Quantum Light from a Whispering-Gallery-Mode Disk Resonator

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Optical parametric down-conversion has proven to be a valuable source of nonclassical light. The process is inherently able to produce twin-beam correlations along with individual intensity squeezing of either parametric beam, when pumped far above threshold. Here, we present for the first time the direct observation of intensity squeezing of −1.2 dB of each of the individual parametric beams in parametric down-conversion by use of a high quality whispering-gallery-mode disk resonator. In addition, we observed twin-beam quantum correlations of −2.7 dB with this cavity. Such resonators feature strong optical confinement and offer tunable coupling to an external optical field. This work exemplifies the potential of crystalline whispering-gallery-mode resonators for the generation of quantum light. The simplicity of this device makes the application of quantum light in various fields highly feasible.

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When going back and forth on a swing, the oscillatory movement of the legs is coupled to the oscillation of the swing—a well-known parametric process. Similarly, optical parametric down-conversion (PDC) links an optical field to its subharmonic mediated by a dielectric medium. In this optical process, one pump photon (p) is converted to two subharmonic photons, called the signal (s) and the idler (i). The signal and idler are therefore strongly correlated in photon number. Experimentally, these two-mode correlations were proven to be quantum correlations [1,2]. Further improvement made this optical parametric process a state-of-the-art squeezing source [3–5], along with squeezing in optical fibers [6,7].

About 20 years ago, a theoretical analysis additionally predicted the generation of an entangled signal and idler, which are simultaneously individually squeezed in intensity, when pumping the process high above the optical parametric oscillation (OPO) threshold [8,9]. However, this could not be shown directly so far, as the above-threshold squeezing was obscured by classical noise of the pump laser source and by relaxation oscillations occurring in triply resonant cavities [10,11]. In our approach, we were able to overcome these obstacles by utilizing a compact and robust whispering-gallery-mode (WGM) resonator, which is a valuable tool in many areas in photonics [12–14]. In these resonators, light is guided by continuous total internal reflection. This lifts the need for high reflectivity coatings, thereby limiting the quality of WGM resonators by internal material loss only [15]. The evanescent field overlap with an external optical element allows for continuously variable coupling to the WGM resonator. So far, the extreme properties of WGM resonators could not be leveraged to generate quantum light from nonlinear effects of the resonator material, as, e.g., thermorefractive noise plays a detrimental role [16,17]. By using the high second-order nonlinearity in a crystalline WGM resonator, we succeeded in showing quantum two-mode correlations in the down-converted beams. In addition, the extremely low threshold enabled us to directly show the squeezing of a single parametric beam.

Cavity properties—in particular, of high quality cavities, such as our WGM resonator—significantly influence the quantum and classical properties of the OPOs [18]. In the following, we discuss theoretical implications of a triply resonant degenerate OPO pumped coherently above threshold (see [9]). The resonator is characterized by the total pump and parametric cavity linewidth γ_p and γ, respectively. The total cavity linewidth γ is equal to the sum of the internal loss rate α and the coupling rate γ_0 (we refer to energy loss rates).

The OPO threshold pump power P_{th}, for instance, depends cubically on the total cavity linewidth of the three fields involved, as $P_{\text{th}} \propto \gamma_p \gamma^2$. We expect an extremely low threshold, as WGM resonators have very high quality factors. Hence, only a low optical pump power P is required, where shot noise limited lasers are easily available. Moreover, we are able to continuously tune the threshold by varying the coupling rate.

The twin-beam intensity correlations between the signal and idler are affected by the presence of a cavity as well. These correlations are proportional to the variance of the difference $\mathcal{V}_-$ of the field amplitude quadratures $\hat{X}_s$ and $\hat{X}_i$ of both beams. The variance of the quadrature sum $\mathcal{V}_+$ yields the total noise of the parametric beams. The theory for these variances normalized to the shot noise limit (SNL) predicts [9]

$$
\begin{aligned}
\mathcal{V}_{-} &= 1 - \frac{\gamma_0}{\gamma} \frac{\eta}{1 + \Omega^2}, \\
\mathcal{V}_{+} &= 1 + \frac{\gamma_0}{\gamma} \frac{\eta}{(\sigma - 1)^2 + \Omega^2}.
\end{aligned} \tag{1}
$$

