Cavity enhanced telecom heralded single photons for spin-wave solid state quantum memories

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Keywords: photon pair source, cavity enhanced SPDC, heralded single photons, narrowband photons, solid state quantum memory

Abstract

We report on a source of heralded narrowband (~3 MHz) single photons compatible with solid-state spin-wave quantum memories based on praseodymium doped crystals. Widely non-degenerate narrow-band photon pairs are generated using cavity enhanced down conversion. One photon from the pair is at telecom wavelengths and serves as heralding signal, while the heralded single photon is at 606 nm, resonant with an optical transition of Pr³⁺:Y₂SiO₅. The source offers a heralding efficiency of 28% and a generation rate exceeding 2000 pairs mW⁻¹ in a single-mode. The single photon nature of the heralded field is confirmed by a direct antibunching measurement, with a measured antibunching parameter down to 0.010(4). Moreover, we investigate in detail photon cross- and autocorrelation functions proving non-classical correlations between the two photons. The results presented in this paper offer prospects for the demonstration of single photon spin-wave storage in an on-demand solid state quantum memory, heralded by a telecom photon.

1. Introduction

The process of spontaneous parametric down conversion (SPDC) finds widespread application in quantum physics as a source for quantum states of light. In SPDC a pump photon inside an optically nonlinear medium can split spontaneously into a pair of photons (called signal and idler), each of them at lower energy than the pump. The detection of one photon signals the presence of its partner photon. Hence, SPDC can be used as a source of heralded single photons [1]. The spectral bandwidth of the SPDC process is typically in the order of hundreds of gigahertz. This results in the generation of broadband photons. However, for many applications in quantum optics, where the photons are interfaced with atomic transitions, much smaller line widths are necessary. One possibility to decrease the spectral bandwidth of the emitted photons is to use spectral filters after the crystal [4–6]. However, this requires a very efficient SPDC process to achieve high pair rates and it is technically challenging to build low loss narrowband filtering systems in the MHz range. Another approach is to insert the nonlinear material in a resonator [7]. In such an optical parametric oscillator (OPO), operated below its oscillation threshold, photons can only be created resonant to the cavity modes. The spectra of the photons resemble the spectrum of the cavity. Additionally, the resonance allows for a spectral enhancement of the creation rate. With this technique, photon sources with line widths in the MHz range could be demonstrated with various designs [8–18]. Interfacing such photons effectively with atomic transitions has also been demonstrated [19–24].

Such sources are of interest for the development of quantum repeaters [25, 26]. One interesting scenario is that one photon of the source is stored in a quantum memory while its partner photon is used to distribute entanglement to the neighboring node [27]. To bridge large distances the second photon should be at a low loss telecommunication wavelength. To that end frequency non-degenerate sources have been developed [5, 6, 12].

Rare earth doped solids cooled to cryogenic temperatures offer suitable atomic transitions to serve as photonic quantum memories. Single photons have been stored as collective optical excitations for pre-determined durations <5 µs in the excited state of such crystals [20–22, 28, 29] using the atomic frequency...
combs [30]. Longer storage times and on demand read-out may be reached by transferring the optical excitations to long-lived collective spin-excitations (known as spin-waves) [31–33]. Among rare-earth doped materials, praseodymium doped crystals (e.g. Pr$^{3+}$:Y$_2$SiO$_5$) offer exceptional properties for quantum storage. Classical images have been stored for durations on the minute time scale [34], and high storage and retrieval efficiencies have been demonstrated for weak coherent states [35] and classical light [36]. Praseodymium ions also offer an electronic structure with three ground state levels allowing spin wave storage. The spin-wave storage of time-bin qubits encoded in weak coherent pulses at single photon level has been shown with praseodymium [32] and europium ions [33]. However, the storage of external quantum states of light in a spin state of a rare earth is still a remaining challenge and could not be demonstrated, yet.

A photon pair source compatible with praseodymium doped memories has been demonstrated by our group recently [12], using cavity enhanced down conversion. The source featured a limited heralding efficiency of 6% and a low detected count rate of around 3 coincidences per mW of pump power, in single mode operation [20]. This prevented us to measure directly the single photon character of the emitted light, by measurement of the autocorrelation signal. While the source was sufficient to demonstrate quantum storage in the excited state of the Pr$^{3+}$:Y$_2$SiO$_5$ crystal [20], its performances were not sufficient for spin-wave storage. While storage in the excited state is an almost noiseless process, this is not the case for storage in the spin state. The control pulses necessary for the transfer between excited and spin state have macroscopic power and thus can contribute to noise by leakage by scattering in the crystal or via interactions with atoms that are not perfectly optically pumped and then decay coherently or incoherently. Hence a major requirement for the storage in a three level system is to reduce the number of unnecessary transfer pulses. As such control pulses will be sent conditioned on the detection of a herald, a high heralding efficiency is mandatory. In both excited state and spin-state storage, the photons interact with the same atomic levels. Hence there is no difference in the requirements for the photon bandwidth.

