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CONTROL OF LASER RADIATION PARAMETERS

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Nonlinear phase shift and frequency jumps in second-harmonic generation in a dual-cavity laser

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***Abstract*. Frequency characteristics and the stability of lasing regimes are considered for a solid-state laser with intracavity second-harmonic generation and feedback at the frequency of the second harmonic. Three stationary states differing from each other by a nonlinear phase shift related to second-harmonic generation are investigated. It is shown that jumps in the frequency of laser radiation may be induced by transitions between the considered stationary states as parameters of a dual-cavity laser are smoothly tuned.**

Keywords: *intracavity second-harmonic generation, dual cavity, nonlinear frequency shift, solid-state laser.*

1. Introduction

Methods of intracavity frequency doubling are widely used to improve the efficiency of nonlinear optical processes [1 – 21]. Previous studies on intracavity second-harmonic generation (SHG) can be divided into two groups. The first group includes research on frequency doubling [1 – 13] and self-doubling [14 – 17] inside a laser cavity. The second group includes studies devoted to SHG with laser radiation injected into an extracavity resonator containing a nonlinear crystal [18 – 21]. The influence of intracavity SHG on the dynamics of lasing was investigated as a part of the first group of these studies. Such investigations were mainly performed for systems where the second harmonic leaves the cavity after a single pass (no feedback is introduced at the frequency of the second harmonic 2ω). The dynamics of SHG inside a laser cavity in the absence of feedback at the frequency 2ω was thoroughly analysed in [2 – 9]. In particular, periodic regimes of antiphase self-modulation and dynamic chaos in multimode lasers have been examined and the ways to stabilise stationary lasing regimes have been considered.

Second-harmonic generation inside a laser cavity in the presence of feedback at the frequency 2ω was theoretically investigated in [10 – 13]. In particular, the stability of stationary lasing in a single-mode laser was analysed [10, 11, 13], parameters of squeezed light were investigated [11, 12], and the influence of the cavity Q-factor at the frequency 2ω on the SHG efficiency was studied [13]. This paper is devoted to a theoretical analysis of the nonlinear phase shift, frequency characteristics, and the stability of lasing regimes of a solid-state laser with intracavity SHG in the presence of feedback in the second harmonic.

The propagation of waves through a nonlinear crystal may be accompanied by the appearance of additional nonlinear phase shifts $\varphi_{1,2}^{\mathrm{NL}}$ at the frequencies ω and 2ω related to the conversion of radiation frequency (nonlinear phase shifts). In the phase-matched regime, the nonlinear phase shift at the fundamental frequency is written as

$$\varphi_1^{\mathrm{NL}} = \int_0^l \chi\sqrt{I_2}\cos\psi \mathrm{d}x, \tag{1}$$

where χ is the nonlinearity coefficient; l is the length of the nonlinear element; I_2 is the second-harmonic intensity; $\psi = 2\varphi_1 - \varphi_2$; $\varphi_{1,2}$ are the phases of the fundamental wave and the second harmonic in the nonlinear crystal. Nonlinear phase shifts $\varphi_{1,2}^{\mathrm{NL}}$ vanish in phase-matched intracavity SHG in the absence of feedback at the frequency 2ω, since $\cos\psi = 0$ in this case. As shown below, such phase shifts may arise in the presence of a feedback.

2. The system of equations and stationary lasing regimes

Consider single-mode lasing in a solid-state laser with intracavity SHG. We assume that the laser cavity includes mirrors with high reflection coefficients at the fundamental frequency and the frequency of the second harmonic. Such a system will be called a dual cavity. The system of rate equations governing the dynamics of lasing in such a system can be written as

$$\frac{\mathrm{d}a_1}{\mathrm{d}t} = \frac{a_1}{2T_{\mathrm{c}}}\left[k_1(N-1) - \sqrt{\varepsilon}a_2\sin\psi\right], \tag{2}$$

$$\frac{\mathrm{d}a_2}{\mathrm{d}t} = -\frac{k_2}{2T_{\mathrm{c}}}a_2 + \frac{\sqrt{\varepsilon}}{2T_{\mathrm{c}}}a_1^2\sin\psi, \tag{3}$$

$$\frac{\mathrm{d}\psi}{\mathrm{d}t} = \frac{\sqrt{\varepsilon}}{2T_{\mathrm{c}}}\left(\frac{a_1^2}{a_2} - 2a_2\right)\cos\psi + \omega_{2\mathrm{c}} - 2\omega_{1\mathrm{c}}, \tag{4}$$

$$\frac{\mathrm{d}N}{\mathrm{d}t} = \frac{1}{T_1}\left[(1+\eta) - N(1+a_1^2)\right]. \tag{5}$$

