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# Numerical Simulation of Ultrashort Pulse Generation by the DFB Dye Laser with Traveling-Wave Pumping

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**Abstract**—A system of equations for pulse generation by a DFB dye laser with spatially modulated traveling-wave pumping is obtained and investigated using numerical simulation. Analytical expressions are derived for the coefficients of amplification and coupling of counterpropagating pulses. It is shown that the traveling-wave pumping in a range of optimum parameters allows practically unidirectional generation of a single ultrashort pulse (USP) to be provided. Its duration is several times as short as the one in the case of a conventional DFB laser with pumping in the steady-state regime. The results obtained agree qualitatively with the available experimental data on kinetics of pulse generation by the dye DFB laser with traveling-wave pumping.

## 1. INTRODUCTION

Dye lasers with light-induced distributed feedback (DFB) allow generation of frequency-tunable single picosecond and subpicosecond ultrashort pulses (USP) to be obtained [1–3]. The DFB dye lasers are at a disadvantage in relation to conventional mode-locked lasers in terms of limiting duration values of their output pulses. However, owing to the design advantages and possibility of realization of stable generation of single USPs, the dye DFB lasers are most promising sources of narrow-band emission employed for purposes of superhigh resolution spectroscopy.

In the conventional scheme of a DFB laser the output USP duration is limited by the transit time $T = L/v$ of emission of the active medium with length $L$ [4] (here, $v = c/n$ is the velocity of light in the active medium, $n$ is the refractive index). To enhance power characteristics and to decrease the duration of the output pulse, a modified DFB-laser scheme with subpicosecond pumping in the traveling-wave regime is used [5]. In such a scheme (see Fig. 1) a side pumping pulse is shaped in such a way that its amplitude front is tilted with respect to the phase front (this pulse resembles a translationally scanned beam). This is achieved by slewing the amplitude front of the pumping pulse by an angle $\Theta$, making use of diffraction grating *1*. After the beam splitting by diffraction grating 2 and reflection from quartz block *3* the radiation of pumping is directed toward the side surface of cell *4* filled with a dye. As a result, the periodically modulated (due to the interference of convergent beams) radiation of pumping induces a light-induced grating responsible for DFB in the active medium. In this case, the spatial position of rulings on the light-induced grating remains stationary, while an intensity envelope of the pumping pulse moves along the active medium at the velocity $v_0 = c/\tan\Theta$. In the scheme under consideration two conditions are fulfilled, namely, a small length of light-induced DFB structure and a relatively large region of amplification of the generation pulse copropagating with pumping. Apparently, at $\Theta \longrightarrow 0$ ($v_0 \longrightarrow \infty$) this geometry corresponds to the traditional DFB-laser scheme with pumping in the steady-state regime.

