Cavity Enhanced Spontaneous Parametric Down-Conversion for the Prolongation of Correlation Time between Conjugate Photons

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By resonating both fields of the spontaneous parametric down-conversion with an optical cavity, the bandwidth of the fields is reduced to a level that can be resolved by photodetectors. This enables us to directly measure the distribution of the time intervals between the two down-converted photons. The correlation time is found inversely proportional to the down-conversion bandwidth. In the meantime, the signal levels of the fields are greatly enhanced by the resonance. Such a two-photon source will find wide application in quantum information processing.

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Ever since the pioneering work of Burnham and Weinberg [1] and more recent development by Mandel and co-workers [2,3], the entangled two-photon state from spontaneous parametric down-conversion has been widely used in the demonstration of fundamental phenomena such as quantum interference [4–7] and nonlocality [8,9]. It is also proven to be a good source for quantum communication [10,11]. All these are possible thanks to the extremely short correlation time between the two photons produced in the wide band spontaneous parametric down-conversion [1,2].

However, it is precisely this property of such two-photon states that makes it difficult to apply them in the quantum interference of independent sources, which gives rise to more interesting phenomena such as single photon and multiphoton nonlocality [12,13] and quantum state teleportation [14]. The reason is that this type of interference requires the fields to be coherent within the detection time [15], which is usually set by the response time of the detectors. While current detector technology has produced fast optical detector with a response time of order of 10 psec [2], it is still too long compared to the sub-psec correlation time between the two photons generated from single-pass spontaneous down-conversion. Because of this, no direct measurement is available for this short correlation time although an attempt was made first by Burnham and Weinberg [1] and later by Friberg *et al.* [2]. So far, only indirect measurement has been made via interference [16] and harmonic generation [17].

It should be pointed out that the slow detector problem can be tackled with an ultrashort pump field for the parametric down-conversion [18]. Indeed, such a strategy was used in recent experiments of quantum state teleportation [19] and swapping [20] and generation of three-photon entangled state [21]. However, complicated medium dispersion and phase matching involved with ultrashort pulses seriously deteriorate the temporal mode match, leading to a reduced quantum interference effect (small visibility) [19–21]. Passive filtering is needed to clean up the temporal mode.

On the other hand, because the correlation time is directly proportional to the inverse of the bandwidth of the detected fields from down-conversion [3], we can increase the correlation time by optical filtering, as confirmed indirectly by the interference technique [16]. Unfortunately, the spontaneous nature of the process makes the conversion probability directly proportional to the possible number of modes (channels) or the bandwidth of the down-conversion. Therefore, simply filtering the spectrum of down-conversion for longer correlation time will inevitably reduce the conversion rate or the signal level. To make the correlation time comparable to the response time of commonly available avalanche photodiodes ($\sim$0.3 nsec), we have to cut the bandwidth and the signal level by 4 to 5 orders of magnitude.

The problem of signal reduction due to filtering can be solved with active filtering by placing the down-conversion source inside an optical cavity. The resonance property limits the bandwidth to that of the cavity and in the meantime effectively makes the interaction length longer by bouncing the light back and forth inside the cavity and hence increases the signal level. The decrease of signal level due to bandwidth reduction can therefore be compensated by the increase in interaction length. In this Letter, we report a successful implementation of the active filtering scheme and a direct measurement of the time interval distribution between the two photons from parametric down-conversion. The active filtering scheme consists of a type-I parametric down-converter inside a high finesse optical cavity on resonance at the degenerate frequency. We find that the enhancement factor in the signal level due to resonance is given approximately by the square of the number of bounces of light in the cavity and that the correlation time is inversely proportional to the detected bandwidth of down-conversion fields.

