Transversely Pumped Counterpropagating Optical Parametric Oscillation and Amplification

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We show that the second-order nonlinear optical material in a waveguide can be used to achieve transversely pumped counterpropagating optical parametric oscillation and amplification based on our proposed novel interaction configuration. By changing the incident angle of the pump beam, one can tune the output frequency through a large range. Although the vertical and horizontal cavities are not required to establish the oscillation, they can be used to reduce the threshold pump power by several orders of magnitude. For the optimum semiconductor structure, the threshold pump power for the oscillation can be as low as $\sim 100\ \mu$W.

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Backward parametric oscillation was proposed [1] without mirror feedback in the parallel propagation configuration. However, to the best of our knowledge, such a fundamental nonlinear process has not been observed due to the lack of appropriate materials for achieving phase matching. In conventional nonlinear materials such as KTP, conventional phase matching (CPM) was achieved by using the birefringent effect of the materials. In the birefringence-free materials, quasiphase matching (QPM) can be achieved by spatially modulating the second-order susceptibility [2]. Following Ref. [2], recently surface-emitting green light was obtained [3] based on second-harmonic generation (SHG) in the waveguide consisting of multilayers. The conversion efficiency of SHG was observed to increase by two orders of magnitude after a vertical cavity was included [4]. Most recently, SHG was observed [5] in asymmetric coupled quantum-well domain (ACQWD) structure following Ref. [6]. SHG in the mid-infrared domain was found to be feasible [7]. Correlators [8], spectrometers [8], and phase detectors [9] were demonstrated. Optical power limiting [10], optical phase-conjugation [10], self-phase modulation [11], and directional coupling [12] are feasible.

In this Letter, for the first time to the best of our knowledge, we show that the second-order optical nonlinear material in a waveguide can be used to achieve transversely pumped counterpropagating optical parametric oscillation (TPCOPO) and amplification (TPCOPA) in a novel interaction configuration. The proposed configuration does not require any cavity to establish the oscillation, and thus is fundamentally different from the collinear propagation configuration [13]. Based on the estimates made below, one can observe the mirrorless optical parametric oscillation with the pulsed lasers currently available. We believe, therefore, that the proposed TPCOPO will make a dramatic impact on fundamental nonlinear optics.

The proposed phase-matching condition is a combination of the CPM along the waveguide and the QPM normal to the waveguide, achieved by spatially modulating the second-order susceptibility [14]. The proposed transverse-pumping geometry is a versatile geometry for optical parametric oscillation in almost any second-order nonlinear medium. This geometry allows us to generate efficient degenerate parametric oscillation (i.e., the frequencies of the signal and idler are the same) for a large wavelength range of the pump wave (i.e., broadband frequency dividers). In addition, by changing the incident angle of the pump beam, we can tune the output frequency over a large range.

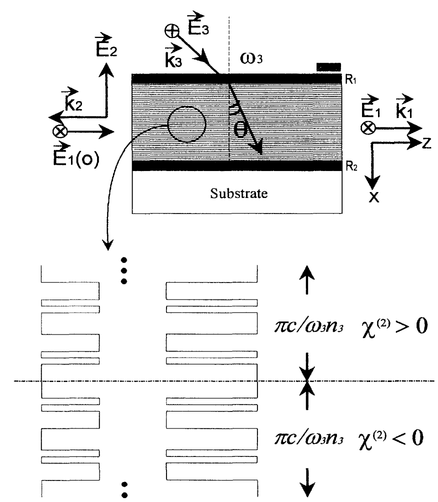

FIG. 1. The wave propagation configuration for TPCOPO and TPCOPA. Two dielectric mirrors with the reflectivities of R_1 and R_2 form a vertical cavity. For the oscillation, $E_1(0) = 0$. For the amplification, $E_1(0) \neq 0$. Inset: the band diagram of the ACQWD structure. The polarizations of the three waves are chosen based on $\chi^{(2)}_{yxy}$.