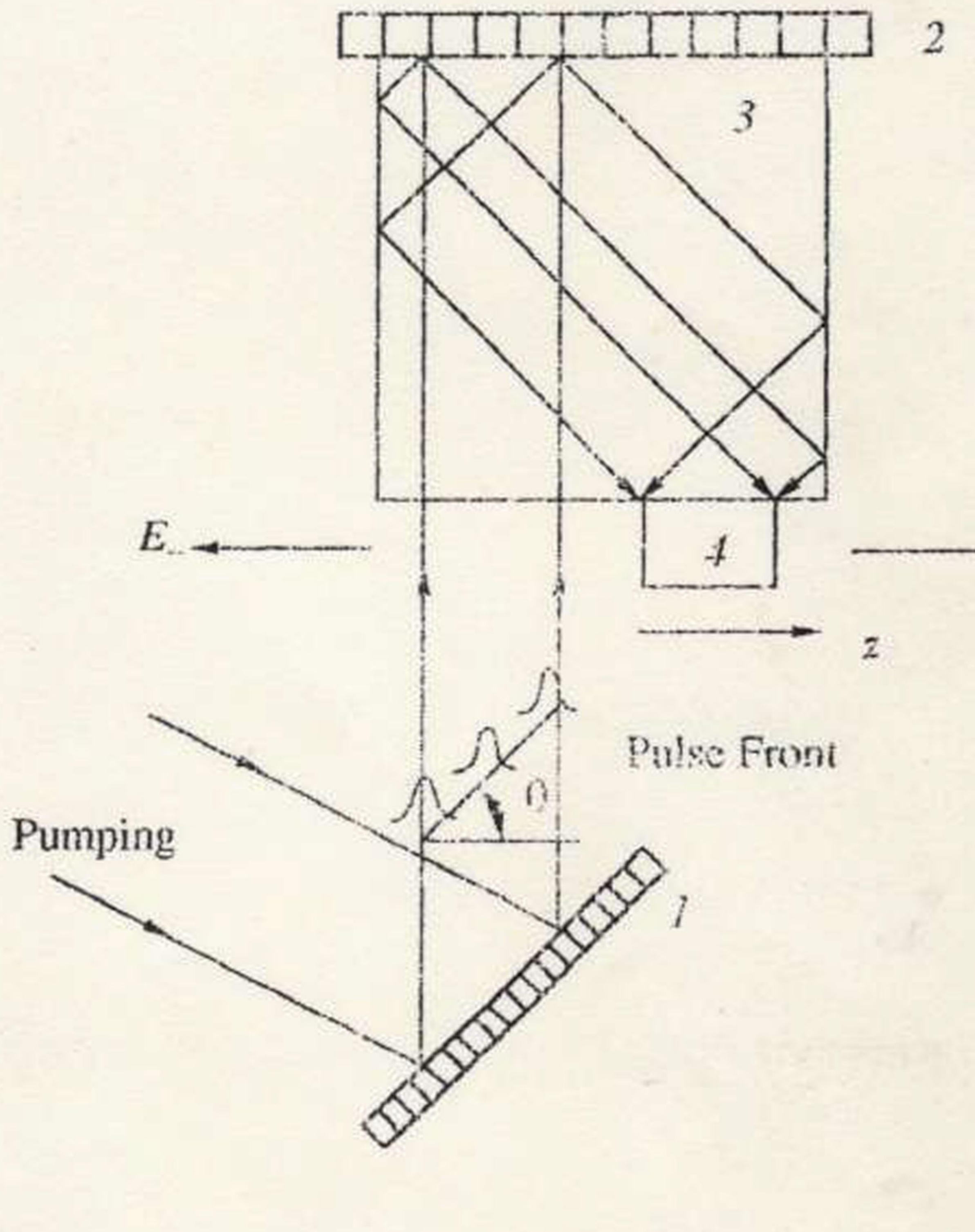


Fig. 1. The scheme of the DFB laser with the pulse pump in the traveling-wave regime [5]: (*1*, *2*) diffraction gratings, (*3*) quartz block, (*4*) dye cell.

This communication is devoted to the investigation of characteristic features of kinetics of pulse generation by a dye DFB laser with transverse pulse pumping in the traveling-wave regime by means of numerical simulation and general analytical estimations. Results of the numerical analysis are qualitatively compared with known experimental data, and an estimation is made of potentialities of such a DFB laser.

## 2. BASIC EQUATIONS

The active medium of the DFB laser is modeled by an ensemble of four-level molecules. It is assumed that the frequency of pulse generation $\omega$ coincides with the Bragg frequency, here DFB is caused by a gain grating induced by a space-modulated traveling pumping pulse, i.e., $\omega = \omega_b = \omega_0$, where $\omega_0$ is the central frequency of an amplification contour of the active medium. Respectively, a pumping rate of the DFB laser is specified as

