Two-Color Bright Squeezed Vacuum

Ivan N. Agafonov, Maria V. Chekhova, and Gerd Leuchs

1Department of Physics, M.V.Lomonosoe Moscow State University, Lenninskie Gory, 119992 Moscow, Russia
2Max Planck Institute for the Science of Light, Günther-Scharowsky-Straße 1/Bau 24, 91058 Erlangen, Germany
3University Erlangen-Nürnberg, Staudtstrasse 7/B2, 91058 Erlangen, Germany

In a strongly pumped non-degenerate traveling-wave OPA, we produce two-color bright squeezed vacuum with up to millions of photons per pulse. Our approach to registering this macroscopic quantum state is direct detection of a large number of transverse and longitudinal modes, which is achieved by making the detection time and area much larger than the coherence time and area, respectively. Using this approach, we obtain a record value of twin-beam squeezing for direct detection of bright squeezed vacuum. This makes direct detection of macroscopic squeezed vacuum a practical tool for quantum information applications.

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We are now witnessing a growth of interest in macroscopic quantum systems [1]. Although the subject of macroscopic superpositions is causing scientific debates ever since the formulation of the famous Schrödinger-cat paradox, nowadays the interest in macroscopic quantum systems is also motivated by applications: gravitational wave detection [2], quantum memory [3], super-resolution [4], etc. An interesting perspective for applying macroscopic nonclassical states of light in quantum technologies stems from the fact that such states, in principle, can provide stronger interactions with matter than microscopic (single-photon and few-photon) states.

An evident example of a macroscopic quantum state of light is two-mode bright squeezed vacuum (SV) generated at the output of a traveling-wave optical parametric amplifier (OPA) [5]. At weak pumping such an OPA produces biphoton light, which is usually characterized in terms of Glauber’s normalized correlation functions and studied using single-photon coincidence counting technique. At strong pumping, two-photon correlations become smeared by the presence of higher photon numbers and coincidence counting gets inefficient since normalized Glauber’s correlation functions approach unity. This does not mean that quantum correlations disappear; they just need a different type of measurement to be revealed [6]. This other kind of measurement can be based on the photon-number difference. Since signal and idler beams are ‘twins’ [7], their photon-number difference does not fluctuate, and its variance is ideally zero. When such beams are detected, the difference of the output photocurrents of the two detectors fluctuates only due to non-unity quantum efficiencies of the detectors, electronic noise, or possible optical losses.

Twin-beam squeezing for two-mode SV is being studied using optical homodyne technique for two decades (see, e.g., Refs [8–10]). The method is powerful since it provides phase information. At the same time, mixing with local oscillator turns SV into a different state, namely a squeezed coherent state. Obtaining the ‘unperturbed’ properties of SV requires direct detection. The latter has been applied to bright SV rather recently [11–13], with the appearance of sensitive CCD cameras and the technique of charge-integrating detection. All these pioneering works showed rather low squeezing (less than 3dB), which disappeared at large photon numbers (more than 10 photons per mode). As we will show below, most probable reason for this behavior is not selecting enough transverse modes. There was, however, one very early paper on the direct detection of weak two-mode SV [14] cleverly using multimode collection and obtaining a certain degree of squeezing.

Measurements based on the photocurrent subtraction are principally different from the ones based on coincidence counting (or photocurrent multiplication). In particular, while coincidence measurement requires as few modes selected as possible, measurement of the difference-signal variance requires a large number of detected modes. (Alternatively, a single Schmidt mode can be selected in both signal and idler beams [15, 16].) Recently, 3dB of polarization squeezing was observed via direct detection of SV by collecting a large number of frequency and angular modes [17], although with the values of gain as small as 0.3. In this work, we measure two-mode squeezing for two-color SV generated via collinear nondegenerate type-I parametric down-conversion at gain values of up to four.

The effect of two-mode squeezing consists of the suppression of fluctuations in the photon-number difference for two light modes (beams) below the classical limit, which is given by the mean sum photon number in these beams. Indeed, it is easy to show that for two single-mode beams, labeled 1, 2, with mean photon numbers equal to $N$, the variance of photon-number difference is

$$\text{Var}(\hat{N}_1 - \hat{N}_2) = N^2(g_{11}^{(2)} + g_{22}^{(2)} - 2g_{12}^{(2)}) + 2N, \quad (1)$$

