Two-photon pattern in a second-order interference

H. H. Arnaut[1] and G. A. Barbosa[2,*]
[1]Departamento de Física, Universidade Federal de Minas Gerais, Belo Horizonte 30161-970, MG, Brazil
[2]Avenida Portugal 1558, Belo Horizonte 31550-000, MG, Brazil
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An interferometric arrangement utilizing a pair of twin photons, generated by parametric down-conversion, is devised to show second-order interferences revealing de Broglie's wavelength of the photon pair as a whole and the Pancharatnam phase depending on the number of photons. This proposed experiment utilizes the phenomenon ''induced coherence without stimulated emission'' to generate a coherent superposition of states distinct from the usual superposition with beam splitters, where the photons are *independently* scattered. In this experiment interference occurs due to recombination of photon trajectories on a beam splitter and also through frustration or enhancement of the down-conversion processes. [S1050-2947(98)04211-5]

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I. INTRODUCTION

Recently Jacobson *et al.* [1] showed that interferometric measurements of *de Broglie's* wavelength of a composite object depends fundamentally on the interaction between the object and the beam splitters of the interferometer. This interaction determines the quantum state of the scattered beams and, in principle, different beam splitter interactions can be *simulated* if one knows how to modify the wave function of these beams appropriately.

Two-photon particle sources are easily produced by pumping nonlinear $\chi^{(2)}$ crystals with ultraviolet laser beams [2] in the process of *parametric down-conversion luminescence* (PDC). The availability of these particles suggests that their *de Broglie's* wavelength could be measured to check the lack of energy interaction between the constituent photons as well as to have the transition from a two-photon to a one-photon particle followed. This transition can be obtained in an interferometer designed to allow careful control of the interference processes.

This work explores a viable interferometric arrangement to measure the associated wavelength of controlled superpositions of one- or two-photon states and discuss the obtained results. This arrangement includes the possibility of introducing a geometric phase (Pancharatnam's phase [3]) in a trajectory, depending on the number of photons present. Two basic interferometric elements are used on the trajectories defining interfering possibilities: lossless—passive and linear—beam splitters (constant transmissivity and reflectivity) and a nonlinear two-mode converter.

The beam splitter exposes the inherent lack of photon-photon interaction at optical frequencies due to the fact that, whenever a two-photon wave packet is incident upon it, the output state shows a binomial probability of occurrence of single photons in either side of the beam splitter, as an *independent* scattering process for each constituting photon. On the other hand, the phenomenon of *induced coherence without stimulated emission* is used in the two-mode converter to produce enhancement or inhibition of photons, as an active quantum ''beam splitter'' able to interact on a two-photon state as a whole. The induced coherence without stimulated emission phenomenon has no classical analog as has been demonstrated by Mandel and collaborators [4]. Spontaneous emission in this case is several orders of magnitude *larger* than any contribution coming from stimulated emission and, in this sense, the situation analyzed is purely quantum mechanical and cannot be explained by a classical theory. A longitudinal multimode theory of PDC is applied to this analysis.

II. PRINCIPLE OF THE PROPOSED EXPERIMENT

The proposed interferometer is outlined in Fig. 1. The two crystals NL1 (centered at $\mathbf{R}=\mathbf{R}_0$) and NL2 (centered at $\mathbf{R}=\mathbf{0}$) have nonlinear susceptibilities $\chi^{(2)}$ and are optically pumped by the same laser beam of midfrequency $\overline{\omega_p}$. Down-converted fields of type I (horizontally polarized) on trajectories (0,1) and (4,5) are generated at NL1 and NL2, respectively. The conjugated beams emerging from NL1 are mixed by the beam splitter BS1 and then superposed to the down-converted beams emerging from NL2. This superposition process on NL2 is a key step in the sense that it erases the signature of the absence of photon-photon interaction, i.e., the binomial statistics of the scattered beams produced at BS1. The superposed beams are again mixed by the beam splitter BS2 so that the detection system (detectors D_6 and D_7) *cannot* distinguish from which crystal a photon was generated [5].

