Photon bunching in parametric down-conversion with continuous wave excitation

Bibiane Blauensteiner,[1, *] Isabelle Herbauts,[1, *] Stefano Bettelli,[1, *] Andreas Poppe,[2] and Hannes Hübel[1, †]

[1]*Quantum Optics, Quantum Nanophysics and Quantum Information, Faculty of Physics, University of Vienna, Boltzmanngasse 5, 1090 Vienna, Austria*

[2]*Austrian Research Centers GmbH - ARC, Donau-City-Str. 1, 1220 Vienna, Austria*

(Dated: June 8, 2009)

The first direct measurement of photon bunching ($g_{(2)}$ correlation function) in one output arm of a spontaneous-parametric-down-conversion source operated with a continuous pump laser in the single-photon regime is demonstrated. The result is in agreement with the statistics of a thermal field of the same coherence length, and shows the feasibility of investigating photon statistics with compact cw-pumped sources. Implications for entanglement-based quantum cryptography are discussed.

I. INTRODUCTION

Light sources based on *spontaneous parametric down-conversion* (SPDC) [1], in which photons from a pump laser are sporadically converted into pairs of photons at lower frequency in a non-linear crystal, have become an essential component in the toolbox of quantum optics laboratories in the last decade. These sources allow state-of-the-art experiments on entanglement and quantum information, due to the ability to supply photons that are highly correlated in time, energy and polarisation.

SPDC-based systems have been exploited as a basis for entangled-photons schemes [2] as well as approximations of single-photon sources in the so-called heralded-photon schemes [3]. Such setups are however limited by the possibility that additional pairs are created within the window determined by the pump-pulse duration or by the electronic gate width of the single photon detectors (for continuous sources). In multi-photon interference experiments, additional pairs in general affect the purity of the investigated state [4]. In quantum cryptography [5] instead, multi-photon pairs may constitute a security threat, since their information-carrying degree of freedom (typically polarisation) might be correlated [6].

The topic of SPDC-pair emission and correlation, deeply related to the statistical properties of the down-converted field, is therefore not only interesting for fundamental quantum optics experiments, but also concerns applied quantum information. Significant theoretical and experimental investigations have already been conducted on this subject, but a direct demonstration that the field of one arm of an SPDC source is thermal is still missing for the special case of a source in *single-photon regime* operated with a *continuous pump*. Such evidence is presented in this communication by measuring the second-order correlation function with a Hanbury Brown-Twiss setup [7] and comparing it to a prediction based both on the detector jitter and a first-order correlation measurement with a Michelson interferometer. This approach relies entirely on experimentally observable quantities and no assumptions are needed.

The intrinsic difficulty of the experiment stems from the fact that resolving times are much larger than the coherence time τ_c of the down-converted field, so that all relevant quantities are averaged out and the statistics becomes similar to a Poissonian one, with the true (thermal) statistics often overlooked.

The remaining part of the introduction presents a review of the theoretical background (Sec.I A) and existing experimental evidence (Sec.I B) concerning multi-pair emission, as well as a discussion of the potential security risk for quantum cryptography (Sec.I C). The experimental setup, data acquisition (Sec.II), and quantitative expectation model (Sec.III) are then described, followed by a discussion of experimental findings (Sec.IV).

A. Field statistics and theoretical description of SPDC

A direct characterisation of the *temporal* statistical properties of a generic quantum electromagnetic field is provided by the (normalised) correlation functions, defined as [8, ch.6]

$$g_{(1)}(t, t+\tau) = \frac{\langle \hat{E}^\dagger(t+\tau)\hat{E}(t)\rangle}{\langle \hat{E}^\dagger(t)\hat{E}(t)\rangle} = \frac{G_{(1)}(t, t+\tau)}{G_{(1)}(t,t)} \quad (1a)$$

$$g_{(2)}(t, t+\tau) = \frac{\langle \hat{E}^\dagger(t)\,\hat{E}^\dagger(t+\tau)\,\hat{E}(t+\tau)\,\hat{E}(t)\rangle}{\langle \hat{E}^\dagger(t)\,\hat{E}(t)\rangle\,\langle \hat{E}^\dagger(t+\tau)\,\hat{E}(t+\tau)\rangle} \quad (1b)$$

$$= \frac{G_{(2)}(t, t+\tau)}{G_{(1)}(t,t)G_{(1)}(t+\tau,t+\tau)}, \quad (1c)$$

where $\hat{E}$ are field amplitude operators. These quantities are, respectively, field and intensity correlations at times

*These authors contributed equally to the work.

