Macroscopic Hong–Ou–Mandel interference

T Sh Iskhakov1, K Yu Spasibko1,2, M V Chekhova1,2,3,4 and G Leuchs1,3

1 Max-Planck-Institute for the Science of Light, Erlangen, Germany
2 Faculty of Physics, M V Lomonosov Moscow State University, Moscow, Russia
3 University of Erlangen-Nürnberg, Staudtstrasse 7/B2, D-91058 Erlangen, Germany
E-mail: maria.chekhova@mpl.mpg.de

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Abstract. We report on a Hong–Ou–Mandel interference experiment for twin beams with photon numbers per mode as large as 106 generated via high-gain parametric down conversion (PDC). The standard technique of coincidence counting leads in this case to a dip with a very low visibility. By measuring, instead of coincidence counting rate, the variance of the photon-number difference, we observe an extremely well-pronounced peak. From the shape of the peak, one can infer information about the spectral properties of the PDC radiation, including the number of frequency/temporal modes.

4 Author to whom any correspondence should be addressed.

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1. Introduction

One of the most fascinating manifestations of quantum interference is the Hong–Ou–Mandel (HOM) dip, an effect that started a whole new direction in quantum optics. First reported in 1987 [1, 2], it consists of the interference of two single-photon wavepackets on a 50% beamsplitter (BS). When indistinguishable photons simultaneously arrive at different input ports of a BS, they always depart from the same port, due to the destructive interference of the probability amplitudes to be both reflected or both transmitted [3]. The effect is usually observed as a ‘dip’ in coincidences, its width determined by the coherence time of the photons and not affected by the group-velocity dispersion [4]. HOM interference is widely used for the identity tests of single-photon states [5], including ones generated by a quantum dot and a laser [6], conditionally prepared via parametric down-conversion (PDC) and a highly attenuated coherent state [7], produced via two consequent emissions of a single atom [8] or molecule [9] or by two different molecules [10]. Other important applications are the implementation of a photon Fock-state filter [11] and quantum-optical coherence tomography [12]. It is worth mentioning that if, as in most cases, the photons are generated via PDC, the level of coincidence rate in the ‘dip’ is not zero but is given by the accidental coincidences rate [13]. For this reason, the visibility of the dip reduces down to 33% at high-gain PDC [14] and becomes vanishingly low if many modes are involved.

HOM effect is known to have a classical analogue. It can be observed for coherent beams, although with only 33% visibility: the coincidence counting rate drops twice in the dip, to the level of half the accidental rate of coincidences [13]. In other words, the second-order normalized correlation function is equal to 1 outside of the dip and to 0 in the dip. This effect has been recently used for testing the fidelity of quantum memories [15].

In this paper we expand the boundaries of the HOM effect applications and observe destructive HOM interference for bright twin beams (containing, on the average, up to $8 \times 10^5$ photons per mode) generated via high-gain PDC in a traveling-wave optical parametrical amplifier. Namely, instead of two single-photon states we send to the BS two macroscopic states containing exactly equal numbers of quanta. We suggest an alternative approach of observing the HOM effect, based on the measurement of the normalized variance of the difference signal at the output ports of the BS. This technique is robust against the multi-mode detection [16].

Under certain conditions, the interference can be constructive see [3].
and allows one to observe the interference with a high visibility. We also study this effect for classical beams, coherent and thermal ones, and see a similar behavior.

The paper is organized as follows. Section 2 gives a simple qualitative description of the effect; section 3 provides its rigorous theory for various types of light: twin-beam squeezed vacuum, thermal light and coherent light. The experiment is described in section 4, and section 5 summarizes and discusses the results.

2. The effect

If a pair of indistinguishable photons arrives at the input ports of a balanced (50\%) BS, it always leaves through a single output port (figure 1). For a flux of photon pairs at the input, there is no correlation between the counts in the two output ports but instead, photon superbunching can be observed in each output port. A much more subtle effect originating from this superbunching is the appearance of a ‘dip’ in the counting rate of a single detector placed in one of the output ports [18].