Distances between the nonlinear converters and the beam splitters can be piezoelectrically controlled. An arrangement to introduce a Pancharatnam phase θ on the trajectory 3 is represented by $U(\theta)$. This arrangement is composed of the following succession of wave plates: a quarter-wave plate with the fast axis at 45° from horizontal followed by a half-wave plate that can be set at an arbitrary angle θ relative to the first plate axis and a last quarter-wave plate also set with the fast axis at 45° from the horizontal. A horizontally polarized photon reaching the first plate will be circularly right-polarized and passing by the half-wave plate will be set left-polarized and will emerge again horizontally polarized after the last quarter-wave plate. In Poincaré's sphere of polarization [6] this will be represented by a closed circuit on its surface defining a sector of angle 2θ.

This experimental setup is able to produce either two-photon or one-photon states through controlled variations of

*Electronic address: gbarbosa@gold.com.br

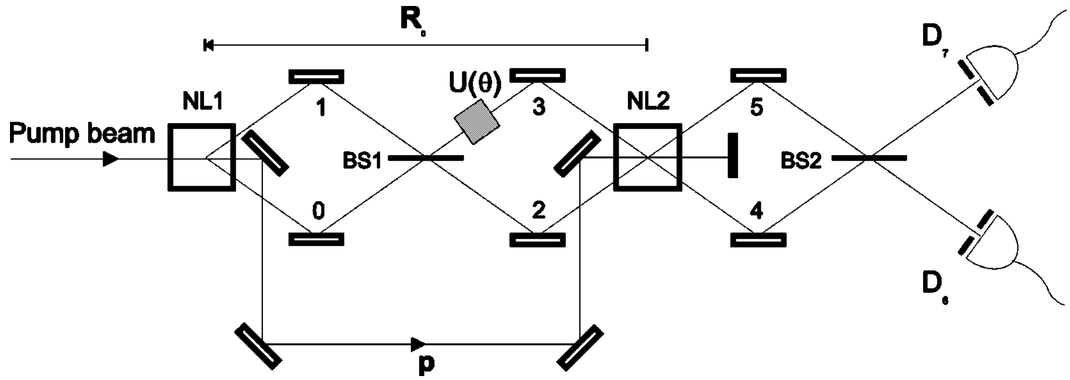

FIG. 1. Two nonlinear crystals NL1 and NL2 are optically pumped by the same laser beam. Down-converted fields propagate on trajectories (0,1), (2,3), and (4,5). The conjugated beams emerging from NL1 are mixed by the beam splitter BS1 and then superposed to the down-converted beams emerging from NL2. The superposed beams are again mixed by the beam splitter BS2 so that the detection system (detectors D_6 and D_7) *cannot* distinguish from which crystal a photon was generated. A Pancharatnam phase θ on the trajectory 3 can be introduced by the set of wave plates represented by $U(\theta)$.

wave-packet superpositions at the output of beam splitter BS1. Pancharatnam's phase θ can be introduced through $U(\theta)$ on path 3 and the superposition of states generated from down-converters NL1 and NL2 on BS2 makes it possible to obtain second-order interference patterns of one-photon and two-photon states at the detection system.

The analysis of these possibilities given by this experimental arrangement demands a study of the complete quantum state produced at beam splitter BS2. This quantum state should contain all information that can be obtained with the photon counting system, such as dependence on controlled variations of the photon trajectories and on Panchratnam's angle θ.

III. QUANTUM STATE OF THE DOWN-CONVERTED FIELD

The pump beam will be represented classically by a complex analytical signal $V(t)=\Sigma_{\omega_p} v(\omega_p)\exp(-i\omega_p t)$ and such that its mean intensity $\langle I_p\rangle=\langle V^*(t)V(t)\rangle$ is in units of photons per second (approximation of narrow bandwidth detection). The phase-matching conditions define basic directions for the down-converted modes. These directions will be experimentally specified by pinholes placed along the photon trajectories and in such a way that transverse effects due to the finite size of the pinholes can be neglected. Under this usual approximation of considering only longitudinal variations, the first-order approximation for the unitary time evolution operator, in the interaction picture, reads [7]

