Analysis of the possibility of analog detectors calibration by exploiting Stimulated Parametric Down Conversion

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Abstract: Spontaneous parametric down conversion (SPDC) has been largely exploited as a tool for absolute calibration of photon-counting detectors, i.e. detectors registering very small photon fluxes. In [4] we derived a method for absolute calibration of analog detectors using SPDC emission at higher photon fluxes, where the beam is seen as a continuum by the detector. Nevertheless intrinsic limitations appear when high-gain regime of SPDC is required to reach even larger photon fluxes. Here we show that stimulated parametric down conversion allow one to avoid this limitation, since stimulated photon fluxes are increased by the presence of the seed beam.

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OCIS codes: (270.4180) Multiphoton processes; (120.1880) Detection; (120.3940) Metrology.

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1. **Introduction**

Prompted by the necessity of a precise calibration of photo-detectors both in the analog and photon-counting regime [1, 2, 3], we presented recently [4] a detailed theoretical analysis on the possibility to use the correlations of spontaneous parametric down conversion (SPDC) light [5, 6] for calibrating analog detectors, as an extension of the technique developed in photon counting regime [7, 8, 9, 10, 11, 12, 13, 14, 15]. However, while the method is suitable for analog calibration at relatively low gain, which means a photon flux lower than $10^{10}$ photon/s, in higher-gain regime it is limited by the difficulty of collecting the same correlated modes in the two branches [4].

In this paper we present a detailed theoretical analysis of a scheme based on the stimulated PDC (on the other hand a discussion of the uncertainty budget is left to a specifically addressed paper [16]). We show that this scheme allow to overcome the problem mentioned above. Indeed, in this case the photon fluxes can be varied by varying the power of the coherent seed beam, without increasing the parametric gain $G$. On the other hand, the SPDC non-classical correlation at single-photon level which enables the absolute calibration of counting detectors survives, in some form, when a coherent seed is injected and the photon flux becomes macroscopic.

2. **Analog detection and quantum efficiency**

The schematic set-up for calibration of photodetectors by using PDC is shown in Fig. 1. We model the photodetection process in the analog regime as a random pulse train [17]

$$i(t) = \sum_n q_n f(t - t_n),$$

i.e. a very large number of discrete events at random times of occurrence $t_n$. The pulse shape $f(t)$ is determined by the transit time of charge carriers. We assume that $f(t)$ is a fixed function with characteristic width $\tau_p$ and a unit area. The pulse amplitude $q_n$ is a random variable in order to account for a possible current gain by avalanche multiplication. The statistical nature of the multiplication process gives an additional contribution to the current fluctuations [18].

In an ideal instantaneous photocell response, without avalanche gain, all values $q_n$ are equal to the charge $e$ of a single electron and $f(t) \sim \delta(t)$. In the case of ideal quantum efficiency, since the probability density of observing a photon at time $t$ at detector $D_j$ ($j = 1, 2$) is related to the quantum mean value $\langle F_j(t) \rangle$ of the photon flux, we calculate the average current output of $D_j$...
Fig. 1. Scheme for absolute calibration of analog detectors by using stimulated PDC. The quantum efficiency of detector 2, collecting the seed beam and the stimulated emission, is estimated by the ratio between the cross-correlation function of the photocurrents and the auto-correlation of the photocurrent $i_1$.

as

\[
\langle i_j \rangle = \sum_n \langle q_{jn} f(t - t_n) \rangle = \int dt_n \langle q_j \rangle f(t - t_n) \langle \hat{F}_j(t_n) \rangle
\]

where the factor $\langle q_j \rangle$ is the average charge produced in a detection event. We have assumed the response function for the two detectors to be the same, $f_1(t) = f_2(t) = f(t)$.