Here, $a_{1,2} = (I_{1,2}/I_{\mathrm{s}})^{1/2}$ are the normalised field amplitudes inside the cavity at the fundamental frequency and the frequency of the second harmonic; $\varphi_{1,2}$ and $I_{1,2}$ are the

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phases and intensities of these fields; I_s is the saturation intensity of the active medium; $k_{1,2}$ is the magnitude of linear losses in the dual cavity; T_c is the cavity single-pass time; $\varepsilon = (\chi l)^2 I_s$ is the nonlinearity parameter; T_1 is the relaxation time of population inversion; ω_{2c} is the cavity eigenfrequency for the second harmonic; $2\omega_{1c}$ is twice the cavity frequency for fundamental radiation; N is the ratio of the population inversion to its threshold value; $1+\eta$ is the ratio of the pump power to its threshold value.

The system of Eqns (2)–(5) is derived in the Appendix. When deriving these equations, we made several assumptions determining the applicability range of this system. We consider SHG for type-I phase matching (ooe interaction). We assume that phase matching is achieved in the nonlinear crystal [see inequality (A5) in the Appendix], and the relevant phase-matching condition determines the range of admissible wave-vector mismatches $\Delta\boldsymbol{\kappa} = \boldsymbol{\kappa}_2 - 2\boldsymbol{\kappa}_1$ for the waves involved in the interaction. Deriving these equations, we also assumed that the total magnitudes of losses per cavity single pass are low for both wavelengths.

We assume also that the change in the phase difference $\Delta\psi$ between the interacting waves is small within a cavity single pass ($|\Delta\psi| \ll \pi$). This assumption imposes restrictions on the tolerable nonlinear phase shifts $\varphi_{1,2}^{\rm NL}$ ($|\varphi_{1,2}^{\rm NL}| \ll \pi$). The phase-matching condition [inequality (A5) in the Appendix] limits the admissible range of mismatches of cavity eigenfrequencies $\omega_{2c} - 2\omega_{1c}$. The mismatch of eigenfrequencies in a dual cavity is due to the dispersion of intracavity optical elements and the inequality of phase shifts $\delta_{1,2}$ arising in the reflection of light waves from cavity mirrors [see Eqns (A6) of the Appendix].

The relative detuning of the fundamental frequency from the centre of the gain line for a single-mode solid-state laser is small. Therefore, this detuning is neglected in Eqns (2)–(5). In a particular case of an inertialess active medium (a class-A laser), when the population inversion N can be adiabatically excluded [$T_1 \ll T_c$, $N = (1+\eta)/(1+a_1^2)$], Eqns (2)–(5) are reduced to the equations employed in [11, 12].

We start with a particular case when the cavity eigenfrequency for the second harmonic ω_{2c} is equal to twice the cavity frequency for fundamental radiation, $\omega_{2c} = 2\omega_{1c}$. In such a situation, we can find three stationary solutions to the system of Eqns (2)–(5). Solution 1 corresponds to $\cos\psi = 0$. The field amplitudes inside the cavity in this case are given by

$$a_1^2 = \frac{-B_0 + \left(B_0^2 + 4A_0 k_1 \eta\right)^{1/2}}{2A_0}, \quad a_2^2 = \frac{\varepsilon a_1^4}{k_2^2}, \tag{6}$$

where $A_0 = \varepsilon/k_2$ and $B_0 = k_1 + A_0$.

Examination of small perturbations of the stationary solution (6) and analysis of the relevant characteristic equation show that this solution is stable when the inequality

$$\varepsilon\left(\eta - \frac{k_2}{2k_1}\right) \leqslant \left(\frac{k_2}{4}\right)^2 \left(2 + \frac{k_2}{k_1}\right). \tag{7}$$

is satisfied.