Here, $\Omega = \nu_{\mathrm{det}}/\gamma$ is the measurement sideband frequency ν_{det} normalized to the parametric cavity bandwidth. The pump parameter is $\sigma = \sqrt{P/P_{\mathrm{th}}}$, and the detection efficiency η includes all optical losses in the signal (or idler) channel. The difference signal $\mathcal{V}_{-}$ is independent of σ and shows twin-beam quantum correlations, as $\mathcal{V}_{-} < 1$. The sum signal $\mathcal{V}_{+}$, on the contrary, depends on the pump parameter. It shows strong excess noise ($\mathcal{V}_{+} \gg 1$) near the threshold and drops to the SNL for high pump powers. Both variances depend on the ratio of parametric coupling rates to total losses. With a high quality WGM OPO, we are able to vary this ratio between coupling and internal losses for optimizing the correlations while maintaining a reasonably low threshold. Thus, shot noise limited or even nonclassical pump fields are accessible.

The individual intensity fluctuations of the signal and equivalently the idler field are proportional to the normalized amplitude quadrature variance [9]

$$
\mathcal{V}(\hat{X}_{s,i}) = 1 - \frac{\eta}{2} \frac{\gamma_0}{\gamma} \frac{\sigma(\sigma - 2)}{(1 + \Omega^2)[\Omega^2 + (\sigma - 1)^2]}. \tag{2}
$$

Contrary to the twin-beam correlations $\mathcal{V}_{-}$, the single-beam intensity variance $\mathcal{V}(\hat{X}_{s,i})$ strongly depends on the pump parameter and drops from excess noise near the threshold to an ideal squeezing limit of 0.5 (-3 dB) far above the threshold. The SNL is reached at 4 times the threshold. Analogous to the twin-beam correlations, the individual squeezing depends on the ratio of coupling rates to the total loss. To reach a high pump parameter, however, a low threshold is needed, implying a high quality resonator. Additionally, one has to assure a coherent pump field for the observation of squeezing, easily accessible at low pump powers. In this respect, the quality of a WGM resonator gives an additional significant advantage.

Triply resonant OPOs are known to undergo relaxation oscillations [11,19]. These induce classical fluctuations in the ac power spectrum of each parametric field at a specific frequency and can disturb the measurement of squeezing. The relaxation oscillation frequency ν_R normalized to the parametric linewidth γ is given by [19]

$$
\nu_N = \frac{\nu_R}{\gamma} = \sqrt{\frac{\gamma_p}{2\gamma}} \sqrt{\sigma - \left(1 + \frac{\gamma_p}{4\gamma}\right)}. \tag{3}
$$

The relaxation oscillations have a threshold at $\sigma = 1 + \gamma_p/(4\gamma)$ and depend on the pump parameter σ. One possible approach to circumventing the relaxation oscillations is to make its threshold very high by having a large ratio of γ_p/γ, i.e., working with a cavity resonant only for the parametric fields. However, we cannot tune both loss rates independently with the present coupling technique. Alternatively, one can shift the relaxation oscillation frequency out of the cavity bandwidth ($\nu_N \gg 1$), by increasing the pump power. For the cavity used in this experiment, we estimated the normalized relaxation oscillations versus σ. The relaxation oscillations lie within our cavity bandwidth only for pump parameters σ between approximately 2.5 and 2.8. Thus, we are able to easily shift these oscillations out of the cavity bandwidth.

For our experiment we used a WGM disk resonator, which possesses axial symmetry. To assure phase matching for PDC all along the circumference of the resonator, the crystal should be uniaxial, with the optical axis along the symmetry axis of the cavity. Hence, our WGM disk is made from a 5% MgO-doped z-cut lithium niobate wafer. Natural type I phase matching in the crystal can be achieved by tuning the refractive indices of the disk via temperature and the electro-optical effect, as discussed in Ref. [20]. Our experimental setup is shown in Fig. 1. We drive our triply resonant WGM OPO above threshold with a continuous wave frequency doubled Nd:YAG laser at 532 nm. As our WGM OPO shows very good stability, no active locking is required. A movable diamond prism (antireflection coated for 532 and 1064 nm) is placed in the vicinity of the WGM resonator, overlapping with its evanescent field. This allows for continuously variable coupling of the pump field to the resonator. The outcoupling is provided by the same prism. The outcoupled pump and strongly nondegenerate signal and idler fields are separated with a dispersion prism and directed to the detection setup, carefully assuring minimal losses for the PDC light. Continuously scanning the pump frequency, we observed the WGM spectrum and selected the pump mode most efficient for PDC. We applied a phase matching temperature near 94 °C. The minimal incoupled pump threshold power for PDC was measured to be 6.7 μW, which is