At the same time, the auto-correlation and the cross-correlation functions for the currents can be expressed as

\[
\langle i_j(t) i_k(t + \tau) \rangle = \sum_{n,m} \langle q_{jn} q_{km} f(t - t_n) f(t - t_m + \tau) \rangle
\]

\[
= \int \int dt_n dt_m \langle q_{jn} q_{km} f(t - t_n) f(t - t_m + \tau) \rangle \langle \hat{F}_j(t_n) \hat{F}_k(t_m) \rangle,
\]

respectively for $j = k$ where $\langle \hat{F}_j(t_n) \hat{F}_j(t_m) \rangle$ is the auto-correlation function of the photon flux at detector $j$, and for $j \neq k$ where $\langle \hat{F}_j(t_n) \hat{F}_k(t_m) \rangle$ is the cross-correlation between the fluxes incident on the two different detectors. By the substitution $\hat{F}_j \equiv \langle \hat{F}_j \rangle + \delta \hat{F}_j$, it is convenient to express them as

\[
\langle \hat{F}_j(t_n) \hat{F}_k(t_m) \rangle = \langle \hat{F}_j \rangle \langle \hat{F}_k \rangle + \langle \hat{F}_j \rangle \delta(t_n - t_m) \delta_{jk} + \langle \delta \hat{F}_j(t_n) \delta \hat{F}_k(t_m) \rangle.
\]

The second contribution, proportional to the photon flux when $j = k$, represents the intrinsic and unavoidable component of the fluctuation that does not depend on the specific property of the field since it generates from the commutation relation of the quantum fields in the free
with conjugate frequencies, linking the fields interaction started. The coefficients approximation, only pairs of modes with opposite transverse wave vectors, process are the input-output transformations the energy and momentum conservation hold. The equations describing the down conversion

Here we consider a type I PDC process. In the limit of monochromatic and plane-wave pump approximation, only pairs of modes with opposite transverse wave vectors, \( \omega_1 = \omega_{\text{pump}}/2 - \Omega \) and \( \omega_2 = \omega_{\text{pump}}/2 + \Omega \), are coupled such that the energy and momentum conservation hold. The equations describing the down conversion process are the input-output transformations

\[
\begin{align*}
\hat{a}_1^{\text{out}}(\mathbf{q}, \Omega) &= U_1(\mathbf{q}, \Omega) \hat{a}_1^{\text{in}}(\mathbf{q}, \Omega) + V_1(\mathbf{q}, \Omega) \hat{a}_2^{\text{in}}(\mathbf{q}, -\Omega) \\
\hat{a}_2^{\text{out}}(\mathbf{q}, \Omega) &= U_2(\mathbf{q}, \Omega) \hat{a}_2^{\text{in}}(\mathbf{q}, \Omega) + V_2(\mathbf{q}, \Omega) \hat{a}_1^{\text{in}}(\mathbf{q}, -\Omega)
\end{align*}
\]  

where we introduced the convolution of the response function of detectors \( \mathcal{F}(\tau) = \int dt \int f(t) f(t + \tau) \).

3. Correlation functions of stimulated PDC

Here we consider a type I PDC process. In the limit of monochromatic and plane-wave pump approximation, only pairs of modes with opposite transverse wave vectors, \( \mathbf{q} \) and \( -\mathbf{q} \), are coupled such that the energy and momentum conservation hold. The equations describing the down conversion process are the input-output transformations

\[
\begin{align*}
\langle i_j \rangle &= \eta_i \langle j_1 \rangle \langle f_j \rangle \\
\langle i_j \rangle_k(t + \tau) &= \langle i_j \rangle_k + \eta_j \langle j_2 \rangle \mathcal{F}(\tau) \langle f_j \rangle \delta_{jk} \\
&+ \eta_j \eta_k \langle j_2 j_k \rangle \int dt \int dt_m f(t - t_n) f(t - t_m + \tau) \langle \delta \hat{F}_j(t_n) \delta \hat{F}_k(t_m) \rangle,
\end{align*}
\]  

which guarantee the conservation of the free-field commutation relations

\[
\begin{align*}
|U_k(\mathbf{q}, \Omega)|^2 - |V_k(\mathbf{q}, \Omega)|^2 &= 1, \quad (k = 1, 2) \\
U_1(\mathbf{q}, \Omega) V_2(-\mathbf{q}, -\Omega) &= U_2(-\mathbf{q}, -\Omega) V_1(\mathbf{q}, \Omega)
\end{align*}
\]  