The lasing frequency corresponding to solution 1 coincides with the frequency of the cavity mode, since $\varphi_1^{\rm NL} = 0$ in this case and the nonlinear frequency shift proportional to $\varphi_1^{\rm NL}$ vanishes.

When the stability of solution 1 is violated, i.e., the inequality

$$\varepsilon\left(\eta - \frac{k_2}{2k_1}\right) > \left(\frac{k_2}{4}\right)^2 \left(2 + \frac{k_2}{k_1}\right), \tag{8}$$

is met, two other stable solutions exist. For these solutions, $\cos\psi \neq 0$, and, consequently, nonlinear phase shifts $\varphi_{1,2}^{\rm NL}$ and frequency shifts arise. For the second (a) and third (b) solutions, the field amplitudes are given by

$$a_2^2 = \left(\eta - \frac{k_2}{2k_1}\right) \Big/ \left(2 + \frac{k_2}{k_1}\right), \quad a_1^2 = 2a_2^2. \tag{9}$$

The corresponding lasing frequencies are written as

$$\omega_{\rm a,b} = \omega_{1c} \pm \omega_{\rm NL}, \tag{10}$$

where

$$\omega_{\rm NL} = \frac{1}{T_c}\left[\frac{\varepsilon(\eta - k_2/2k_1)}{2 + k_2/k_1} - \left(\frac{k_2}{4}\right)^2\right]^{1/2} \tag{11}$$

is the nonlinear shift of the lasing frequency. The existence of two solutions defined by Eqns (9)–(11) is due to the fact that Eqns (2)–(4) for $\sin\psi$ have two roots equal in their absolute values, but opposite in sign.

Fig. 1 presents the dependence of the frequency shift $\Delta f = (\omega_{\rm a,b} - \omega_{1c})/2\pi$ on the excess of the pump power over the threshold power. This dependence was calculated for a system with the following parameters: $\varepsilon = 5\times10^{-5}$, $T_c = 0.2$ ns, $k_1 = 0.01$, and $k_2 = 0.01$. These characteristics correspond (except for the value of k_2) to parameters of a YAG : Nd laser with intracavity SHG in a KTP crystal with a length $l = 5$ mm employed in experiments [3–5].

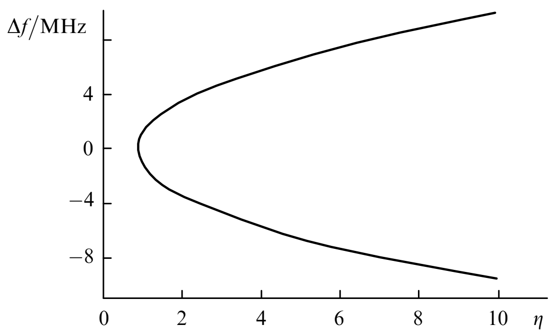

Figure 1. Dependence of the shift of the lasing frequency $\Delta f = (\omega_{\rm a,b} - \omega_{1c})/2\pi$ on the excess of the pump power over the threshold power η for $\varepsilon = 5\times10^{-5}$, $T = 0.2$ ns, $k_1 = 0.01$, and $k_2 = 0.01$.

Nonlinear phase and frequency shifts arise also for second-harmonic radiation. Using Eqns (2)–(5), we can demonstrate that the frequency of the second harmonic is determined by the formula

$$\omega_2 = 2\omega_{1c} + \frac{\sqrt{\varepsilon}}{T_c}\cos\psi.$$

One can see from this formula, the frequency of the second harmonic is equal to twice the lasing frequency:

$$\omega_{\rm L} = \omega_{1c} + \frac{\sqrt{\varepsilon}}{2T_c}\cos\psi.$$

Now, let us consider the influence of mismatch of cavity frequencies $\omega_{2c} - 2\omega_{1c}$ on the frequency shift and the stability of lasing.

3. Jumps of the lasing frequency induced by varying the mismatch of cavity frequencies $\omega_{2c} - 2\omega_{1c}$

Amplitude and frequency characteristics of radiation produced by a dual-cavity laser with intracavity SHG are highly sensitive to the mismatch of cavity eigenfrequencies $\Delta = \omega_{2c} - 2\omega_{1c}$. Applying Eqns (2)–(5), we can easily derive formulas governing stationary lasing in such a system with $\Delta \neq 0$:

$$a_1^2 = \frac{-B + \left(B^2 + 4Ak_1\eta\right)^{1/2}}{2A}, \quad a_2 = \frac{\sqrt{\varepsilon}}{k_2} a_1^2 \sin\psi, \qquad (12)$$

$$\Delta \equiv \omega_{2c} - 2\omega_{1c} = \frac{\sqrt{\varepsilon}}{2T_c}\left(2a_2 - \frac{a_1^2}{a_2}\right)\cos\psi, \qquad (13)$$

where $A = (\varepsilon/k_2)\sin^2\psi$ and $B = k_1 + A$.