A full theoretical description of the intensity correlations at the output of a BS, with the consideration of both classical (coherent and thermal) and quantum (two-photon and squeezed vacuum) states has been given by Klyshko [13] who has formulated ‘a conservation law for the sum of correlations and fluctuations’. Indeed, the sum of second-order intensity autocorrelation functions for outputs 1, 2, \( G_{11} \equiv \langle N_1^2 \rangle \) and \( G_{22} \equiv \langle N_2^2 \rangle \), and the doubled second-order intensity cross-correlation function \( G_{12} \equiv \langle N_1 N_2 \rangle \) is the same before and after the BS:

\[
G_{11} + G_{22} + 2G_{12} = \text{const.} \tag{2.1}
\]

Hence, the drop of the correlation \( G_{12} \) at the output of the BS is accompanied by the increase in the fluctuations \( (G_{11}, G_{22}) \).

What happens if there are many photon pairs at the input? In practice, this case can be realized with squeezed-vacuum twin beams generated through high-gain PDC. It is known that for high-gain PDC, measurement of correlation functions like \( G_{12} \) reveals a high background,
caused by accidental coincidences [16]. At the same time, the reduced level of correlations can be observed in this case by measuring the variance of the photon-number difference [16], \( \text{Var}(N_1 - N_2) \), for the output BS channels. If the delay between the beams is large, the beams ‘do not notice’ each other, pass through the BS without interaction, and the variance of the difference photon number corresponds to the shot-noise level, which is given by the sum of mean photon numbers, \( \langle N_1 + N_2 \rangle \). If the beams interfere, there will be increased fluctuations in each beam, but reduced correlations between them (figure 1); this will obviously lead to the increased noise in the photon-number difference. The variance in this case will be much higher than the shot-noise level; hence, if the variance is measured as a function of the group delay between the input beams, a peak will be observed. The exact calculation is given in the next section.

3. Theory

3.1. Twin beams

Consider first bright squeezed vacuum twin beams at the input of the BS. The twin beams differ by polarization and are obtained via type-II frequency-degenerate collinear PDC. Our theoretical treatment is based on the model of [19], describing the quantum state evolution in the Heisenberg picture within the undepleted pump approximation. The spatial degrees of freedom are ignored assuming that a single spatial mode is selected in both output beams. The Hamiltonian describing collinear type-II optical parametric amplifier (OPA) with the signal and idler frequencies \( \pm \Omega_1 \) tuned around the degenerate frequency \( \omega_p \), with \( \omega_p \) denoting the pump wavelength, is

\[
\hat{H} \propto \Gamma(\Omega)[a_s^T(\Omega)a_i^T(-\Omega)] + \text{h.c.},
\]

where \( \Gamma(\Omega) \) is the spectral parametric gain coefficient and \( a_s^T, a_i^T \) are the photon creation operators in the signal and idler beams. Taking into account the collinear geometry, calculations are done within the plane-wave approximation. The operators at the output of the OPA are given by the Bogolyubov transformations [20]

\[
a_s(\Omega) = U(\Omega, t)a_s^\text{vac}(\Omega) + V(\Omega, t)a_i^\text{vac}(\Omega),
\]

\[
a_i(\Omega) = U(\Omega, t)a_i^\text{vac}(\Omega) + V(\Omega, t)a_s^\text{vac}(\Omega),
\]

where \( a_s^\text{vac}, a_i^\text{vac}, a_s^\text{vac}^T, a_i^\text{vac}^T \) are the photon annihilation and creation operators at the vacuum inputs of the OPA. The gain functions \( U(\Omega, t) \) and \( V(\Omega, t) \) are given by

\[
U(\Omega, t) = \cosh \Gamma(\Omega, t) + i \frac{\Delta(\Omega)l_c}{2\Gamma(\Omega, t)} \sinh \Gamma(\Omega, t),
\]

\[
V(\Omega, t) = \frac{G(t)}{2\Gamma(\Omega, t)} \sinh \Gamma(\Omega, t)
\]

with

\[
\Gamma(\Omega, t) = \sqrt{G(t)^2 - \frac{\Delta(\Omega)^2 l_c^2}{4}},
\]