$$\begin{aligned}\hat{U}(t,0)=1&+\frac{\eta_1(\delta\omega)^{3/2}}{\sqrt{2\pi}}\sum_{\omega}\sum_{\omega'}\sum_{\omega_p}\phi_1(\omega_p,\omega,\omega')v(\omega_p)\\&\times\frac{\sin(\Omega t/2)}{(\Omega/2)}e^{i\Omega t/2}e^{i(\mathbf{K}_p-\mathbf{K}-\mathbf{K}')\cdot\mathbf{R}_0}\hat{a}_0^\dagger(\omega)\hat{a}_1^\dagger(\omega')\\&+\frac{\eta_2(\delta\omega)^{3/2}}{\sqrt{2\pi}}\sum_{\omega}\sum_{\omega'}\sum_{\omega_p}\phi_2(\omega_p,\omega,\omega')v(\omega_p)\\&\times\frac{\sin(\Omega t/2)}{(\Omega/2)}e^{i\Omega t/2}\hat{a}_4^\dagger(\omega)\hat{a}_5^\dagger(\omega').\end{aligned}\tag{1}$$

Here $\Omega=\omega+\omega'-\omega_p$, η_j $(j=1,2)$ is the quantum efficiency of the crystal j and $\phi_j(\omega_p,\omega,\omega')$ is the spectral function of the down-converted light, symmetric in ω and ω', peaked at $\omega=\omega'=\overline{\omega_p}/2=\overline{\omega}$ and normalized so that

$$2\pi\delta\omega\sum_{\omega}|\phi_j(\overline{\omega_p},\overline{\omega_p}-\omega,\omega)|^2=1.\tag{2}$$

Similar crystals will be considered such that the index j in the spectral function can be dropped out.

In the case where the time interval between two successive down-conversions is long compared to the coherence time T_{DC} of the down-converted beams, the first-order approximation (1) for the unitary time evolution operator is sufficient to describe the main features and characteristics of the system [8].

The following identifications are introduced after consideration of the transformations on the beam splitters:

$$v(\omega_p)e^{i\mathbf{K}_p\cdot\mathbf{R}_0}=v(\omega_p)e^{-i\omega\tau_p},\tag{3}$$

$$\hat{a}_0^\dagger(\omega)e^{-i\mathbf{K}\cdot\mathbf{R}_0}=[\mathcal{T}\hat{a}_4^\dagger(\omega)e^{i\theta}e^{i\omega\tau_3}+\mathcal{R}\hat{a}_5^\dagger(\omega)e^{i\omega\tau_2}]e^{i\omega\tau_0},\tag{4}$$

$$\hat{a}_1^\dagger(\omega)e^{-i\mathbf{K}\cdot\mathbf{R}_0}=[\mathcal{R}\hat{a}_4^\dagger(\omega)e^{i\theta}e^{i\omega\tau_3}+\mathcal{T}\hat{a}_5^\dagger(\omega)e^{i\omega\tau_2}]e^{i\omega\tau_1},\tag{5}$$

where $\mathcal{R}$, $\mathcal{T}$ are the complex reflectivity and transmissivity [9] of the beam splitter BS1, and $c\tau_\alpha$ is the length of the arm labeled α of the interferometer outlined in Fig. 1. Substituting Eqs. (3), (4), and (5) in Eq. (1) and considering that the down-converted modes are in the vacuum state at time $t=0$, we find the final form for the state of the field at a later time t:

$$
\begin{aligned}
|\psi(t)\rangle=&|\mathrm{vac}\rangle+\frac{\eta_2(\delta\omega)^{3/2}}{\sqrt{2\pi}}\sum_{\omega}\sum_{\omega'}\sum_{\omega_p}\phi(\omega_p,\omega,\omega')v(\omega_p)\frac{\sin(\Omega t/2)}{(\Omega/2)}e^{i\Omega t/2}|1_\omega\rangle_4|1_{\omega'}\rangle_5\\
&+\frac{\eta_1(\delta\omega)^{3/2}}{\sqrt{2\pi}}\sum_{\omega}\sum_{\omega'}\sum_{\omega_p}\phi(\omega_p,\omega,\omega')v(\omega_p)\frac{\sin(\Omega t/2)}{(\Omega/2)}e^{i\Omega t/2}e^{i\omega\tau_0}e^{i\omega'\tau_1}e^{-i\omega_p\tau_p}[(\mathcal{T}\mathcal{R}e^{2i\theta}e^{i(\omega+\omega')\tau_3}|1_\omega\rangle_4|1_{\omega'}\rangle_4\\
&+\mathcal{T}\mathcal{R}e^{i(\omega+\omega')\tau_2}|1_\omega\rangle_5|1_{\omega'}\rangle_5)(1-\delta_{\omega,\omega'})+\sqrt{2}(\mathcal{T}\mathcal{R}e^{2i\theta}e^{i(\omega+\omega')\tau_3}|2_\omega\rangle_4+\mathcal{T}\mathcal{R}e^{i(\omega+\omega')\tau_2}|2_\omega\rangle_5)\delta_{\omega,\omega'}\\
&+\mathcal{T}^2e^{i\theta}e^{i\omega\tau_3}e^{i\omega'\tau_2}|1_\omega\rangle_4|1_{\omega'}\rangle_5+\mathcal{R}^2e^{i\theta}e^{i\omega\tau_2}e^{i\omega'\tau_3}|1_{\omega'}\rangle_4|1_\omega\rangle_5]
\end{aligned}
\tag{6}
$$

IV. RATES OF PHOTON DETECTION

The state $|\psi(t)\rangle$ given by Eq. (6) expresses the input state at the beam splitter BS2. Photon detection rates for the down-converted light on the output side of BS2 can be easily obtained from the electric field operators $\hat{E}_6^{(+)}(t)$ and $\hat{E}_7^{(+)}(t)$ expanded in terms of the annihilation operators at the input of BS2 [10]

$$
\begin{aligned}
\hat{E}_6^{(+)}(t)&=\left(\frac{\delta\omega}{2\pi}\right)^{1/2}\sum_{\omega_6}[R\hat{a}_4(\omega_6)e^{i\omega_6\tau_4}+T\hat{a}_5(\omega_6)e^{i\omega_6\tau_5}]e^{-i\omega_6 t},\\
\hat{E}_7^{(+)}(t)&=\left(\frac{\delta\omega}{2\pi}\right)^{1/2}\sum_{\omega_7}[T\hat{a}_4(\omega_7)e^{i\omega_7\tau_4}+R\hat{a}_5(\omega_7)e^{i\omega_7\tau_5}]e^{-i\omega_7 t},
\end{aligned}
\tag{7}
$$

where R, T are the complex reflectivity and transmissivity of the beam splitter BS2. The rate of photon detection by detector D_6 at time t is proportional to

$$
R_6(t)=\langle\langle\psi(t)|\hat{E}_6^{(-)}(t)\hat{E}_6^{(+)}(t)|\psi(t)\rangle\rangle_{\mathrm{clas.}},
\tag{8}
$$

where the external angular brackets stand for the average over the classical ensemble of laser amplitudes. The rate of photon detection by detector D_7 can readily be obtained from $R_6(t)$ by the substitution $T\to R$, $R\to T$. Equations (6) and (7) give directly $R_6(t)$:

$$
\begin{aligned}
R_6(t)=\Bigg\langle\Bigg|&\frac{\eta_2(\delta\omega)^2}{2\pi}\sum_{\omega}\sum_{\omega_p}\sum_{\omega'}v(\omega_p)\phi(\omega_p,\Omega+\omega_p-\omega,\omega)\frac{\sin(\Omega t/2)}{(\Omega/2)}[e^{-i\Omega(t/2-\tau_4)}Re^{i(\omega_p-\omega)(\tau_4-t)}|1_\omega\rangle_5\\
&+e^{-i\Omega(t/2-\tau_5)}Te^{i(\omega_p-\omega)(\tau_5-t)}|1_\omega\rangle_4]+\frac{\eta_1(\delta\omega)^2}{2\pi}\sum_{\omega_p}\sum_{\omega}\sum_{\omega'}v(\omega_p)\phi(\omega_p,\Omega+\omega_p-\omega,\omega)\frac{\sin(\Omega t/2)}{(\Omega/2)}\\
&\times[\mathcal{T}\mathcal{R}Re^{2i\theta}e^{i\omega(\tau_1-\tau_0-\tau_4+t)}e^{i\omega_p(\tau_0+\tau_3+\tau_4-\tau_p-t)}e^{-i\Omega(t/2-\tau_0-\tau_3-\tau_4)}|1_\omega\rangle_4\\
&+\mathcal{T}\mathcal{R}Re^{2i\theta}e^{i\omega(\tau_0-\tau_1-\tau_4+t)}e^{i\omega_p(\tau_1+\tau_3+\tau_4-\tau_p-t)}e^{-i\Omega(t/2-\tau_1-\tau_3-\tau_4)}|1_\omega\rangle_4\\
&+\mathcal{T}\mathcal{R}Te^{i\omega(\tau_1-\tau_0-\tau_5+t)}e^{i\omega_p(\tau_0+\tau_2+\tau_5-\tau_p-t)}e^{-i\Omega(t/2-\tau_0-\tau_2-\tau_5)}|1_\omega\rangle_5\\
&+\mathcal{T}\mathcal{R}Te^{i\omega(\tau_0-\tau_1-\tau_5+t)}e^{i\omega_p(\tau_1+\tau_2+\tau_5-\tau_p-t)}e^{-i\Omega(t/2-\tau_1-\tau_2-\tau_5)}|1_\omega\rangle_5\\
&+\mathcal{T}^2Re^{i\theta}e^{i\omega(\tau_1+\tau_2-\tau_0-\tau_3-\tau_4+t)}e^{i\omega_p(\tau_0+\tau_3+\tau_4-\tau_p-t)}e^{-i\Omega(t/2-\tau_0-\tau_3-\tau_4)}|1_\omega\rangle_5\\
&+\mathcal{T}^2Te^{i\theta}e^{i\omega(\tau_0+\tau_3-\tau_1-\tau_2-\tau_5+t)}e^{i\omega_p(\tau_1+\tau_2+\tau_5-\tau_p-t)}e^{-i\Omega(t/2-\tau_1-\tau_2-\tau_5)}|1_\omega\rangle_4\\
&+\mathcal{R}^2Re^{i\theta}e^{i\omega(\tau_0+\tau_2-\tau_1-\tau_3-\tau_4+t)}e^{i\omega_p(\tau_1+\tau_3+\tau_4-\tau_p-t)}e^{-i\Omega(t/2-\tau_1-\tau_3-\tau_4)}|1_\omega\rangle_5\\
&+\mathcal{R}^2Te^{i\theta}e^{i\omega(\tau_1+\tau_3-\tau_0-\tau_2-\tau_5+t)}e^{i\omega_p(\tau_0+\tau_2+\tau_5-\tau_p-t)}e^{-i\Omega(t/2-\tau_0-\tau_2-\tau_5)}|1_\omega\rangle_4]\Bigg|^2\Bigg\rangle_{\mathrm{clas.}}.
\end{aligned}
\tag{9}
$$

Due to the continuum character of the PDC light, terms $\delta\omega\Sigma$ summed over ω' can be approximate to an integral in Ω and, observing that as t becomes much longer than T_{DC} the dominant contribution for the Ω integral comes from $\Omega\ll\Delta\omega$, where $\Delta\omega$ is the bandwidth of the downconverted light, one may write $\phi(\omega_p,\Omega+\omega_p-\omega,\omega)\simeq\phi(\omega_p,\omega_p-\omega,\omega)$ and approximate the Ω integral by the standard Dirichlet integral

$$
\frac{1}{2\pi}\int_{-\infty}^{\infty}\frac{\sin(\Omega t/2)}{\Omega/2}e^{\pm i\Omega(t/2-\tau)}d\Omega=\begin{cases}1 & \text{if}\quad 0<\tau<t\\ 0 & \text{if}\quad \tau<0\quad\text{or}\quad\tau>t\end{cases}
\tag{10}
$$