These formulas directly describe the parametric dependence of a_1^2 and a_2^2 on the mismatch Δ. Defining the parameter ψ, we can use Eqns (12) and (13) to calculate a_1^2, a_2^2, and Δ. The normalised second-harmonic intensity a_2^2 as a function of Δ is shown in Fig. 2.

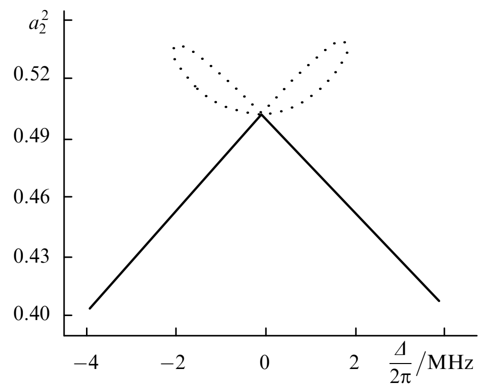

Figure 2. The normalised second-harmonic intensity a_2^2 as a function of the mismatch of dual-cavity eigenfrequencies $\Delta/2\pi$ for $\eta = 2$ and $k_1 = k_2 = 5 \times 10^{-3}$. The other parameters are the same as in Fig. 1 (the unstable solution is shown by the dotted line).

The stability of stationary lasing in the case of $\Delta \neq 0$ was investigated with the use of the Rauss–Gurwitz criterion. The dependence $a_2^2(\Delta)$ within the stability range is shown by a solid curve, while the dotted line shows this dependence in the area of instability. One can see from Fig. 2 that the maximum intensity of second-harmonic emission is achieved with $\Delta = 0$. The intensity a_2^2 monotonically decreases with the growth in $|\Delta|$.

Two solutions differing from each other by the sign of the nonlinear frequency shift may exist within some range of mismatches of cavity frequencies $\Delta = \omega_{2c} - 2\omega_{1c}$. This result is illustrated by Fig. 3, which presents the lasing frequency shift

$$\omega_L - \omega_{1c} = \frac{\sqrt{\varepsilon}}{2T_c}\cos\psi$$

as a function of the mismatch Δ. For $\Delta \neq 0$, the dependence of $\omega_L - \omega_{1c}$ on Δ in the stability range is shown by a solid curve, while the dotted line represents unstable states. One can see from Fig. 3 that the sign reversal of the frequency mismatch switches lasing regimes for parameters characteristic of solid-state lasers. The lasing regimes in the considered system differ from each other by the sign of the nonlinear frequency shift $\omega_L - \omega_{1c}$. As a result, when the mismatch Δ is smoothly tuned, the lasing frequency may change in a jumpwise manner. This jump in the lasing frequency is equal to $2\omega_{NL}$, where ω_{NL} is given by Eqn (11).

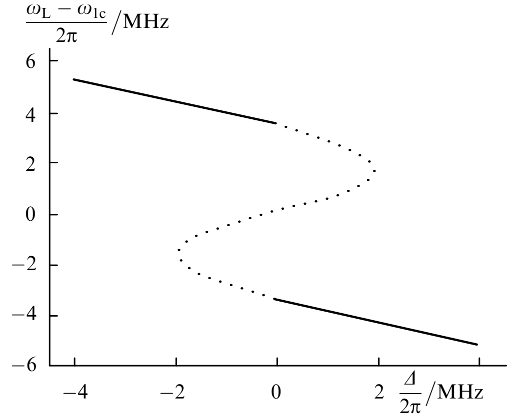

Figure 3. The shift of the lasing frequency $(\omega_L - \omega_{1c})/2\pi$ as a function of the mismatch of dual-cavity eigenfrequencies $\Delta/2\pi$. The laser parameters are the same as in Fig. 2 (the unstable solution is shown by the dotted line).