†Corresponding author: hannes.huebel@gmail.com

t and $t + \tau$ at the same spatial location. Since for large time delays realistic fields are uncorrelated, $g_{(2)}(\infty) = 1$. The behaviour for finite delays depends however on the actual statistics of the field. If $g_{(2)}$ increases around $\tau = 0$, the field is said to be bunched.

A chaotic field in an interval short with respect to its coherence time τ_c will show thermal statistics, that is the distribution P_n of the number of photons is

$$P_n = \frac{\nu^n}{(\nu + 1)^{(n+1)}}, \qquad (2)$$

where ν is the average number of photons in the interval. Chaotic light, for which $g_{(2)}(0) = 2$, is therefore bunched. In contrast, the output of an ideal mono-mode laser displays Poissonian behaviour, that is $g_{(2)}(\tau) = 1$. Hence, in a short time interval, the probability of counting two photons for chaotic light is twice as large as that expected for Poissonian statistics.

With the picture of an SPDC process as a spontaneous emission of pairs from the splitting of pump photons, it could be naïvely believed that the statistics of pairs is similar to that of the pump field. It is however known [9] that linear interaction in a resonant medium affects incident light in such a way that it becomes a mixture of the initial "signal" field and an additional thermal field from spontaneous emission. If the incident field is absent (or already thermal), the output field is therefore exactly thermal. In analogy, since the SPDC process corresponds to optical amplification without input signal (as both signal and idler fields are initially in their vacuum states), the statistics of pairs needs not mimic that of the pump, and thermal behaviour is to be expected.

As for incoherent light sources, the quantum interpretation of SPDC bunching relies on constructive two-photon interference between probability amplitudes for the paths of two indistinguishable photons [10, 11]. It has been shown that this process can equivalently be interpreted as a pair emission stimulated by another down-converted pair [11, ch.7].

That the statistics of one arm of a mono-mode SPDC source is thermal was theoretically first proven in 1987 by Yurke and Potasek [12], who derived from the pure state of the combined single-mode down-converted fields the density operator of the mixed state of one field only by tracing out either signal or idler field. In this model, the quantum state $|\phi\rangle$ of the combined signal-idler fields in the Fock state expansion reads

$$e^{i(\eta\, a^\dagger b^\dagger + \eta^* ab)}|\emptyset\rangle = \mathrm{sech}|\eta| \sum_{n=0}^{\infty} \left(\frac{\eta}{|\eta|}\tanh|\eta| \right)^n |n\rangle_{si}$$
$$= (1 - \tfrac{1}{2}|\eta|^2)|\emptyset\rangle + \eta|1\rangle_{si} + \eta^2|2\rangle_{si} + O(|\eta|^3), \quad (3)$$

where η is a parameter proportional to the pump amplitude, $a^\dagger$ ($b^\dagger$) is the creation operator for the idler (signal) mode with fixed polarisation, and $|n\rangle_{si} = \frac{1}{n!}(a^\dagger b^\dagger)^n|\emptyset\rangle$ is a Fock state with n pairs. By tracing over one field, say i, the dominant signal-idler correlation is suppressed, and

a mixed density matrix is obtained for the residual field that displays exactly thermal statistics:

$$\rho_s = \sum_{n=0}^{\infty} P_n |n\rangle_s\langle n| = \sum_{n=0}^{\infty} \frac{\nu^n}{(\nu + 1)^{(n+1)}}|n\rangle_s\langle n|, \quad (4)$$

where $\nu = \sinh^2|\eta|$ and P_n is the probability of finding exactly n signal photons. The SPDC process is nevertheless inherently quantum mechanical, and its "thermality" is not an effect of the trace operation, but originates directly from the nature of the pair production process [12].

Since SPDC production is far from being mono-mode, Yurke and Potasek's analysis was not realistic; however, the thermal nature of a multi-mode SPDC process was later confirmed by Tapster and Rarity [13], and in a more general way by Ou, Rhee and Wang [4, 14], who analysed the case of sources with *pulsed pumping*. In their derivation, what is actually measured is the pulse integral of $G_{(1)}$ and $G_{(2)}$, *i.e.*, the probabilities of one or two photons in the pulse. With time resolution limited to the whole pulse, one must resort to test the "bunching excess", which was shown [4] to be essentially the ratio of the coherence time τ_c of the photons in the inspected arm and the duration τ_p of the pulse, when $\tau_c \ll \tau_p$. Shorter pump pulses then increase the excess, but, contrary to intuition, the limit of a very short pulse is not sufficient alone to guarantee a full bunching peak ($g_{(2)}(0) = 2$). In [11, ch.7] it is shown that this condition can only be achieved with the help of narrow spectral filtering.