Despite the sometimes heard opinion that accidental coincidences are of purely technical origin and can always be avoided, they are inevitable for high-gain PDC. Indeed, at high-gain PDC the rate of coincidences becomes practically equal to the rate of accidental coincidences, regardless of the experimental setup features.
where $l_c$ is the length of the nonlinear crystal and $G(t)$ depends on the pump field envelope $A_0(t)$ as $G(t) = \sigma A_0(t) l_c$, with the coefficient $\sigma$ taking into account the properties of the crystal. Because the frequency spectrum of type-II PDC is narrow (in our case, 1.3 nm, see the experimental part), the phase mismatch $\Delta$ is defined only by the temporal walk-off between the signal and idler waves at the degenerate frequency, $\Delta(\Omega) = \left(\frac{\partial k_s}{\partial \omega} \omega = \omega_p/2 - \frac{\partial k_i}{\partial \omega} \omega = \omega_p/2\right)\Omega$. Assuming that the signal and idler beams arrive at the BS with the relative delay $\tau$, the BS transformations take the form

$$a_1(t) = \frac{a_{s}(t - \tau) + a_{i}(t)}{\sqrt{2}}, \quad a_2(t) = \frac{-a_{s}(t - \tau) + a_{i}(t)}{\sqrt{2}}.$$ (3.5)

The numbers of photons per pulse at each output of the BS $N_1, N_2$ are found by integrating the combination $a_{1,2}^\dagger(t)a_{1,2}(t)$ over time. We further calculate the variance of their difference, $\text{Var}(N_-) \equiv \langle (N_1 - N_2)^2 \rangle - \langle N_1 - N_2 \rangle^2$. It is convenient to normalize this variance to the mean sum of the photon numbers, $\langle N_+ \rangle \equiv \langle N_1 + N_2 \rangle$, which yields

$$\frac{\text{Var}(N_-(\tau))}{\langle N_+ \rangle} = 1 + \frac{1}{\int_0^\infty d\Omega |V(\Omega, 0)|^2} \int_0^\infty d\Omega |V(\Omega, \tau)|^2 |\text{Re}([U^*(\Omega, 0)]^2 e^{2i\Omega \tau})|.$$ (3.6)

At delays $\tau$ much larger than the pump pulse duration, both terms in the brackets vanish, and the normalized variance corresponds to the shot-noise level. At $\tau$ less than the pulse duration but exceeding the coherence time of the twin beams (assumed to be much less than the former), only the first term in the brackets survives, yielding the normalized variance scaling as the mean photon number, $\langle N_{1,2} \rangle = \int_0^\infty d\Omega |V(\Omega, 0)|^2$. Finally, at zero delay both terms in the brackets are non-zero and, at high gain, are equal to each other, which yields twice the value of $\langle N_{1,2} \rangle$. From these considerations, it is clear that the HOM effect measured for the variance of the difference photon number results in a double peak: the narrow ‘top’ has a typical width given by the coherence time of the PDC radiation while the broader ‘pedestal’ width depends on the PDC pulse duration. The exact shape of the peak can be only calculated numerically. Figure 2 shows this shape calculated for two different maximal values of the parametric gain, $G_{\text{max}} = 1.5$ and 7.5, with the PDC assumed to be type II from a beta-barium borate (BBO) crystal of length 10 mm pumped by 18 ps-long pulses at 355 nm.
If the detection efficiency $\eta$ is less than unity, the height and the visibility of the peak will be reduced. Indeed, by denoting the numbers of detected photons in channels 1, 2 by $S_{1,2} \equiv \eta N_{1,2}$, we obtain the variance of their difference $S_- \equiv S_1 - S_2$ normalized to the mean value of their sum $S_+ \equiv S_1 + S_2$ in the form
\[
\frac{\text{Var}(S_- (\tau))}{\langle S_+ \rangle} - 1 = \eta \left( \frac{\text{Var}(N_- (\tau))}{\langle N_+ \rangle} - 1 \right).
\]
(3.7)

Therefore, the background of the normalized variance is given by a unity (the shot-noise level) regardless of the quantum efficiency, while the height of the peak above the background scales as $\eta$.