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Consider a second-order nonlinear optical material sandwiched between the top and bottom dielectric mirrors with the reflectivities R_1 and R_2, respectively, to form a vertical Fabry-Pérot cavity. The pump wave at the frequency ω_3 propagates onto the top mirror with an angle, as shown in Fig. 1. If the power of the pump wave outside the cavity is $P_3^{(o)}$, the power of one of the two traveling waves forming the pump quasistanding wave inside the cavity is $P_3^{(i)} = [(1 - R_1)/(1 - \sqrt{R_1 R_2})^2] P_3^{(o)}$, when the total thickness of the material is $m\lambda_3 \cos\theta / 2n_3$, where m is an integer, λ_3 is the pump wavelength in vacuum, n_3 is the refractive index at λ_3, and $\theta \ll 1$ is the small angle between the propagation direction of the pump wave inside the vertical cavity and the surface normal. Assuming that the nonlinear material is thin, we neglect the depletion of the pump wave. Two oscillating waves at the frequencies ω_1 and ω_2 are generated and counterpropagate along the waveguide formed by the mirror or extra confinement layers (see Fig. 1). Assuming that $\hbar\omega_3$ is smaller than the effective bandgap of the material, we can neglect the absorption of the pump and oscillating waves. The equations governing these two waves in the quasi-CW regime can be shown to be [1,10]

$$\nabla^2 E_{1,2} - \frac{n_{1,2}^2}{c^2}\frac{\partial^2 E_{1,2}}{\partial t^2} = \frac{1}{c^2}\frac{\partial^2}{\partial t^2}[\chi^{(2)} E_{2,1}(E_3 + E_3')], \quad (1)$$

where E_1 and E_2 are the electric fields of the signal and idler waves that propagate in the $+z$ and $-z$ directions, respectively, E_3 is the electric field of the pump quasistanding wave inside the cavity, E_3' is the electric field of the quasistanding wave at ω_3 as a result of the back conversion, $n_{1,2}$ are the refractive indices at ω_1 and ω_2 ($\omega_1 \geq \omega_2$), and $\chi^{(2)}$ is the second-order susceptibility of the nonlinear material. Here, we omit the subscripts of $\chi^{(2)}$ that correspond to the polarizations of three waves. Therefore, all the electric fields in Eq. (1) are scalars. $E_{1,2}$ can be expressed as the products of the amplitudes that vary slowly in space ($A_{1,2}$) and normalized waveguide modes [$\psi_{1,2}(x)$]:

$$E_{1,2} = A_{1,2}(z)\,\psi_{1,2}(x)\,e^{i(\pm\beta_{1,2}z - \omega_{1,2}t)} + \text{c.c.}, \quad (2)$$

where $\beta_{1,2} = (\omega_{1,2}/c)\,n_{1,2} = (2\pi/\lambda_{1,2})\,n_{1,2}$. E_3 can be expressed as the product of a constant amplitude (A_3) and a normalized factor describing the spatial dependence of the quasistanding wave along its propagation direction [$\psi_3(x)$]:

$$E_3 = A_3\psi_3(x)\,e^{i(\beta_3 \sin\theta z - \omega_3 t)} + \text{c.c.}, \quad (3)$$

where $\beta_3 = (\omega_3/c)\,n_3 = (2\pi/\lambda_3)\,n_3$. All of the ψ_i satisfy the conditions $\int_{-\infty}^{+\infty} \psi_i^2\, dx = 1$, for $i = 1$, 2, and 3. E_3' can be expressed as

$$E_3' = A_3'(z)\,\psi_3(x)\,e^{i(\beta_3 \sin\theta z - \omega_3 t)} + \text{c.c.}, \quad (4)$$

where A_3' is the amplitude that varies slowly in z. It can be determined in the same way as in Ref. [10]:

$$A_3' = ig_3 A_1 A_2, \; g_3 = 2\omega_3 d\chi_0^{(2)}/n_3 c(1 - R_1R_2)\, d_{\text{eff}}^{1/2}, \quad (5)$$

where d_{eff} is the effective waveguide thickness

$$d_{\text{eff}} = [\chi_0^{(2)}]^2 \Bigg/ \left[\int_{-\infty}^{+\infty} \chi^{(2)}(x)\psi_1(x)\psi_2(x)\psi_3(x)\, dx\right]^2. \quad (6)$$