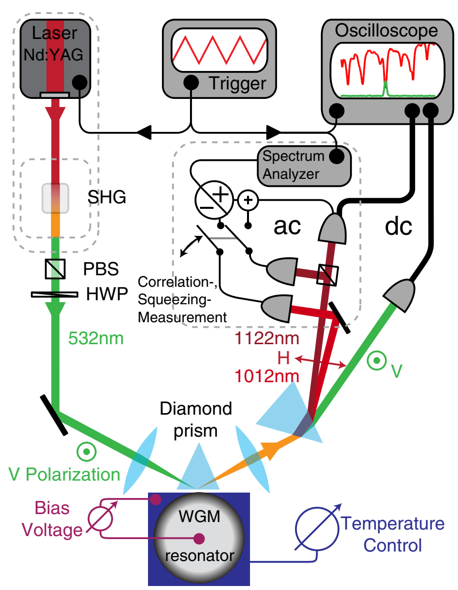

FIG. 1 (color online). PDC setup with WGM resonator (disk radius 1.9 mm, rim radius 0.25 mm, height 0.5 mm).

expectedly extremely low compared to state of the art experiments with 300 μW [21] and gives us the necessary flexibility for the quantum measurements. The classical properties of our OPO are discussed in Ref. [22].

For coupling the pump field critically to the resonator ($\gamma_{p,0}/\gamma_p = 0.5$), the total cavity bandwidth of the pump field is observed to be $\gamma_p = 30$ MHz, and its coupling rate is $\gamma_{p,0} = 15$ MHz. Moreover, we found $\gamma_0/\gamma = 0.22$ for the parametric light, by using the OPO output power dependence on the pump parameter (see [22]). This indicates that the signal and idler couple more weakly to the resonator than to the pump field. This is surprising, as the evanescent coupling is stronger for a fixed spacing, the longer the wavelength. However, we assume that the parametric WGMs excited in our OPO are nonequatorial and deeper-lying modes than the pump mode. By adjusting the phase matching conditions to excite parametric equatorial modes, we would expect stronger outcoupling and hence more squeezing.

For the quantum measurements, we use two photodetectors with an optical noise equivalent power of less than 1.5 μW at a measurement frequency of 3.2 MHz (resulting in a normalized cavity frequency of $\Omega = 0.6$). The dc part of the photocurrent is used as a monitor and for shot noise calibration. We simultaneously monitor the sum and the difference ac fluctuations of these two detectors with two spectrum analyzers configured for zero span around 3.2 MHz, a resolution bandwidth of 300 kHz, and a video bandwidth of 10 kHz. These settings result in a time resolution of around 0.1 ms, while the pump laser sweep time is 50 ms. We postprocessed the spectrum analyzer data with a running average of 50. The detection setup is linear in ac and dc over more than 17 dB. The common mode rejection ratio between sum and difference channels is at least 20 dB (i.e., the detectors are balanced to within 1%). All squeezing measurements performed were verified by the linear behavior of attenuation measurements.

The PDC process generates photons by pairs, leading to quantum correlations in the intensity of the signal and idler [1]. These two-mode correlations can be observed by focusing the signal and idler on two detectors and analyzing the fluctuations of the balanced sum and difference photocurrents (see Fig. 1 with switches right). When the difference fluctuations are below the SNL, the twin beams are quantum correlated. The total quantum efficiency of the measurement is $\eta = (87 \pm 4)\%$. For this twin-beam measurement we set the threshold of our resonator slightly below the fixed pump power by adjusting the coupling. In Fig. 2(a), we present a sweep of the pump frequency through the pump WGM resonance. We see three PDC channels, each corresponding to a pair of signal and idler WGMs incidentally phase matched to the pump WGM at a particular frequency detuning [20]: one at the center of the pump resonance and two on the wings. The difference signal in Fig. 2(a) is reliably below the SNL for the central channel, displaying -2.7 ± 0.4 dB of two-mode squeezing (corrected for electronic noise). The sum signal greatly exceeds the SNL in this near-threshold measurement. In a separate measurement we observed that it approaches the SNL for high pump powers, as predicted in Eq. (1).