where $g_{11}^{(2)}, g_{22}^{(2)}$ are normalized second-order intensity correlation functions for beams 1 and 2, respectively, and $g_{12}^{(2)}$ is their second-order cross-correlation function. For
For two-mode SV, \( g_1^{(2)} = g_2^{(2)} = 2 \) (both twin beams have thermal statistics), \( g_1^{(2)} = 2 + 1/N \), and, according to Eq. (1), NRF=0. Our setup was a traveling-wave OPA based on two type-I 2-mm BBO crystals with the optic axes oriented in the same plane but symmetrically with respect to the pump direction, in order to eliminate spatial walkoff. As a pump, we used the third harmonic of a Nd:YAG laser (wavelength \( \lambda_p = 355 \text{ nm} \)) with pulse duration 17 ps, repetition rate 1 kHz, and energy per pulse up to 0.2 mJ. The fundamental and second-harmonic radiation of the laser was eliminated using a prism. The crystals were oriented to produce frequency-nondegenerate parametric down-conversion (PDC) at wavelengths 635 nm (signal) and 805 nm (idler). After the crystals, the pump radiation was cut off by two dichroic mirrors with high transmission (\( T > 97\% \)) at the down-converted wavelengths. The residual pump radiation was suppressed by two Glan prisms, the first one selecting the pump polarization (H) before the crystals and the second one transmitting the SV polarization (V). The pump power could be varied by rotating a waveplate before the first Glan prism. Signal and idler beams were then separated using a dichroic beamsplitter and focussed, by lenses with 5 cm focal lengths, on two charge-integrating detectors based on Hamamatsu S3883 p-i-n diodes \cite{17,20}. Quantum efficiencies of the detectors at wavelengths 635 nm and 805 nm were 85% and 95%, respectively. The angular bandwidth selected was determined by the sizes of two iris apertures (about 1 cm) placed at a distance of 36 cm from the crystals. (In order to properly match frequency and angular modes in the signal and idler channels, the apertures had different diameters, see below.) The number of transverse modes \( m_t \) was given by the squared ratio of the signal (idler) aperture size and the transverse coherence length of signal (idler) radiation in the plane of the aperture (about 2 mm), yielding \( m_t \approx 25 \). The number of longitudinal modes was given by the ratio of the pulse duration and PDC coherence time (corresponding to the spectral bandwidth selected by the apertures, about 13 nm), yielding \( m_l \approx 150 \). The overall detection efficiencies at wavelengths 805 nm and 635 nm, determined by 11 AR coated surfaces, 3 dichroic coatings and detectors’ quantum efficiencies, were estimated to be 77% and 70%, respectively, resulting in the best achievable NRF of 0.33. Registration of the detectors’ output signals and their calibration was performed the same way as described in Ref. \cite{17}. However, in comparison with the results of \cite{17} we now had much larger photon numbers, on the order of \( 10^6 \) photons per pulse, and the shot-noise level (about \( 10^3 \) electrons) much exceeded the electronic noise level (180 electrons). Still, in the results shown below the electronic noise was subtracted.

The most important point about direct detection of SV is that the detection time and detection area should be much larger than the coherence time and coherence area, respectively. In other words, the number of detected modes should be large, as it was shown in \cite{18,21}. In particular, the detection aperture size should much exceed the transverse coherence length of down-converted radiation. In order to study the dependence of squeezing on the number of registered modes we made the measurement for different sizes of the apertures. First, signal and idler aperture diameters \( D_s, D_i \) were varied simultaneously, so that their ratio satisfied the relation (set by the transverse phase matching condition for PDC)

\[
D_i / D_s = \lambda_i^{max} / \lambda_s^{min},
\]

where \( \lambda_i^{max}, \lambda_s^{min} \) are maximal idler and minimal signal wavelengths selected, respectively (Fig.2a). Then, the idler aperture had a fixed diameter and the signal aperture diameter was varied (Fig.2b). From Fig.2a we see that, according to the predictions of \cite{21}, squeezing improves with the increase of the aperture diameters, i.e., with the number of angular modes selected. This can be understood by recalling that squeezing is sensitive to losses, so that collection of a limited number of modes leads to the existence of uncorrelated photons of the beam, which is similar to the loss of a certain amount of photons. Ideally, to observe maximal squeezing one should detect all frequency-angular spectrum of PDC. But, as it was shown in Ref. \cite{18}, in practice a high degree of squeezing can be obtained even for a limited number of modes provided that it is high.

Fig.2b shows that the optimal squeezing is obtained when aperture sizes are matched according to condition \cite{23}. If the aperture sizes are not matched, squeezing disappears and even turns into anti-squeezing. This happens because both signal and idler beams are thermal ones, and they consist of strongly fluctuating 'speckles’. There is of course perfect correlation between the
speckles in the two beams. If the aperture sizes are not matched, then in one of the beams there are 'unmatched' speckles whose excess intensity fluctuations are not compensated and lead to an increase in the NRF. A straightforward calculation shows that for two sets of $k$ independent thermal modes with equal photon numbers $N$, noise reduction factor does not depend on $k$ and is only determined by the mean photon number per mode: \(\text{NRF} = \frac{N}{N} + 1\). In other words, difference-photocurrent measurement on a large number of independent thermal modes will provide NRF as high as for a single thermal mode. This is yet another demonstration of the huge difference between NRF measurement and photocount coincidence detection where thermal beams ‘poissonise’ if a large number of modes is collected. The discrepancy between the experimental NRF value of 4 dB and the calculated value of 5.2 dB might be due to imperfect alignment (apertures shape deviating from round etc).