Within this picture it is possible to take into account the quantum efficiency by the following substitutions:

\[
\begin{align*}
\langle \hat{F}_j(t) \rangle &\rightarrow \eta_j \langle \hat{F}_j(t) \rangle \\
\langle \hat{F}_j(t) \hat{F}_k(t') \rangle &\rightarrow \eta_j \eta_k \langle \hat{F}_j(t) \hat{F}_k(t') \rangle.
\end{align*}
\]  

Thus, \( \langle \hat{F}_j(t) \rangle \) being time independent, according to Eq. (1) we obtain:

\[
\langle i_j \rangle = \eta_i \langle j_1 \rangle \langle f_j \rangle.
\]  

Eq. (2) becomes

\[
\begin{align*}
\langle i_j(t) \rangle_k(t + \tau) &= \langle i_j \rangle_k + \eta_j \langle j_2 \rangle \mathcal{F}(\tau) \langle f_j \rangle \delta_{jk} \\
&+ \eta_j \eta_k \langle j_2 j_k \rangle \int dt \int dt_m f(t - t_n) f(t - t_m + \tau) \langle \delta \hat{F}_j(t_n) \delta \hat{F}_k(t_m) \rangle,
\end{align*}
\]  

where we introduced the convolution of the response function of detectors \( \mathcal{F}(\tau) = \int dt f(t) f(t + \tau) \).
Here, $U_k$ and $V_k$ define the strength of the process and at the same time the range of transverse momentum and frequencies in which it takes place and thus are named gain functions. In fact, the finite length of the crystal introduces a partial relaxation of phase matching condition concerning the third component of the momenta. By selecting a certain frequency, the transverse momentum uncertainty, i.e. the angular dispersion of the emission direction, is not null ($\Delta q \sim (l \tan \vartheta)^{-1}$) for the non-collinear PDC, where $l$ is the crystal length and $\vartheta$ is the central angle of propagation with respect to the pump direction). On the contrary, by fixing the transverse momentum $q$, or equivalently a direction of propagation $\vartheta$, the spectral bandwidth $\Delta \Omega$ is proportional to $l^{-1/2}$, and typically $1/\Delta \Omega \sim 10^{-13}$ for type I.

Hereinafter we will focus on the properties of the far field, observed at the focal plane of a thin lens of focal length $f$, placed at distance $f$ from the crystal. The spatial distribution of the far field is, in this case, the Fourier transform of the field just at the output face of the crystal. This special imaging configuration is convenient to show the basic of calculation, nonetheless the validity of the results is more general. Thus, any transverse mode $q$ is associated with a single point $x$ in the detection (focal) plane according to the geometric transformation $q \rightarrow 2\pi x/(\lambda f)$. The far field operator in the space-temporal domain is therefore

$$\hat{B}_k(x,t) = \sum_{\Omega} e^{-i\Omega t} \hat{a}^{\text{out}}_k \left( \frac{2\pi}{\lambda f} x, \Omega \right).$$

(9)

The mean value of the operator $\hat{I}_k(x,t) \equiv \hat{B}_k^\dag(x,t) \hat{B}_k(x,t)$ is

$$\langle \hat{I}_k(x,t) \rangle = \sum_{\Omega\Omega'} e^{-i(\Omega'-\Omega)t} \langle \hat{a}^{\text{out}}_k \left( \frac{2\pi}{\lambda f} x, \Omega \right) \hat{a}^{\text{out}}_k \left( \frac{2\pi}{\lambda f} x, \Omega' \right) \rangle,$$

(10)

representing the intensity profile of the emission in the detection plane. This mean value has to be calculated over the following initial state, in which the field 1 is in the vacuum state whereas the field 2 is in a multi-mode coherent state:

$$|\psi_m\rangle = |0\rangle_1 \bigotimes_{q,\Omega} |\alpha(q,\Omega)\rangle_2,$$

