

# Realistic $g^{(2)}$ measurement of a PDC source with single photon detectors in the presence of background

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Received 28 August 2010, accepted 30 September 2010

Published online 21 February 2011

**Keywords** quantum optics, nonlinear optics

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The  $g^{(2)}$  correlation function of a light beam depends on its photon statistics. In quantum optics, this effect can be used to determine the effective mode number of a type II parametric downconversion source.

However, the outcome is very sensitive against background light from other sources. We consider the experimental challenges of this approach and introduce a simple model to account for influence of the background.

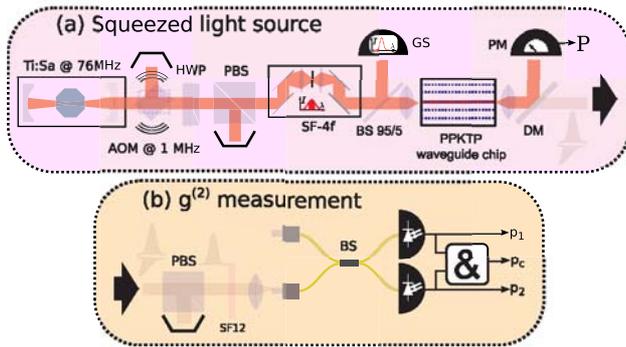
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Many modern quantum optical applications require sources of nonclassical states of light, most prominently single photons and squeezed light. In solid states physics, recent efforts concentrate on single photon sources based on semi-conductor quantum dots and color centers in crystal structures. In nonlinear optics, parametric downconversion (PDC) sources are an established standard to probabilistically generate photon pairs. In the type II PDC process, a pump photon decays inside a  $\chi^{(2)}$ -nonlinear medium into one horizontally polarized signal and one vertically polarized idler photon. The idler is used to herald the advent of the signal photon, thus creating a heralded single photon source by post-selection. Recent work has shown that PDC source engineering [1,2] is capable of producing spectrally separable two-photon states  $|1\rangle_s \otimes |1\rangle_i$ , allowing for the preparation pure heralded single photons [3]. Going beyond the single photon pair approximation, PDC in general can be understood as a source of squeezed states of light [4,5]. First observed by Slusher et al. [6] in 1985, squeezed states originally garnered interest for their noise reduction in the quadrature observables  $\hat{X}, \hat{Y}$  below the classical shot noise level, applicable in quantum-enhanced interferometry [7]. Until now, PDC experiments rely on spatial [3] or narrow spectral [8] filtering of multimode

[9,5] squeezers to approximate single mode bi-photonic states, with severe loss of source brightness. In [10], we introduce a source of two-mode squeezed vacuum states based on type II PDC in a single mode PP-KTP waveguide and spectral engineering [1–3,11], to avoid narrow-band filtering entirely. The strict photon number correlation between both output modes, in conjunction with efficient photon number resolving detection, promises high fidelity preparation of pure heralded Fock states. The two mode character of such a source can be demonstrated by measuring the  $g^{(2)}$  second order correlation of one of the source's output modes.

In this article, we demonstrate how to deal with experimental imperfections inherent to our  $g^{(2)}$  measurement experiment, namely a spectrally ultrabroad flux of background photons from the PDC process as well as dark count events of the InGaAs avalanche photo diodes (APDs) used as single photon detectors.

It has been shown early on in the experimental exploration of squeezing that PDC produces squeezed states of light [4]. In photon number representation, a two-mode



**Figure 1** Experimental setup: (a) Squeezed light source: A PP-KTP waveguide, spatially single mode at 1550 nm, is pumped with a mode locked Ti:Sa laser emitting ultrafast pulses with 8 nm FWHM. An acousto-optic modulator (AOM) reduces full repetition rate of 76 MHz to 1 MHz, a HWP+PBS combination controls pump beam power. We adjust pump spectral width with a 4f spectral filter setup (SF-4f) and monitor it with a grating spectrometer (GS). Pump light coupled through the waveguide is then separated from the generated signal and idler beams with a dichroic mirror (DM) and its power measured (PM). (b)  $g^{(2)}$  measurement: Background light is removed from the signal beam with a 12 nm FWHM spectral filter (SF12), then split at a 50/50 BS and each output arm fed into APDs. Single, coincidence and trigger event rates are recorded.

squeezed vacuum state has the form

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

where  $a$  and  $b$  are two orthogonal modes,  $\hat{S}_{a,b}$  is the two-mode squeezing operator, and  $\hat{H}_{a,b}$  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  $b$ . The photon number correlation between both modes allows for the creation of heralded Fock states using a photon number resolving detector, in the simplest case for pure heralded single photons with binary detectors. However, the underlying bilinear effective Hamiltonian  $\hat{H}_{a,b} = \zeta \hat{a}^\dagger \hat{b}^\dagger + \text{h. c.}$  describes only a special case of PDC.