$$W(z, t) = W_0[1 + \eta\cos\beta z], \tag{1}$$

where $W_0(z, t) = \overline{W}_0 \exp\{-4(t - z/v_0)^2/T_p^2\}$ is the amplitude of the Gaussian pumping rate with duration $T_p$; $\beta$ and $\eta \le 1$ is the spatial frequency of modulation and the parameter of visibility of an interference picture of pumping; $\overline{W}_0 = I_0(\sigma_p/\hbar\omega_p)$; $\sigma_p$, $I_0$, $\omega_p$ is the absorption cross section, peak intensity, and pumping pulse frequency, respectively. In the assumption that the duration of the pumping pulse $T_p$ is longer than the relaxation time $T_0$ of vibrational sublevels of the ground and excited state of molecules ($T_0 \sim 10^{-12}$ s), for the inverse population of a single molecule the following equation [7] is valid:

$$\frac{\partial N}{\partial t} + \left(\frac{1}{T_1} + W_p + W\right)N = W_p, \tag{2}$$

where $W = (cn\sigma_0/2\hbar\omega_0)|E|^2$ is the probability of the induced transition under the action of the generation radiation with the amplitude $E(\omega_0)$, $\sigma_0$ is the section of the induced transition at frequency $\omega_0$, $T_1$ is the lifetime of the excited electronic state. The resonant polarization of the active medium is determined by the expression

$$P = -iN_a\frac{cn\sigma_0}{4\pi\omega_0}NE(\omega_0), \tag{3}$$

where $N_a$ is the concentration of dye molecules.

A solution of the wave equation

$$\frac{\partial^2 E}{\partial z^2} - \frac{n^2}{c^2}\frac{\partial^2 E}{\partial t^2} = \frac{4\pi}{c^2}\frac{\partial^2 P}{\partial t^2} \tag{4}$$

for the active medium with periodically modulated parameters at $\omega_0 = \omega_b = c\beta/2n$ can be sought in the form of two couterpropagating pulses

$$E(\omega_0) = [A_+(z, t)\exp\{i\beta z/2\} + A_-(z, t)\exp\{-i\beta z/2\}]\exp\{-i\omega_0 t\} \tag{5}$$

with real amplitudes $A_\pm(z, t)$. In accordance with (1) and (5), for performing spatial averaging in (4) it is convenient to represent the inverse population $N$, being a fast-oscillating function of the $z$-coordinate, as a superposition of spatial harmonics with multiple frequencies $\beta_m = m\beta$ [7, 8],

$$N = N_0 + 2\sum_{m=1}^{\infty} N_m\cos m\beta z, \tag{6}$$

where $N_m(z, t)$ are the slowly changing amplitudes. After the substitution of (1), (5), and (6) into (2) and (4) and some simple transformations, a system of equations describing pulse generation by the DFB laser can be written as

$$\pm\frac{\partial E_\pm}{\partial z} + \frac{1}{v}\frac{\partial E_\pm}{\partial t} = G_0[N_0(E_\pm + F) + N_1 E_\mp], \tag{7}$$

$$\frac{N_m}{\partial t} + (T_1^{-1} + W_0 + E_+^2 + E_-^2)N_m + \left(\frac{\eta}{2}W_0 + E_+E_-\right)(N_{m+1} + N_{|m-1|}) = W_0\left(\delta_{0m} + \frac{\eta}{2}\delta_{1m}\right), \quad m = 0, 1, 2, \ldots \tag{8}$$

with the following boundary and initial conditions:

$$E_+(0, t) = E_-(L, t) = 0, \quad N_m(t = -\infty) = 0. \tag{9}$$

Here, $G_0N_0$ is the amplitude gain, $G_0N_1$ is the coupling coefficient of counterpropagating pulses due to the diffraction by the grating of the inverse population of the first order; $G_0 = N_a\sigma_0/2$, $E_\pm = \sqrt{cn\sigma_0/2\pi\hbar\omega_0}A_\pm$, $\delta_{jm}$ is the Kronecker symbol. The function $F$ of sources of spontaneous emission is inserted into wave equations (7). The power of these sources with allowance for the normalization of the amplitudes of generation pulses is determined by the relation [9]