SPDC sources have also been implemented with continuous pumping, using *continuous wave* (cw) lasers, which are both more practical and more affordable than femto-second (fs) pulsed SPDC sources. The field emitted by cw sources, and its statistical properties, is still relatively unexplored. Theoretically, it is usually assumed that resolving times are so much larger than the coherence time of the down-converted field that all relevant quantities are not accessible. It has been stated [15, 16, 17, 18] that the statistics of SPDC as a multi-mode process, *i.e.* when the pulse duration of the laser is much longer than the coherence time of the produced photons, will lead to Poissonian statistics. Careful experimental design and data analysis, as shown in this paper, reveal however the inherent thermal statistics, hidden by an apparent Poissonian behaviour.

B. Current experimental evidence

Photon statistics in SPDC processes can be directly accessed using a Hanbury Brown-Twiss setup by looking at photon emission in a single output arm only, *i.e.* by measuring the $g_{(2)}$ of the signal or idler arm. In such an investigation however, jitter and arrival-time discretisation smear the result. Due to the aforementioned disproportion between the time resolution of available single-photon detectors and the SPDC coherence time, most experimental investigations have exploited fs-pulsed sources

and strong filtering to increase artificially the coherence time of the measured field. Positive evidence in single-photon regime was first reported by Tapster and Rarity [13], who compared quantitative predictions based on a specific experimental hardware with measurements performed with various filters; with the narrowest one, a very convincing bunching peak of 1.85 could be demonstrated. Six years later, de Riedmatten *et al.* [19] (almost) replicated the experiment, with similar shortcomings and results, concluding that the amount of bunching depends only on the ratio τ_c/τ_p. In both cases an apparent transition from thermal to Poissonian statistics was observed when τ_c was varied below τ_p. Söderholm *et al.* [17] studied the dependence of single- and double-click rates on pump power in the two cases $\tau_c \sim 1.75\tau_p$ and $\tau_c \sim 0.6\tau_p$, and found good agreement with a thermal model in the first case, and an intermediate behaviour in the second case. On the other hand, Mori *et al.* [18] studied a source with very long pump pulses (40ps) and broad filtering leading to a rather short coherence time τ_c ($\sim$ 120fs), and failed to detect any significant bunching. Similarly, in a recent article, Avenhaus *et al.* [15], using 60ps laser pulses, observed Poissonian statistics for the SPDC process (τ_c not stated).

In an experiment where, contrary to the single-photon regime, several photons per pulse are created, Paleari *et al.* [20] were able to demonstrate, with careful fitting of the measured distribution, a thermal behaviour, although their ratio τ_p/τ_c was rather unfavourable (in the range 10-20). Vasilyev *et al.* [21] used a strongly filtered optical homodyne tomography to extract the photon-number distribution of a SPDC process and compare it with a thermal one, finding almost perfect agreement[32].

SPDC sources have also been implemented using *continuous-wave* (cw) pumping; this alternative approach is still relatively unexplored due to the technical difficulty of measuring very short correlation times. Larchuk *et al.* [16] tried to investigate sources in the single-photon regime with a single detector and delay lines but failed to obtain any evidence because the dead time, about 1μs, makes the apparatus blind to the interesting region $\tau \lesssim$ 1ps. Super-Poissonian behaviour was found by Zhang *et al.* [22] by directly observing the photocurrents of orthogonally polarised twin beams from a continuously driven KTP crystal in a cavity.

Measurements of the marginal, *i.e.* one-arm, SPDC field are not to be confused with measurements of the better known statistics of the whole bi-photon state, as well as conditioned statistics, which have been experimentally confirmed, among others, by [23, 24, 25, 26].

C. Quantum key distribution (QKD) and SPDC bunching

Sources based on continuously pumped SPDC are often employed in quantum cryptographic devices, generating the multi-photon pair state of Eq.3. Correlations between multi-photon pairs constitute a potential threat to the security of QKD as the exchange is then susceptible to a photon-number-splitting attack [27]. The case of *entangled-photon* sources [28] is still relatively unexplored, and its solution depends on a thorough understanding of the extent to which multiple SPDC pairs are correlated.