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Appendix

To derive Eqns (2)–(5), we shall consider a distributed model of a unidirectional ring laser. In the plane-wave approximation, the intracavity optical fields E_m at the fundamental frequency ($m = 1$) and at the frequency of the second harmonic ($m = 2$) can be written in the following way:

$$E_m = \mathrm{Re}\{e_m(x,t)\exp[\mathrm{i}(m\omega t - \boldsymbol{\kappa}_m x)]\}, \qquad (A1)$$

where, $m\omega$ and $\boldsymbol{\kappa}_m$ are the optical frequencies and the wave vectors of the fields.

Slowly varying complex amplitudes $e_{1,2}(x,t)$ and the density of population inversion $\Delta(x,t)$ are governed by the following system of equations (see, e.g., [22]):

$$\frac{\partial e_1}{\partial x} + \frac{n_1}{c}\frac{\partial e_1}{\partial t} = -\mathrm{i}\sigma_1 e_1^* e_2 + \frac{\sigma_a}{2}\Delta e_1 + \left[\mathrm{i}(\omega - \omega_{1c})\frac{n_1}{c} - \frac{\alpha_1}{2}\right]e_1, \quad (A2)$$

$$\frac{\partial e_2}{\partial x} + \frac{n_2}{c}\frac{\partial e_2}{\partial t} = -\mathrm{i}\sigma_2 e_1^2 + \left[\mathrm{i}(2\omega - \omega_{2c})\frac{n_2}{c} - \frac{\alpha_2}{2}\right]e_2, \qquad (A3)$$

$$\frac{\partial \Delta}{\partial t} = \frac{\Delta_0 - \Delta(1 + |e_1|^2/I_s)}{T_1}, \qquad (A4)$$

where $n_{1,2}(x)$ and $\alpha_{1,2}(x)$ are the refractive indices and the distributed absorption coefficients at the frequencies ω and 2ω, respectively; $\Delta_0(x,t)$ is the density distribution of population inversion in the absence of saturation; and I_s is the saturation intensity of the active medium. Equations (A2) and (A3) are written with an assumption that phase-matching conditions are satisfied in the nonlinear crystal [1].

With a nonzero wave-vector mismatch $\Delta\boldsymbol{\kappa} = \boldsymbol{\kappa}_2 - 2\boldsymbol{\kappa}_1$, the phase-matching condition is met when the mismatch $\Delta\boldsymbol{\kappa}$ satisfies the inequality

$$|\Delta\boldsymbol{\kappa}|l \ll 1. \tag{A5}$$

We assume that a nonlinear crystal with a length l is located within an area $x_0 < x < x_0 + l$, and the active crystal with a length l_a is located within the area $x_a < x < x_a + l_a$. In a particular case when $x_0 = x_a$ and $l = l_a$, the considered system of equations describes a laser with a nonlinear active medium. The ends of crystals (or the ends of a nonlinear active element) are assumed to be totally antireflective and ignore the reflection of light waves from these ends. As light passes from one medium to another, wave amplitudes change. If the refractive indices of the active and nonlinear crystals considerably differ from each other, then the amplitudes of the waves in these crystals also considerably differ from one another. To exclude such effects, we assume that the active element and the nonlinear crystal have close refractive indices.

The boundary conditions for intracavity optical fields $E_{1,2}$ on the output coupler are written as

$$\tilde{R}_{1,2}E_{1,2}(L,t) = E_{1,2}(0,t),$$

where $\tilde{R}_{1,2} = R_{1,2}\exp(\mathrm{i}\delta_{1,2})$ are the complex reflection coefficients; $R_{1,2}$ are the moduli of these coefficients; $\delta_{1,2}$ are the phase shifts introduced by reflection; L is the cavity perimeter. We assume for simplicity that only one of the cavity mirrors has reflection coefficients $R_{1,2}$ different from unity, while the reflection coefficients of all the other mirrors are equal to unity. Neglecting the nonlinear polarisability of the crystal and the population inversion in the active medium, we can find the wave vectors of intracavity fields $\boldsymbol{\kappa}_{1,2}$ and the cavity eigenfrequencies $\omega_{1c,2c}$. The following expressions can be derived for $\omega_{1c,2c}$:

$$\omega_{1c} = \frac{2\pi m + \delta_1}{T_{c1}}, \quad \omega_{2c} = \frac{4\pi m + \delta_2}{T_{c2}}, \tag{A6}$$

where m is the index of the axial mode and $T_{c1,c2} = c^{-1}\times\int_0^L n_{1,2}(x)\mathrm{d}x$ are the cavity single-pass times for the waves at the fundamental and second-harmonic frequencies.