$$F^2 = \frac{2n^2c\sigma_0}{\lambda^4}\Delta\Omega\Delta\lambda, \tag{10}$$

where $\lambda$, $\Delta\Omega$, and $\Delta\lambda$ are the wavelength, solid angle, and the spectral interval of radiation, respectively. Equations (7) and (8) do not take into account the possible absorption of generation radiation in making the transition from the first excited singlet state of molecules to the second state. They follow from more general equations derived in [10] with allowance for the finiteness of the relaxation time of vibrational sublevels of the electronic dye molecule states at $T_0 \longrightarrow 0$.

System of equations (8) has an exact analytical quadrature solution [11], from which the following formulas are obtained for the pertinent amplitudes of harmonics of the inverse population:

$$N_0 = \int_{-\infty}^{t} W_0 e^{\Gamma}[I_0(\gamma) + \eta I_1(\gamma)]dt',$$
$$N_1 = \int_{-\infty}^{t} W_0 e^{\Gamma}\left[\frac{\eta}{2}I_0(\gamma) + I_1(\gamma) + \frac{\eta}{2}I_2(\gamma)\right]dt', \tag{11}$$

where $I_m(\gamma)$ is the modified Bessel function of the $m$th order,

$$\Gamma(t, t') = -\int_{t'}^{t}\left[W_0 + \frac{1}{T_1} + E_+^2 + E_-^2\right]dt'',$$

$$\gamma(t, t') = -\int_{t'}^{t}[\eta W_0 + 2E_+E_-]dt''.$$

Note that for the numerical simulation of the above problem the expansion of the inverse population $N$ into a harmonic series is more preferable than the representation of the corresponding harmonics in the form of integral Fourier coefficients (11).

From the viewpoint of stable single USP generation the DFB laser is in an intermediate position between the conventional DFB laser and the superluminescence laser with spatially homogeneous traveling-wave pumping. As is shown in [9], with the fulfillment of one of the inequalities $\tau_p = T_p v/L > 1$ or $F^2\exp\{2G_0N_0L\} \ll 1$, the DFB laser with pumping in the traveling-wave regime approaches (in terms of parameters) the conventional DFB laser with pumping in the steady-state regime ($\Theta = 0$). This is attributable to the fact that, when the first condition is fulfilled, the active medium is excited practically in the same manner as in the case of steady-state pumping, but, when the second condition is fulfilled ($\tau_p \ll 1$), the generation pulse, propagating in step with pumping, changes the inverse population insignificantly. Correspondingly, the basic portion of the stored energy is yielded up after the pumping pulse passed through the active medium of the DFB laser.

In the other limiting case, when the inverse inequalities $\tau_p \ll 1$ and $F^2\exp\{2G_0N_0L\} \gg 1$ are simultaneously fulfilled, the DFB laser is similar to the superluminescence laser with spatially homogeneous transverse pumping in the traveling-wave regime. Here, owing to considerable saturation of the inverse population by the generation pulse copropagating with pumping, there occurs the effective "erasure" of the dynamical gain grating, and the DFB influence becomes insignificant, which must be accompanied by considerable broadening of the spectrum of generation radiation. Correspondingly, the characteristic features of kinetics of pulse generation by the DFB laser, which are due to pumping in the traveling-wave regime, will be most pronounced in the intermediate range of parameters.