In this paper we report on the improved generation of telecom heralded single photons compatible with a Pr quantum memory, using cavity enhanced down-conversion. We report an order of magnitude improvement in detected coincidence count rate, and a significant improvement in heralding efficiency, compared to previous realizations of this source [12, 20]. This allows us to measure directly the single photon nature of the heralded photons through autocorrelation measurements, with measured antibunching parameters down to 0.01. With the help of cross- and unconditioned autocorrelation measurements we demonstrate the violation of the Cauchy–Schwarz inequality. We complete the characterization with a detailed analysis of the non-classical state of our photons for different pump powers and an investigation of the spectral modes formed by the double resonance of the cavity. With these improved performances, our source should now be suitable for spin wave storage in a Pr$^{3+}$:Y$_2$SiO$_5$ quantum memory [32]. To our knowledge our source is at the moment the only one compatible with Pr solid-state quantum memories and telecom heralding, making it a unique device for implementing quantum repeater technology. It is worth mentioning that there exist other techniques producing correlated photons compatible with memories, apart from SPDC. In the so-called DLCZ protocol [37] quantum correlations between a photon and its delayed (stored) partner are established via Raman transitions in atomic ensembles. This way of combining memory and source could recently also be applied to solid state systems [38, 39]. A similar technique based on rephasing of spontaneous emission (RASE, [40]) has recently been applied to Pr based solid-state memories [41], demonstrating continuous variable entanglement between a spin-wave solid state quantum memory and a light field. However, none of these realizations features telecom compatible photons.

The outline of this report is as follows. We first introduce the experimental setup of the source and discuss the improvements compared to the previous version. It follows a presentation of temporal and spectral parameters of the photons and results for the important figures of merit, the heralding efficiency and the coincidence rates. Afterwards we investigate the non-classical correlations of the photons by cross correlation as well as unconditioned and conditioned auto correlation measurements. We compare the results with theoretical models and check for consistency.

2. Source

A sketch of the experimental setup can be found in figure 1. The photon pairs are generated in a 2 cm long periodically poled lithium niobate crystal (poling period $\Lambda = 16.5$ $\mu$m). We use type-I phase matching where both signal and idler are polarized along the crystal $Y$-axis. The crystal is temperature stabilized and positioned in the center between the two curved mirrors ($R = 100$ mm) of a four mirror bow-tie type ring cavity (free spectral range, FSR = 423 MHz). Three mirrors have a highly reflective coating for signal and idler, one of the plane mirrors is used as output coupler. This makes the resonator doubly resonant and allows to exploit the clustering effect. It has been described in detail for OPOs as laser [42] and SPDC sources [15, 43]. In summary, it implies the
fact that the photons can only be generated in resonance with the cavity where the spectrum can be described by an Airy function. However, as the cavity contains a crystal, it has a different optical length for signal and idler wavelength. Hence the spectrum of allowed modes can be described by the product of the Airy functions for signal and idler, having slightly different periodicity, and the phase matching function. The simultaneous resonance for both wavelengths is thus only fulfilled for a few modes close together, forming a cluster. The benefit of clustering is a drastic reduction in the number of generated modes which facilitates the filtering of a single mode. For application with our quantum memory, the wavelength of the signal photons is 606 nm. With the help of a reference laser we actively stabilize our source to the resonance frequency of the QM \[12\]. We use a pump wavelength around 426 nm, resulting in idler photons around 1436 nm in the telecom E-band. In order to ensure that the idler photon is also resonant with the cavity, we use classical light at 1436 nm generated by difference frequency generation between the pump laser at 426 and the 606 nm laser. The frequency of this 1436 nm light is actively stabilized to be on resonance by active feed-back on the pump frequency. We use a chopped lock scheme, where we alternate between locking and measurement period at a rate of 30 Hz (see \[12\] for more details). The full locking scheme allows stable operation of the source for several hours, enough for detailed investigations of the photon properties or experiments implementing the quantum memory. For example the data in figure 2 is composed of four measurements, the longest lasting 4.6 h without interruption. Behind the OPO the photons are split by a dichroic mirror. The idler photons are then sent through an actively stabilized Fabry–Perot filter cavity (linewidth ca. 80 MHz, FSR 17 GHz) before being coupled to a single mode fiber. This additional filter cavity in the idler arm extracts a single mode out of the OPO spectrum and was used in all measurements presented in this paper. The signal photons pass an etalon filter (linewidth 4.25 GHz, FSR 100 GHz) suppressing eventual side clusters and are then coupled to a polarization maintaining single mode fiber. The transmission losses between the cavity output and the detectors are 71\% for the signal arm, and 35\% for the idler arm, including 50\% for the idler filter cavity. We detect the telecom photons with an InGaAs single photon counting module (SPCM, IdQuantique) and the signal photons with a silicon SPCM (Perkin Elmer) with detection efficiencies of 10\% and 62\%, respectively. All detectors are fiber coupled.

In the first design \[12\] of our source, a low transmission out-coupling mirror prevented the photons to escape the cavity with efficiency higher than 30\%. By reducing the reflectivity of the output mirror to 97\%, we increased the escape efficiency to ca. 56\% for the signal photons and 74\% for idler photons. To calculate the escape efficiency we measure the linewidth and the free spectral range of the cavity to calculate the finesse. As the
transmission of the output coupler is known, we can use this value to estimate the internal round trip loss of the cavity (2.4% for signal, 1.1% for idler). These losses are dominated by absorption in the nonlinear crystal which differs significantly for our signal and idler wavelengths [44, 45]. The escape efficiency can be estimated with these internal losses $L_{\text{int}}$ and the transmission of the output coupler $T_{\text{out}}$. For a detailed analysis see appendix A.

