Highly Efficient Single-Pass Source of Pulsed Single-Mode Twin Beams of Light

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We report the realization of a bright ultrafast type II parametric down-conversion source of twin beams free of any spatiotemporal correlations in a periodically poled $KTiOPO_4$ (PP-KTP) waveguide. From a robust, single-pass setup it emits pulsed two-mode squeezed vacuum states: photon-number entangled pairs of single-mode pulses or, in terms of continuous variables quantum optics, pulsed Einstein-Podolsky-Rosen states in the telecom wavelength regime. We verify the single-mode character of our source by measuring Glauber correlation functions $g^{(2)}$ and demonstrate with a pump energy as low as 75 pJ per pump pulse a mean photon number of 2.5.

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The main obstacle to the real-world deployment of wide area quantum communication networks is the limited distance of guaranteed security between communication partners. In order overcome it, quantum repeaters [1] are needed to counter the security-degrading effects of transmission losses. For continuous variable (CV) quantum communication, these protocols heavily rely on the concatenation of non-Gaussian states and squeezed Gaussian states [2], namely, EPR states produced by parametric down-conversion (PDC) combined with photon counting [3]. In general though, PDC does not produce single-mode (SM) but multimode (MM) EPR states, requiring additional post-processing for optimal fidelity. Their MM structure is intrinsic to their generation process [4], and only direct manipulation of that process allows for the production of SM states.

For the generation of photon pairs, PDC sources have become an established standard: Inside a $\chi^{(2)}$-nonlinear medium, a pump photon decays into one signal and one idler photon. Recent works have shown that PDC source engineering [5,6] is capable of producing spectrally separable two-photon states $|1\rangle_s \otimes |1\rangle_i$, allowing for the preparation pure heralded single photons [7]. Going beyond the single photon pair approximation, PDC in general can be understood as a source of squeezed states of light [8,9]. First observed by Slusher *et al.* [10] in 1985, squeezed states originally garnered interest for the noise reduction in their quadrature observables $\hat{X}$, $\hat{Y}$ below the classical shot noise level, applicable in quantum-enhanced interferometry [11]. With the availability of mode locked lasers, MM pulsed squeezed states [12] became accessible [13]. Measuring with detectors incapable of resolving this MM structure, such as avalanche photo diodes (APD) implementing non-Gaussian operations [14], introduces mixedness which degrades the quantum features of the state [15]. Spectral engineering [5–7,16] made it possible to generate ultrafast photon pairs without spectral correlations in one spectral broadband mode [17]. Until now, PDC experiments relied on spatial [7] or narrow spectral [18] filtering of MM [9,12] squeezers to approximate SM biphotonic states, with severe loss of source brightness. In recent years, waveguide PDC sources [19] have become more and more popular as a means of achieving higher brightness [20] in a single-pass configuration, as well as for their easy integrability into miniaturized quantum optical experiments.

In this Letter, we demonstrate a waveguided single-pass type II PDC source of ultrafast SM EPR states of unprecedenced brightness in the telecom wavelength regime. For low pump powers ($\langle n\rangle \ll 1$), it doubles as a source of pure heralded single photons. By utilizing a SM PP-KTP waveguide, the output states can be used without narrow spatial filtering, and spectral engineering lets us avoid narrow spectral filtering, boosting source brightness considerably. The ultrafast, broadband nature of the pump beam makes spectral broadband modes a natural choice to describe our system. Our source emits pairs of spectrally broadband SM pulses. This we corroborated by a measurement of spectral separability, a $g^{(2)}$ measurement of one of the output arms to ensure the expected photon statistics, and a $\sinh^2$ gain in mean photon number. In contrast to MM PDC sources, the generated photons are not scattered over a number of broadband modes but concentrated in one mode. The resulting squeezing of several broadband modes cannot be trivially combined into one mode, as this would amount to entanglement distillation, which has been shown to be impossible using Gaussian operations [2]. But spectral engineering and a SM waveguide allow us to efficiently emit all output light into one spatiospectral mode, thus leading to a mean photon number $\langle n\rangle = 2.5$ per mode when pumping with picosecond pulses of only 75 pJ.

