Chirped quasi-phase-matching with Gauss sums for production of biphotons

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Abstract
We study the theory of linearly chirped biphoton wave packets produced in two basic quasi-phase-matching configurations: chirped photonic-like crystals and chirped aperiodically poled crystals. The novelty is that these structures are considered as definite assembles of nonlinear layers that leads to a detailed description of spontaneous parametric down-conversion processes through the discrete Gauss sums. In this approach, the spectra of biphoton wave packets for each of the two presented quasi-phase-matching configurations display new important properties that were not obtained and explained in the standard phenomenological approach.

(Some figures may appear in colour only in the online journal)

1. Introduction

One of the challenges in the fields of quantum optics and quantum communications is the engineering of new sources of nonclassical light, such as biphotons with controllable spectral and temporal properties. The standard method for the generation of biphotons is spontaneous parametric down conversion (SPDC) [1–3], where photon pairs are generated under the action of a strong pump field interacting with a nonlinear crystal. The photon pairs are also generated in four-wave mixing [4, 5], for atomic systems in the bichromatic field [6, 7] and for cooled and trapped ions [8]. In particular, the preparation of biphoton wave packets with small correlation times between two photons, as well as with broad frequency spectra, is of considerable interest. Correlated photon pairs with broad spectra are mainly generated by SPDC using the method of quasi-phase-matching (QPM) that provides a feasible alternative to the conventional phase matching for many optical parametric process applications. Several methods have been suggested and experimentally implemented in this direction for generation of biphotons, including SPDC in periodically poled crystals or photonic crystals. The applications have been made for the synthesis of twin-photon states by manipulating the overall group delay mismatches between interacting waves in multilayered structures by using compensating dispersive effects [15–18], for the production of entangled three-photon states in cascaded parametric processes [19–21].

It has recently been experimentally demonstrated that biphotons with broad spectra and ultrashort correlation time at a high emission rate can be generated using chirped QPM nonlinear crystals involving nonlinear grating in the spatial coordinate along the direction of pump propagation with a nonuniform poling period [22–25]. In this case, phase matching is achieved along the crystal in such a way that the phase-matching conditions vary longitudinally by the use of QPM which leads to a broadband biphoton.

The theory of linearly chirped biphoton wave packets is usually based on a simple model allowing an analytical description of linearly chirped periodically poled crystals or optical fibres (see [22–25]). In this model, a phase-matching function of three-wave interaction \( \exp[i\Delta k(z)z] \) with a QPM grating with the linearly varying spatial frequency \( \Delta k(z) = \Delta k - \alpha z \) has been considered. The interaction leads to the biphoton spectral density in the form of theerfi(x) error function.
function. Such a phenomenological approach qualitatively explains the basic properties of chirped QPM, in this number: extremely wide spectral bandwidths of biphoton wave packets and the sharp temporal separation between the signal and the idler photons comprising a pair. However, this approach has been unable to describe some important details of chirped three-wave interactions as will be demonstrated in sections 3 and 4.

In this paper, we present a more detailed description of chirped structures for generation of two-photon light via SPDC from a cw pump in another way: we consider these structures as the definite assembly of nonlinear segments. This approach provides more flexibility for designing QPM gratings and allows us to engineer the controllable phase relation between the various spectral components. On the whole, it becomes possible to control the frequencies and bandwidths of the signal and idler photons by varying the number of layers equally with the local poling period along the length of the crystal. In this approach, the total amplitude of two-photon generation is calculated as the superposition of partial amplitudes corresponding to each layer with the local spatial chirp frequency. As frequencies of signal and idler photons are varied with different layers, a much broader range of photon frequencies is formed due to the superposition. Such detailed analysis of chirped QPM allows one to consider the interference and transition effects in SPDC stimulated by definite numbers of layers. In this approach, the QPM function of SPDC is expressed through the so-called Gauss sums instead of the continuous error function that appears in the phenomenological approach. It should be mentioned that during recent years an impressive number of experiments implementing Gauss sums has been seen

\[ S_N(\zeta) = \sum_{m=-M}^{M} W_m e^{2\pi i (m+\zeta)} \]  

(1)
in physical systems, particularly, for number factorization schemes. These systems range from nuclear magnetic resonance methods [26, 27] via cold atoms [28] and Bose–Einstein condensates [29], tailored ultrashort laser pulses [30, 31] to classical light in a multi-path Michelson interferometer [32] (for an introduction to this field, see [33, 34]). In addition to these results, here we present new quantum systems for implementation of Gauss sums to factor numbers.