3. Temporal and spectral characterization

As a result of changing the output coupling, the finesse of the resonator was reduced to 184 for idler and 114 for signal, and the bandwidth of the photons became slightly larger than in [12]. To investigate the new bandwidth of the photons we use the second order cross-correlation function $G^{(2)}$. We record the detection times of both detectors with fast time stamping electronics (Signadyne). In the post processing we correlate the time delays between the detection times of idler photons (start events) and signal photons (stop events). A resulting histogram is shown in figure 2.

Its temporal shape is a result of the cross correlation of the temporal shapes of both photons. The spectrum of each photon should resemble the Lorentzian shape of the cavity mode spectrum corresponding to an exponential time structure. Hence we see in figure 2 the rising exponential of the idler and the falling exponential of the signal wave packet. We fit the histogram with the following function:

$$G \propto \exp [-2\pi \cdot \Delta \nu_{c} \cdot t] \cdot \Theta(t) + \exp [2\pi \cdot \Delta \nu_{c} \cdot t] \cdot \Theta(-t) + c_{0}$$

which directly results in the bandwidths $\Delta \nu_{c}$ of the photons, typically 3.7 MHz for signal and 2.3 MHz for idler. We use the heaviside function $\Theta(t - t_{0})$ to distinguish between the idler ($t < t_{0}$) and signal ($t > t_{0}$). The smaller value for the idler photons can be explained by a higher cavity finesse due to slightly lower intra-cavity losses. We further define the full width at half maximum (FWHM) of the measured $G^{(2)}$-function as the correlation time $\tau_{c}$ of the photon pair ($\tau_{c} = 78 \text{ ns}$). It is connected to the bandwidth of the individual photons via $\tau_{c} = \frac{\ln 2}{2\pi \cdot \Delta \nu}$, which is also the linewidth of the heralded single photon.

The non-degenerate double resonance leads to a clustering effect containing three clusters with several modes each [12]. We investigate the spectrum of the dominant cluster in the center of the phase matching envelope and show the purity on the longitudinal mode in the idler arm. To this end the signal photons pass a scanning Fabry–Perot cavity (FPI in figure 1) and the detection events are recorded as well as the scanning trigger. The resulting histogram is shown in figure 3. The gray curve shows the detection histogram as a function of the cavity length, i.e. relative frequency. Additionally, we recorded the heralding photons which were filtered to a single spectral mode as discussed above. The black curve shows the spectral distribution of the heralded signal photons. We see that the envelope of the main cluster contains seven to nine modes and has a width (FWHM) of roughly 1.9 GHz. This is a slightly broader cluster than obtained in [12], due to the lower finesse. The measured cluster width is also larger than the estimation that can be obtained using Sellmeier equations and the product of Airy functions (720 MHz). However, the heralded detection shows that only one longitudinal mode is transmitted by the idler cavity. The detection events clearly agglomerate at the position of a single longitudinal mode.

In the context of connecting the source with our quantum memory, the memory itself will act as a spectral filter for the signal photons [20]. This means that no additional filtering is necessary there to suppress unwanted modes. All transitions necessary to operate the memory are in a range of a few tens of MHz around the central

Figure 3. The gray curve shows the spectrum of the signal photons, demonstrating the mode cluster. The black curve shows the spectral distribution of the signal photons heralded by a single mode idler photon.
frequency of the photons. The result is that all modes, except for the central one in the memory operation range, will be absorbed by the inhomogeneously broadened transition of the memory crystal. For this reason the measurements with no single mode filtering in the signal arm are not overestimating the performance of the source.

4. Heralding efficiency and coincidence count rate

While the finesse decreased with the higher transmission output coupler, the overall cavity-extracting efficiency of the photons increased. In addition, we increased the transmission of the signal photons between the cavity and the single photon detector. The noise in the idler mode was also reduced by the use of a new bandpass filter and a new photon detector with a low dark count rate. We use the heralding efficiency \( \eta_{H} = \frac{p_{s,i}}{p_{i\text{det},s}} \), defined via the coincidence probability, the probability to detect a herald, and the detection efficiency of the signal SPCM, as a figure of merit for generating heralded single photons. The improvements resulted in increased heralding efficiencies \( \eta_{H} > 28\% \) at the signal photon detector, as shown in figure 4. This is a significant increase, compared to our previous results in single mode configuration [20]. As outlined in the introduction a high heralding efficiency is mandatory for storage in the spin state. It has been shown that the condition for quantum storage in quantum memories is that \( \eta_{H} > \mu_{1} \), where \( \mu_{1} \) is the minimum input number of photons to achieve a signal-to-noise ratio of 1 after the memory [32, 33]. In state of the art demonstrations of solid-state spin-wave quantum memories with weak-coherent states, \( \mu_{1} \) ranges between 0.03 and 0.11 [32, 33, 46]. Current memory demonstrations use weak pulses, longer than our single photons (>260 ns). However, we also performed such tests with 100 ns weak coherent pulses, comparable to the heralded single photon duration. In this case we still observe \( \mu_{1} < 0.1 \). As the resulting heralding efficiency is still above the \( \mu_{1} \)-value our source fulfills the requirement for the spin wave storage.