It has been shown early on in the experimental exploration of squeezing that PDC produces squeezed vacuum states of light [8]. In photon number representation, a two-mode squeezed vacuum state or SM EPR state has the form

$$|\psi\rangle = \hat{S}_{a,b}|0\rangle = e^{\imath \hat{H}_{a,b}}|0\rangle = \sqrt{1-|\lambda|^2}\sum_n \lambda^n |n,n\rangle, \quad (1)$$

where a and b are two orthogonal modes, $\hat{S}_{a,b}$ is the two-mode squeezing operator, and $\hat{H}_{a,b} = \zeta \hat{a}^\dagger \hat{b}^\dagger + \text{H.c.}$ is its effective Hamiltonian. It is a coherent superposition of strictly photon number correlated Fock states, and exhibits thermal photon statistics in both modes a and b. Its mean photon number $\langle n\rangle$ in each output beam is a measure of how much two-mode squeezing is generated $s = \frac{20}{\ln(10)} a\sinh(\sqrt{\langle n\rangle})$. The photon number correlation between both modes allows for heralding pure single photons with binary detectors. However, the underlying bilinear effective Hamiltonian $\hat{H}_{a,b}$ describes only a special case of PDC.

In general, the effective PDC Hamiltonian has a richer spatiospectral structure, and emits a continuum of momentum modes. Waveguided PDC, due to the boundary conditions of the guiding structure, exhibits a discrete spatial mode spectrum. In principle, any combination of pump, signal and idler waveguide modes will contribute to the overall PDC process. Their coupling strength is determined by their "overlap" integral [21], and multiple nonzero coupling coefficients between existing modes will result in a complex spectral structure of the output state. But using a SM waveguide allows us to restrict our analysis to one spatial mode for each beam, and we find

$$\hat{H}_{\text{PDC}} = \zeta \int d\omega_1 \int d\omega_2 f(\omega_1,\omega_2)\hat{a}^\dagger(\omega_1)\hat{b}^\dagger(\omega_2) + \text{H.c.}, \quad (2)$$

which generates a generalized version of the two-mode squeezed vacuum in Eq. (1) with spectrally correlated output beams. The coupling constant ζ determines the strength of this interaction, while spectral correlations between photons of the pairs produced are governed by the normalized joint spectral amplitude $f(\omega_1,\omega_2)$.

By applying a Schmidt decomposition to the joint amplitude $f(\omega_1,\omega_2) = \sum_k c_k \varphi_k(\omega_1)\psi_k(\omega_2)$, we obtain two orthonormal basis sets of spectral amplitude functions $\{\varphi_k(\omega_1)\}$ and $\{\psi_k(\omega_2)\}$ and a set of weighting coefficients $\{c_k\}$ with $\sum_k |c_k|^2 = 1$. For ultrafast pumped type II PDC, the φ_k, ψ_k are in good approximation to the Hermite functions [5,17]. Now the PDC Hamiltonian can be expressed in terms of broadband modes

$$\hat{H}_{\text{PDC}} = \sum_k \hat{H}_k = \zeta\sum_k c_k(\hat{A}_k^\dagger \hat{B}_k^\dagger + \hat{A}_k\hat{B}_k). \quad (3)$$

Each broadband mode operator $\hat{A}_k$, $\hat{B}_k$ describes a temporal pulse mode, or equivalently, an ultrafast spectral mode. It is defined as superposition of monochromatic creation/annihilation operators $\hat{a}(\omega)$, $\hat{b}(\omega)$ operators weighted with a function from the Schmidt basis: $\hat{A}_k^\dagger := \int d\omega \varphi_k(\omega)\hat{a}^\dagger(\omega)$ and $\hat{B}_k^\dagger := \int d\omega \psi_k(\omega)\hat{b}^\dagger(\omega)$. The effective Hamiltonians $\hat{H}_k$ do not interact with each other (since $[\hat{H}_k, \hat{H}_l] = 0$), and thus the PDC squeezing operator represents in fact an ensemble of independent two-mode squeezing operators $\hat{S}_{a,b} = e^{\imath \hat{H}_{\text{PDC}}} = \hat{S}_{A_0,B_0} \otimes \hat{S}_{A_1,B_1} \otimes \ldots$ where the coefficients c_k determine the relative strength of all squeezers as well as spectral correlation between signal and idler beams. This correlation is characterized by the source's effective mode number $K = \frac{1}{\sum_k |c_k|^4}$ [6,22]. We note that as the overall mean photon number n is shared between all modes, the amount of two-mode-squeezing has to be considered for each mode separately. For $c_0 = 1$ and all other $c_k = 0$, K assumes its minimum value of 1, and the PDC process can be described as a two-mode squeezer according to Eq. (1), and also optimal squeezing performance can be expected.