### 3.2. Classical beams

Consider now that at the input of the BS, instead of twin beams there are classically correlated thermal beams. In order to observe the classical analogue of the HOM effect, a random phase $\phi$ should be introduced into one of the beams to ‘erase’ the usual first-order interference \[13\]. (For the case of twin beams, there is no first-order interference due to their uncertain relative phase.) In the experiment we consider here, this phase should vary from pulse to pulse. Then the fields at the input of the BS are $E(t - \tau)$ and $E(t)e^{i\phi}$, where $\tau$ is the delay introduced before the BS. The fields at the output ports of the BS are
\[
E_{1,2}(t) = [E(t - \tau) \pm E(t)e^{i\phi}] / \sqrt{2}.
\]
(3.8)

The output intensities are
\[
I_{1,2}(t) = [|E(t - \tau)|^2 + |E(t)|^2 \pm (E^*(t - \tau)E(t)e^{i\phi} + \text{c.c.})] / 2.
\]
(3.9)

The value measured in experiment is the variance of the difference $I_1(t) - I_2(t)$ integrated over an optical pulse. The integral of the intensity difference is
\[
N_-(\tau) \equiv \int (I_1(t) - I_2(t)) \, dt = \int dt (E^*(t - \tau)E(t)e^{i\phi} + \text{c.c.})
\]
(3.10)

and the variance is found as
\[
\text{Var}(N_- (\tau)) = \left\langle \left[ \int dt (E^*(t - \tau)E(t)e^{i\phi} + \text{c.c.}) \right]^2 \right\rangle
\]
(3.11)

with the averaging over the ensemble of pulses. Due to the averaging, the terms containing the random phase $\phi$ in the exponential vanish. After taking into account that for thermal fields
\[
\langle E_{1}E_{2}E_{3}E_{4} \rangle = \langle E_{1}^{*}E_{3}\rangle\langle E_{2}^{*}E_{4}\rangle + \langle E_{1}^{*}E_{4}\rangle\langle E_{2}^{*}E_{3}\rangle,
\]
we obtain
\[
\text{Var}(N_- (\tau)) = 2 \left[ \int dt \sqrt{I_0(t - \tau)I_0(t)} \right]^2 |g^{(1)}(\tau)|^2
\]
\[+ 2 \int dt \int dt' \sqrt{I_0(t - \tau)I_0(t)I_0(t' - \tau)I_0(t')} |g^{(1)}(t - t')|^2,
\]
(3.12)

where $I_0(t) \equiv |\langle E(t) \rangle|^2$ is the initial pulse intensity profile and $g^{(1)}(\tau)$ is the first-order normalized correlation function of the input beams. In the case where the pulse width is much larger than the $g^{(1)}(\tau)$ width (not Fourier-limited pulse), the expression simplifies to
\[
\text{Var}(N_- (\tau)) \approx 2 \left[ \int dt I_0(t) \right]^2 |g^{(1)}(\tau)|^2 + 2 \int dt I_0(t - \tau)I_0(t).
\]
(3.13)
After normalization of equation (3.13) to $\langle N_+ \rangle = 2 \int dt I_0(t)$, the shape of a narrow peak on top of a broader background is reproduced. It looks almost as in figure 2 but without the shot-noise term, which cannot appear in the classical picture:

$$\frac{\text{Var}(N_-(\tau))}{\langle N_+ \rangle} \approx \int dt I_0(t) |g^{(1)}(\tau)|^2 + \int dt I_0(t-\tau) I_0(t) \int dt I_0(t).$$ (3.14)

At zero delay, $\text{Var}(N_-(0)) = \langle N_+ \rangle^2$.

Thus, the expected effect for correlated thermal beams is the same as for twin beams. More recent research [21] shows that the difference can be observed by measuring the photon-number probability distribution at the BS output provided that the losses are very low and the detectors resolve photon numbers.