We also measured the coincidence rate for different pump powers. For the detected rate we observe a slope of 34 Hz mW\(^{-1}\). Compared to our previous results for single mode operation [12] this is an increase by one order of magnitude. Correcting for the measured transmission in the signal and idler arms, we find a creation rate of photon pairs of around 2200 pairs s\(^{-1}\) mW\(^{-1}\) behind the cavity output mirror for the central mode of the cluster. This gives a spectral brightness at the output of the cavity of approximately 800 s\(^{-1}\) mW\(^{-1}\) MHz\(^{-1}\).

As the coincidences follow an exponential function while the noise is equally distributed in time, the coincidence rate will depend on the integration window of the coincidences (the \( g_{s,i}^{(2)} \) function will be introduced in the following section). In figure 4(b) we show the dependence of such parameters on the size of the coincidence window, while the window is always centered around the coincidence peak. Obviously coincidence rate and heralding efficiency saturate. For this reason, in this paper when we integrate over a certain window, we use a window width of 400 ns, illustrated in figure 2.

5. Cross-correlations measurements

To prove the non-classicality and the single photon character of the signal photons three types of correlation measurements are performed and compared with each other. A first indication for the non-classicality of the generated pairs is given by the normalized cross correlation function \( g_{s,i}^{(2)} \). It is defined as \( g_{s,i}^{(2)} = \frac{p_{s,i}}{p_{i} p_{s}} \), where \( p_{s,i} \) describes the probability for a coincidence detection of a signal and an idler photon, and \( p_{i} \) are the detection

Figure 4. (a) Source characterization for different pump powers: detected coincidences (upward triangles) and heralding efficiency (downward triangles); size of integration window \( \Delta t = 400 \) ns. (b) Dependence of coincidence rate (upward triangles), heralding efficiency (downward triangles), and cross correlation value (circles) on the choice of the coincidence window width (the \( g_{s,i}^{(2)} \) values were divided by a factor of 5 to use the same scale for all measures); pump power 1 mW. Error bars are partly hidden by the markers.
probabilities for single signal and idler events, respectively. We measure the correlation value with histograms, as the one shown in figure 2. The coincidence rate outside the coincidence peak is a result of accidental coincidences between uncorrelated photons or detector dark counts. We further integrate the coincidence rate in an interval of 400 ns around the correlation peak, shown in figure 2. The ratio of this signal of interest and the background is the $g^{(2)}_{s,i}$-value. The results for different pump powers are illustrated in figure 5. We measure a maximum value of $g^{(2)}_{s,i}$ of 161 ± 38, for a pump power of 125 μW. The value of $g^{(2)}_{s,i}$ is then decreasing for higher pump powers, as expected for an SPDC process. For very low pump powers on the other hand, the rate of detected photons becomes comparable to the dark count rates of the detectors, effectively reducing the $g^{(2)}_{s,i}$-value. The dependence of $g^{(2)}_{s,i}$ on the pump power, including noise, was theoretically described in [47]. We compare our data with this model, including our detection efficiencies, the dark count rate of our detectors and the inferred photon creation probability $p$. The photon creation probability per time window (400 ns) is calculated from our measured spectral brightness, corrected for the photon bandwidth and escape efficiencies. We find $p \approx 2.2 \times 10^{-3}$ mW$^{-1}$. To describe the data we use for the dark count probability the average value of both detectors. The model then gives results for $g^{(2)}_{s,i}(0)$-values. For comparison with our measured data we need to modify the values for an integration window of 400 ns. Assuming an exponential decay of the correlation function (with correlation time $\tau_c$) and a background constant in time we find the relation

$$\frac{g^{(2)}_{s,i}(\Delta \tau)}{g^{(2)}_{s,i}(0)} - 1 = \frac{\tau_c}{\Delta \tau} \left[ 1 - \exp \left( -\frac{\Delta \tau}{\tau_c} \right) \right]$$

(2)

In our case, we find $g^{(2)}_{s,i}(\Delta \tau) \approx g^{(2)}_{s,i}(0)/5.16$. The result of the modeling is shown as blue solid line in figure 5 describing well our experimental findings. In the figure we additionally plot the function $1/(p \cdot P_{\text{pump}})$ (blue dashed line), describing a noiseless cross correlation measurement.

