidelity due to filtering and single-mode spectral purity, which is directly related to the HOM visibility for independent sources of genuine single photons [88]. Clearly, the impact of multipair emission as purity increases shows that maintaining high rates and high raw fidelity can be challenging, which is an important consideration when one is designing N -photon applications.

APPENDIX C: TIME-BIN VERSUS POLARIZATION ENTANGLEMENT DISTRIBUTION IN NOISY ENVIRONMENTS

Here we explore some differences between time-bin encoding and polarization encoding in noisy environments, including how time-bin qubits are measured.

1. Basis dependence in measuring time-bin qubits in high-noise environments

a. X and Y bases

For time-bin entanglement encoding where the X and Y bases are measured with use of a 1-bit delay interferometer [71], the maximum and minimum coincidence counts based on the model in the main text for the middle bin are given by

$$C_{\text{max}}^{\text{xy}} = \frac{1}{4} \mu_{\text{both}} \eta_s \eta_i \left(\frac{1}{2} + \frac{1}{2} V_{\text{int}} \right) + S_s^{\text{xy}} S_i^{\text{xy}}, \quad (\text{C1})$$

$$C_{\text{min}}^{\text{xy}} = S_s^{\text{xy}} S_i^{\text{xy}} + \frac{1}{4} \mu_{\text{both}} \eta_s \eta_i \left(\frac{1}{2} - \frac{1}{2} V_{\text{int}} \right), \quad (\text{C2})$$

with the singles count probabilities

$$S_j^{\text{xy}} = \frac{1}{2} \mu_j \eta_j + \left(\frac{1}{2} \eta_{r_j} \alpha_{\text{pol}} \mathcal{R}_j \Delta \lambda_j + d_j \right) \Delta T_j. \quad (\text{C3})$$

Compared with a correlated pair source, the coincidence rate is reduced by 1/4. This is because both signal and idler photons split and recombine at the first and second beam splitters of the interferometer, but half of the counts are lost to the two side bins. Because of the random arrival time for the SpRS photons, the noise is also reduced by 50% in each detector, since photons will randomly split and recombine at each splitter. This means that at the output, the noise is equal in all time intervals but with 50% of the input intensity. For unpolarized noise, applying polarization can then reduce the noise by another 50% since $\alpha_{\text{pol}} = 0.5$.

b. Optimal and nonoptimal measurements in the Z basis

In our experiment, where the Z basis is measured with use of the first and third bins of the interferometer output, the single counts are 1/4 of a correlated pair source, since these photons split at both beam splitters in the interferometer but do not recombine as for the middle bin. In this case, the coincidence probability is 1/16 of that of a correlated pair source. However, the noise photon count rate is equal across all three bins, indicating that the first and third bins are more impacted by noise than the X and Y bases with this Z -basis measurement scheme.

We model the maximum ($C_{t_1 t_1}$ or $C_{t_3 t_3}$) and minimum ($C_{t_1 t_3}$ or $C_{t_3 t_1}$) coincidences for the Z basis by the equations

$$C_{\text{max}} = \frac{1}{16} \mu_{\text{both}} \eta_s \eta_i + S_s^z S_i^z, \quad (\text{C4})$$

$$C_{\text{min}} = S_s^z S_i^z, \quad (\text{C5})$$

and the singles count probability is given by

$$S_j^z = \frac{1}{4} \mu_j \eta_j + \left(\frac{1}{2} \eta_{r_j} \alpha_{\text{pol}} \mathcal{R}_j \Delta \lambda_j + d_j \right) \Delta T_j. \quad (\text{C6})$$

Thus, it is clear that when one uses the first and third bins to measure the Z basis, one will have an unequal degradation of fidelity when projecting onto qubits on the equator or poles of the Bloch sphere. This would, for example, lead to the reconstruction of a density matrix that resembles a nonmaximally entangled state when quantum state tomography is performed [71], even if the input Bell state is ideal when noise is ignored.

There are some ways in which this can be addressed. For example, a beam splitter or switch could be placed in front of the interferometer [Fig. 10(a)], where some fraction of the photons are sent toward another detector without an interferometer to perform Z -basis measurements. In this case, the Z basis would not suffer an additional source of loss compared with the noise and would maintain the same fidelity as the X and Y bases. Thus, it appears that use of an interferometer for all three bases is not optimal for time-bin encoding in high-noise environments. For applications in which only two bases are required, the X and Y bases would be optimal to achieve the highest fidelity or tolerance to background noise.

2. Importance of polarization filtering and comparison with polarization encoding

In this section, we compare time-bin entanglement with polarization entanglement and examine the benefit of polarization filtering of unpolarized noise photons.