We study the theory of linearly chirped biphoton wave packets produced for two basic QPM configurations: chirped layered photonic crystals and aperiodically poled layered crystals. In the phenomenological approach, both structures are described equally taking a linearly varying spatial frequency of the form \( \Delta k(\zeta) = \Delta k - az \). However, detailed calculations of two-photon spectral amplitudes in our approach (based on the consideration of nonlinear media as an assemble of nonlinear layers) for both the chirped structures lead to radically different basic equations (9) and (10) and, hence, to different results for two-photon spectra. These spectra have been separately analysed in sections 3 and 4.

The rest of the paper is organized as follows. In section 2, the brief description of SPDC in multilayered media is presented. Section 3 is devoted to the chirp in layered photonic-like crystals. The production of broadband biphotons in aperiodically poled layered crystals is described in section 4. We summarize our results in section 5.

2. SPDC in multilayered media: brief description

In this section, we briefly describe the generation of two-photon light via SPDC in one-dimensional \( \chi^{(2)} \) media consisting of layers with different coefficients of nonlinearity and refractive indices [15, 18, 19]. We consider collinear, type-II QPM configurations for generation of photons at the same frequency. Since the generated photons are collinear, their directions cannot be used to distinguish the two photons. Thus, we assume that photons in pairs have different polarizations, but omit the polarization indices below. The three-wave interaction Hamiltonian is expressed as the sum of interactions in each layer in terms of the electric fields for the \( m \)th layer,

\[ H(t) = \sum_{m} \int_{z_m}^{z_{m+1}} dz \chi^{(2)}(z)E_{out}^m(z,t)E_{in}E_{in}^* + \text{h.c.} \]  

(2)

Here, \( \chi^{(2)}(z) \) is the second-order susceptibility, \( \omega_{out}(z) \) represents the classical laser field at the frequency \( \omega_0 \), and \( E_{in}^m(z,t) \) and \( E_{in}^m(z,t) \) represent the positive-frequency parts of the fields of the subharmonics centred at the frequencies \( \omega = \frac{\omega_0}{2} + \Omega \) and \( \omega = \frac{\omega_0}{2} - \Omega \). The two-photon state can be written as

\[ |\psi\rangle = \int d\omega_1 d\omega_2 \Phi(\omega_1, \omega_2) a^+(\omega_1) a^+(\omega_2)|0\rangle, \]  

(3)

where \( a^+(\omega) \) and \( a^+(\omega) \) are the creation photon operators for modes with frequencies \( \omega \) and \( \omega \), \( \omega = \omega_t + \omega_i \), \( |0\rangle \) is a vacuum state of the signal and the idler fields, and \( \Phi(\omega_t, \omega_i) \) is the spectral amplitude of two-photon radiation. The two-photon amplitude for multilayered structures has been presented in many papers [15, 18], (see also [19–21]). It has been calculated as the sum of definite partial amplitudes in the general form given by the product of the pump envelope function \( E_0(\omega_0) \) and the phase-matching function \( F(\Delta k) \):

\[ \Phi(\omega_t, \omega_i) = -\frac{2\pi i}{h} E_0(\omega_t + \omega_i) F(\Delta k), \]  

(4)

\[ F = \sum_{m} \ln \chi_m F_m, \]

\[ F_m = e^{-i(\psi_m + \Delta k_m \phi)} \frac{\sin((\Delta k_m \phi)}{2}, \]

(5)

Here, \( \chi_m \) is the second-order susceptibility in the \( m \)th layer, \( \ln = \sum_{m+1} - \sum_m \) is the length and \( \Delta k_m = k^n_m - k^n_m - k^n_i \) is the phase mismatch vector for the \( m \)th layer, \( k^n_m(z, \omega) = \frac{\lambda}{2\pi} n_m(z, \omega), n_m(\omega) \) is the corresponding refractive index of the medium at the given frequency that describes the effects of dispersion on the properties of photon pairs. In this case, the probability of twin-photon generation is calculated as
has been demonstrated that one-dimensional photonic crystals for the chirp configuration. In the scheme shown in [40].