### 6. Measurements of unconditional auto-correlation

The signal-idler second-order cross correlation is a quantity easy to measure, but it does not prove the quantumness of the generated light without further assumptions. Hence we performed additional measurements to unambiguously show the quantum character of the correlations. First we measured the unconditional second order autocorrelation functions for signal ($g^{(2)}_{s,s}$) and idler ($g^{(2)}_{i,i}$), shown in figure 6. To this end we split the particular photons in a 50/50 fiber beam splitter followed by single photon detectors at both outputs. Again we correlate the detection events in the post processing of the time tags, using the data from both the signal detectors only. For an ideal two-mode squeezed state we expect values of $g^{(2)}_{s,s}(0) = 2$, which is not the case for our results. For comparison with other measurements we further define the $g^{(2)}_{s,i}$-values via integration over a $\Delta \tau = 400$ ns interval, analog to the cross correlation value. This eventually results in measured values of $g^{(2)}_{s,s}(\Delta \tau) = 1.10(1)$ and $g^{(2)}_{i,i}(\Delta \tau) = 1.32(5)$. Thanks to the use of a low dark count detector this is the first time we observe a bunching peak for the idler photon autocorrelation (see [20] for a comparison). From the fits to the histograms we determine the width of the peaks as 138 ns and 184 ns (FWHM; signal and idler, respectively), which is twice the time constant as for the cross correlation, as expected (see appendix B). From the height of the peaks we deduce values of $g^{(2)}_{s,s}(0) = 1.18(4)$ for the signal and $g^{(2)}_{i,i}(0) = 1.5(2)$ for the idler photons.

Two reasons can explain the decrease of the measured correlation values: the presence of several spectral modes and noise in the detection process. The measurements, especially for idler, are affected by detector noise. We can estimate this influence with the following basic model that is described in the supplement of [20]. The detected count rate $N_{s,i,B} = S_{s,i} + B_{s,i,B}$ comprises the signal of interest $S_{s,i}$ and the noise $B_{s,i,B}$ for both detectors $A, B$. At time delays much larger than the correlation time we expect a coincidence rate proportional to $N_A \cdot N_{B,s,i}$.
while it is proportional to $N_A N_B + S_A S_B$ at zero delay, due to the bunching of the SPDC photons. The resulting $g^{(2)}_{\text{s,i}}(0)$-value is then: $g^{(2)}_{\text{s,i}}(0) = 1 + \frac{S_A S_B}{N_A N_B}$. We measured dark count rates of $B_A = 30$ Hz and $B_B = 50$ Hz for the signal measurement as well as $B_A' = 18$ Hz and $B_B' = 192$ Hz for the idler measurements. For the autocorrelation measurement of idler we gated the detectors in phase with the chopper, i.e. the cavity locking cycle, to avoid recording noise when no photons are produced. This is necessary because one of the two telecom detectors has a rather high dark count rate, significantly affecting the measurement outcome. Taking only these detector dark counts into account we find upper limits of $g^{(2)}_{\text{i,s,noise}}(0) \leq 1.91$ and $g^{(2)}_{\text{i,l,noise}}(0) \leq 1.36$.

These numbers describe the peak values assuming a single mode and uncorrelated noise. For comparison with our other results we estimate the value that is expected integrating over a 400 ns window (according to an equation similar to equation (2)) to be $g^{(2)}_{\text{s,s,noise}}(\Delta \tau) \leq 1.57$.

Apart from detection noise we also have to take the effect of multiple modes of the signal photons into account. It has been shown [48, 49] that the presence of several modes results in a reduced value following $g^{(2)}_{\text{s,i}}(0) = 1 + \frac{1}{N}$, where $N$ is the number of contributing modes. In the signal branch there is no narrowband filtering installed, which results in the contribution of all modes of the central cluster (see spectrum in figure 3). By determining the total intensity as the sum of detected counts in all modes and comparing it to the detections in the central mode, we infer a number of $N = 3.9$ effective modes (assuming all modes have the same intensity). This corresponds to $g^{(2)}_{\text{s,i,central,noise}}(\Delta \tau) = 1.15$, close to the measured value.

We can now use the measured data to prove the quantum character of the correlations. The Cauchy–Schwarz inequality $R = \frac{g^{(2)}_{\text{s,i}}^2}{g^{(2)}_{\text{s,i,noise}} g^{(2)}_{\text{i,s,noise}}} \leq 1$ is bound for classical correlations. When we insert the results for cross-correlation ($g^{(2)}_{\text{s,i}}(\Delta \tau) = 70(1)$ at 1 mW) and auto-correlation measurements we find $R = 2711 \pm 247$ for an integration window of 400 ns. Hence we violate the classical boundary by more than 10 error margins which is a proof of non-classical correlations.

7. Heralded narrow-band single-photon source

We next check the single photon character of the generated light. This is done by measuring the autocorrelation conditioned on the detection of a heralding photon. We use a gate window of 400 ns around the detection time of the idler photons (analog to previous measurements) to herald the signal photons. We then use a method introduced by Fasel et al [50] to generate histograms as shown in figure 7. The histogram illustrates the triple coincidences sorted by the number of heralding events between succeeding detections at different signal detectors. The ratio of events in the central bin (bin 0, coincidence of the signal detectors in the same heralding window) divided by the mean value of the outer bins corresponds to the conditioned autocorrelation $g^{(2)}_{\text{s,s,conditioned}} = 0.03(5)$. This value is considerably below the classical threshold $g^{(2)}_{\text{s,s,conditioned}} \geq 1$.