The idea of placing the parametric down-conversion source inside an optical cavity is not new. Yurke [22] first realized the role of cavity in the generation of squeezed state and, more recently, Aiello *et al.* [23] discussed further the situation in the context of cavity QED. Such a device is usually referred to as optical parametric

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oscillator (OPO). It is found that the maximum amount of squeezing is produced when the system is operated close to but under the threshold of oscillation [22,24]. So far, some of the largest squeezing was observed in this system [25,26]. However, the application here is quite different from the case of squeezed states in that we operate the device far below threshold. In this way, mainly spontaneous emission occurs for two-photon generation. The chance of stimulated emission for the production of four or more photons is negligible. The generation of a squeezed state, on the other hand, relies on stimulated emission, thus requiring close to threshold operation. Even so, the theory of OPO below threshold for a squeezed state applies equally well to the case far below threshold for a two-photon state. There are a number of versions [22,24,27] of the theory which all make the same prediction for the generation of squeezed states. We adopt here the one formulated by Collett and Gardiner [24]. The output operator of a degenerate OPO on resonance is related to the inputs as follows [Eq. (46) of Ref. [24]]:

$$\hat{a}_{\rm out}(\omega_0 + \omega) = G_1(\omega)\hat{a}_{\rm in}(\omega_0 + \omega) + g_1(\omega)\hat{a}_{\rm in}^\dagger(\omega_0 - \omega) + G_2(\omega)\hat{b}_{\rm in}(\omega_0 + \omega) + g_2(\omega)\hat{b}_{\rm in}^\dagger(\omega_0 - \omega), \quad (1)$$

with

$$G_1(\omega) = \frac{\gamma_1 - \gamma_2 + 2i\omega}{\gamma_1 + \gamma_2 - 2i\omega}, \qquad g_1(\omega) = \frac{4\epsilon\gamma_1}{(\gamma_1 + \gamma_2 - 2i\omega)^2},$$
$$G_2(\omega) = \frac{2\sqrt{\gamma_1\gamma_2}}{\gamma_1 + \gamma_2 - 2i\omega}, \qquad g_2(\omega) = \frac{4\epsilon\sqrt{\gamma_1\gamma_2}}{(\gamma_1 + \gamma_2 - 2i\omega)^2}.$$

Here ϵ is the single-pass parametric gain amplitude and is proportional to the pump amplitude and the nonlinear coefficient. We dropped the $|\epsilon|^2$ term in the denominator of Eq. (1) in the transition from Eq. (46) of Ref. [24] because the OPO is operated far below threshold. ω_0 is the degenerate frequency of the OPO. $\hat{b}_{\rm in}$ represents the unwanted vacuum mode coupled in due to loss in the system. γ_1, γ_2 are the coupling constants (also known as decay constants) for $\hat{a}_{\rm in}$ and $\hat{b}_{\rm in}$, respectively.

First, let us look at the enhancement effect in down-conversion due to resonance. For this, we calculate from Eq. (1) the spectrum of the field $S(\omega)$ defined by

$$\langle \hat{a}_{\rm out}^\dagger(\omega_0 + \omega)a_{\rm out}(\omega_0 + \omega')\rangle \equiv S(\omega)\delta(\omega + \omega'). \quad (2)$$

The result is

$$S(\omega) = |g_1(\omega)|^2 + |g_2(\omega)|^2 = \frac{16|\epsilon|^2\gamma_1(\gamma_1 + \gamma_2)}{[(\gamma_1 + \gamma_2)^2 + 4\omega^2]^2}. \quad (3)$$

Therefore the full width at half height (FWHH) of the spectrum from the down-converted field is simply $\Delta\omega = 0.64(\gamma_1 + \gamma_2)$. The overall signal level is then

$$R_{\rm resonance} = \frac{1}{2\pi}\int_{-\infty}^{\infty} d\omega\, S(\omega) = |r|^2\mathcal{F}^2/\pi\,\Delta t\,\mathcal{F}_0 \approx |r|^2\mathcal{F}^3\,\Delta\omega/4\pi\mathcal{F}_0, \quad (4)$$