The generation of polarization-entangled photon pairs [35, 36]. It has been considered only the zero- and first-order terms: 

\[ n_j = n_0 - m\beta_j, \quad m \leq z < (m+1)l, \]

where \( \beta_0 \) and \( \beta_j (j = s, i) \) are chirp parameters for the refractive indices of the pump and subharmonic waves, correspondingly. In this case, the phase-mismatch vector for the \( m \)th layer has a linear chirp of the following form 

\[ \Delta k_m = \Delta k - \alpha (m-1)l, \]

where \( \Delta k = k_0(\omega_0) - k_s(\omega_j) - k_i(\omega_i) \) is the phase mismatch vector at the first layer, and \( \alpha = \frac{\beta_0 - \beta_j}{2} \) is the spatial chirp parameter.

Taking into account (7), we present the general expression (5) as

\[ F(\Delta k) = i e^{-\frac{\omega k^2}{2}} \sum_{m=1}^{N} F_m, \]

\[ F_m = e^{-i\left(m\Delta k - \alpha(m-1)l\right)} \times \sin\left(\frac{\Delta k - \alpha (m-1)l}{2}\right). \]

where \( \sin(x) = \frac{\sin(x)}{x} \).

It is easy to see that in the frame of this presentation, the phase-matching function is given by the Gauss sum, but is not determined by the error functions that appear in the phenomenological approach [13, 23]. Below, we calculate these sums considering the phase-mismatch function as 

\[ \Delta k = \Delta k_0 + \Omega D, \]

where \( \Delta k_0 = k_0 - k_s(\omega_j) - k_i(\omega_i) \) is temporal walk-off between signal and idler modes, assuming that \( \Delta k_0 \) satisfies the QPM condition. In order to illustrate the broadening of the spectrum and the other features of obtained phase-matching function, we have calculated the squared amplitude of the phase-matching function \( |F(\Omega)|^2 \) on the basis of formula (6) and the wave vector dispersion expansion as

\[ |F(\Omega)|^2 = I^2 \chi_0^2 \sum_{m=1}^{N} \sin^2 \left(\frac{\Omega + \Delta k_0 - \alpha (m-1)l}{2}\right) \times \sin\left(\frac{\Omega + \Delta k_0 - \alpha (m-1)l}{2}\right) \times \cos\left(p(D\Omega + \Delta k_0)l - \alpha p (m-1 - \frac{p}{2})^2\right). \]
If the number of domains is increasing, the chirp is displayed as a broadening of the biphoton spectra as it is depicted in figures 1(c) and (d) for \( N = 20 \) and \( N = 80 \) layers and the phase-matching conditions \( \Delta k_0 = 10\alpha l \) and \( \Delta k_0 = 40\alpha l \), respectively. One can see that the spectrum of the biphoton field is quite broad (from 900 to 1300 nm in wavelength) and has nearly rectangular shape already for \( N = 80 \) (see figure 1(d)). For further increasing of the number of domains, the interference picture becomes smooth and the results qualitatively coincide with the analogous results obtained on the basis of a simple model [22, 23]. These power spectra also qualitatively coincide with the experiment (see [22, figure 2(c)]).

It should be noted that the length of biphoton spectral shape does not depend on the number of layers for the fixed total length of nonlinear media but is only determined by the chirping parameter. We illustrate this statement in figures 1(e) and (f) by consideration of the other chirping parameter \( \alpha = 600 \text{ cm}^{-2} \) that is two times less than the previous one, for the case of \( N = 5 \) and \( N = 80 \) layers, correspondingly. As we see, increasing the chirping parameter leads to the increase in the intervals between the spectral peaks for the case of a small number of layers. For a case of a large number of layers, this increase leads to the increase in the range of the wavelength spectral broadening.

At the end of this section, we have selected some important details in the production of chirped biphoton wave packets that were not explained in the phenomenological approach: (a) it has been demonstrated that for the case of a small number of layers, the biphoton spectra consist of resolved spectral peaks corresponding to the number of domains. The frequencies of the spectral lines can be controlled by the chirp parameter and temporal walk-off between signal and idler modes. This process leads to the production of a few various spectrally nondegenerate photon pairs that might be applied in a new type of experiment in the area of photon entanglement and (b) there are critical ranges of the number of domains when the spectral lines are not resolved and the chirp leads to the broadening of the biphoton spectra. In the regime of spectral broadening, the width of biphoton spectra does not depend on the number of layers for a fixed total length of nonlinear media but is only determined by the chirping parameter.