We assume that $\chi^{(2)}$ is modulated along the x axis with a spatial period λ_3/n_3 and an amplitude $\chi_0^{(2)}$. As will be shown below, this spatial modulation results in the QPM. Substituting Eqs. (2)–(4) into Eq. (1), multiplying them by $\psi_{1,2}(x)$, and integrating over x, we obtain

$$\frac{dA_{1,2}}{dz} = \pm ig_{1,2}(A_3 + A_3')\,A_{2,1}^* e^{i(\beta_3 \sin\theta + \beta_2 - \beta_1)z}, \quad (7)$$

where $g_{1,2} = \omega_{1,2}\chi_0^{(2)}/2n_{1,2}c d_{\text{eff}}^{1/2}$ is the coupling coefficient among the oscillating waves and pump wave. Based on Eq. (7), we can readily show that $|A_1|^2 + (\omega_1 n_2/\omega_2 n_1)\,|A_2|^2 = |a|^2$, i.e., constant. This is the modified Manley-Rowe relation [1]. We assume that CPM condition along the z axis is satisfied:

$$\beta_3 \sin\theta + \beta_2 - \beta_1 = 0. \quad (8)$$

Consider the just-above-threshold oscillation [15], $g_3|A_1A_2| \ll |A_3|$. Equation (7) then reduces to

$$dA_{1,2}/dz = \pm ig_{1,2}A_3A_{2,1}^*. \quad (9)$$

The solutions of Eq. (9) can be written as

$$A_1 = a_1 \exp[i(g_1g_2)^{1/2}A_3z] + a_2 \exp[-i(g_1g_2)^{1/2}A_3z], \quad (10)$$

$$A_2 = (g_2/g_1)^{1/2}\{a_1^* \exp[-i(g_1g_2)^{1/2}A_3z] - a_2^* \exp[i(g_1g_2)^{1/2}A_3z]\}, \quad (11)$$

where a_1 and a_2 are complex constants. For the TPCOPO: $A_1(0) = 0$ and $A_2(L) = 0$ where L is the waveguide length, we obtain the threshold condition for the TPCOPO [16]:

$$(g_1g_2)^{1/2}(A_3)_{\text{th}}L = \pi/2. \quad (12)$$

The threshold pump power outside the vertical cavity can be determined to be

$$P_{\text{th}} = \frac{n_1n_2n_3\lambda_1\lambda_2}{4\eta_0[\chi_0^{(2)}]^2}\,\frac{(1 - \sqrt{R_1R_2})^2}{1 - R_1}\,\frac{W}{L}\,\frac{d_{\text{eff}}}{d}, \quad (13)$$

where W and d are the width and thickness of the waveguide, respectively, and η_0 is the vacuum impedance.

Equations (10) and (11) reduce to

$$A_1 = 2ia\sin(\pi z/2L), \quad A_2 = 2(g_2/g_1)^{1/2}a^* \cos(\pi z/2L), \quad (14)$$

where $a = a_1 = -a_2$. We can see that without mirror feedback these amplitudes start from zero at their respective input planes and reach maxima at the respective output planes (see Fig. 2). This is similar to the backward parametric oscillation [1], but fundamentally different from the conventional optical parametric oscillation, which cannot occur without the mirror feedback [13].

Consider the most fundamental interaction configuration for TPCOPO: *no cavity at all* (see Fig. 3). Assuming $\lambda_3 \approx 0.9\ \mu$m, $R_1 = R_2 \approx 0$, $d_{\text{eff}}/d \approx 10$, $W/L \approx 1/10$, and $\chi_0^{(2)} \approx 10^{-10}$ m/V, we estimate the threshold pump power for the TPCOPO to be $P_{\text{th}} \approx 8.2$ MW. Such a power can be achieved by available pulsed lasers.