If the input beams are coherent, already at the stage of integration in equation (3.10) we get

$$N_-(\tau) = g^{(1)}(\tau) \int dt \sqrt{I_0(t-\tau) I_0(t)} e^{i\phi} + \text{c.c.},$$ (3.15)

which yields, after averaging over the random phase, the normalized variance in the form

$$\frac{\text{Var}(N_-(\tau))}{\langle N_+ \rangle} = \left[\int dt \sqrt{I_0(t-\tau) I_0(t)}\right]^2 |g^{(1)}(\tau)|^2.$$ (3.16)

This dependence represents a peak with the shape mainly determined by the first-order correlation function and the height twice as small as for the thermal light: $\text{Var}(N_-(0)) = \langle N_+ \rangle^2 / 2$. If the pulse duration is much larger than the coherence time, the shape exactly coincides with the first-order correlation function.

It is worth mentioning that if the input thermal or coherent beams are independent, the result will be the same as derived above for the case of delays much exceeding the coherence time. Note that in our consideration, the random relative phase of the beams is constant within the pulse duration; therefore, at small delays the beams are not fully independent.

4. Experiment

In experiment, we have first studied the HOM effect for squeezed-vacuum twin beams obtained via high-gain PDC. The experimental setup (figure 3) contained three main parts. In the state generation part, horizontally polarized third harmonic of a pulsed Nd:YAG laser with the wavelength 355 nm, pulse duration 18 ps and repetition rate 1 kHz was used as a pump. The intensity was changed by rotating a half-wave plate ($\lambda_p/2$) in front of a Glan-laser polarizer (GP). Bright twin-beam polarization squeezed vacuum state was generated in a frequency-degenerate traveling-wave OPA based on two 5 mm width crystals BBO cut for type-II phase matching. The effect of the spatial walk-off was reduced by placing two crystals with the optical axis in the horizontal plane tilted oppositely with respect to the pump. The OPA was tuned to generate twin beams in the collinear direction. A high parametric gain was achieved by reducing the pump beam diameter to 180 $\mu$m with the help of a telescope consisting of a convex lens ($F = 50$ cm) followed by a concave lens ($F = -7.5$ cm) at a distance of 42.5 cm. Due to the group velocity difference, the ordinarily polarized pulse of the signal radiation was always delayed relative to the extraordinarily polarized idler pulse. After the crystals the pump was eliminated by two dichroic mirrors DM$_{1,2}$ with high reflection at 710 nm and high transmission.
Figure 3. Experimental layout for the macroscopic HOM interference: Nd:YAG (3\(\omega\)): third harmonic of the Nd:YAG laser; \(\lambda_p/2\): half-wave plate for 355 nm; GP: Glan prism; M: mirror; BBO: two 5 mm type-II beta barium borate crystals with the optical axes tilted oppositely with respect to the pump; DM\(_{1,2}\): dichroic mirrors with high transmission for the pump and 99.5\% reflection for the down-converted radiation; RG: red glass filter RG630 with anti-reflective (AR) coating at 710 nm; \(L_s\): lens with the focal length 0.75 m; PBS\(_{1,2}\): polarization BS; \(\lambda_s/4\): quarter-wave plate for 710 nm; M\(_{1,2}\): high-reflection mirrors at 710 nm; A: 3 mm aperture placed in the focal plane of the lens \(L_s\); \(\lambda_s/2\): half-wave plate for 710 nm; L: lens; D\(_{1,2}\): detectors. Bottom left: the dependence of the PDC signal on the pump power.

at 355 nm. The rest of the pump was absorbed by the red glass filter RG630 with AR coating in the spectral range of the PDC radiation.

The delay between the signal and idler beams was changed in the delay line with the polarization BS (PBS\(_1\)) at the input by scanning the position of the mirror M\(_1\) with respect to the mirror M\(_2\). The linear polarization of the signal and idler beams was rotated by 90° after passing twice through the quarter wave plates (\(\lambda_s/4\)) oriented at 45° to the PBS\(_1\). Therefore both beams left the interferometer from the same output port and were directed to the registration part of the setup.