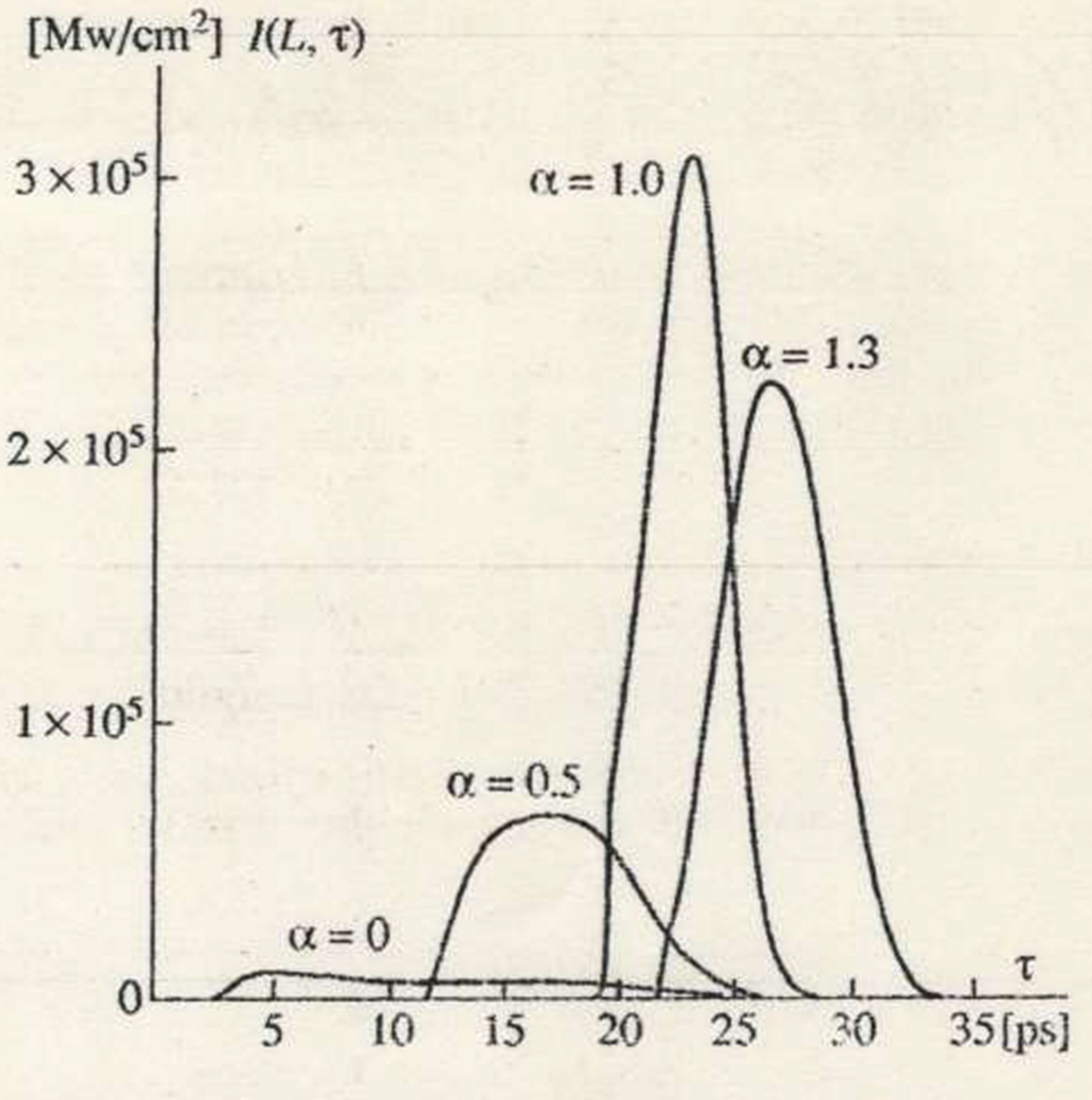


**Fig. 2.** The generation pulse intensity $I(t)$ versus the parameter $\alpha$ at $H_0 = Lv^{-1}\overline{W}_0 = 60$ and $G_0L = 50$, which corresponds to $I_0 = 4.5$ Mw/cm$^2$ and $N_a = 1.4 \times 10^{18}$ cm$^{-3}$.