To realize the TPCOPO and TPCOPA experimentally, we consider II-VI and III-V semiconductor structures. $\chi^{(2)}$ for the bulk materials originates from the center-inversion asymmetry. The nonzero elements of the $\chi^{(2)}$

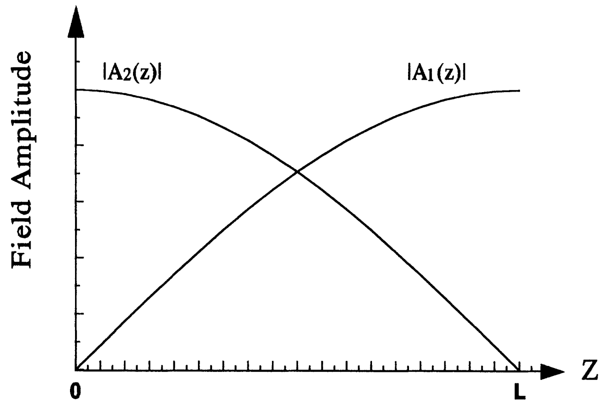

FIG. 2. The dependence of the electric field amplitudes of the parametric oscillation on the propagation distance for the pump power, which is slightly larger than the threshold pump power.

tensor are $\chi^{(2)}_{xyz}$, $\chi^{(2)}_{yzx}$, and $\chi^{(2)}_{zxy}$. On the other hand, in the asymmetric QW structures, the additional center-inversion symmetry is removed because of the asymmetric structure along the x axis. Additional nonzero elements are $\chi^{(2)}_{yxy}$ and $\chi^{(2)}_{xyy}$ [5,6].

Consider the TPCOPO pumped by a single short laser pulse. The vertical dielectric mirrors consist of the quarter-wave stacks. For the $GaAs/Al_{0.8}Ga_{0.2}As$ multilayer structure, $\chi^{(2)}$ has a spatially oscillating component with the period being λ_3/n_3. This is similar to the situation for SHG [3]. For $\lambda_3 \approx 0.9$ μm, $R_1 = R_2 \approx 99.9\%$ [17], $d_{\rm eff}/d \approx 10$, $W/L \approx 1/10$, $\chi_0^{(2)} \approx 10^{-10}$ m/V, $n_{1,2} \approx 3.3$, and $n_3 \approx 3.5$, we obtain $P_{\rm th} \approx 8.2$ kW. For the ZnSe/ZnS multilayer structure: $\lambda_3 \approx 0.49$ μm, $n_{1,2} \approx 2.5$, and $n_3 \approx 2.7$, we obtain $P_{\rm th} \approx 1.1$ kW. For the GaAs/AlAs ACQWD structure, $\lambda_3 \approx 10$ μm, $\chi_0^{(2)} \approx 30 \times 10^{-9}$ m/V [7], $n_{1,2} \approx 3.2$, and $n_3 \approx 3.3$, we obtain $P_{\rm th} \approx 10$ W. To illustrate the concept of the QPM, we have shown this structure in the inset of Fig. 1. The adjacent domains mirror each other. As a result, $\chi^{(2)}_{yxy}$ changes sign from one domain to the adjacent one. For the ZnSe/ZnS ACQWD structure: $n_{1,2} \approx 2.4$, $n_3 \approx 2.6$, $\chi_0^{(2)} \approx 6 \times 10^{-9}$ m/V, and $\lambda_3 \approx 5$ μm, we obtain $P_{\rm th} \approx 28$ W. All of these powers can be readily achieved with available pulsed lasers.

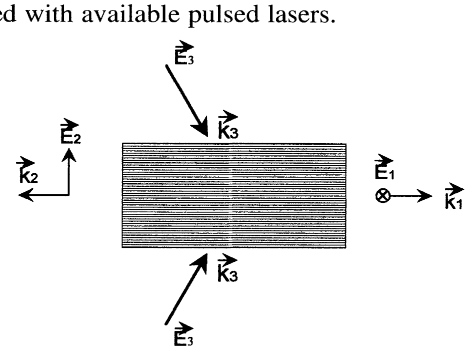

FIG. 3. The wave propagation configuration for the TPCOPO without any cavity.