In the latter, a narrow angular spectrum (with the width 4 mrad) was restricted by a 3 mm iris aperture (A) placed in the focal plane of a 75 cm collecting lens (\(L_s\)). The HOM interference was observed at the output of the polarization BS (PBS\(_2\)) placed after a halfwave plate (\(\lambda_s/2\)) which rotated the polarization basis of the PBS\(_2\) by 45°. In this case, one can show that the
trajectories of two multi-photon states arriving simultaneously at the polarization BS became indistinguishable and the interference occurred. After the BS the radiation was focused by two lenses (L) on the analogue detectors (D$_1$, D$_2$) based on pin-diodes [22]. The electric pulses from the detectors had a fixed duration (8 μs) and their amplitudes were proportional to the numbers of photons at the input. The pulses were time-integrated with the help of a 10 Msample s$^{-1}$ 14-bit analogue-to-digital converter card. The card was triggered by the synchronization pulses of the laser. The measured signals per pulse $S_{1,2}$ at the output of D$_{1,2}$ were used for the calculation of the normalized second-order intensity correlation function $g^{(2)} = \frac{\langle S_1 S_2 \rangle}{\langle S_1 \rangle \langle S_2 \rangle}$ and the normalized variance of the difference signal. In the experiment, the variance of the difference signal, the second-order intensity correlation function, and the mean values of the signals were calculated for each position of the mirror M$_1$ by averaging the data over 30 000 pulses.

At the beginning, the brightness of the generated state was evaluated from the parametric gain ($G$) measurement. The intensity of the PDC radiation was measured as a function of the pump power (see bottom-left part of figure 3). The measurement was done after narrowband spectral selection (0.2 nm around the wavelength of 710 nm) by means of a spectrometer HORIBA Jobin Yvon Micro HR (not shown) and angular selection (4 mrad around the collinear direction) by means of an aperture. The red circles represent the experimental results and the blue line is the fit plotted according to the dependence $N = \sinh^2 G$. The curve clearly demonstrates that the process was strongly nonlinear. The gain reached the value of $G = 7.5$ at 55 mW pump power, which is equivalent to $N_{\text{mode}} \sim 8 \times 10^5$ photons per mode.

Figure 4(a) shows the interference dip in the second-order intensity correlation function $g^{(2)}$ plotted versus the delay time. According to the definition of the interference visibility $V = \frac{\max - \min}{\max + \min}$, the obtained value of $V = 0.022$ was very small. The value of the $g^{(2)}$ at the edges was used for the estimation of the number of detected modes. In this case, the relative delay was large enough and there was no interference at the BS. Hence, the initial correlation between the signal and idler beams given by the second-order intensity correlation function for a single mode $g^{(2)} = 2 + \frac{1}{N_{\text{mode}}}$ was conserved. Assuming that due to the multi-mode detection the initial normalized second-order intensity correlation function $g^{(2)}$ is transformed into $g_m^{(2)} = 1 + \frac{g^{(2)} - 1}{mN_{\text{mode}}}$ [17], the number of detected modes was found to be $m = 10$.

The normalized variance of the difference signal is displayed as a function of the time delay in figure 4(b). As predicted by the theory, a narrow peak is observed on a broader pedestal.
The width of the narrow peak is determined by the coherence time of the PDC radiation and its amplitude is proportional to $2(N_{\text{mode}} + 1)$. The broader pedestal, with the amplitude proportional to $N_{\text{mode}} + 1$, has the width determined by the pulse duration of the PDC radiation. One should mention here that due to the strong nonlinearity of the parametric process the pulse width of the PDC radiation is considerably shorter than the pulse width of the pump. As it was discussed above, the signal and idler pulses arriving to the BS in sequence were split independently and their excess fluctuations were eliminated as a result of subtraction. This is why the value of the normalized variance of the difference signal at the edges corresponds to the shot noise level.

As a result, the peak manifests a visibility $V = 0.999$ exceeding all previous records of HOM interference, including the one achieved recently with two oppositely chirped classical optical pulses in an experiment on sum frequency generation [23].