(11)

where

$$\alpha(q,\Omega) = \alpha(q) \delta_{\Omega_0,\Omega}$$

is the complex parameter associated to this state. The seed beam is therefore represented as a coherent state of field 2 with fixed frequency $\Omega_0$ and a certain distribution of the transverse momentum $q$. Here we assume that, for the modes that are stimulated by the seed, the spontaneous component of the emission is negligible with respect to the stimulated one. This corresponds to the assumption $|\hat{V}_2(q,\Omega_0)|^2 \ll 1$ and, at the same time, $|\alpha(q)|^2 \gg 1$, that is a typical experimental situation in which a few millimeters length crystal is pumped by a continuous pump and the seed has a power just around the microwatt or more. The intensities (10) of the two stimulated beams after the crystal can be evaluated by using Eq. (7), with the substitution $q \rightarrow 2\pi x/(\lambda f)$. This leads to

$$\langle \hat{I}_1(x,t) \rangle \approx |\hat{V}_1(x,-\Omega_0)|^2 |\alpha(-x)|^2$$

(12)
\[ \langle \tilde{I}_2(x,t) \rangle \approx \frac{\bar{U}_2(x,\Omega_0)}{\bar{V}_2(x,\Omega_0)} |\bar{\alpha}(x)|^2 \]
\[ = \left( \bar{V}_2(x,\Omega_0)^2 + 1 \right) |\bar{\alpha}(x)|^2 \]  
(13)

where we defined
\[ \bar{U}_k(x,\Omega) = U_k \left( \frac{2\pi}{\lambda f} x, \Omega \right) , \]
\[ \bar{V}_k(x,\Omega) = V_k \left( \frac{2\pi}{\lambda f} x, \Omega \right) , \]
\[ \bar{\alpha}(x) = \alpha \left( \frac{2\pi}{\lambda f} x \right) . \]  
(14)

We used the first property of the gain functions in (8) into the last line of Eq. (13), in order to show explicitly the two contributions to the beam 2, one given by the original coherent seed and the other from the stimulated emission, proportional to the parametric gain \( G = \max |F_k(x,\Omega)|^2 \). We notice that, aside from the spontaneous emission that is neglected, the generated beams conserve the original momentum distribution of the seed but weighted according to the gain function.

The normal-ordered correlation function of the intensities fluctuation is defined as
\[ \langle : \delta \tilde{I}_k(x,t) \delta \tilde{I}_j(x',t') : \rangle = \langle \tilde{I}_k(x,t) \tilde{I}_j(x',t') \rangle - \langle \tilde{I}_k(x,t) \rangle \langle \tilde{I}_j(x',t') \rangle \]  
(15)

The calculation leads to the following normal-ordered mean values:
\[ \langle : \delta \tilde{I}_1(x,t) \delta \tilde{I}_1(x',t') : \rangle \approx 0 , \]  
(16)

\[ \langle : \delta \tilde{I}_2(x,t) \delta \tilde{I}_2(x',t') : \rangle \approx \sum_{\Omega} e^{-i(\Omega_0 - \Omega)(t-t')} |\bar{\alpha}(x)|^2 \]
\[ \bar{V}_2(x,\Omega)^2 |\bar{U}_2(x,\Omega_0)|^2 \delta_{xx'} \]
\[ + \text{c.c.} \]  
(17)

Fluctuations of the stimulated field 1 are negligible, because they have the same order of magnitude of the spontaneous emission. Concerning the cross-correlation we have:
\[ \langle : \delta \tilde{I}_1(x,t) \delta \tilde{I}_2(x',t') : \rangle \approx \sum_{\Omega} e^{-i(\Omega_0 + \Omega)(t-t')} |\bar{\alpha}(-x)|^2 \]
\[ \bar{U}_2(-x,\Omega_0) \bar{U}_2^*(x,-\Omega) \]
\[ \bar{V}_1(x,-\Omega_0) \bar{V}_1^*(x,\Omega) \delta_{xx'} \]
\[ + \text{c.c.} \]  
(18)

Let us consider two detectors, \( D_1 \) and \( D_2 \), to register the photons crossing the regions \( R_1 \) and \( R_2 \) centered on the two correlated bright beams in the detection plane. Let the two areas be much larger than the dimension of the seed beam on the detection plane. This simply means to collect all the bright components of the emission, corresponding to the modes stimulated by the seed.