For $W \approx 1$ mm, the corresponding threshold pump intensities inside the vertical cavity for the four structures above are $I_{\rm th}^{(i)} \approx 82, 11, 0.1$, and 0.28 MW/cm^2. The maximum increase of the temperature of the material is $(\Delta T)_{\max} \approx 2 I_{\rm th}^{(i)} \tau_p \alpha_3 / c_p$, where c_p is the heat capacity of the material, τ_p is the duration of the laser pulse, and α_3 is the below-band-edge absorption coefficient at ω_3. For $c_p \approx 1.8$ J/cm^3 K, $\tau_p \approx 1$ ns, $\alpha_3 \approx 10$ cm^{-1}, and $I_{\rm th}^{(i)} \approx 82$ MW/cm^2, we obtain $(\Delta T)_{\max} \approx 0.91$ K. Therefore, the effect of the laser heating for all the materials considered above is negligible.

Consider a horizontal cavity along the z axis, formed by two dielectric mirrors on the both ends $z = 0$ and $z = L$ with reflectivities $R_3, R_4 \sim 1$, and the degenerate TPCOPO, for which $\omega_1 = \omega_2 = \omega_3/2$. Assuming

$$A_1 = 2ia \sin(gA_3L + \pi/4 - \gamma),$$
$$A_2 = 2a^* \cos(gA_3L + \pi/4 - \gamma), \qquad (15)$$

where $g = g_1 = g_2$ and $\gamma \ll 1$, and using the boundary conditions $A_1(0) = \sqrt{R_3}A_2(0)$, $A_2(L) = \sqrt{R_4}A_1(L)$, we obtain the threshold condition for the oscillation [16]:

$$g(A_3)_{\rm th}L = (1 - \sqrt{R_3R_4})/2. \qquad (16)$$

If $R_3 \approx R_4 \approx 99\%$, the threshold pump powers for the oscillation in the four structures discussed above are determined to be 83, 11, 0.1, and 0.28 mW, respectively. If $W \approx 10$ μm [18], and for $I_{\rm th}^{(i)} \approx 8.3$ MW/cm^2, we obtain $(\Delta T)_{\max} \approx 0.092$ K; i.e., the effect of the laser heating of the materials is negligible.

One can tune the frequencies of the signal and idler by changing the incident angle of the pump wave (θ). Based on Eq. (8) and energy conservation ($\omega_3 = \omega_1 + \omega_2$), the maximum and minimum wavelengths are

$$\lambda_{\min} = [(n_1 + n_2)/(n_2 + 1)]\lambda_3,$$
$$\lambda_{\max} = [(n_1 + n_2)/(n_2 - 1)]\lambda_3. \qquad (17)$$

For $GaAs/Al_{0.8}Ga_{0.2}As$ we find $\lambda_{\min} \approx 1.4$ μm and $\lambda_{\max} \approx 2.6$ μm. For ZnSe/ZnS, with $\lambda_3 \approx 0.49$ μm, we obtain $\lambda_{\min} \approx 0.7$ μm and $\lambda_{\max} \approx 1.6$ μm, while for $\lambda_3 \approx 5$ μm, we get $\lambda_{\min} \approx 7.1$ μm and $\lambda_{\max} \approx 17$ μm. Finally, for GaAs/AlAs, we get $\lambda_{\min} \approx 15$ μm and $\lambda_{\max} \approx 29$ μm. If one uses a 2 μm laser beam to pump the TPCOPO based on GaAs/AlGaAs or InGaAs/GaAs, one can generate tunable (3.5–5.5 μm) and high-power (40 W) pulsed output.

One can achieve a large wavelength range of the pump wave for the efficient TPCOPO. It can be shown based on Eq. (6) that

$$d_{\rm eff}(\lambda_3) \approx d_{\rm eff}(\lambda_0)[\Delta\beta_3 d/\sin(\Delta\beta_3 d)]^2, \qquad (18)$$

where $\Delta\beta_3 = 2\pi(\lambda_3^{-1} - \lambda_0^{-1})n_3$, and λ_0/n_3 is the spatial period of the modulation of $\chi^{(2)}$. Thus, if $\lambda_3 = \lambda_0$, the QPM condition is satisfied. For $n_3 \approx 3.5$, $d \approx 1$ μm, and $\lambda_0 \approx 0.9$ μm, the usable bandwidth of the pump wavelength is $\approx$1300 Å [15].