In general, the effective PDC Hamiltonian has a richer spatio-spectral structure, but using a single mode waveguided approach allows us to neglect the spatial structure, 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  $\zeta$  determines the strength

of this interaction, while spectral correlations between photons of the pairs produced are governed by the normalized joint spectral amplitude

$$f(\omega_1, \omega_2) \propto \alpha(\omega_1 + \omega_2) \Phi(\omega_1, \omega_2) \quad (3)$$

where  $\alpha(\omega)$  is the spectral amplitude of the pump beam and  $\Phi(\omega_1, \omega_2)$  is the phasematching function that depends on the nonlinear medium's dispersion properties.

For low pump powers, PDC is in good approximation a probabilistic source of photon pairs. By applying a Schmidt decomposition to the pairs' joint amplitude [12]  $f(\omega_1, \omega_2) = \sum_k c_k \varphi_k(\omega_1) \psi_k(\omega_2)$ , we obtain two orthonormal basis sets  $\varphi_k(\omega_1)$  and  $\psi_k(\omega_2)$  and a set of weighting coefficients  $c_k$  with  $\sum_k |c_k|^2 = 1$ . Now the PDC Hamiltonian can be expressed in terms of broadband modes

$$\hat{H}_{\text{PDC}} = \sum_k \hat{H}_k = \zeta \sum_k c_k \left( \hat{A}_k^\dagger \hat{B}_k^\dagger + \hat{A}_k \hat{B}_k \right). \quad (4)$$

Each broadband mode operator  $\hat{A}_k, \hat{B}_k$  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  $\hat{B}_k^\dagger := \int d\omega \psi_k(\omega) \hat{b}^\dagger(\omega)$ . Since the effective Hamiltonians  $\hat{H}_k$  do not interact with each other (i. e.  $[\hat{H}_k, \hat{H}_l] = 0$ ), we see that the PDC time evolution operator is in fact an ensemble of independent two-mode squeezing operators  $\hat{U}_{\text{PDC}} = e^{i\hat{H}_{\text{PDC}}} = \hat{S}_{A_0, B_0} \otimes \hat{S}_{A_1, B_1} \otimes \dots$  where the coefficients  $c_k$  determine the relative strength of all squeezers as well as spectral correlation between signal and idler beams. The strength of correlation between the output beams of such a source is characterized by its effective mode number  $K = \frac{1}{\sum_k |c_k|^4}$ . 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).

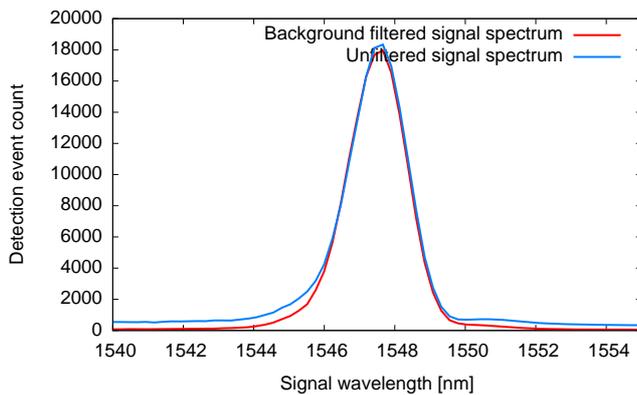
The second order correlation function  $g^{(2)}$  can be used to discriminate between beams with thermal ( $g^{(2)} = 2$ ) and Poissonian photon statistics ( $g^{(2)} = 1$ ) from a squeezing source [8]. As has been noted above, type II PDC can in general be seen as an ensemble of two-mode squeezers, 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 [13]. If, on the other hand, 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  $\zeta$ , we can find a simple connection to the effective mode number:  $g^{(2)} = 1 + \sum |c_k|^4 = 1 + \frac{1}{K}$ . In the presence of additional background events, we measure again a convolution of different photon statistics. This will always reduce the experimentally obtained value of  $g^{(2)}$  further towards 1.

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 the effective mode number  $K$  by simply adjusting the spectral width of the pump pulses [1,2].

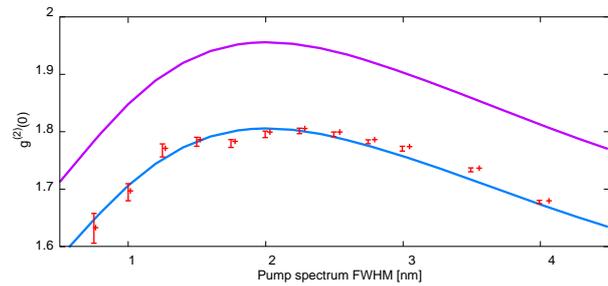
The setup in Fig. 1 (a) illustrates the PDC source: Ultrafast pump pulses at 768 nm are prepared with a TiSa mode locked laser system, spectrally filtered with a variable bandpass filter 4f setup, and then used to pump a type II PDC process within the PP-KTP waveguide with a poling period of 104  $\mu\text{m}$  by AdvR Inc. Central wavelengths of signal and idler beam were 1544 nm and 1528 nm, respectively. Fig. 1 (b) illustrates the  $g^{(2)}$  measurement: Idler is discarded, and the signal beam split by a 50/50 beamsplitter. The 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

$$g^{(2)} = \frac{\langle \hat{a}^\dagger \hat{a}^\dagger \hat{a} \hat{a} \rangle}{\langle \hat{a}^\dagger \hat{a} \rangle^2} \approx \frac{p_c}{p_1 p_2} \quad (5)$$