where $r \equiv \epsilon\,\Delta t$ is the single-pass gain parameter with Δt as the round-trip time, and $\mathcal{F} \equiv 2\pi/(\gamma_1 + \gamma_2)\,\Delta t \approx 4/\Delta\omega\,\Delta t$ is the finesse of the cavity (which is of the order of the number of bounces of light before it leaves the cavity) and can be measured directly, and $\mathcal{F}_0 \equiv 2\pi/\gamma_1\Delta t$ is the same quantity without the loss ($\gamma_2 = 0$). To find the enhancement factor, we need the signal rate without the cavity. In the single-pass case, we simply have $g_1(\omega) = r\eta(\omega)$ and $g_2 = 0$. Here $\eta(\omega)$ is the gain spectrum of single-pass spontaneous down-conversion determined by phase-matching condition with normalization $\eta(0) = 1$. In the experiment, we usually have an interference filter (IF) in front of the detector. The bandwidth $\Delta\omega_{\rm IF}$ of IF is normally smaller than that of down-conversion so that $\eta(\omega) = 1$ for ω within $\Delta\Omega_{\rm IF}$. Hence, the signal rate without the cavity is

$$R_{\text{single-pass}} = |r|^2\Delta\omega_{\rm IF}/2\pi, \quad (5)$$

and the average enhancement factor per mode is

$$B \equiv \frac{R_{\rm resonance}/\Delta\omega}{R_{\text{single-pass}}/\Delta\omega_{\rm IF}} = \mathcal{F}^3/2\mathcal{F}_0, \quad (6)$$

or roughly the square of the number of bounces of light before it leaves the cavity consistent with the two-photon nature of parametric down-conversion. The loss of the system will reduce the effect by a factor of $\mathcal{F}/\mathcal{F}_0$.

To find the correlation time between the two down-converted photons, we calculate the intensity correlation function defined as

$$\Gamma^{(2,2)}(\tau) = \langle \hat{E}^{(-)}(t)\hat{E}^{(-)}(t + \tau)\hat{E}^{(+)}(t + \tau)\hat{E}^{(+)}(t)\rangle, \quad (7)$$

with

$$\hat{E}^{(+)}(t) = [\hat{E}^{(-)}(t)]^\dagger = \frac{1}{\sqrt{2\pi}}\int d\omega\, \hat{a}(\omega)e^{-i\omega t}. \quad (8)$$

From Eqs. (1), (7), and (8) with some calculation, we find that

$$\Gamma^{(2,2)}(\tau) = |\epsilon|^2(\mathcal{F}/\mathcal{F}_0)^2 e^{-|\tau|(\gamma_1+\gamma_2)}. \quad (9)$$

If we define the correlation time T_c between the two photons as the FWHH of the distribution in Eq. (10), then T_c is inversely proportional to the bandwidth of the down-converted field:

$$T_c = 1.39/(\gamma_1 + \gamma_2) = 0.89/\Delta\omega. \quad (10)$$

So, exactly as for the single-pass case [3], the correlation time in the active filtering scheme is inversely proportional to the bandwidth of down-conversion.

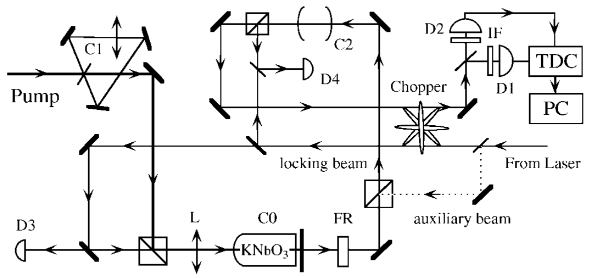

FIG. 1. Experimental layout.

The outline of the experiment arrangement is shown in Fig. 1. The pump field for the OPO is produced by frequency doubling of a single mode cw Ti:sapphire laser operating around 855 nm. The harmonic field from frequency doubling is then used to pump a 4 mm long $KbNO_3$ crystal for type-I parametric down-conversion. The crystal is polished at one end with a 7 mm curvature and at the other with a flat surface. The curved side is optically coated so that it is highly reflective at 855 nm ($R > 99.99\%$) and relatively high in transmission for 427 nm ($T > 70\%$). The flat face is antireflection coated for both 855 and 427 nm. A flat output coupler with a measured 1.5% transmissivity at 855 nm is placed closely (with a 0.5 mm gap) to the flat side of the crystal forming a semimonolithic standing wave cavity C0. Such a compact geometry is designed for an optimum bandwidth of down-conversion taking into consideration both signal level and bandwidth. The output coupler has a relatively high transmission ($T > 80\%$) for 427 nm so that the pump field interacts only once with the nonlinear medium. This eliminates complications involved in resonating the cavity at two wavelengths (855 and 427 nm). It is important to have a good mode match of the pump field to the TEM_{00} mode of the OPO cavity, not only because it can increase the pump efficiency but also because it can inhibit the excitation of the complicated transverse spatial modes of the OPO. The mode match is done with the aid of an auxiliary cavity C1 [26].