We performed such measurements of the conditioned autocorrelation for various pump powers. The results are shown in figure 5 (red diamonds). For all accessible pump powers the $g^{(2)}_{\text{s,s,conditioned}}$-value is well below the classical threshold, with a minimum of 0.01(4). Our source is therefore well suited to generate narrowband heralded single photons with high fidelity. The approximately linear dependence of the $g^{(2)}_{\text{s,s,conditioned}}$-values on pump power was also observed earlier [50]. The inflection of the curve and increase at very low pump powers indicates a regime where the detected photon rate is comparable to the dark count rate. The single photon correlation is then dominated by noise coincidences.
We expect the relation $g^{(2)}_{i:s,s} = \frac{g^{(2)}_{i:s} g^{(2)}_{s,s}}{g^{(2)}_{i,i}}$ \([21]\) between conditioned and unconditioned autocorrelation. Hence we can give an estimation with the help of the independently measured values. However, since we condition the detection of the signal photons to the idler, we project the signal photons to a single mode too. Thus we assume for the unconditioned autocorrelation value of the signal photons the theoretical value for a single-mode, corrected for dark counts, and retrieve a prediction of $g^{(2)}_{i:s,s} = 0.03$ for 1 mW of pump power. The red solid line in figure 5 is calculated with these autocorrelation values while for the cross correlation values the model curve (blue solid line in the same figure) was used. This is in good agreement with experimental results. It should be mentioned that the $g^{(2)}_{i:i}$-value here was measured with a low and a high dark count detector, while for the conditioned autocorrelation only the low dark count detector was used for heralding. Additionally, $g^{(2)}_{i:i}$ was measured only at a high pump power of 4.3 mW. Hence this value is underestimating the quality of the photons.

We summarize the different correlation results, and values derived by different methods in table 1 for comparison.

### 8. Conclusion

In summary, we developed a source of telecom-heralded single photons with high heralding efficiency and a high creation rate of narrowband photons compatible with a solid-state quantum memory based on Pr$^{3+}$: Y$_2$SiO$_5$. We proved the single photon character of the emission by direct measurement of the heralded autocorrelation function, with antibunching parameters as low as 0.01. In addition, we reported significant improvements in terms of count rate, quality of correlations, and heralding efficiency (up to 28%), compared to the state of the art. This source is well suited for advanced experiments with solid state quantum memories. In particular, it meets the requirements to demonstrate single-photon spin-wave storage in a solid-state quantum memory, which would be an important milestone for the use of solid-state quantum memories in quantum networks and quantum repeaters architectures.

**Acknowledgments**

We acknowledge financial support by the ERC starting grant QuLIMA, by the Spanish Ministry of Economy and Competitiveness (MINECO) and Fondo Europeo de Desarrollo Regional (FEDER) through Grant No. FIS2015-
Appendix A. Estimation of the escape efficiency

In general we have two approaches to estimate the escape efficiency of the photons. We can use the heralding efficiency and calculate backwards to the cavity output or we can use the parameters of the cavity to directly estimate the escape efficiency. Here we want to give a more detailed description of these two ways together with an error estimation.

In figure 4 we can read a heralding efficiency of \( \eta_h = 28\% \). The heralding efficiency is defined via the probability of detecting a coincidence \( p_{si} \), a heralding photon \( p_{pi} \) and the detection efficiency \( \eta_{det} \) by \( \eta_h = \frac{p_{si}}{p_{pi} \eta_{det}} \). This number can be affected by the measurement accuracy of the detector resulting in \( \eta_h = 28(1)\% \). From measurements of the idler spectrum, that look similar to the one shown in figure 3, we know that there is a floor of uncorrelated photons present. From the signal to noise ratio, visible in the spectrum, we can estimate that about 10% of the detected photons are not correlated with a signal photon. This results in an underestimation of the heralding efficiency, as \( p_{pi} \) is corrupted by noise. To estimate the escape efficiency, the heralding efficiency is divided by the transmission losses between cavity and detector, i.e. \( \eta_{esc}^{calc} = \frac{\eta_h}{\eta_c} \). For the transmission we found a value of 71(3)%. The error accounts for e.g. variations in the fiber coupling and accuracy of the power meter we used, and is found by repetitive measurements. This finally results in \( \eta_{esc}^{calc} = 39(2)\% \) or \( \eta_{esc}^{calc} = 44(2)\% \), if we take the uncorrelated noise into account.

On the other hand, looking at the performance of the cavity we take the uncorrelated noise into account.

Within the error bars we see an agreement between these two methods, resulting in an escape efficiency of \( \eta_{esc} = 56(13)\% \). The uncertainty on the re-fluence of the photons. We can use the heralding efficiency only depends on the internal losses \( L_{int} \) and the output coupling:

\[
\eta_{esc} = \frac{1 - R_{oc}}{1 - R_{oc} + L_{int}}.
\]

The manufacturer of the mirrors gives a reflectivity of \( R_{oc} = 0.9999 \). We solve the above equations numerically and find \( L_{int} = 2.4(7)\% \) and \( \eta_{esc} = 56(13)\% \). The uncertainty on the reflectivity of the output coupling mirror has a big influence on the error bar of the escape efficiency. We also remark that the internal losses include the absorption in the crystal (ca. 0.01 cm\(^{-1}\) [44]; corresponding to 2% loss for our case) and losses at the crystal facets (AR coating \( R < 1\% \)).