**Figure 2** Signal beam spectrum with background events in unfiltered spectrum (blue line) and dark count events only in background-filtered spectrum (red line)

Our waveguide source supports more PDC processes than the phase-matched type II process we are utilizing, most prominently the type I processes slow  $\rightarrow$  slow, slow and fast  $\rightarrow$  fast, fast. Being not phase-matched in the telecom wavelength range, we detect them as a low intensity background in a wide spectral range. Also production im-



**Figure 3**  $g^{(2)}$  values from setup 1(b) for a variable pump FWHM. (red) Experimental values. (blue) Theory curve according to Eq. (6) with  $R = \frac{1}{20}$ . (violet) Background corrected theory curve according to Eq. (5)

perfections in periodic poling structure or waveguide diameter add to this. Other fluorescence processes besides PDC, caused by color centers or faults in the nonlinear crystal, are an unlikely source, since one would expect distinctly different arrival time for their photons, due to long decay times compared to the time frames of our PDC process. We were however unable, with the equipment at our disposal allowing for timing resolution of ca. 500 ps, to observe any significant arrival time difference between signal beam photons and background photons. Whatever the photon statistics of these processes are, in combination with the thermal statistics of either beam of a perfect two-mode squeezed vacuum state they result in a “less thermal” distribution, i. e.  $g^{(2)}$  will drop below a value of 2 for a two-mode squeezer source with background noise.

Detector dark counts are most easily dealt with: We apply the least possible bias voltage to our APDs and use the smallest possible gate width of 2.5 ns to arrive at a dark count probability of  $5 \times 10^{-5}$ .

Background photons degrade our experimental results most dramatically. However, most of them can be removed from the signal beam by using a suitable bandpass filter. After introducing the SF12 filter in setup 1(b), our photon count dropped by a factor of five and we were able to see  $g^{(2)}$  values significantly different from one with a photon detection event probability of  $2 \times 10^{-2}$ .

This leaves the background photons that are transmitted by the bandpass filter SF12. Assuming identically efficient single photon detectors ( $p = p_1 = p_2$ ), we add the background photons as a statistically independent source of detector events with probability  $q$  at each detector. Coincidence detection events from these uncorrelated background photons happen with probability  $q^2$ . We substitute  $p \rightarrow p + q - pq$  and  $p_c \rightarrow p_c + q^2 - p_c q^2$  and find

$$g^{(2)} \approx \frac{p_c + q^2 - p_c q^2}{(p + q - pq)^2} \quad (6)$$

In order to estimate  $q$ , we compare the unfiltered spectral distribution of the signal beam to the same distribution with and without a spectral filter SF12 applied (Fig. 2). These

spectra have been measured with a fiber spectrometer [13, 10]. We see that the filtered spectrum quickly falls to the level of detector dark counts outside the main peak, while the unfiltered beam maintains a background event level 450 counts over the dark count level of 50 counts over a wide range. We can expect those to be present in the main peak in both instances, and the ratio  $R = \frac{q}{p}$  of detection events to background events equals the ratio of the areas of the background level “under the main peak” and the main peak itself. From the graph in Fig. 2, we estimate  $R = \frac{1}{20}$  and thus find with Eq. (6) the theory curve in Fig. 3 in excellent agreement with our experimental data. We use the background corrected curve to predict  $g^{(2)} = 1.95$  in the absence of background at an optimal pump FWHM of 1.95 nm. From the relationship between photon statistics and effective mode number [10] we determine the mode number  $K = \frac{1}{g^{(2)} - 1} = 1.05$  of our source.

In conclusion, we have demonstrated the degrading impact of background events in the  $g^{(2)}$  measurement used to determine the effective mode number of our PDC source, as well as effective counter measures: By driving our APDs with low bias voltage and short gate width, we reduce detector dark counts. By simply employing a spectral filter, we dispose of most of the background photon events. To negate the effects of residual background light in the same spectral range, we develop a simple model for  $g^{(2)}$  measurements in the presence of an independent background source of events, and thus can infer the measurement outcome in the absence of background.

**Acknowledgements** This work was supported by the EC under the grant agreement CORNER (FP7-ICT-213681).

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