The observed shape is well described by the theoretical dependences (3.6), (3.7) for the gain value $G = 7.5$, with the only fitting parameter being $\eta$. The obtained value of $\eta$ allows us to estimate the losses as 97%. This huge value can be partly explained by the losses in the optical channel (30%) but is mainly caused by the spatial filtering of the signal and idler beams. The beams have different angular spectra because of the anisotropy and therefore have to be restricted by the aperture $A$, small enough to select the area where they overlap. This leads to more than 90% of losses.

Because the background, given by the first term on the right-hand side of equation (3.6), is equal to a unity, we conclude that the brighter the state the higher visibility can be obtained. We would like to stress that the visibility is high only at high gain; at low gain the peak height is comparable with the background. The ratio between the widths of the broad pedestal and the central narrow peak gives the effective number of the longitudinal modes, $m_{\text{long}} = 8$. It is important to mention here that for pure states this ratio is an estimate for the degree of entanglement in the time/frequency domain [24].

Figure 5 shows the full width at half maximum (FWHM) of the central peak as a function of the parametric gain. A 20% narrowing of the peak was observed by changing the parametric gain in the range of $5.5 < G < 7.5$. The red circles are the experimental data, in agreement with the theoretical prediction (the blue line). As it was shown in [25], at high gain most of the photons are generated at the output part of the crystal, which leads to the broadening of the spectrum and therefore, to the reduction of the coherence time of the PDC radiation. If the HOM
Figure 6. The peak obtained for coherent pulses of the second harmonic of Nd:YAG laser with the wavelength 532 nm.

effect is used for resolving time intervals, as, for instance, in quantum coherence tomography, then the resolution will increase with the parametric gain.

Note that because of the huge losses, photon-number correlations of the squeezed-vacuum state are reduced to the classical level. Therefore the state at the BS input is not different from a phase-mixed thermal state considered in section 3.2. However, according to the theory, the shapes of the peaks observed for twin beams and for thermal beams should be exactly the same. In both cases, the shape contains information about the pulse duration (width of the pedestal peak) as well as the coherence time (width of the narrow peak). The ratio of the two yields the number of the temporal (longitudinal) modes.

The shape of the peak should be different for the case of coherent beams at the input. According to equation (3.16), it should be given by a product of the first-order correlation function and the convolution of the field pulse shape with itself. In the case of a Gaussian pulse with the FWHM $T$ and the coherence time $\tau_c$, the shape of the peak is also Gaussian:

$$\frac{\text{Var}(N_-(\tau))}{\langle N_+ \rangle} = \frac{1}{2} \langle N_+ \rangle \exp \left[ -2 \ln 2 \frac{\tau^2}{\Delta \tau^2} \right],$$  

(4.1)

where $1/\Delta \tau^2 \equiv 1/\tau_c^2 + 1/T^2$.

This result was tested in experiment by feeding into the setup the radiation of the Nd:YAG laser second harmonic. The obtained shape of the peak is shown in figure 6. As predicted by the theory, a single peak was observed with nearly Gaussian shape. The coherence time value was obtained from the fit as 5.8 ps.

5. Conclusion

To summarize, we have observed the HOM interference of macroscopic states containing on the average $8 \times 10^5$ photons per mode. By measuring the normalized variance of the difference signal we were able to achieve the visibility of 0.999. This is an evident advantage over the correlation function measurement where the visibility is strongly reduced in the presence of high gain and many temporal modes. The experimental results provide the information on the temporal mode structure of the interfering fields. It is shown that the width of the interference
peak reduces with the growth of the parametric gain leading to an increase in the time resolution. This can be very useful for applications involving time measurements, such as quantum-optical coherence tomography.

The effect is shown to be similar for twin beams and classically correlated thermal beams. In both cases, from the shape of the dip one can simultaneously determine the coherence time and the pulse duration. This result provides a simple method for the characterization of not Fourier-limited pulses with thermal statistics. In practice, it can be used for the characterization of beams with excess fluctuations. In particular, it can be used for the measurement of pulse spreading after dispersive optical elements.

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