In our waveguided source pumped by an ultrafast pulsed laser beam we can manipulate spectral correlations of the photon pair joint spectra, thus the coefficients c_k, and as a result minimize K by simply adjusting the spectral width of the pump pulses [5,6,22].

We verified this by measuring the joint spectral intensity (JSI) of generated photon pairs at different spectral pump widths, to show the dependence of bi-photon frequency correlations on the pump width. The setup in Fig. 1(a) illustrates the PDC source: Ultrafast pump pulses at 768 nm are prepared with a Ti:sapphire mode locked laser system, spectrally filtered with a variable bandpass filter $4f$ setup, and then used to pump a type II PDC process within the PP-KTP waveguide with a poling period of 104 μm and 4 μm $\times$ 6 μm size. Its length is 10 mm but an effective length of 8 mm is used to correctly predict the measurement results in Figs. 2 and 3, since manufacturing imperfections in poling period and waveguide diameter lead to a widened phasematching distribution $\Phi(\omega_1,\omega_2)$ as if from a shorter waveguide. The generated photon pairs

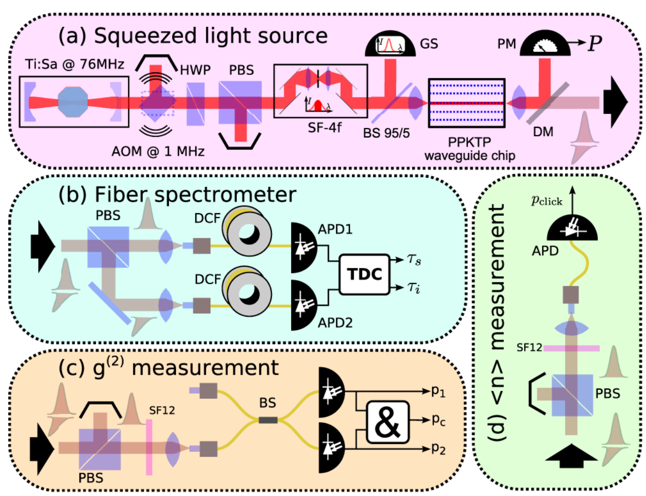

FIG. 1 (color online). Experimental setup: (a) PP-KTP waveguide source of two-mode squeezed vacuum states. (b) Fiber spectrometer for JSI measurement. (c) $g^{(2)}$ measurement setup. (d) Mean photon number $\langle n\rangle$ measurement.

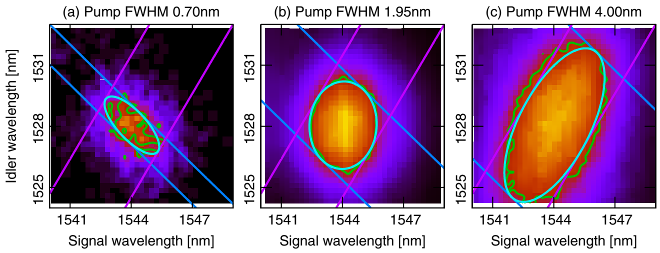

FIG. 2 (color online). Two-photon spectral intensities from the setup 1(b) with pump width below, equal to and above photon pair separability width at 1.95 nm FWHM. Green: 50% intensity. Violet: phase matching width. Blue: pump width. Bright blue: theoretical 50% intensity.