It follows that for bright PDC, mismatch of the aperture diameters will be more noticeable at large numbers of photons per mode, i.e., at large gain. In fact, whether two-mode squeezing can be still observed for SV at high gain is an important question. Previous works \([11, 22]\) showed an increase of NRF at high values of the gain. At the same time, in theory, NRF for SV should only be given by losses and not depend on the gain. The growth of NRF with the increase of parametric gain was explained in Refs. \([11, 22]\) by the dependence of the mode size on the gain, and, as a consequence, the reduction of the selected number of modes with the gain growth. However, this should have little effect if a large number of modes are collected; also, this does not explain the appearance of anti-squeezing.

To study the dependence of squeezing on the gain (Fig.3a), we measured NRF versus the pump power. The aperture sizes were matched in accordance with Eq.3. The gain was estimated by fitting the dependence of the PDC output photon number $N$ on the pump power (shown in the inset to Fig.3a) with the formula

\[
N = m \sinh^2 \gamma + N_0,
\]

where $\gamma$ is the parametric gain, scaling as square root of the pump power, $m = m_m m_l$, and $N_0$ is the noise (fluorescence) linear in the pump, measured separately. In all dependencies, NRF indeed grew with the increase in the gain. However, this growth was very sensitive to alignment. For the best alignment we could achieve, a growth of NRF is only noticeable at the end of the dependence, where $\gamma \approx 2$ (Fig.3a).

To pass to higher-gain PDC, we focused the pump into the crystals by means of a 5:1 telescope. This way we achieved gain values of up to 4.5, where the increase of NRF was noticeable (Fig.3b). Still, even at $\gamma = 4$, corresponding to about 700 photons per mode, squeezing was still observed. At high gain, the sensitivity of squeezing to misalignment was even more evident. For instance, Fig.4a shows NRF as a function of the mean photon number per mode $\gamma = m_m m_l^2 \Gamma$ for the best alignment we could achieve (squares) and for the idler aperture (whose diameter was 10.2 mm) displaced by 500 $\mu$m (circles). Clearly, this 5% displacement of the aperture leads to a dramatic increase in NRF. Fig.4b shows NRF as a function of the aperture shift measured at the gain value 3.7.
High sensitivity of NRF to misalignment is of the same origin as its sensitivity to the aperture size mismatch. Indeed, if the mode sets collected by the signal and idler detectors are not matched, there are uncompensated (independent) intensity fluctuations in the signal and idler channels. They should lead to NRF growth, scaling linearly with the mean photon number. Our calculation shows that, provided that in each channel there are \( m \) matched PDC modes and \( k \) unmatched ones, all having the same mean photon number \( \bar{N} \), NRF is increased by

\[
\Delta = \frac{k(\bar{N} + 1)}{m + k}.
\]  

For small displacements of a circular aperture, this should lead to a linear dependence of NRF on \( \bar{N} \), which we used to fit the experimental data in Fig. 4a.

It follows that observation of twin-beam squeezing in high-gain PDC is the more sensitive to mode matching the higher the gain. This explains a linear growth of NRF with the increase of \( \bar{N} \) observed in [11]. Similar effect is used in quantum metrology to measure spatial displacements with high precision [24], and determines the resolution of quantum imaging [25].

In conclusion, we have generated and tested via direct detection two-color squeezed vacuum containing up to a thousand photons per mode, and up to a million of photons per pulse. By collecting a large number of angular and frequency modes in both channels, considerable amount of two-mode squeezing (4dB) has been obtained, the best result ever achieved for the direct detection of squeezed vacuum. We have also studied the behavior of squeezing on the angular bandwidths selected in signal and idler channels and showed that the effect is extremely sensitive to the sizes and alignment of the angle-selecting apertures. On the other hand, if a large number of conjugated angular modes is properly selected, squeezing is not much sensitive to the growth of the parametric gain. In particular, we observed nearly constant amount of squeezing up to the gain \( \Gamma = 2 \), corresponding to 13 photons per mode. Squeezing, although to a smaller degree, was observed up to 900 photons per mode.

Our result is an important step in the investigation of mesoscopic and macroscopic quantum systems. We show that difference-photocurrent measurements can replace coincidence-based measurements in the study of macroscopic squeezed vacuum. Pairwise correlations are still evident provided that losses are sufficiently small and proper mode matching is ensured.

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