Taking the pump beam as quasimonochromatic ($\Delta\omega_p \ll \overline{\omega}_p$) and assuming that the spectral function of the down-converted light is slowly varying in ω_p, the spectral function can be approximate as $\phi(\omega_p,\omega_p-\omega,\omega)\simeq\phi(\overline{\omega}_p,\overline{\omega}_p-\omega,\omega)$. Equation (9) then reads

$$\begin{aligned}R_6(t)=&\langle I_p\rangle(|\eta|^2+|\eta_1|^2)+2\langle I_p\rangle\mathrm{Re}\{\eta_1\eta_2^*[\mathcal{TRR}T^*\gamma(\tau_p+\tau_5-\tau_0-\tau_3-\tau_4)\mu(\tau_4+\tau_0-\tau_5-\tau_1)e^{2i\theta}\\&+\mathcal{TRR}T^*\gamma(\tau_p+\tau_5-\tau_1-\tau_3-\tau_4)\mu(\tau_4+\tau_1-\tau_5-\tau_0)e^{2i\theta}+\mathcal{T}^2TT^*\gamma(\tau_p-\tau_1-\tau_2)\mu(\tau_2+\tau_1-\tau_3-\tau_0)e^{i\theta}\\&+\mathcal{R}^2TT^*\gamma(\tau_p-\tau_0-\tau_2)\mu(\tau_2+\tau_0-\tau_3-\tau_1)e^{i\theta}+\mathcal{TR}TR^*\gamma(\tau_p+\tau_4-\tau_0-\tau_2-\tau_5)\mu(\tau_5+\tau_0-\tau_4-\tau_1)\\&+\mathcal{TR}TR^*\gamma(\tau_p+\tau_4-\tau_1-\tau_2-\tau_5)\mu(\tau_5+\tau_1-\tau_4-\tau_0)+\mathcal{T}^2RR^*\gamma(\tau_p-\tau_0-\tau_3)\mu(\tau_0+\tau_3-\tau_1-\tau_2)e^{i\theta}\\&+\mathcal{R}^2RR^*\gamma(\tau_p-\tau_1-\tau_3)\mu(\tau_1+\tau_3-\tau_0-\tau_2)e^{i\theta}]\},\end{aligned}\tag{11}$$

where

$$\mu(\tau)=2\pi\delta\omega\sum_\omega|\phi(\overline{\omega}_p,\omega,\overline{\omega}_p-\omega)|^2e^{-i\omega\tau}\tag{12}$$

and

$$\gamma_p(\tau'-\tau)=\frac{\langle V^*(t+\tau)V(t+\tau')\rangle_{\text{clas.}}}{\langle I_p\rangle}\tag{13}$$

are the normalized self-correlation functions of the down-converted and the pump beams, respectively. Finally, suppose that variations $\delta\tau_\alpha$ due to changes in the optical path travel time τ_α are much smaller than the second-order coherence time of the downconverted light and the pump beam, it follows that

$$\mu(\tau+\delta\tau)\simeq\mu(\tau)e^{-i\overline{\omega}\delta\tau}\tag{14}$$

and

$$\gamma_p(\tau'+\delta\tau'-\tau-\delta\tau)\simeq\gamma_p(\tau'-\tau)e^{-i\overline{\omega}_p(\delta\tau'-\delta\tau)}.\tag{15}$$

These usual approximations simplify Eq. (11), giving

$$\begin{aligned}R_6(t)\simeq&\langle I_p\rangle(|\eta_2|^2+|\eta_1|^2)+2\langle I_p\rangle|\eta_1\eta_2\vartheta|\cos[\overline{\omega}(\delta\tau_0+\delta\tau_1\\&+\delta\tau_4-\delta\tau_5)+\overline{\omega}_p(\delta\tau_3-\delta\tau_p)+2\theta+\varphi]\\&+2\langle I_p\rangle|\eta_1\eta_2\vartheta'|\cos[\overline{\omega}(\delta\tau_0+\delta\tau_1+\delta\tau_5-\delta\tau_4)\\&+\overline{\omega}_p(\delta\tau_2-\delta\tau_p)+\varphi']\\&+2\langle I_p\rangle|\eta_1\eta_2\vartheta''|\cos[\overline{\omega}(\delta\tau_0+\delta\tau_1+\delta\tau_2+\delta\tau_3)\\&-\overline{\omega}_p\delta\tau_p+\theta+\varphi''],\end{aligned}\tag{16}$$