are analyzed in a fiber spectrometer [23] [Fig. 1(b)]: After separating signal and idler photons by polarization, they independently travel through long dispersive fibers, and are detected by a pair of idQuantique id201 avalanche photo diodes (APDs). Because of the chromatic dispersion of the fibers, the photons' group velocity and arrival time at the APDs depend on their wavelength. Thus we are able to determine the spectral intensity distribution of a stream of single photons from its arrival time spread. For a spectral pump FWHM of 0.70, 1.95, and 4.0 nm, we observe in Fig. 2 negative spectral correlations, an uncorrelated spectrum, and positive spectral correlations between signal and idler photons, respectively. For a fully uncorrelated state with just one contributing Schmidt mode pair φ_0, ψ_0 the marginal spectra of the biphoton amplitude are given by $|\varphi_0|^2$, $|\psi_0|^2$. In Fig. 3 (left), their shapes according to Fig. 2 (middle) are in good agreement with the Gaussian marginal distributions of their theory curve, indicating that the dominant Schmidt functions of the generated state are, up to phase, Gaussians as well. We have demonstrated control over spectral entanglement between signal and idler by filtering the pump spectrum, and found minimal spectral correlations of photon pairs around 1.95 nm pump FWHM.

To prove the genuine two-mode squeezer character of our source, an uncorrelated JSI is necessary but not sufficient. It is proportional to the modulus square of the complex joint amplitude $|f(\omega_1, \omega_2)|^2$ of the photon pair, so all phase information is lost in an intensity measurement. In order to detect phase entanglement between signal and idler, we need to measure an additional quantity sensitive to the source's mode number K, which is unity only in the absence of entanglement on the photon pair level, and larger otherwise.

The second order correlation function $g^{(2)}$ can be used to discriminate between beams with thermal ($g^{(2)} = 2$) and Poissonian photon statistics ($g^{(2)} = 1$) from a PDC source [18]. As has been noted above, type II PDC can in general be understood as an ensemble of two-mode squeezers, each of them emitting two beams with thermal photon statistics. In our waveguided type II setup, all broadband modes, A_k or B_k, share one polarization mode, a and b or, respectively. A detector with a spectral response function much wider than the characteristic width of the broadband modes cannot resolve them. It "sees" a convolution of the thermal photon statistics of all broadband modes, and in the limit of a large number of modes, this is a Poissonian distribution [24]. But if there is only one mode per polarization to begin with (which is only true for a two-mode squeezer), the detector receives a thermal distribution of photon numbers. Therefore, with the assumption that PDC emits a pure state, we can infer from $g^{(2)} = 2$ measured in either output beam a two-mode squeezer source. Indeed, for low pump power and thus low coupling strength ζ, we can find a simple connection between the $g^{(2)}$ correlation function on the one hand, and the broadband mode structure of our source and the effective mode number K on the other: $g^{(2)} = 1 + \sum |c_k|^4 = 1 + \frac{1}{K}$.

Figure 1(c) illustrates the $g^{(2)}$ measurement: Idler is discarded, and the signal beam is split by a 50/50 beamsplitter. Its output modes are fed into APDs, single (p_1, p_2) and coincidence (p_c) click probabilities for different spectral pump widths are recorded. When using binary detectors far from saturation, rather than intensity measurements, one finds $g^{(2)} = \frac{\langle \hat{a}^\dagger \hat{a}^\dagger \hat{a} \hat{a} \rangle}{\langle \hat{a}^\dagger \hat{a} \rangle^2} \approx \frac{p_c}{p_1 p_2}$ As has been demonstrated, frequency correlations between signal and idler beam and thus squeezer mode number K can be controlled by manipulation of the pump width. In Fig. 3 (middle) measurement results show a maximum $g^{(2)}$ value at 1.95 nm pump FWHM, in accordance with Fig. 2. When departing from the optimum pump width, $g^{(2)}$ decreases as predicted. Because of uncorrelated, residual background events from waveguide material fluorescence and detector dark counts that make up 5% of the total single event counts, we obtain a maximum of $g^{(2)} = 1.8$, and $g^{(2)} = 1.95$ after background correction. This highlights the next-to-perfect SM EPR states our source emits, and the degree of control we exact over the mode number and photon statistics of the system. Note that the two-mode character is shown with respect to frequency as well as