We measured quantum correlations for different pump powers in two coupling regimes [see Fig. 2(b)]. The theoretical estimate of -0.6 dB for a critically coupled pump, where our signal and idler modes are undercoupled ($\gamma_0/\gamma = 0.22$), is close to the measured value for weak coupling. Stronger coupling leads to an increased output rate relative to the loss rate. This increases the threshold but also the two-mode squeezing. This result is consistent with our theoretical estimate yielding -2.0 dB of quantum correlations for signal and idler WGMs critically coupled ($\gamma_0/\gamma = 0.5$). As stated in Eq. (1), the correlations are independent of the pump power in both regimes within the measurement uncertainty.

On the contrary, the intensity fluctuations of a *single* parametric beam are expected to depend strongly on the pump power, as discussed. Investigations of squeezing high above threshold now are feasible. We measured the signal (or idler) intensity fluctuations along with the SNL with a balanced self-homodyning setup, by using the sum and difference photocurrents (see the circuit in Fig. 1).

In the single-beam measurement the overall detection efficiency is $(73 \pm 4)\%$. We varied the coupling to

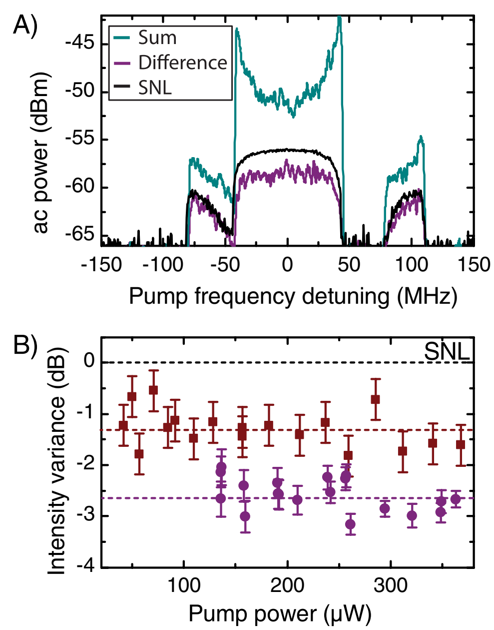

FIG. 2 (color online). Twin-beam measurement. (a) Variances of the signal and idler photocurrent sum and difference compared to the SNL vs laser frequency detuning from the pump WGM center. (b) Twin-beam squeezing vs pump power for stronger coupling (circle) and weaker coupling (square). Error bars are determined from the raw data, as seen in (a). Horizontal lines indicate the SNL and the mean squeezing values, for weak coupling -1.3 dB and for strong coupling -2.7 dB.

facilitate a very low threshold. In Fig. 3(a), the pump frequency is swept through the resonance while the intensity noise and the SNL are recorded. Again, multiple PDC channels are excited. The central channel shows -1.2 ± 0.4 dB of squeezing (corrected for electronic noise) when pumped with approximately 400 μW. However, we chose a different conversion channel and coupling to be able to measure the squeezing in the full range near and far above the threshold [see Fig. 3(b)], as we were limited by the pump power and signal to noise ratio of the detectors. Figure 3(b) shows strong excess noise just above the threshold. As the pump power increases, the noise decreases and finally falls below the SNL, at which point the field becomes squeezed. A theoretical fit for the normalized intensity variance from Eq. (2) yields a threshold power of 12.3 μW and converges to 0.9 shot noise units (SNU) ($-$ 0.5 dB) for high pump powers. This is in good agreement with the theory for critically coupled parametric WGMs ($\gamma = 2\gamma_0$), where the squeezing limit is $V_{\mathrm{crit},s}(\hat{X}_1) = 0.84$ SNU ($-$ 0.7 dB). The squeezing was not impeded by relaxation oscillations, as their frequency depends very strongly on the pump power, and lies within the cavity bandwidth only for a very narrow range of pump power [see Fig. 3(b)]. This has enabled the first direct observation of sub-Poissonian photon statistics in a single OPO beam above threshold.