## 3. RESULTS OF NUMERICAL SIMULATION

For the numerical solution of problem (7) through (9), use was made of the characteristic finite-difference method. Parameters of the problem corresponded to the rhodamine C ethanol solution DFB laser with pumping by a Gaussian pulse of duration $T_p = 5$ ps ($\lambda_p = 527$ nm) and duration emitted at the wavelength $\lambda_p = 605$ nm [12]: $n = 1.36$, $\sigma_p = 1.5 \times 10^{-16}$ cm$^2$, $\sigma_0 = 1.4 \times 10^{-16}$ cm$^2$, $T_1 = 4$ ns, and $L = 0.5$ cm. Results of the numerical experiment at $\eta = 1$ are given in Figs. 2–4.

Figure 2 shows intensity envelopes of the generation pulse[1] copropagating with pumping at the output boundary of the active medium, $z = L$, as a function of the parameter of velocities "synchronization" $\alpha = v/v_0$. The curve for $\alpha = 0$ corresponds to the case of the traditional DFB-laser scheme with steady-state pulsed pumping ($v_0 \longrightarrow \infty$). In this case, the intensity envelopes of generation pulses at the left ($z = 0$) and right ($z = L$) boundaries of the active medium coincide, and their duration is comparable with the time of radiation propagation along the active medium ($T = L/v \approx 23$ ps). As is seen, the traveling-wave regime of pumping ($\alpha \neq 0$) considerably increases the peak intensity of generation pulse $I(t)$ and, at the same time decreases its dura-

[1] All results below refer to the generation pulse copropagating with pumping.

