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Diode amplifier of modulated optical beam power

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Abstract. **Analytical relations are obtained between characteristics of modulated light at the output and input of an optical diode power amplifier operating in the highly saturated gain regime. It is shown that a diode amplifier may act as an amplitude-to-phase modulation converter with a rather large bandwidth (~10 GHz). The low sensitivity of the output power of the amplifier to the input beam power and its high energy efficiency allow it to be used as a building block of a high-power multielement laser system with coherent summation of a large number of optical beams.**

Keywords: *diode optical amplifier, amplitude/phase modulation.*

1. Introduction

The development of high-speed free-space mobile optical communications requires bright sources of modulated optical beams. One optimal type of such sources is a diode laser or a single-crystal integrated diode heterostructure in which a semiconductor power amplifier is the output stage of the integrated structure (see e.g. Refs [1, 2]).

Recent work [3, 4] has shown that, using an optical scheme with a single-frequency diode laser and diode amplifier/modulator, one can obtain modulated optical beams ~100 mW in average power, capable of information transfer at a rate of ~20 Gb s^{-1}. This average power is, however, not quite sufficient for free-space communications. The guiding target here is a power level of 1 W and above, characteristic of, e.g., intersatellite communication [5]. Such average power is difficult to reach through direct modulation of a diode laser source because of both the necessity for a large amplitude of the modulated electrical signal and the associated degradation of the optical quality of the beam.

In this context, it is of interest to use an optical power amplifier. A modulated optical beam with a low or moderate average power, e.g. ~100 mW or less, is fed to its input, and the output power reaches 1 W or more. The amplifier is energised by a direct current sufficient for maintaining a constant average output optical power.

It is worth noting here that, to ensure efficient electrical-to-optical power conversion, the amplifier should operate deep in saturation. This inevitably entails changes in the modulation state of a beam when it travels through the amplifier, due to its optical nonlinearity. The purpose of this work was to find out how and to what extent the beam modulation changes in this process.

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2. Analysis of modulated beam amplification

The propagation of a modulated optical beam through an amplifier will be analysed here in terms of theory developed previously [4]. The amplifier configuration is schematised in Fig. 1. The amplifier is thought to be a laser diode based on a ridge heterostructure with an optical waveguide that supports only one, fundamental transverse mode. The diode facets are nonreflective, e.g. owing to an antireflection coating, and are inclined at an angle θ to the waveguide axis, like in a previous study [6]. Thus, the laser diode under consideration is a single-pass amplifier, or a travelling wave amplifier.

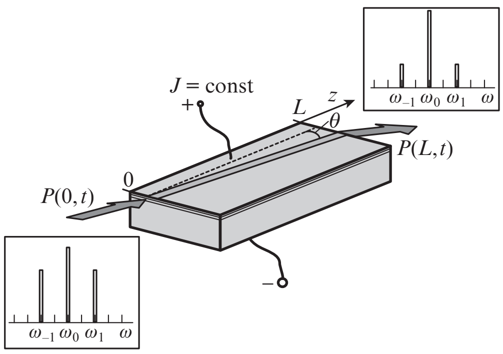

Figure 1. Simplified schematic of the amplifier: θ is the angle between the optical axis of the amplifier and the normal to the diode facet.

To analyse the operation of such an amplifier, consider first the amplification of a quasi-monochromatic wave harmonically modulated in the microwave range, with an optical carrier frequency ω_0 and side frequencies ω_1 and ω_{-1}, which are related by

$$\omega_1 = \omega_0 + \Omega,\ \ \omega_{-1} = \omega_0 - \Omega. \tag{1}$$

The frequencies ω_{-1} and ω_1 correspond to the Stokes and anti-Stokes components, respectively. For a beam propagating along the z axis of the amplifier, the field amplitude $\mathcal{E}$ as a function of time t can be represented as

$$\mathcal{E}(z,t) = \frac{1}{2}\{E_0(z)\exp[\mathrm{i}(k_0 z - \omega_0 t)][1 + V_{+1}(z)\exp[\mathrm{i}(qz - \Omega t)] + V_{-1}(z)\exp[-\mathrm{i}(qz - \Omega t)]] + \text{c.c.}\}, \quad (2)$$

where $E_0(z)$ is the 'slow' complex-valued amplitude of the wave with a carrier frequency ω_0; k_0 is its waveguide propagation constant; and V_{+1} and V_{-1} are the 'slow' relative amplitudes of the waves at the side frequencies. Their propagation constants k_1 and k_{-1} satisfy the relations

$$k_{\pm 1} = k_0 \pm q. \quad (3)$$

Here, $k_0 = (\omega_0/c)n$ and $q = (\Omega/c)n_{\text{gr}}$, where n and n_{gr} are the effective waveguide phase and group indices of refraction and c is the speed of light.