In entangled-photon sources, the state, to first order, corresponds to a pair completely entangled in the information-carrying degree of freedom (polarisation); *e.g.* for type-II SPDC $|\psi^-\rangle = (|H_s, V_i\rangle - |V_s, H_i\rangle)/\sqrt{2}$. When the second order of the expansion (Eq.3) is considered (two pairs), two limiting cases emerge: *sister* pairs are in well distinguishable temporal modes, and are therefore completely uncorrelated in polarisation, with all polarisation combinations equally likely,

$$|\text{sisters}\rangle = \frac{|2H, 2V\rangle + |2V, 2H\rangle - \sqrt{2}|HV, VH\rangle}{2}, \quad (5)$$

while *twin* pairs are in the same temporal mode (within τ_c), and

$$|\text{twins}\rangle = \frac{|2H, 2V\rangle + |2V, 2H\rangle - |HV, VH\rangle}{\sqrt{3}}. \quad (6)$$

For twin states it has been shown that the probability that the two signal photons have the same polarisation is $\frac{2}{3}$ [6]; therefore, whereas no kind of photon-number-splitting attack is possible with sister states, twin states are potentially dangerous for quantum key distribution. The situation is analogous for type-I SPDC.

The works of Tsujino *et al.* [29] and Ou [30] have shown that polarisation correlation and one-arm bunching are two sides of the same phenomenon: the peculiar properties of the state space of identical particles imply that, for delays shorter than τ_c, signal photons with the same polarisation occur more frequently than expected for classically uncorrelated objects (as differently polarised photons).

Summarising, the study of one-arm $g_{(2)}$ in entangled-photon sources [33] gives information about the maximum possible extent of the security threat posed to quantum key distribution by higher orders of SPDC.

II. EXPERIMENTAL SETUP AND DATA ACQUISITION

In our SPDC source [34] (Fig.1), a 532nm cw laser pumps a single 30mm temperature-stabilised nonlinear crystal (periodically-poled KTP) for type-I downconversion (532nm $\rightarrow$ 810nm + 1550nm). The signal and idler photons have fixed polarisations, central wavelengths of respectively 810nm and 1550nm, and are emitted collinearly. The signal coherence time, as measured with a Michelson interferometer (Fig.3), is $\tau_c \sim$ 2.8ps FWHM, corresponding to a bandwidth of less than 1nm.

Signal and idler photons are separated according to their frequency, and coupled into single-mode optical fibres. Signal photons (at 810nm) are then recollimated and sent to a non-polarising balanced beam splitter. The resulting beams are collected into multi-mode fibres. The coupling efficiency is greater than 90% as the light is coupled from single-mode fibres into multi-mode ones. The multi-mode fibres direct the beams into a *PerkinElmer* SPCM-AQ4C single-photon detector array. The parameters of each detector are as follows: quantum efficiency at 810nm $\sim 50\%$, dark-count rate < 500Hz, dead time ~ 50ns, and detector saturation level above 1MHz. All detector clicks are recorded by a time-tagging unit (TTU) (*TT8, smart systems Division, ARC*) with a common time basis and an intrinsic resolution of $\tau_{TT} = 82.2$ps. Time tags are then processed to define coincidences.

One of the two fibres is a 100m spool, so that it operates as a delay line of approximately 500ns (this delay is removed by software during data analysis). The purpose of the delay line is to make photons that happen to come with a picosecond delay to impinge on detectors at sufficiently different times to avoid electronic cross-talk and increase sensitivity.

The timing jitter of the overall detection unit was characterised by measuring correlations in arrival times of signal-idler photon-pairs from an additional degenerate SPDC source at 810nm (405nm $\rightarrow$ 810nm + 810nm). This auxiliary photon-pair source is based on a Sagnac configuration [31]. For the present measurement of jitter, it was operated to yield $\sim$ 1MHz detected single rates, below detector saturation. The coincidence rate of the photon-pair source was approximately 100kHz. The combined jitter of the detection unit is estimated to be $\tau_j = 640$ps $\gg \tau_c$. This estimated FWHM is however not used in further experimental analysis, and no assumption on the shape of the jitter is needed. Instead, the fully normalised jitter curve, as seen in the cross-correlation in Fig.2, is used to characterise the detection unit, the only extracted parameter being the value of the area under the jitter curve (611ps). Several measurements of the jitter histogram were taken, and the deviation on the jitter area is estimated to be less than 3%.

A cross-correlation histogram is obtained from recorded detector clicks. A coincidence from detector D_1 to detector D_2 with delay $\Delta t = t_2 - t_1$ is counted if D_1 clicked at time t_1, and D_2 clicked at time $t_2 > t_1$, irrespective of any other click. Coincidences with the role of detectors reversed are defined accordingly, and the two functions are joined by arbitrarily defining the delays Δt from D_2 to D_1 as negative.