4. Chirp in aperiodically poled crystals

Recently, aperiodic poling has been used for various applications. It has been provided not only for compensation of a natural phase mismatch but has also allowed one to tailor the properties of emitted photon pairs using nonlinear domains with variable lengths (chirped periodical poling). Domains of different lengths in an ordered structure allow an efficient nonlinear interaction in an ultra-wide spectral region extending typically over several hundreds of nanometres [9–14, 41, 42]. It has been shown that photon pairs generated in such structures can possess quantum temporal correlations at the timescale of femtoseconds. Other applications of ultra-wideband biphotons include nonclassical metrology [43] and large bandwidth quantum information processing [44, 45].

We consider aperiodically poled crystals as a multilayered structure consisting of \( N \) layers with a variation of lengths: the length of each layer is larger from the previous by a chirp parameter \( \zeta \). So the length of the \( n \)th layer is given by the following expression \( l_n = l_0 + (n - 1)\zeta \), where \( l_0 \) is the length of the first layer. We consider the quadratic nonlinearity \( \chi_n \), having a constant \( \chi_n = (-1)^{(n-1)}\chi_0 \) value within each \( n \)th layer. This structure can be divided into \( N/2 \) domains such that each of them will consist of two layers with reversed crystal axes. So the poling period \( \Lambda \), which is the length of the domain, races by \( 2\xi l \) for any next domain, which means that it depends on the coordinate. Thus, it can be easily obtained that the relation between chirp parameters \( \zeta \) and \( \alpha \) can be written as follows: \( \zeta = \frac{\alpha l}{\Delta k} \).

According to the formulae (5), we obtain two-photon spectral amplitude \( F(\Delta k) \) for the case of aperiodically poled crystal in the following form:

\[
F(\Delta k) = \frac{\chi_0}{\Delta k} \exp \left(-i\frac{\Delta kl_0}{2}\right) \times \sum_{m=1}^{M} (-1)^m \exp \left(-i\Delta k \left(ml_0 + \frac{(m - 1)^2\xi}{2}\right)\right) \times \sin \left(\frac{\Delta k(l_0 + (m - 1)\xi)}{2}\right),
\]

through the discrete Gauss sum, where \( \Delta k \) is the phase mismatch function. In this case, the probability of biphoton
generation as a function of the frequency $\Omega$ is proportional to

$$|F(\Omega)|^2 = \frac{\chi_0^2}{(D\Omega + \Delta k_0)^2} \times \left[ \left( \sum_{m=1}^{N} \sin^2 \left( \frac{(D\Omega + \Delta k_0)(l_0 + \zeta (m - 1))}{2} \right) \right) + 2 \sum_{m=1}^{N-1} \sum_{n=1}^{N-m} (-1)^n \times \sin \left( \frac{(D\Omega + \Delta k_0)(l_0 + \zeta (m + n - 1))}{2} \right) \times \sin \left( \frac{(D\Omega + \Delta k_0)(l_0 + \zeta (m - 1))}{2} \right) \times \cos \left( p(D\Omega + \Delta k_0)(l_0 + \zeta (2m + p - 2)) \right) \right].$$

(11)

Here, $\Delta k_0 = k_0 (\frac{\zeta}{m}) + k_0 (\frac{n}{N}) - k_0$ is the phase mismatch vector at the central idler and signal frequencies $\omega_1 = \omega_s = \frac{\omega_0}{2}$, where $\omega_0$ is the laser pump frequency. In this formula, we have considered also $\Delta k = \Delta k_0 + D\Omega$.

It is easy to realize that the biphoton spectra for two systems under consideration are essentially different in form. The main difference is that biphoton spectra for aperiodically poled structures do not consist of separated spectral peaks in the case of small layers contrary to the case of photonic crystal. If the chirp parameter $\zeta$ is equal to zero, the spectrum (11) is reduced to the spectral function of biphotons for the purely periodically poled configuration that includes segments of length $l_0$ with positive $\chi_0$ and negative $-\chi_0$ susceptibilities that alternate one with the other

$$|F(\Omega)| = l_0 \chi_0 \times \sin \left( \frac{N(D\Omega + \Delta k_0 l_0)}{2} \right) \times \sin \left( \frac{(D\Omega + \Delta k_0)(l_0 + \zeta)}{2} \right).$$

(12)

In the approximation $N \gg 1$, this probability reads $|F(\Omega)|^2 \sim l_0^2 \chi_0^2 \sin^2 \left( \frac{N(D\Omega + \Delta k_0)}{2} \right)$.