We move now from the discrete modes representation of the electromagnetic field, in which it is artificially enclosed in a fictitious cube of side \( L \), to a continuous set of modes for \( L \rightarrow \infty \).
Thus, in the following, the \( \Omega \)-sum and the Kronecker function \( \delta_{xx'} \) in Eqs. (17) and (18) are substituted respectively with the integral over \( \Omega \) and the Dirac function \( \delta(x-x') \). The fact that the second-order correlation function has an unphysical spatial behaviour, dominated by a delta function, comes from the starting assumption of a plane-wave pump, having infinite transverse dimension. In a realistic case the delta function is replaced by the Fourier transform of the pump transverse profile.

The quantum mean values of the photon fluxes \( F_j \equiv \int_R \hat{I}_j(x,t) \, dx \) reaching detectors 1 and 2 can be obtained integrating equations (12) and (13) respectively on \( R_1 \) and \( R_2 \). By applying relations (8) we have

\[
\langle \delta F_1(t) \delta F_1(t') \rangle \approx \int_{R_1} dx \mid \tilde{V}_1(x,-\Omega_0) \mid^2 |\tilde{\alpha}(-x)|^2, \tag{19}
\]

\[
\langle \delta F_2(t) \delta F_2(t') \rangle \approx \int_{R_2} dx \left( |\tilde{V}_2(x,\Omega_0)|^2 + 1 \right) |\tilde{\alpha}(x)|^2. \tag{20}
\]

At the same time integrating Eqs. (16) and (17) over \( R_1 \times R_1 \) and \( R_2 \times R_2 \) respectively, we have the following auto-correlation functions for the photon fluxes

\[
\langle \delta F_1(t) \delta F_1(t') \rangle \approx 0 \tag{21}
\]

\[
\langle \delta F_2(t) \delta F_2(t') \rangle \approx \sum_{\Omega} e^{-i(\Omega_0-\Omega)(t-t')} \int_{R_2} dx |\tilde{V}_2(x,\Omega)|^2 |\tilde{\alpha}(x)|^2 + c.c., \tag{22}
\]

Here and in the following, we have assumed that in the limit of small \( G \) the gain functions \( |\tilde{U}_j(x,\Omega)| \sim 1 \), according to the relations (3). Then, integrating Eq. (18) over \( R_1 \times R_2 \) gives the following cross-correlation:

\[
\langle \delta F_1(t) \delta F_2(t') \rangle \approx \sum_{\Omega} e^{-i(\Omega_0+\Omega)(t-t')} \int_{R_1} dx |\tilde{\alpha}(-x)|^2 \tilde{V}_1(x,-\Omega_0) \tilde{V}_1^*(x,\Omega) + c.c. \tag{23}
\]

Although the seed beam is monochromatic with frequency \( \Omega_0 \), the temporal width of both auto-correlation and cross-correlation in a fixed point \( x \) still depends on the spectral bandwidth \( \Delta \Omega \) of the gain functions. This explains why they originate from the correlation between the monochromatic seed and the spontaneous emission that is broadband. Therefore, the coherence time of the stimulated emission is still of the order of \( \tau_{coh} = 1/\Delta \Omega \), analogously to the spontaneous case.