Consider TPCOPA. Assuming that there is only one input, $A_1(0) \neq 0$, see the inset of Fig. 4(a), and based on Eqs. (10) and (11), we obtain

$$|A_1(L)| = |A_1(0)|\,|\sec[(g_1 g_2)^{1/2} A_3 L]|\,, \quad (19)$$
$$|A_2(0)| = |A_1(0)|\,(g_2/g_1)^{1/2} |\tan[(g_1 g_2)^{1/2} A_3 L]|\,, \quad (20)$$

i.e., the input signal is amplified if $(g_1 g_2)^{1/2} A_3 L > 0$ [see Fig. 4(a)]. This indicates that there is no threshold for amplifying the transmitted beam. One can see from Eq. (20) that $|A_2(0)| > |A_1(0)|$ if $(g_1 g_2)^{1/2} A_3 L > \pi/4$. In general, if $\omega_1 \neq \omega_2$, since there is no input at $z = L$ for A_2, the generation of $A_2(0)$ from the input $A_1(0)$ is a frequency conversion process. For the degenerate TPCOPA ($\omega_1 = \omega_2$, $g_1 = g_2 = g$), however, $A_1(0)$ can be regarded as an input for A_2. In this regard, there is a threshold for amplifying the reflected beam [see Fig. 4(a)]. The threshold pump is 2 kW, 0.27 kW, 2.5 W, and 6.9 W, respectively, for the four structures considered above. We can eliminate the threshold if the end $z = L$ is coated with the dielectric layers whose reflectivity (R_4) is close to 100% [see the inset of Fig. 4(b)]. Based on $A_1(0) = a_{\text{in}}$, $A_1(L) = A_2(L)$, and Eqs. (10) and (11), we can obtain

$$a_{\text{out}} = A_2(0) = a_{\text{in}}[\sec(2gA_3L) + i\tan(2gA_3L)]\,. \quad (21)$$

There is no threshold for amplifying the reflected beam [see Fig. 4(b)].

In conclusion, we have proposed to use the second-order nonlinear optical material in a waveguide to achieve TPCOPO and TPCOPA based on a novel interaction configuration. If one changes the incident angle of the pump wave, one can tune the output frequency of the TPCOPO in a broad range. Although the vertical and horizontal cavities are not required to establish the oscillation, they can be used to reduce the threshold pump power by several orders of magnitude. For the optimum semiconductor structure, the threshold pump power for the oscillation can be as low as $\sim 100\ \mu$W. TPCOPO can be used to frequency-divide laser radiation in a large frequency range. The TPCOPO and TPCOPA can be also implemented in the conventional second-order nonlinear materials such as KTP, $LiNbO_3$, or nonlinear optical polymer materials.

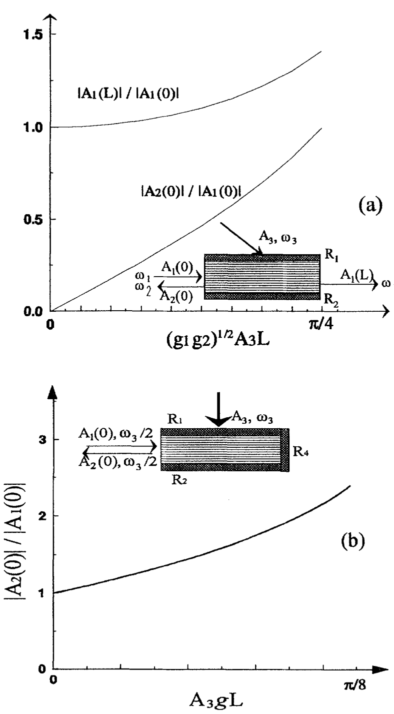

FIG. 4. The ratio between the output electric field amplitude for the transmitted or reflected wave and the input electric field amplitude vs the dimensionless pumping amplitude $[(g_1 g_2)^{1/2} A_3 L]$ for (a) no horizontal mirrors and (b) one horizontal mirror at the right end. Insets: the wave propagation configurations.

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