The coincidence density for uncorrelated photons, *e.g.* when Δt is large with respect to any coherence time, is the product of the individual experimental detection rates λ_1 and λ_2 of D_1 and D_2 (which include the effects of dead times and dark counts). Therefore, every bin of the coincidence histogram will contain, on average, and in the case of uncorrelated photons, $C = \lambda_1 \lambda_2 \tau_{TT} T$ counts, where T is the duration of a data-taking run. Since
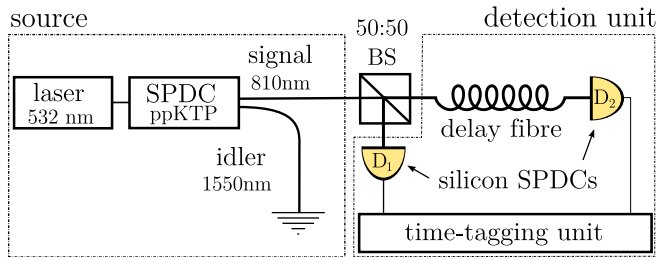

FIG. 1: (colour online) The experimental setup consists of a SPDC source of photon pairs, a balanced beam splitter and two single-photon silicon detectors, combined in a Hanbury Brown and Twiss [7] configuration.

C contains both single rates λ_1 and λ_2, it is directly proportional to the square of the pump power.

III. MODEL OF EXPECTED PHOTON BUNCHING AND MEASUREMENTS

Since the jitter τ_j of the detection unit is much larger than the coherence time τ_c of the down-converted light, any $g_{(2)}$ peak will be strongly smeared. A rough estimate for the residual peak height is given by the ratio of the coherence time to the jitter $\tau_c/\tau_j \sim 4 \times 10^{-3}$.

For a more accurate calculation, we model how a theoretical thermal bunching peak $(g_{(2)}(0) = 2)$ will be transformed due to our instruments and then compare it with the experimental result. SPDC-bunching is characterised by the same relation between $g_{(2)}$ and the first order coherence $g_{(1)}$ that holds for chaotic light [8], namely:

$$g_{(2)}(\tau) = 1 + |g_{(1)}(\tau)|^2 \qquad (7)$$

The $g_{(1)}$ of the signal beam was measured in a Michelson interferometer and can be seen in Fig.3. The setup for this measurement consisted of a Michelson interferometer with a motorised movable mirror and some limiting apertures in both arms to increase visibilities. The maximum visibility of the interferogram was measured to be 90%. The mirror is set to move at a constant speed of 0.002 mm/s over a range of 3mm, and detection counts were recorded over integration windows of 2 ms. At unequal path lengths, the detector count rate at the output of the interferometer was in the range of 200kHz. The interference fringes in Fig.3 were obtained from the raw data by subtracting the background rates (around 5kHz), and are displayed with an averaging of 30 data points to reduce fluctuations away from the centre of the interferogram.

The envelope of the interferogram as shown in Fig.3, that is $|g_{(1)}|$, is calculated from the visibility of the interference fringes in the raw data (see inset in Fig.3). An estimate of the coherence time $\tau_c \sim 2.8$ps is obtained from the graph. After rescaling the data points of the $|g_{(1)}|$ envelope to compensate for the reduced visibility (90%) in the interferogram, an area of 2.17ps was found for the
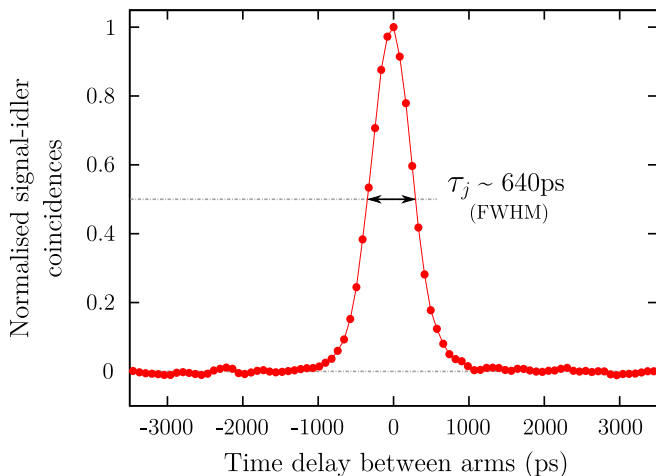
FIG. 2: (colour online) Second-order correlation function between the idler and signal arms of an *auxiliary* degenerate SPDC source at 810nm (405nm → 810nm + 810nm), rescaled to unity. Since the time correlation between photons of the same pair is tight (sub-picoseconds), the FWHM, approximately 640ps, is essentially an estimate of the combined jitter of the two silicon detectors and the time-tagging unit.
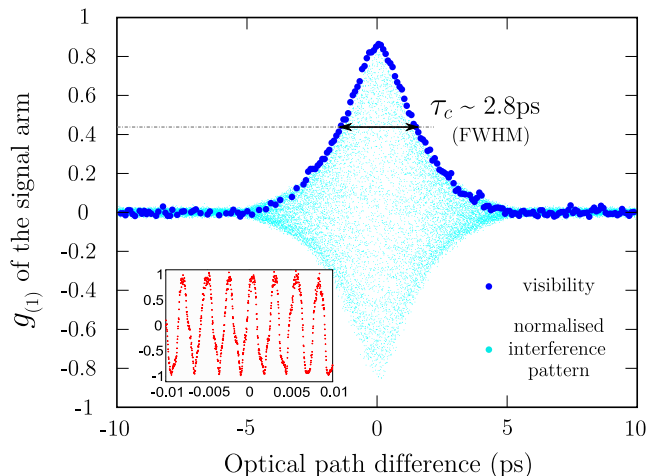