Within the error bars we see an agreement between these two methods, resulting in an escape efficiency between 40% and 50% for the signal photons. The escape efficiency of the idler photons is considerably higher due to the lower absorption in the crystal.

Appendix B. Width of autocorrelation function

The temporal shape of the input wave packet with a Lorentzian spectrum of width \( \Delta \nu \) can be described by

\[
f(t) = \exp(-2\pi \cdot \Delta \nu \cdot t) \cdot \Theta(t),
\]

where \( \Theta(t) \) is the Heaviside step function. The autocorrelation of such a signal can be calculated as follows:

\[
\begin{align*}
\alpha(\tau) &= \int_{-\infty}^{+\infty} f(t) \cdot f(t + \tau) \, dt \\
&= \int_{-\infty}^{+\infty} \exp[-2\pi \cdot \Delta \nu \cdot t] \cdot \Theta(t) \cdot \exp[-2\pi \cdot \Delta \nu \cdot (t + \tau)] \cdot \Theta(t + \tau) \, dt \\
&= \int_{-\infty}^{+\infty} \exp[-2\pi \cdot \Delta \nu \cdot (2t + \tau)] \cdot \Theta(t) \cdot \Theta(t + \tau) \, dt.
\end{align*}
\]
To solve the integral we can distinguish two cases for $\tau$:

$$
\tau \geq 0 \quad \Rightarrow \quad \Theta(t) \cdot \Theta(t + \tau) = \Theta(t),
$$

$$
a(\tau \geq 0) = \int_{-\infty}^{\infty} \exp[-2\pi \cdot \Delta \nu \cdot (2t + \tau)] \cdot \Theta(t) \, dt
= \exp[-2\pi \cdot \Delta \nu \cdot \tau] \cdot \int_{0}^{\infty} \exp[-2\pi \cdot \Delta \nu \cdot 2t] \, dt
= \exp[-2\pi \cdot \Delta \nu \cdot \tau] / 4\pi \Delta \nu,
$$

$$
\tau < 0 \quad \Rightarrow \quad \Theta(t) \cdot \Theta(t + \tau) = \Theta(t + \tau),
$$

$$
a(\tau < 0) = \int_{-\infty}^{\infty} \exp[-2\pi \cdot \Delta \nu \cdot (2t + \tau)] \cdot \Theta(t + \tau) \, dt
= \exp[-2\pi \cdot \Delta \nu \cdot \tau] \cdot \int_{-\infty}^{\infty} \exp[-2\pi \cdot \Delta \nu \cdot 2t] \, dt
= \exp[2\pi \cdot \Delta \nu \cdot \tau] / 4\pi \Delta \nu,
$$

which means in total:

$$a(\tau) = \frac{\exp(-2\pi \cdot \Delta \nu \cdot |\tau|)}{4\pi \Delta \nu}.
$$

The result for the autocorrelation is a symmetric function where both sides have a decay time given by the spectral width $\Delta \nu$. Hence the full width of the correlation function is twice the width we infer from the cross correlation measurement for signal and idler.