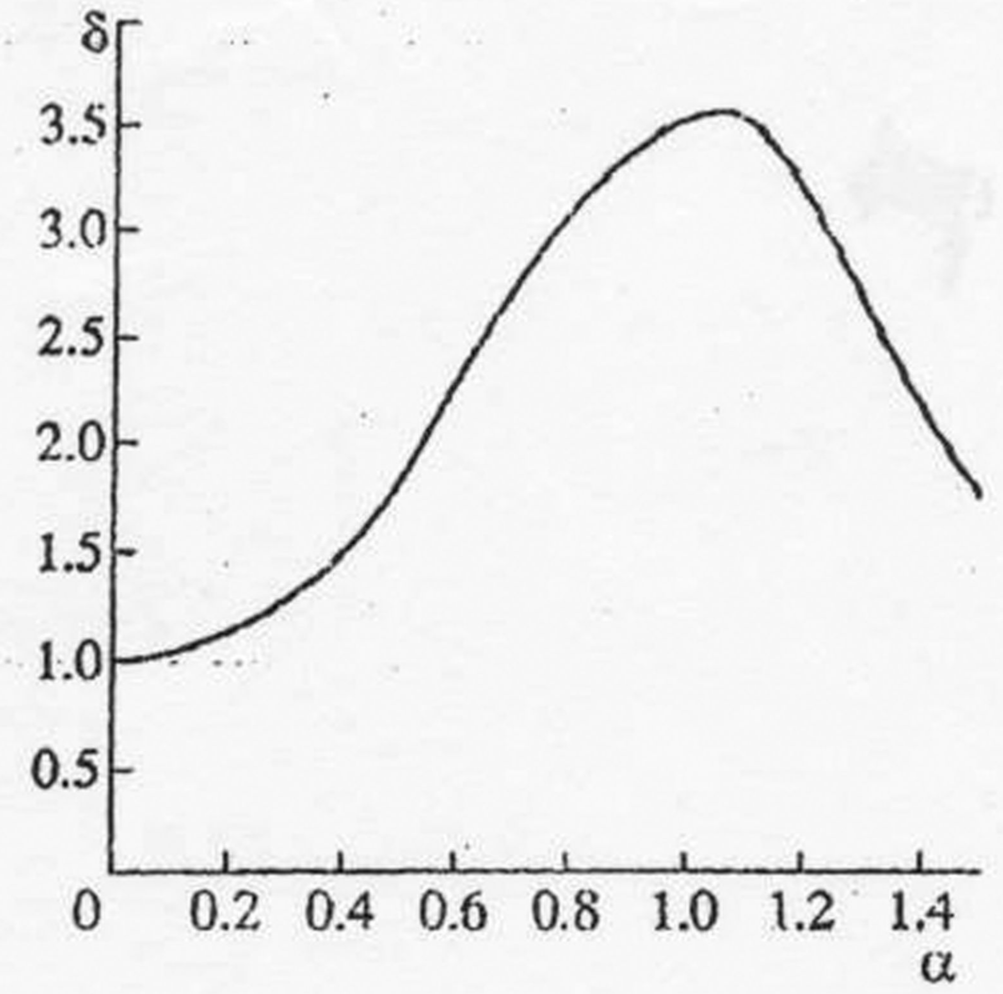


**Fig. 3.** The dependence of the compression degree δ on α for the generation pulse at $H_0 = 60$ and $G_0L = 50$.

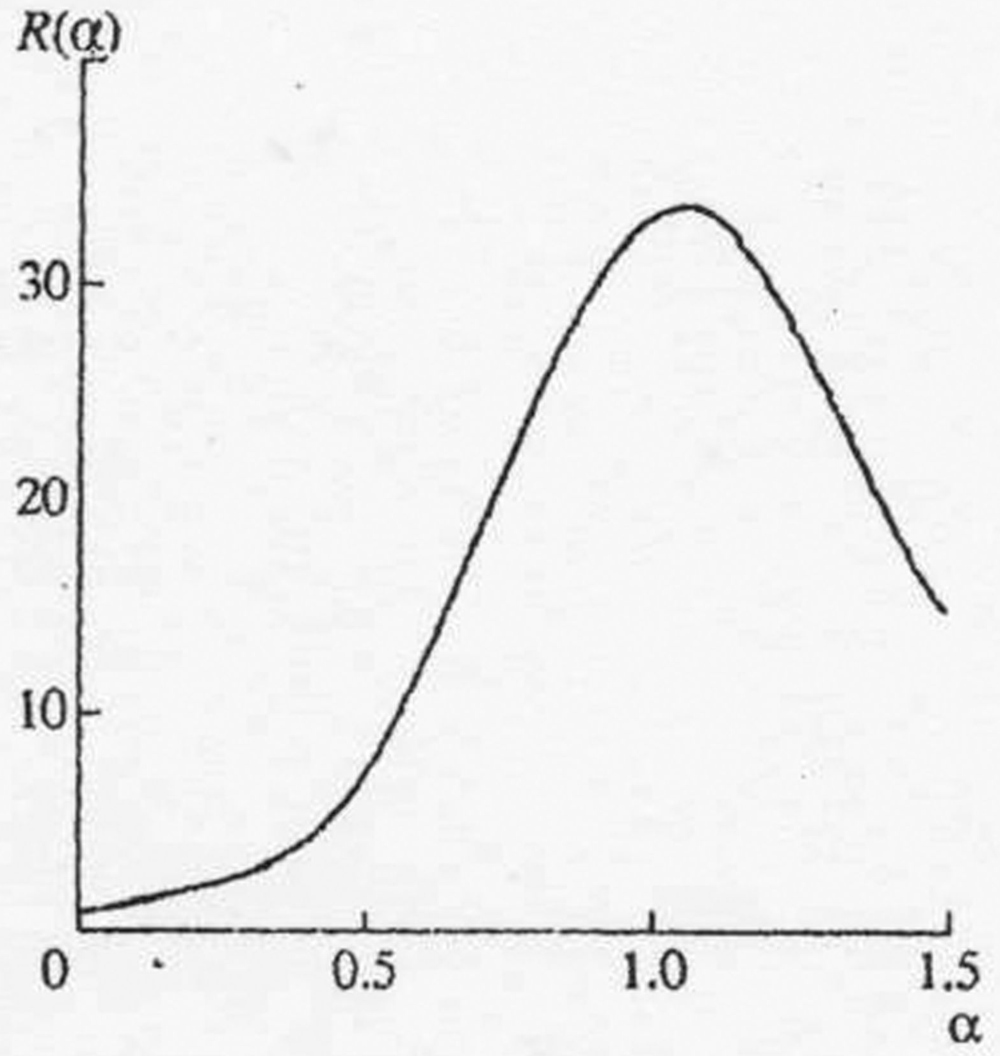


**Fig. 4.** The dependence of the normalized peak intensity $R(\alpha) = I_{max}(\alpha)/I_{max}$ $(\alpha = 0)$ on α for the generation pulse at $H_0 = 60$ and $G_0L = 50$.

tion $T_g$. In this case, in keeping with changing the velocity $v_0$ of a traveling wave of pumping, the generation pulse is displaced, with the growth of α, toward large values of $t$. In the case of complete "synchronization" of the velocities ($\alpha = 1$), the form of the pulse is close to the Gaussian form with the maximum intensity and shortest duration $T_g \approx 0.2T \leq T_p$. Further increasing the parameter $\alpha > 1$ (see the curve corresponding to $\alpha = 1.3$) causes a decrease in the intensity and an increase in the duration of the generation pulse. Note that for $\alpha \leq 1$ the generation pulse counterpropagating with the traveling wave of pumping has the considerably lower peak intensity and, respectively, longer duration.

Figure 3 illustrates a relative variation of the generation pulse duration as a function of the parameter α. As is seen in Fig. 3, in the given range of parameters, owing to the traveling-wave pump, the modified DFB-laser scheme provides for the generation of pulses, the duration of which is several times as short (at $\alpha = 1$, $\delta(\alpha) > 3$) as the one in the case of the conventional DFB laser operated in a steady-state pumping regime [6]. Here, the peak intensity of the generation pulse intensity increases by more than an order of magnitude (see Fig. 4), and the regime of the unidirectional pulse generation is realized.