Since we analyse a mode whose transverse profile remains unchanged as it propagates in the amplifier and is amplified, the particular character of the profile is unimportant in this case and is missing in (2). The field intensity $\mathcal{E}(z,t)$ is here taken to mean its value on the optical axis of the amplifier. In solving the waveguide problem, the characteristics of the transverse mode profile, n and n_{gr}, and the optical confinement factor Γ, which relates the mode and material gain coefficients to each other, are determined separately (see e.g. Refs [7, 8]).

The dimensionless functions $V_{+1}(z)$ and $V_{-1}(z)$ characterise the modulation depth, and the relationship between them specifies the type of modulation: at $V_{+1} = V_{-1}^*$, the signal is amplitude-modulated, and $V_{+1} = -V_{-1}^*$ corresponds to a phase-modulated wave. The other cases correspond to a mixed, amplitude/phase modulation. In the case of linear operation of the amplifier, neglecting the refractive index and gain dispersion we find that the spectral components in (2) propagate and are amplified in the same way and independently of each other. The relationship between their slow amplitudes remains completely unchanged and, hence, the modulation parameters also persist. In the case of nonlinear amplification, the situation is fundamentally different.

We consider modulation frequencies Ω at which the intraband carrier dynamics have a quasi-equilibrium character, i.e. $\Omega\tau_2 \leqslant 1$, where τ_2 is the intraband relaxation time. It is reasonable to assume, for example, that τ_2 is far less than 10^{-13} s [9], so Ω is limited from above by frequencies of the order of hundreds of gigahertz.

The theory developed in Ref. [4] is valid in the case of moderate modulation depths, i.e. for $|V_{\pm 1}| < 0.5$. Here, the optical power is also thought to be concentrated in the central spectral component (in the carrier). Therefore, the expression for its amplitude $E_0(z)$ has the form [4]

$$E_0(z) = |E_0(z)|\exp[\mathrm{i}\varphi(z)],$$

where

$$|E_0(z)|^2 = (8\pi I_s/cn)u(z);$$

$$\varphi(z) = \varphi(0) - \frac{R}{2}\left[\alpha z + \ln\left(\frac{u(z)}{u_0}\right)\right]; \quad (4)$$

$u(z) = I(z)/I_s$; $I_s = \hbar\omega_0/(\sigma\tau)$; $u_0 = u(0)$; $I(z)$ and I_s are the carrier wave and saturation intensities, respectively; $u(z)$ is the normalised (dimensionless) carrier wave intensity; σ is the stimulated transition cross section (differential gain); τ is the spontaneous lifetime; R is the amplitude–phase coupling factor, whose waveguide value was determined e.g. in Ref. [10]; and α is the background loss coefficient.

As shown earlier [4], the coordinate dependence of the normalised intensity, $u(z)$, can be found by solving the transcendental equation

$$\frac{u(z)}{u_0}\left|\frac{g - u_0}{g - u(z)}\right|^{g+1} = \exp(g\alpha z), \quad (5)$$

where $g = (\Gamma G_0 - \alpha)/\alpha$ is a dimensionless quantity characterising the mode gain; G_0 is the unsaturated material gain in the active layer, related to the pump current J by

$$G_0 = \frac{\sigma\tau}{ed_{\text{a}}W_0 L}(J - J_{\text{tr}}); \quad (6)$$

e is the electron charge; d_{a} is the thickness of the active layer; W_0 is the effective width of the pump region; L is the amplifier length; and J_{tr} is the transparency current for resonance optical losses (which corresponds to concentration N_{tr} at the inversion threshold).

From the above, we have

$$g = \beta(J - J_{\text{tr}}) - 1, \quad (7)$$

where

$$\beta = \frac{\Gamma\sigma\tau}{\alpha ed_{\text{a}}W_0 L}.$$

The power of the optical flux $P(z)$ propagating through the amplifier can be expressed through the normalised intensity $u(z)$:

$$P(z) = dWI_s u(z), \quad (8)$$

where W and d are the in-plane and out-of-plane effective cross-sectional beam dimensions in the amplifier. They can be found together with the effective index of refraction and optical confinement factor in solving the waveguide problem, as pointed out above. In our case,

$$d \cong d_{\text{a}}/\Gamma, \; W \cong \Gamma_{\text{w}}W_0. \quad (9)$$

Here, Γ_{w} is taken to mean the pump efficiency. It differs from unity because of both the spatial spread of the current and the presence of additional shunt paths.