We illustrate the peculiarities of the biphoton spectra depending on $\Omega$ for the aperiodically poled structure as well as for the periodically poled structure in figure 2 for $N = 50$ and $\Delta k_0 = 0$. As we see, in the limit $\zeta = 0$ and $N \gg 1$, only one narrow pick for positive $\Omega$ centred at $\frac{\pi}{m_0}$ occurs. Increasing the chirp parameter broadens the frequency spectrum.

The typical biphoton spectra for aperiodic, chirped poling are shown in figure 3 depending on the signal-field (idler-field) wavelength ($\lambda_s \gg 2\lambda_0$) for a various number of layers and chirping parameters. For comparison of these spectra with analogous results obtained for the case of photonic crystals as well as with the experimental results (see, e.g., [24]), we use the parameters that are suitable for the LiTaO$_3$ aperiodically poled crystal, with $D = 0.3$, the whole length $L = 0.8$ cm and for the laser pump frequency $\lambda_0 = 0.458 \mu m$.

Biphoton spectral density for an aperiodically poled crystal with a number of layers $N = 50$, chirp parameter $\zeta = 1$ and $l_0 = 109.5$, which correspond to a smaller spatial chirp parameter $\alpha = 240$ cm$^{-2}$, is shown in figure 3(a). The analogous result for the case of the crystal with the same number of layers, but with the chirp parameter $\zeta = 2.82$ and $l_0 = 88.09$, which correspond to the spatial chirp parameter $\alpha = 240$ cm$^{-2}$, is shown in figure 3(b).
\[ \alpha = 1200 \text{cm}^{-2}, \] is shown in figure 3(b). From comparison of these two results, we can see that increasing the chirp parameter by two times leads to wide broadening of the signal- and idler-field spectra. Moreover, for each number of layers, there is an optimal value of the chirp parameter for which the spectrum is the broadest. In figures 3(b)–(d), we show the results of forming the biphoto spectra as the number of layers increases provided that the chirping parameter is fixed. For this goal, we consider the number of layers: \( N = 50, N = 100 \) and \( N = 160 \) assuming that the chirping parameters are chosen so that again \( \alpha = 1200 \text{cm}^{-2} \). As we can see from these figures, the spectra for these three cases have an approximately identical form. Wider broadening occurs for a higher number of layers: for the number of layers \( N = 160 \), the spectral width achieves 0.2 \( \mu \text{m} \). Thus, we conclude that for the case of aperiodic chirping the width of the spectrum is strongly dependent on the number of layers despite the refractive index frequency chirp considered in the previous section. We also observe that in the case of aperiodically poled crystals the spectral lines corresponding to definite segments are not resolved as it is demonstrated in the case of photonic crystals.

At the end of this section, we conclude that the biphoto spectra for an aperiodically poled configuration display an asymmetric form depending on the wavelength of the signal (idler) field (see figure 3) contrary to the case of the previous model and phenomenological approach (see figure 2 in [13]). Nevertheless, these calculated spectra have some familiarity with the spectrum obtained in the experiment [24]. It has also been shown that biphoto spectra for aperiodically poled structures in the case of a small number of layers do not consist of separated spectral peaks contrary to the case of photonic-like crystals (section 3). It has also been demonstrated that for the case of aperiodic chirping the width of the spectrum is strongly dependent on the number of layers contrary to both cases of the refractive index frequency chirp and the phenomenological approach. Indeed, the phenomenological approach described in most of the literature does not allow any connections between the shape and the width of the biphoto spectra with the number of layers.

5. Conclusion

In conclusion, we have investigated the production of biphoto wave packet in two chirped QPM structures presented as the ensemble of second-order nonlinear layers. In this approach, the total amplitude of two-photon generation has been calculated as the superposition of partial amplitudes corresponding to each layer with local spatial chirp frequency. In this way, the spectrum of spontaneously generated chirped photons has been calculated through Gauss sums that involve realistic parameters of the three-wave interaction in each layers. The detailed calculations have been performed for two QPM structures: (a) chirped nonlinear one-dimensional multi-layered photonic crystals and (b) aperiodically poled one-dimensional multi-layered crystals. Thus, in comparison with earlier theoretical investigations, the schemes considered in this paper correspond directly to experimentally realizable systems, whereas earlier works used modelling of the \( z \)-dependent mismatch by the continuous error function. As an advantage of the presented approach, we have demonstrated, in sections 3 and 4, some important properties of biphoto wave packets for each of the chirped structures that were not displayed and explained in the phenomenological approach.

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