FIG. 3: (colour online) $g_{(1)}$ of the 810nm arm of the SPDC source as measured in a Michelson interferometer vs. optical path difference. $g_{(1)}$ was calculated from the fringe-visibility of the interferogram. The experimental visibility at zero optical path difference is slightly less than 1. When the maximum of the visibility is rescaled to one, the total area under the *squared* envelope is ~ 2.2ps. The inset shows the visibility fringes around the central part of the interferogram.

squared envelope $|g_{(1)}|^2$. The error on this value can be estimated by considering the area of several interference scans in identical experimental conditions; a conservative estimate is a 4% error on the calculated area for the squared envelope.

To model the washed out $g_{(2)}$ peak, we note that the total number of coincidences in the excess peak is preserved despite the jitter. Hence we expect to observe a smeared peak with the same temporal profile as that in Fig.2, but with an area equal to that of $|g_{(1)}|^2$ [35]. This method does not rely on any *a priori* knowledge of the actual shapes of $g_{(1)}$ or $g_{(2)}$, but only on a ratio of areas. The estimated coherence time τ_c and jitter τ_j of the detection unit from the FWHM of their respective histogram are not used in the calculation, neither is any assumption on curves shapes necessary.

The correlation function $g_{(2)}$ of the field from one arm of a SPDC source was measured using the Hnbury Brown-Twiss-like setup introduced in the previous section. A cross-correlation histogram for positive and negative delays was obtained from detector clicks and the bunching peak emerged in the region $\Delta t \sim 0$ over the plateau of accidental coincidences. The average count number of the plateau was then used to normalise the histogram. Experimental data were collected during a 16 hours data-taking session with an average count rate of ~ 1 MHz in each detector (below saturation), giving a plateau count of $N \simeq 4.4 \times 10^6$ events per bin. The observed rates used for jitter characterisation were set to be identical to the rates in this measurement, since detector jitter is dependent on actual count rates. The double pair term in Eq.3 scales with the square of the pump power, as does the plateau (see the estimate C at the end of

Sec.II), hence the normalised $g_{(2)}$ function is independent of pump power. The same argument can be used to prove the peak to be independent of optical losses in the setup.

In Fig. 4, we show a comparison of the measured data to the expected $g_{(2)}$ of our model. Data points correspond to the experimental $g_{(2)}$ histogram of the marginal SPDC field. The solid line is the expectation for the bunching peak of a field with thermal statistics. This line was obtained by vertically rescaling the jitter function shown in Fig.2 to yield the same area as $|g_{(1)}|^2$. The comparison does not involve any free parameter to be fitted: the theoretical model is only based on the specific and independently measured parameters of the setup.

The peak was shifted laterally by ~ 50ps for better comparison with the experimental data; this is however less than the intrinsic resolution τ_{TT}. Statistical fluctuations of the number of events in a bin of the histogram are of the order of $\sqrt{N}$. The cumulative error of $\pm 5\%$ on the expected $g_{(2)}$ curve, arising from a conservative error estimation from the jitter and $|g_{(1)}|^2$ data, is shown in Fig.4 as a shaded area. Even this 10% error margin fits the data points, and therefore the good match of the predicted peak is not accidental.