The theory presented earlier is for a single mode OPO resonant at degenerate frequency. In the experiment, because the two down-converted fields have the same polarization in the type-I scheme, there are numerous nondegenerate conjugate pairs of down-conversion on resonance simultaneously with the degenerate pair. The nondegenerate pairs ($\omega_0 \pm \Delta\Omega_{FSR}$) will be located in the spectrum on the two sides of the degenerate pair (ω_0) with a spacing of $\Delta\Omega_{FSR}$, the free spectral range of the cavity, and have about the same strength as the degenerate one. So a passive filter is needed to further eliminate them. This is done with another cavity C2 which has a bandwidth larger than that of single mode down-conversion. The resonance condition is achieved by locking C0 and C2 to the laser frequency (which is also the degenerate frequency ω_0 of the down-conversion) with a beam from the laser via photodiodes D3, D4. Part of this beam (auxiliary) is also used to align and mode match C1 to C0 and C0 to C2. However, the locking beam has the same frequency and polarization as the down-converted signal field, creating an enormous background. To eliminate the background, we alternate the periods of cavity locking and signal detection with a mechanical chopper. Because of the rigid and compact structure of C0 and C2, the cavities remain locked even in the period when the locking beam is blocked for signal detection.

In the first part of the experiment, we examine the enhancement effect of parametric down-conversion due to cavity resonance. To have a faithful measurement of the down-conversion signal, we bypass the filter C2 and place the detectors directly at the output of the OPO. Interference filters are placed in front of the detectors [avalanche photodiodes (APD) D1, D2, EG&G SPCM-AQ-121] to eliminate background light mainly from the pump field. We measured the count rate as well as the coincidence rate as a function of the cavity length of C0, as shown in Fig. 2. A strong resonance effect is obvious. At the main peak, we obtain the calibrated single mode count rate $R_{\mathrm{resonance}} = 1.2 \times 10^6/\mathrm{sec}$ at 1 mW pumping. The smaller peaks are higher order transverse cavity modes excited by the imperfectly mode-matched pump. The finesse of the OPO cavity is measured to be $\mathcal{F} = 350$, and the bandwidth of the OPO cavity is $\Delta\omega = 2.8 \times 10^8$ rad/sec. To obtain the enhancement factor per frequency mode for comparison with Eq. (6), we also measured the count rate with the output coupler removed as $R_{\mathrm{single\text{-}pass}} = 10^5/\mathrm{sec}$ at 1 mW. For the single-pass case, the bandwidth is determined by that of interference filter, which is $\Delta\omega_{\mathrm{IF}} = 1.29 \times 10^{12}$ rad/sec ($\Delta\lambda = 0.5$ nm). So the measured enhancement factor is $B_{\exp} = 5.5 \times 10^4$. The theoretical prediction from Eq. (6) gives $B_{\mathrm{th}} = 5.1 \times 10^4$ for measured values of $\mathcal{F} = 350$ and $\mathcal{F}_0 = 420$, which agrees relatively well with the experimental result. The difference may be due to the crude model for the single-pass case.