For polarization encoding, we build on the model in Ref. [20] but modify it to account for the properties of the source's JSA and filtering. For a source pumped in a coherent equal superposition of horizontal (H) and vertical (V)

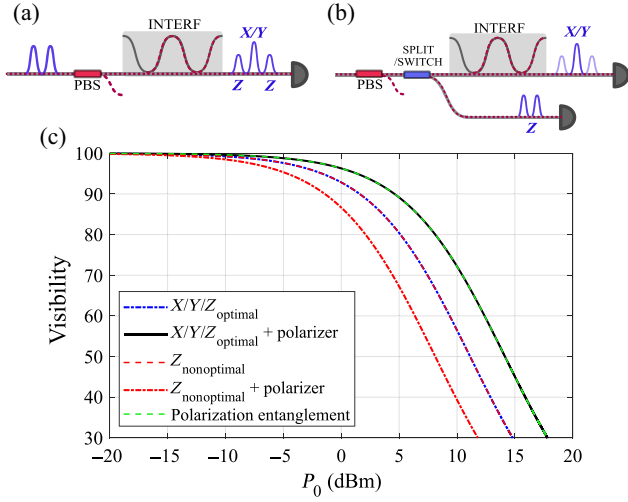


FIG. 10. Comparison between time-bin entanglement distribution and how the receiver is designed for polarization entanglement. (a) Time-bin encoding measured with a 1-bit delay interferometer (INTERF), where all bases are measured with use of the interferometer output. (b) Time-bin encoding with a splitter or switch used to measure the Z basis and interferometer for the X and Y bases, which is the optimal case for the Z basis. (c) Entanglement visibility as a function of classical power for each case assuming ideal sources and FHE. PBS, polarizing beam splitter.

polarizations such that the number of photons generated across the JSI in each qubit degree of freedom is given by $\mu_T \equiv \mu_T^H = \mu_T^V$, the coincidence counts in the low-gain regime can be modeled by

$$C_{\max}^{\text{pol}} = \mu_{\text{both}} \eta_s \eta_i + S_s^{\text{pol}} S_i^{\text{pol}}, \quad (\text{C7})$$

$$C_{\min}^{\text{pol}} = S_s^{\text{pol}} S_i^{\text{pol}}, \quad (\text{C8})$$

where the singles count probability including background noise is given by

$$S_j^{\text{pol}} = \mu_j \eta_j + (\eta_j \alpha_{\text{pol}} \mathcal{R}_j \Delta \lambda_j + d_j) \Delta T_j. \quad (\text{C9})$$

This shows that for unpolarized noise, standard measurements of polarization entanglement intrinsically achieve the same tolerance to noise as including polarizers aligned to a photon pair with definite polarization, since, for example, μ_j^H photons will pass a polarizer at the measurement apparatus projecting onto $|H\rangle$ with 100% probability. For unpolarized noise, the input noise intensity is reduced by $\alpha_{\text{pol}} = 1/2$, similarly to for polarization filtering with time-bin qubits due to the polarizer in the qubit analyzer. Thus, the analyzer naturally applies polarization filtering in the measurement procedure and has a 4 times higher coincidence probability compared with time-bin encoding in the X and Y bases [78] and 16 times higher coincidence probability than the Z-basis measurements using interferometer outputs. Note that the coincidence probabilities for

all bases are the same if we assume unpolarized background noise photons and the generation and alignment of a maximally entangled Bell state.

To demonstrate the differences between these approaches and also evaluate the limits of source and filter designs, we simulate the performance of time-bin entangled photons compared with polarization entangled photons measured with a standard polarization analyzer. We further simulate an ideal source and receiver to see the limits of the tolerable classical power when all key parameters are optimized. We simulate time-bin encoding using a 1-bit delay interferometer to measure all three bases [Fig. 10(a)] and time-bin encoding using a splitter or switch to better measure the Z basis [Fig. 10(b)]. For each case, we compare the performance with and without polarization filtering.

In this simulation, we keep the same filtering used in our time-bin experiment (6.5 GHz, 200 ps) but assume an ideal FHE to examine the limits in an idealized case. We further assume no loss at the source node, which maximizes the SNR at each detector since this loss does not affect the noise photons generated in the long-distance fibers. The mean photon number at the source is optimized to find the highest possible visibility at each classical power according to the background noise level and channel losses. All other parameters are identical to those used in the 1536.5 nm entanglement distribution experiment over 25 km/25 km fibers with the copropagating 1547.72 nm classical signals, representing the best performance that could be achieved in principle with these filter settings.

The results are shown in Fig. 10(c). We find that polarization encoding outperforms the time-bin entanglement distribution without polarization filtering. However, addition of polarizers before the 1-bit delay interferometer makes time-bin encoding in the X and Y bases as resilient to incident noise rates as polarization entanglement. However, the coincidence rate will still be 25% of a polarization entangled state. When the Z basis is measured at the interferometer output (nonoptimal), the fidelity is more impacted and is unable to tolerate as high a classical power, which is what we observed in our experiment. However, if one uses a switch or splitter to measure the Z basis without the interferometer (optimal), this can result in the same tolerance to incident background noise as for the X and Y bases. When polarization filtering is applied, all cases achieve the same visibility for all classical power levels.