4. Quantum efficiency estimation

The auto-correlation of the current fluctuations \( \delta i_1(t) = i_1(t) - \langle i_1 \rangle \) can be calculated by introducing the result in Eq. (21) in (6), obtaining

\[
\langle \delta i_1(t) \delta i_1(t+\tau) \rangle = \eta_1 \langle q_1^2 \rangle \mathcal{F}(\tau) \langle F_1 \rangle. \tag{24}
\]
Under the condition of small gain ($G \ll 1$) and large intensity of the seed ($|\alpha|^2 \gg 1$) the fluctuation of the current at detector $D_1$ is dominated by the shot noise component.

By substituting Eq. (23) in (6), under the same condition, for $j = 1$ and $k = 2$, one has

$$\langle \delta i_1(t) \delta i_2(t + \tau) \rangle = 2 \eta_1 \eta_2 \langle q_1 \rangle \langle q_2 \rangle \mathcal{F}(\tau) \langle F_1 \rangle$$

(25)

where we assumed $\tau_p \gg \tau_{coh}$. We stress that it is the usual situation, the coherence time of SPDC being on the order of picoseconds or less and the typical resolving time of detectors on the order of nanosecond. In this case any fluctuations in the light power are averaged over $\tau_p$.

The cross-correlation function of the fluctuations has the same form as the one obtained in the case of spontaneous down-conversion in our earlier work [4]. This is the evidence that quantum correlations do not disappear when the emission is stimulated. In fact, the down-converted photons are still produced in pairs, although in the beam 2 the photons of the pairs are added to the bright original coherent beam propagating in the same direction. The factor 2, appearing in Eq. (25) can be interpreted as due to the fact that, for any down-converted photon of a pair propagating along direction 2, there is also the original photon of the seed that stimulated the generation of that pair. Formally it emerges from the time integration of two identical contribution in (23), the explicit one and its complex conjugated.

Since we are interested in a relatively large power of incident light, we can at first consider detectors without internal gain, and assume that the charge produced in any detection event is equal to the single electron charge $q$, i.e. $\langle q_k \rangle = q$ and $\langle q_k^2 \rangle = q^2$. Therefore, according to Eqs. (24) and (25), the quantum efficiency can be evaluated as

$$\eta_2 = \frac{1}{2} \frac{\langle \delta i_1(t) \delta i_2(t + \tau) \rangle}{\langle \delta i_1(t) \delta i_1(t + \tau) \rangle}.$$  

(26)

For detectors in which electrons multiplication occurs the statistic fluctuations of this process do not allow to use Eq. (26) for absolute calibration. However, it can be performed by integrating Eq. (25) over time $\tau$ [4]. It corresponds to evaluating the power spectrum of the fluctuations at frequencies around zero, namely much smaller than $1/\tau_p$. We would like to stress that in this case, the assumption $f_1(t) = f_2(t) = f(t)$ is not necessary. Since $\int d\tau \mathcal{F}(\tau) = 1$ we obtain

$$\eta_2 \langle q_2 \rangle = \frac{1}{2} \frac{\int d\tau \langle \delta i_1(t) \delta i_2(t + \tau) \rangle}{\langle i_1 \rangle},$$

(27)

In Eq. (27) the statistic of the electron gain does not play any role.

5. Conclusion

With the purpose to eliminate some problems that appear when working with Spontaneous PDC [4], we have extended our method to higher flux regimes, taking advantage of the stimulated PDC as a source of bright correlated beams. We show that this scheme allows one to overcome these former problem: an important result in view of metrology applications. Indeed this result allows to link the photon counting regime with the photon rates of the traditional optical metrology supplying a single radiometric standard for all these regimes. Finally, we would like to stress once more that equations (26) and (27), between quantum efficiency and current fluctuations, even if identical to the ones obtained in our previous work [4], are here derived for a more general condition of stimulated emission.

Acknowledgments

The Turin group acknowledges the support of Regione Piemonte (ricerca scientifica applicata E14) and San Paolo Foundation. The Russian group acknowledges the support of Russian Foun-
dation for Basic Research (grant # 06-02-16393). Both acknowledge the joint project of Associazione Sviluppo del Piemonte by Grant RFBR-PIEDMONT # 07-02-91581-APS.