In the second part of the experiment, we measure directly the time interval distribution between the two

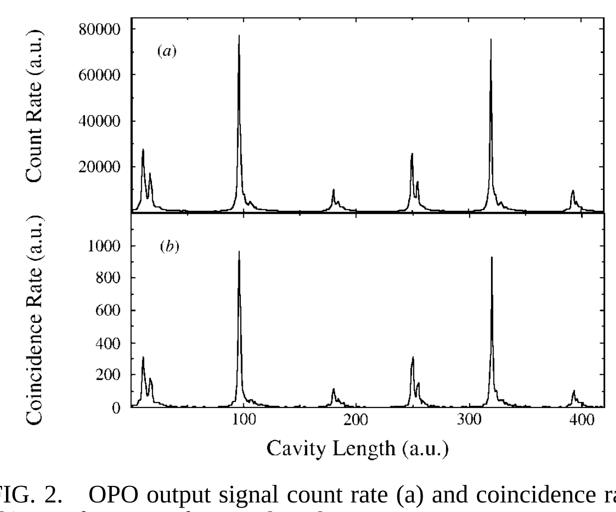

FIG. 2. OPO output signal count rate (a) and coincidence rate (b) as a function of cavity length.

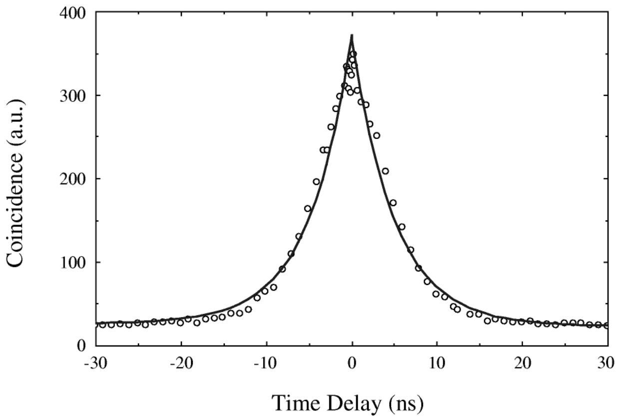

FIG. 3. Directly measured time delay distribution with an exponential fit (solid curve).

photons from parametric down-conversion with a time-to-digital converter (TDC, Lecroy 2228A) and extract out the correlation time from the distribution. Cavity C2 is inserted to eliminate the nondegenerate components for single mode operation. Figure 3 shows a typical plot of the number of coincidence as a function of the difference in arrival times for the two photons. The solid curve is an exponential fit of the data to Eq. (9). As can be seen from Fig. 3, the fit is relatively good. The misfit may come from the modification of the spectrum of the down-converted fields by the filter cavity C2. In fact, we also made a similar measurement without the filter cavity C2, and the exponential curve fits the data very well, but with a quite different time constant due to the multimode nature. Nevertheless, we determine the correlation time T_c as the FWHH from the solid curve in Fig. 3 and make the measurement for a number of the filter bandwidths or equivalently the detected down-conversion bandwidths $\Delta\omega$. Figure 4 plots T_c against $1/\Delta\omega$ and shows a linear dependence as expected, although the slope is 2.3 instead of 0.89 as in Eq. (10), reflecting the modified spectrum by the filter C2. The offset of T_c at the origin is from the finite resolution time of 0.8 ns of the APDs.

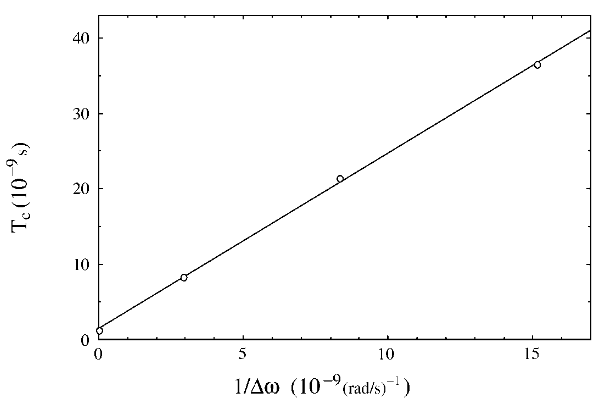

FIG. 4. Correlation time as a function of the inverse of detected bandwidth with a linear fit.

In conclusion, we have substantially reduced the bandwidth of the spontaneous parametric down-conversion without sacrificing the signal level. The correlation time between the two down-converted photons is hence lengthened so that it can be directly measured. Such a two-photon source may find a wide range of applications in quantum information processing.

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