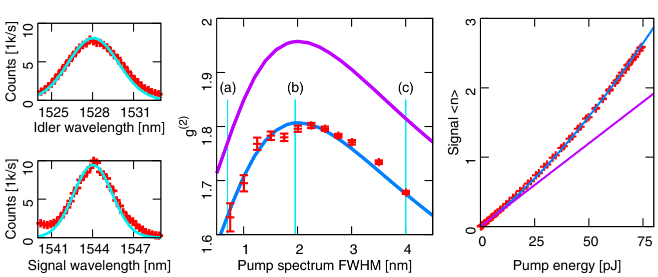

FIG. 3 (color online). Left: Marginal spectra of signal, idler beams from Fig. 2(b). Middle: $g^{(2)}$ values from setup 1(c) (red) with theory curve (blue) and background corrected theory curve (violet); (a), (b) and (c) mark pump FWHM of the JSI from Fig. 2. Right: Mean photon number from setup 1(d) (red) with the theoretical gain of a two-mode squeezer (blue) and the linear gain of a highly MM squeezer (violet).

spatial degrees of freedom. Owing to the waveguide nature of our source, signal and idler beam occupy a single waveguide mode.

Nonlinear waveguides allow for dramatically higher source brightness when compared to bulk sources [20]: Instead of coupling to a continuum of spatial modes, inside a waveguide structure the generated waves couple to a discrete spectrum, and ideally to just one mode, boosting self-seeding of the PDC process and greatly simplifying collection of the output light. Adjusting our setup with a CW laser beam shows collection efficiency of the waveguide output mode into SM fibers up to 80%, indicating in good approximation a Gaussian mode profile. At mean photon numbers of $\langle n \rangle \approx 1$ per mode we will be able to observe the superlinear gain of a two-mode squeezer $\sinh^2(r)$ caused by self-seeding of signal and idler along the waveguide length, further corroborating our source's SM character. With a pump FWHM of 1.95 nm producing separable photon pairs, we measured the mean photon number $\langle n \rangle \approx \frac{p_{\mathrm{click}}}{\eta}$ of the signal beam [Fig. 1(d)] by recording the power dependent APD click probability p_{click}. For binary detectors far from saturation, this is proportional to $\langle n \rangle$, with an overall quantum efficiency η of the setup. The source gain in Fig. 3 (right) exhibits with increasing pump power the departure from the linear gain profile that would be expected for a highly MM squeezer, while it is in very good agreement with the theoretical prediction for a two-mode squeezer gain. Mean photon numbers of up to 2.5 were achieved. Assuming ideal photon collection and detectors, this is equivalent to 11 dB of two-mode squeezing, in a pulsed, single-pass setting, demonstrating the potential of our source for future CV experiments. For an optimized setup we observed an overall detection efficiency of 15%. For a specified APD quantum efficiency of 25% at 1550 nm, this makes a photon collection efficiency into SM fiber of 60%, with our waveguide output facet not antireflection coated.

In conclusion, we have applied spectral engineering to a waveguided PDC source to create a bright, genuinely ultrafast pulsed two-mode squeezer in the telecom wavelength regime with mean photon number as high as 2.5, with only 75 pJ pump pulse energy. In future experiments, this value can be easily scaled up to harvest even higher photon numbers. It features near thermal photon statistics with $g^{(2)} = 1.95$ after background correction, or an effective mode number of $K = 1.05$. A collection efficiency of 60% into SM fibers demonstrates the high spatial mode quality of our waveguide device and shows its potential for inclusion into integrated optical networks. Because of its true two-mode character and brightness, we expect widespread adoption of our source in continuous variable quantum communication, where high squeezing values, purity and low-loss fiber transmission are prerequisite for efficient quantum cryptography [25], teleportation [26,27], and ultimately entanglement distillation [3,14] to overcome transmission losses in wide area quantum communication networks, a vital building block of quantum repeaters [1].

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