The boundary conditions for the complex amplitudes of the fields $e_{1,2}(x,t)$ on the output coupler can be written as

$$R_{1,2}e_{1,2}(L,t) = e_{1,2}(0,t). \tag{A7}$$

Integrating Eqns (A2)–(A4) with respect to x from zero to L, we derive the following differential equations for the complex wave amplitudes and the density of population inversion averaged over the cavity length, $\langle e_{1,2}\rangle = L^{-1}\times\int_0^L e_{1,2}(x,t)\mathrm{d}x$ and $\langle\varDelta\rangle = \int_0^L \varDelta(x,t)\mathrm{d}x/L$:

$$T_c\frac{\mathrm{d}\langle e_1\rangle}{\mathrm{d}t} + e_1(L,t) - e_1(0,t) = -\mathrm{i}\sigma_1 l\langle e_1^*\rangle\langle e_2\rangle + \left[\frac{\sigma_a l_a\langle\varDelta\rangle}{2} + \mathrm{i}T_c(\omega - \omega_{1c}) - \frac{\alpha_1 l}{2}\right]\langle e_1\rangle, \tag{A8}$$

$$T_c\frac{\mathrm{d}\langle e_2\rangle}{\mathrm{d}t} + e_2(L,t) - e_2(0,t) = -\mathrm{i}\sigma_2 l\langle e_1\rangle^2 + \left[\mathrm{i}T_c(2\omega - \omega_{2c}) - \frac{\alpha_2 l}{2}\right]\langle e_2\rangle, \tag{A9}$$

$$\frac{\partial\langle\varDelta\rangle}{\partial t} = \frac{\langle\varDelta\rangle_0 - \langle\varDelta\rangle(1 + |e_1|^2/I_s)}{T_1}, \tag{A10}$$

where $\langle\varDelta\rangle_0 = \int_0^L \varDelta_0(x,t)\mathrm{d}x/L$.

Performing integration in the expressions above, we took into consideration that the complex wave amplitudes in a high-Q cavity only slightly vary within the cavity length. In addition, since we are interested in the mismatches of cavity eigenfrequencies $\omega_{1c} - \omega_{2c}$ that are small as compared with the frequency difference between the adjacent axial modes, the cavity single-pass times $T_{c1,c2}$ in Eqns (A8) and (A9) are replaced by the mean cavity single-pass time $T_c = (T_{c1} + T_{c2})/2$ (the difference in the times T_{c1} and T_{c2} should be taken into consideration only in formulas for cavity eigenfrequencies). Taking into consideration the boundary conditions, we can approximately write the terms $e_{1,2}(L,t) - e_{1,2}(0,t)$ involved in Eqns (A8) and (A9) as $e_{1,2}(L,t) - e_{1,2}(0,t) = (1 - R_{1,2})\langle e_{1,2}\rangle$. Now, let us introduce the coefficients of linear losses,

$$\frac{k_{1,2}}{2} = 1 - R_{1,2} + \int_0^L \frac{\alpha_{1,2}}{2}\mathrm{d}x$$

and assume that the nonlinearity coefficients σ_1 and σ_2 are equal to each other, $\sigma_1 = \sigma_2 = \chi/2$. We represent also $\langle e_{1,2}\rangle$ as $\langle e_{1,2}\rangle = a_{1,2}\sqrt{I_s}\exp(\mathrm{i}\varphi_{1,2})$ and introduce the nonlinearity parameter $\varepsilon = I_s(\chi l)^2$. Then, one can easily verify that Eqns (A8) and (A9) yield Eqns (2)–(4). Setting N in Eqn (A10) equal to the ratio of the mean density of population inversion $\langle\varDelta\rangle$ to its threshold value $\varDelta_{th}$ and introducing the notation $\langle\varDelta_0\rangle/\varDelta_{th} = 1 + \eta$, we also arrive at Eqn (5) for N.

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