The presented results of numerical simulation are obtained at $\bar{F} = F\sqrt{T} \approx 3 \times 10^{-5}$, which corresponds to $\Delta\lambda = 3$ nm and $\Delta\Omega = 10^{-3}$ ster [12].

Note that the basic characteristics of the generated pulse change insignificantly (within 10%).

## 4. CONCLUSION

The results of numerical simulation agree qualitatively with experimental data of [12] where at $\alpha = 1$ a single USP with the duration $T_g < T$ was generated. A quantitative theoretical description of the experiments [5] and [12] involves difficulties due to the absence of a number of values of the parameters necessary in theory. In particular, in [5] there are no data on the length of the DFB structure and intensity of the pumping pulse, while the parameter of "synchronization" of velocities α changes with changing a refractive index $n$ of active medium, which leads to a frequency change in the wavelength of emitted radiation in the band $\Delta\lambda = (576–606)$ nm. In this case the gain cross section $\sigma_0(\lambda)$ undergoes a corresponding change. In particular, for rodamine 6G the detuning $|\Delta\lambda| \approx 15$ nm causes, approximately, a 20% decrease in $\sigma_0(\lambda)$ against the maximum value of $\sigma_0(\lambda_0)$ [13]. In [12] no data are reported on the intensity of the pumping pulse and dye concentration (rhodamine C).

Thus, the analysis made shows that the modified DFB-laser scheme (see Fig. 1) allows single output USPs to be generated, with their profiles being close to the Gaussian profile and duration $T_g$ considerably shorter than $T$ of radiation propagation along the active medium. In this case, practically unidirectional pulse generation is realized with the intensity being appreciably higher than in the traditional DFB-laser scheme with steady-state pumping [6].

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