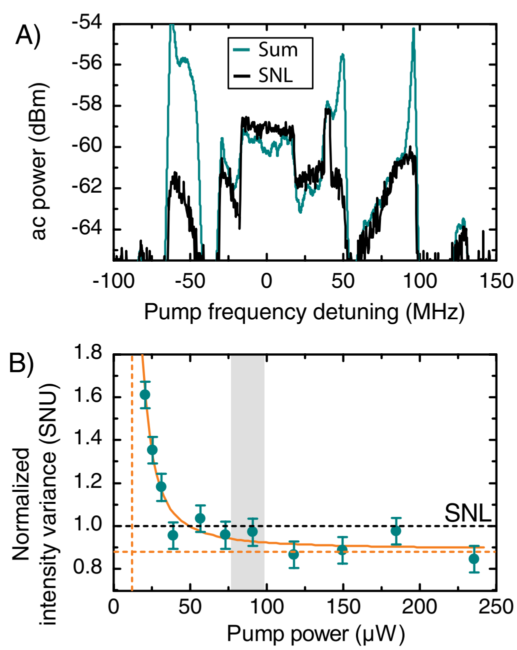

FIG. 3 (color online). Single-beam intensity noise measurement. (a) Intensity noise compared to SNL vs laser frequency detuning from the pump WGM center. (b) Intensity noise in SNU vs pump power (measurement and theoretical fit). Error bars are estimated as the mean variance of two measurements at every point. The horizontal lines indicate the SNL (black) and the asymptotic limit of 0.9 for the normalized intensity variance (orange). The vertical line visualizes the OPO threshold of 12.3 μW. The gray area marks the range of the pump power, where relaxation oscillations may be excited.

This work illustrates the potential of our WGM resonator in the generation of quantum states of light and nonlinear optics in general. The signal and idler beams in our experiment are not only quantum correlated but also individually intensity squeezed. In addition, in the far-below-threshold regime, our WGM resonator is expected to be a highly efficient narrow bandwidth single photon source, compatible with atomic linewidths. High flexibility in wavelength selection for coupling the single photons to atoms could be achieved by using quasi-phase-matching in periodically poled resonators [23]. As crystalline WGM resonators offer high mechanical Q factors as well [24], one could combine optomechanics with nonclassical states of light. Furthermore, studies of dynamical nonclassical effects in second-order nonlinear WGM resonators at high pump powers are intriguing.

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[1] A. Heidmann *et al.*, Phys. Rev. Lett. **59**, 2555 (1987).
[2] L.-A. Wu *et al.*, Phys. Rev. Lett. **57**, 2520 (1986).
[3] J. Laurat *et al.*, Opt. Lett. **30**, 1177 (2005).
[4] M. Mehmet *et al.*, Phys. Rev. A **81**, 013814 (2010).
[5] A. S. Coelho *et al.*, Science **326**, 823 (2009).
[6] R. M. Shelby *et al.*, Phys. Rev. Lett. **57**, 691 (1986).
[7] R. Dong *et al.*, Opt. Lett. **33**, 116 (2008).
[8] M. D. Reid and P. D. Drummond, Phys. Rev. Lett. **60**, 2731 (1988).
[9] C. Fabre *et al.*, J. Phys. (Paris) **50**, 1209 (1989).
[10] Y. Zhang *et al.*, J. Opt. Soc. Am. B **21**, 1044 (2004).
[11] A. Porzio *et al.*, Opt. Commun. **194**, 373 (2001).
[12] A. B. Matsko and V. S. Ilchenko, IEEE J. Quantum Electron. **12**, 1 (2006).
[13] K. J. Vahala, Nature (London) **424**, 839 (2003).
[14] V. S. Ilchenko and A. B. Matsko, IEEE J. Quantum Electron. **12**, 15 (2006).
[15] A. A. Savchenkov *et al.*, Phys. Rev. A **70**, 051804 (2004).
[16] M. L. Gorodetsky and I. S. Grudinin, J. Opt. Soc. Am. B **21**, 697 (2004).
[17] A. B. Matsko *et al.*, J. Opt. Soc. Am. B **24**, 1324 (2007).
[18] A. Aiello, D. Fargion, and E. Cianci, Phys. Rev. A **58**, 2446 (1998).
[19] D.-H. Lee *et al.*, Appl. Phys. B **66**, 747 (1998) (corrected resonance frequency, as it is a driven harmonic oscillator).
[20] J. U. Fürst *et al.*, Phys. Rev. Lett. **104**, 153901 (2010).
[21] K. S. Zhang *et al.*, Phys. Rev. A **64**, 033815 (2001).
[22] J. U. Fürst *et al.*, Phys. Rev. Lett. **105**, 263904 (2010).
[23] See V. S. Ilchenko *et al.*, Phys. Rev. Lett. **92**, 043903 (2004); K. Sasagawa and M. Tsuchiya, Appl. Phys. Express **2**, 122401 (2009); T. Beckmann *et al.*, arXiv:1012.0801v1.
[24] J. Hofer, A. Schliesser, and T. J. Kippenberg, Phys. Rev. A **82**, 031804(R) (2010).