The emerging bunching from the SPDC field is in excellent agreement with the theoretical estimation of a thermal statistics. The peak protrudes 6 standard deviations above the Poissonian limit of $g_{(2)} = 1$ and is therefore experimentally confirmed.
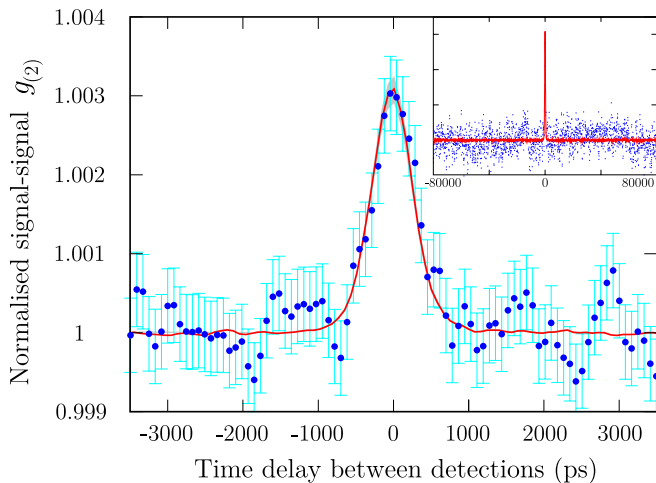
FIG. 4: (colour online) Second-order correlation function (signal-signal $g_{(2)}$) of the 810nm arm of a SPDC source, normalised over a measured plateau of N counts in a bin. The solid line represents the washed-out bunching peak (original height of 2), obtained from the jitter plot rescaled to an area of 2.17ps. Error bars, representing statistical fluctuations of the measured data, are set to $\sqrt{N}$. The shaded grey area, only visible at the very top of the peak, represents the uncertainty in the model prediction. The inset shows an enlarged section of the plateau.

IV. CONCLUSIONS

We experimentally measured the marginal temporal second-order correlation function $g_{(2)}(\tau)$ of SPDC light with continuous wave excitation in the single-photon regime. To our knowledge, the results represent the first direct observation of photon bunching for such a field, and strongly confirm its thermal character. Even with an experimental excess peak $g_{(2)}(0)$ as low as 10^{-3}, our experimental accuracy is sufficient to properly observe the residual bunching peak, which can be modelled solely on the experimental timing jitter and the output of a first-order interferometric measurement ($g_{(1)}$). With much

higher timing precision (e.g. lower detector jitter), it is expected that a full peak with the shape given by Eq.(7) is recovered. The relevant parameter is the ratio between the coherence length τ_c and the measurement resolution (the jitter τ_j in the cw case).

We therefore conclude that thermal photon statistics of the marginal SPDC field can be observed both in the fs-pulsed and in the cw regime, the latter being more practical and more affordable. A multimode excitation does not change the statistics of SPDC but alters only the experimental conditions, so that without a careful analysis an *apparent* Poissonian character is observed (as presented in some works reviewed in section I B), instead of the inherent thermal statistics. We have shown in this paper that even with far-from-ideal detection modules, and without any additional spectral filtering, we could demonstrate as a proof of principle that the true nature of the field remains accessible.

It is hoped that the results presented here stimulate further investigations into the nature of SPDC light and its thermal properties. For cw pumping, there is no clear consensus on the actual state of multi-photon pairs, nor is its extension to entangled states produced by cw pumping well understood. Similar experiments would provide experimental data to confirm that bunching implies polarisation correlations. In addition, cw pumping allows a continuous temporal investigation in a time domain not accessible to experiments with pulsed pumping and hence, with an improved timing resolution, access to the shape of $g_{(2)}$, and not only to its integral (a single data point) as in pulsed setups.

[1] D. Bouwmeester, A. K. Ekert, and A. Zeilinger, eds., *The Physics of Quantum Information* (Springer, 2000), ISBN 978-3540667780, and references therein.

[2] P. G. Kwiat, K. Mattle, H. Weinfurter, A. Zeilinger, A. V. Sergienko, and Y. Shih, Phys. Rev. Lett. **75**, 4337 (1995).

[3] C. K. Hong and L. Mandel, Phys. Rev. A **31**, 2409 (1985), and references therein.

[4] Z.-Y. J. Ou, J.-K. Rhee, and L. J. Wang, Phys. Rev. A **60**, 593 (1999).

[5] V. Scarani, H. Bechmann-Pasquinucci, N. J. Cerf, M. Dušek, N. Lütkenhaus, and M. Peev (2008), accepted in Rev. Mod. Phys., `arXiv:0802.4155 [quant-ph]`.

[6] M. Dušek and K. Brádler, J. Opt. B: Quantum Semiclass. Opt. **4**, 109 (2002), `quant-ph/0011007`.

[7] R. H. Brown and R. Q. Twiss, Nature **177**, 27 (1956).

[8] R. Loudon, *The Quantum Theory of Light* (Oxford University Press, 1983), II ed., ISBN 0-19-851155-8.

[9] S. Carusotto, Phys. Rev. A **11**, 1629 (1975).