References

[1] Grangier P, Roger G and Aspect A 1986 Europhys. Lett. 1 173–9
[2] Afzelius M, Gisin N and de Riedmatten H 2015 Phys. Today 68 42–7
[3] Bussières F, Sangouard N, Afzelius M, de Riedmatten H, Simon C and Tittel W 2013 J. Mod. Opt. 60 1519–37
[4] Kaiser F, Issautier A, Ngah I A, Alibart O, Martin A and Tanzilli S 2013 Laser Phys. Lett. 10 045202
[5] Clausen C, Bussières F, Tiranov A, Herrmann H, Silberhorn C, Sohler W, Afzelius M and Gisin N 2014 New J. Phys. 16 093058
[6] Saglamyurek E, Sinclair N, Slater J A, Heshami K, Oblak D, O’Gradys and Tittel W 2011 Phys. Rev. Lett. 107 250501
[7] Ou Z Y and Lu Y J 1999 Phys. Rev. Lett. 83 2556–9
[8] Kuklewicz C E, Wong F N C and Shapiro J H 2006 Phys. Rev. Lett. 97 223601
[9] Bao X H, Qian Y, Jiang H, Chen Z B, Yang T and Pan J W 2008 Phys. Rev. Lett. 101 190501
[10] Scholz M, Koch L and Benson O 2009 Phys. Rev. Lett. 102 060603
[11] Wolfgramm F, de Icaza Astiz Y A, Beduini F A, Cerè A and Mitchell M W 2011 Phys. Rev. Lett. 106 053602
[12] Fekete J, Rieländer D, Cristiani M and de Riedmatten H 2014 Phys. Rev. Lett. 110 220502
[13] Förtsch M, Fürst J U, Wittmann C, Strekalov D, Aiello A, Chekhova M V, Silberhorn C, Leuchs G and Marquardt C 2013 Nature Commun. 4 14118
[14] Zhou Z Y, Ding D S, Li Y, Wang F Y and Shi B S 2014 J. Opt. Soc. Am. B 31 128–34
[15] Luo K H, Herrmann H, Krapick S, Brecht B, Ricken R, Quiring V, Suche H, Sohler W and Silberhorn C 2015 New J. Phys. 17 073039
[16] Förtsch M et al 2015 Phys. Rev. A 91 023812
[17] Wang J, Lv P Y, Cui J M, Liu B H, Tang S, Huang Y F, Li C F and Guo G C 2013 Phys. Rev. Appl. 4 066011
[18] Afzelius M and Benson O 2016 Appl. Phys. Lett. 108 021111
[19] Zhang H et al 2011 Nat. Photon. 5 628–32
[20] Rieländer D, Kuttler K, Ledingham P M, Gündoğan M, Fekete J, Mazzera M and de Riedmatten H 2014 Phys. Rev. Lett. 112 040504
[21] Clausen C, Usmani I, Bussières F, Sangouard N, Afzelius M, de Riedmatten H and Gisin N 2011 Nature 469 308–11
[22] Saglamyurek E, Sinclair N, Jin I, Slater J A, Oblak D, Bussières F, George M, Ricken R, Sohler W and Tittel W 2011 Nature 469 512–5
[23] Schunk G, Vogl U, Strekalov D V, Förtsch M, Sedlmeir F, Scheweif H G L, Göbelt M, Christiansen S, Leuchs G and Marquardt C 2015 Optica 2 773
[24] Lenhard A, Bock M, Becher C, Kucera S, Brito J, Eich P, Müller P and Eschner J 2015 Phys. Rev. A 92 063827
[25] Briegel H J, Dür W, Cirac J I and Zoller P 1998 Phys. Rev. Lett. 81 5932–5
[26] Sangouard N, Simon C, de Riedmatten H and Gisin N 2011 Rev. Mod. Phys. 83 33–80
[27] Simon C, de Riedmatten H, Afzelius M, Sangouard N, Zbinden H and Gisin N 2007 Phys. Rev. Lett. 98 190503
[28] Zhou Z Q, Hu X Y, Liu X, Chen G, Xu S, Han Y J, Li C F and Guo G C 2013 Phys. Rev. Lett. 110 070502
[29] Tang J S et al 2015 Nat. Commun. 6 6852
[30] Afzelius M, Simon C, de Riedmatten H and Gisin N 2009 Phys. Rev. A 79 052329
[31] Afzelius M, Usmani I, Amari A, Lauritzen B, Walther A, Simon C, Sangouard N, Minár J, de Riedmatten H, Gisin N and Kröll S 2010 Phys. Rev. Lett. 104 040503
[32] Gündoğan M, Ledingham P M, Kuttler K, Mazzera M and de Riedmatten H 2015 Phys. Rev. Lett. 114 230501
[33] Jobez P, Laplanche C, Timoney N, Gisin N, Ferrier A, Goldner P and Afzelius M 2015 Phys. Rev. Lett. 114 230502
[34] Heinze G, Hubrich C and Halflmann T 2013 Phys. Rev. Lett. 111 036001
[35] Hedges M P, Longdell J I, Li Y and Sellers M J 2010 Nature 465 1052–6
[36] Sabooni M, Li Q, Kröll S and Rippe L 2013 Phys. Rev. Lett. 110 133604
[37] Duan L M, Lukin M D, Cirac J I and Zoller P 2001 Nature 414 413–8
[38] Sekatski P, Sangouard N, Gisin N, de Riedmatten H and Afzelius M 2011 Phys. Rev. A 83 053840
[39] Kutluer K, Mazzera M and de Riedmatten H 2016 Manuscript in preparation
[40] Ledingham P M, Naylor W R, Longdell J J, Beavan S E and Sellars M J 2010 Phys. Rev. A 81 012301
[41] Ferguson K R, Beavan S E, Longdell J J and Sellars M J 2016 Phys. Rev. Lett. 117 020501
[42] Eckardt R C, Nabors C D, Kozlovsky W J and Byer R L 1991 J. Opt. Soc. Am. B 8 646–67
[43] Pomarico E, Sanguinetti B, Gisin N, Thew R, Zbinden H, Schreiber G, Thomas A and Sohler W 2009 New J. Phys. 11 113042
[44] Schwesyg J R, Kajiyama M C C, Falk M, Jundt D H, Buse K and Fejer M M 2010 Appl. Phys. B 100 109–15
[45] Waasem N, Fieberg S, Hauser J, Gomes G, Haertle D, Kuhnemann F and Buse K 2013 Rev. Sci. Instrum. 84 023109
[46] Kutluer K, Pascual–Winter M F, Dajczgewand J, Ledingham P M, Mazzera M, Chanelière T and de Riedmatten H 2016 Phys. Rev. A 93 040302
[47] Sekatski P, Sangouard N, Bussières F, Clausen C, Gisin N and Zbinden H 2012 J. Phys. B: At. Mol. Opt. Phys. 45 124016
[48] McNeil K J and Gardiner C W 1983 Phys. Rev. A 28 1560–6
[49] Christ A, Laiho K, Eckstein A, Cassemiro K N and Silberhorn C 2011 New J. Phys. 13 033027
[50] Fasel S, Alibart O, Tanzilli S, Baldo P, Beveratos A, Gisin N and Zbinden H 2004 New J. Phys. 6 163
[51] Chou C W, Polyakov S V, Kuzmich A and Kimble H J 2004 Phys. Rev. Lett. 92 213601