where the symbols ϑ, ϑ', and ϑ'' are given by

$$\begin{aligned}\vartheta=\mathcal{TRR}T^*[&\gamma(\tau_p+\tau_5-\tau_0-\tau_3-\tau_4)\mu(\tau_4+\tau_0-\tau_5-\tau_1)\\&+\gamma(\tau_p+\tau_5-\tau_1-\tau_3-\tau_4)\mu(\tau_4+\tau_1-\tau_5-\tau_0)],\end{aligned}\tag{17}$$

$$\begin{aligned}\vartheta'=\mathcal{TR}TR^*[&\gamma(\tau_p+\tau_4-\tau_0-\tau_2-\tau_5)\mu(\tau_5+\tau_0-\tau_4-\tau_1)\\&+\gamma(\tau_p+\tau_4-\tau_1-\tau_2-\tau_5)\mu(\tau_5+\tau_1-\tau_4-\tau_0)],\end{aligned}\tag{18}$$

$$\begin{aligned}\vartheta''=[&\mathcal{T}^2|T|^2\gamma(\tau_p-\tau_1-\tau_2)\mu(\tau_2+\tau_1-\tau_3-\tau_0)\\&+\mathcal{R}^2|T|^2\gamma(\tau_p-\tau_0-\tau_2)\mu(\tau_2+\tau_0-\tau_3-\tau_1)\\&+\mathcal{T}^2|R|^2\gamma(\tau_p-\tau_0-\tau_3)\mu(\tau_0+\tau_3-\tau_1-\tau_2)\\&+\mathcal{R}^2|R|^2\gamma(\tau_p-\tau_1-\tau_3)\mu(\tau_1+\tau_3-\tau_0-\tau_2)].\end{aligned}\tag{19}$$

In these equations, $\varphi=\arg(\eta_1\eta_2^*\vartheta)$ and so on.

V. CONCLUSIONS

The counting rate $R_6(t)$ is shown to be able to reveal interference patterns corresponding to wavelengths λ_p and $2\lambda_p$, either as separate or in simultaneous patterns. The quantum character of the considered processes establishes these distinct dependencies as signatures of the de Broglie wavelength for a composite and a single particle and cannot be obtained from a classical analysis.

The relevant terms to be considered in $R_6(t)$ are the ones connected with the interferences between the two down-converters, namely, terms proportional to $\eta_1\eta_2$—those with ϑ, ϑ', and ϑ''. The explicit expressions of all terms in $R_6(t)$ allow one to study specific variations of any chosen experimentally controllable parameter. A large number of possibilities then arise. As a guideline one can see, for example, that the term proportional to ϑ'' is due to the amplitude probability for photons leaving BS1 on distinct trajectories. This term vanishes whenever this amplitude probability is zero, that is to say, in cases where $|\mathcal{R}|=|\mathcal{T}|$ and the trajectory lengths 0 and 1 are made equal within the coherent length of the photon wave packets such that creation of two-photon states results and the photon pairs always come together in the same trajectory, either path 2 or 3.

Photon trajectory lengths from NL1 to NL2 have to be also approximately equal to produce a nonzero term and approximately equal to the laser trajectory length from NL1 to NL2. Any differences have to be within the laser coherence length on NL2. These dependencies have been experimentally observed and explained [11–13].

The term in ϑ (analogous reasoning for the ϑ' term) is dependent on the probability amplitude that photons generated in NL1 leave BS1 from port 3 as two-photon states. This term contributes whenever the laser coherence time is long enough to assure coherence on BS2 of the photon pairs emitted by both crystals and any imbalance of the interferometer arms (0,1) do not differ from imbalances of arms (4,5) in more than a coherence length of the down-converted photon packets [14]. This term shows dependence on the two-photon particle as a whole, either depending on changes in path lengths or on Pancharatnam's phase θ. One can also note that, say, if arm 3 is made very different from arm 2 [see Eq. (17)] the term in ϑ may be kept different from zero with a long coherence time laser. Consequently, signatures of photon pairs can be detected in single count experiments. This dependence on photon pairs in a single counting experiment is a novelty made possible by the use of the nonlinear converter NL2 as an effective beam splitter for photon pairs.

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