[10] U. Fano, Am. J. Phys. **29**, 539 (1961).

[11] Z.-Y. J. Ou, *Multi-Photon Quantum Interference* (Springer US, 2007), ISBN 978-0-387-25532-3.

[12] B. Yurke and M. Potasek, Phys. Rev. A **36**, 3464 (1987).

[13] P. R. Tapster and J. G. Rarity, J. Mod. Opt. **45**, 595 (1998), `arXiv:0802.4155v2 [quant-ph]`.

[14] Z.-Y. J. Ou, Quantum Semiclass. Opt. **9**, 599 (1997).

[15] M. Avenhaus, H. B. Coldenstrodt-Ronge, K. Laiho, W. Mauerer, I. A. Walmsley, and C. Silberhorn, Phys. Rev. Lett. **101**, 053601 (2008), `arXiv:0804.0740 [quant-ph]`.

[16] T. S. Larchuk, M. C. Teich, and B. E. A. Saleh, in *Fun-*

damental Problems in Quantum Theory: A Conference Held in Honor of Professor John A. Wheeler (1995), no. 755 in Annals of the New York Academy of Sciences, pp. 680–686.

[17] J. Söderholm, K. Hirano, S. Mori, S. Inoue, and S. Kurimura, in *Proc. of the 8th Int. Symposium on Foundations of Quantum Mechanics in the Light of New Technology*, edited by S. Ishioka and K. Fujikawa (World Scientific, Singapore, 2006), pp. 46–49, ISBN 981-256-858-1, `quant-ph/0511128`.

[18] S. Mori, J. Söderholm, N. Namekata, and S. Inoue, Opt. Commun. **264**, 156 (2006), `quant-ph/0509186`.

[19] H. de Riedmatten, V. Scarani, I. Marcikic, A. Acín, W. Tittel, H. Zbinden, and N. Gisin, J. Mod. Opt. **51**, 1637 (2004), `quant-ph/0310167`.

[20] F. Paleari, A. Andreoni, G. Zambra, and M. Bondani, Opt. Express **12**, 2816 (2004).

[21] M. Vasilyev, S.-K. Choi, R. Kumar, and G. M. D'Ariano, Opt. Lett. **23**, 1393 (1998).

[22] Y. Zhang, K. Kasai, and M. Watanabe, Opt. Lett. **27**, 1244 (2002).

[23] O. Haderka, J. Peřina jr., M. Hamar, and J. Peřina, Phys. Rev. A **71**, 033815 (2005), `quant-ph/0405118`.

[24] E. Waks, B. C. Sanders, E. Diamanti, and Y. Yamamoto, Phys. Rev. A **73**, 033814 (2006), `quant-ph/0602118`.

[25] D. Achilles, C. Silberhorn, and I. A. Walmsley, Phys. Rev. Lett. **97**, 043602 (2006), `quant-ph/0606170`.

[26] M. Tengner and D. Ljunggren (2007), submitted to Phys. Rev. A, `arXiv:0706.2985 [quant-ph]`.

[27] G. Brassard, N. Lütkenhaus, T. Mor, and B. C. Sanders, Phys. Rev. Lett. **85**, 1330 (2000).

[28] A. Treiber, A. Poppe, M. Hentschel, D. Ferrini, T. Loruenser, E. Querasser, T. Matyus, H. Huebel, and A. Zeilinger, New J. Phys. **11**, 045013 (2009).

[29] K. Tsujino, H. F. Hofmann, S. Takeuchi, and K. Sasaki, Phys. Rev. Lett. **92**, 153602 (2004), `quant-ph/0311075`.

[30] Z.-Y. J. Ou, Phys. Rev. A **72**, 053814 (2005).

[31] A. Fedrizzi, T. Herbst, A. Poppe, T. Jenewein, and A. Zeilinger, Opt. Express **15**, 15377 (2007).

[32] Homodyne detection with a 10MHz bandpass filter, hence *extremely* tight filtering.

[33] As opposed to studying $g_{(2)}$ in attenuated lasers and heralded-photon sources.

[34] This polarisation-entangled source was designed for quantum cryptography, and will be the subject of a future publication [M. Hentschel, H. Hübel, A. Poppe and A. Zeilinger, *Three-color Sagnac source of polarization-entangled photon pairs*, 2009, arXiv, to be submitted to Optics Express]. Here the source was employed in unentangled mode, with down-converted fields having fixed linear polarisation.

[35] This analysis does not take into account spatial (transversal) coherence. Due to the use of single mode fibres, the setup was sufficiently optimised for such considerations to have no major implications in the model.