We also see that the tolerable classical power is much higher than in the experimental demonstration, which is due to the increase of the source intensity to overcome the higher background noise, the ideal FHE, and the reduction of the loss before multiplexing with the classical signal. Using these improvements, we predict that it is feasible to reach C-band classical launch powers of around 10 dBm in each fiber while still allowing entanglement-based quantum applications at C-band wavelengths.

APPENDIX D: GAUSSIAN APPROXIMATION FOR MODELING A FILTERED JSA

Here we briefly summarize the Gaussian approximation for the filtered JSA used in simulations of the CAR as a function of pump pulse width and filter bandwidth in the main text. We use the models in Refs. [57,88], where the JSA of a source is represented by the equation [57]

$$f(\omega_s, \omega_i) = N \exp\left(\frac{-(\omega_s - \omega_{s0} + \omega_i - \omega_{i0})^2}{4\sigma_p^2}\right) \times \text{sinc}\left(\frac{([\omega_s - \omega_{s0}] \sin \theta + [\omega_i - \omega_{i0}] \cos \theta)}{2\sigma_{\text{PM}}}\right), \quad (\text{D1})$$

where σ_p is the pump's frequency bandwidth, σ_{PM} is the phase-matching bandwidth, and $\theta = \arctan\left(\frac{k'_p - k'_s}{k'_p - k'_i}\right)$ is the phase-matching angle, with k'_x being the frequency derivative of k for the signal or idler photons. Filters can then be applied, and are modeled by normalized transmission functions $g_j(\omega_j)$ with central frequencies of ω_{j0} . Then each parameter can be obtained by numerical integration.

Alternatively, simple analytical equations can be obtained for each parameter with use of a Gaussian approximation for the pump, phase-matching, and filter functions [57,88]. To do this, the sinc function in the JSA is approximated by $\text{sinc}(x) \approx \exp(-\alpha x^2)$ with $\alpha = 0.193$ such that the two have the same FWHM. Gaussian filters for the signal and idler photons can then be applied with bandwidths $\sigma_j = \sigma_{j\text{FWHM}}/(2\sqrt{2\ln 2})$.

Using this approximation, we can derive the equation for Γ_{both} that relates to the coincidence probability ($\mu_{\text{both}} = \mu_T \Gamma_{\text{both}}$) as [57]

$$\Gamma_{\text{both}} \approx \sqrt{\frac{a_0 b_0 - c^2}{ab - c^2}}, \quad (\text{D2})$$

where

$$\begin{aligned} a &= \frac{\alpha^2 \sin^2 \theta}{\sigma_{\text{PM}}^2} + \frac{1}{\sigma_p^2} + \frac{1}{\sigma_s^2}, \\ b &= \frac{\alpha^2 \cos^2 \theta}{\sigma_{\text{PM}}^2} + \frac{1}{\sigma_p^2} + \frac{1}{\sigma_i^2}, \\ c &= \frac{\alpha^2 \cos \theta \sin \theta}{\sigma_{\text{PM}}^2} + \frac{1}{\sigma_p^2}. \end{aligned} \quad (\text{D3})$$

Setting one of the signal or idler filter bandwidths to infinity gives the marginal probabilities for each individual

photon, giving the equations [57]

$$\begin{aligned} \Gamma_i &\approx \sqrt{\frac{a_0 b_0 - c^2}{a_0 b - c^2}}, \\ \Gamma_s &\approx \sqrt{\frac{a_0 b_0 - c^2}{a b_0 - c^2}}, \end{aligned} \quad (\text{D4})$$

which relate to the single-detector counts by $\mu_s = \mu_T \Gamma_s$ and $\mu_i = \mu_T \Gamma_i$.

The FHE, δ_j , can then be approximated with use of the equations [57]

$$\begin{aligned} \delta_s &= \frac{\Gamma_{\text{both}}}{\Gamma_i} = \sqrt{\frac{a_0 b - c^2}{ab - c^2}}, \\ \delta_i &= \frac{\Gamma_{\text{both}}}{\Gamma_s} = \sqrt{\frac{ab_0 - c^2}{ab - c^2}}. \end{aligned} \quad (\text{D5})$$

In the main text, we also investigate the spectral purity as a function of filtering to relate to the ability to achieve HOM interference with independent photons. The equation for the single-mode purity for a heralded single photon, $\mathcal{P}$, is derived in Refs. [57,88]. Using this approach, the spectral purity can be approximated by the equation [57]:

$$\mathcal{P} \approx \sqrt{\frac{ab - c^2}{ab}}. \quad (\text{D6})$$

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