Thus, Eqns (4)–(9) are sufficient for finding the static characteristics of the amplifier in our model. Setting the input power $P_0 = P(0)$, the material parameters of the amplifier, and the pump current J, we can find the power distribution along the length of the amplifier and the output power $P = P(L)$.

Figures 2 and 3 present numerical calculation results for power P at the amplifier parameters indicated in Table 1. The parameters are typical of state-of-the-art high-power single-transverse-mode semiconductor oscillators. It follows from the data in Figs 2 and 3 that the background optical loss α plays a key role in determining the efficiency and output power limit of the amplifier. To achieve an output power at a level of several watts, this loss should be under 2 cm^{-1}. Clearly, one should remember other limitations on the output power,

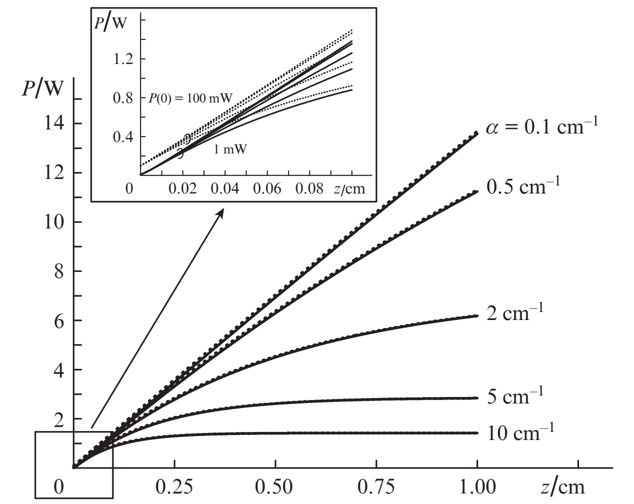

Figure 2. Power (P) distribution of an optical beam along the length (z) of the amplifier at various background loss coefficients, a pump current of the amplifier $J = 10$ A and an input optical beam power of 1 (solid lines) and 100 mW (broken lines). Inset: portions of the $P(z)$ curves on expanded scales.

related to the distortion of spatially single-mode operation (see e.g. Ref. [8]) and optical damage to the laser cavity medium, which have been the subject of several experimental and theoretical studies [11–14].

The data in Fig. 2 can be used to optimise the cavity length at a given α value. To this end, the curves should be regarded as plots of the amplifier output power against length $z = L$ at a constant pump current density.

It can be seen from the data in Fig. 3 that the energy efficiency of the amplifier as a light source can approach that of state-of-the-art diode lasers, i.e. the physical limit ($\sim\hbar\omega_0/e$). Moreover, at high input powers it is essentially independent of the input beam power and is determined primarily by the background loss. Indeed, it can be shown directly from (4)–(9) (see Appendix 1) that, in the strong gain saturation approximation,

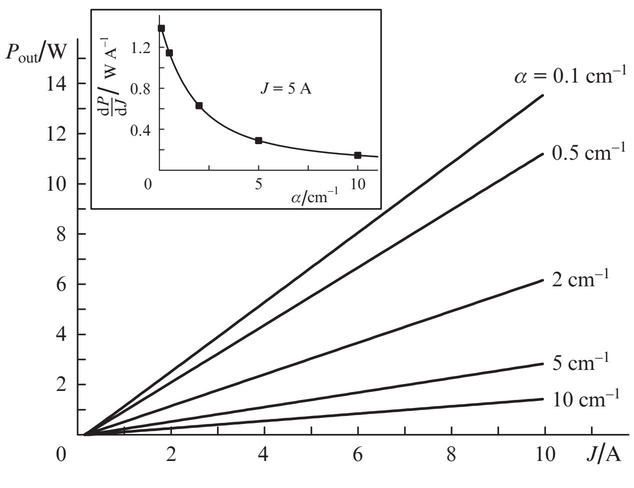

Figure 3. Power–current characteristics of the amplifier at various α values and an input optical beam power $P(0) = 10$ mW. Inset: amplifier efficiency $\mathrm{d}P/\mathrm{d}J$ as a function of α: the solid squares represent numerical calculation using Eqns (5)–(9) and the solid line represents the results obtained using the approximate relation (10).

Table 1. Parameters used in assessing characteristics of a power amplifier.

Parameter	Notation	Value
Stimulated transition cross section (differential gain)	σ/cm^2	10^{-15}
Spontaneous carrier recombination time	τ/ns	1.0
Optical confinement factor	Γ	0.01
Thickness of the active region	d_a/nm	8.0
Background loss coefficient in the waveguide	α/cm^{-1}	$0.1<\alpha<10$
Transparency carrier concentration in the active region	$N_{tr}/10^{18}$ cm^{-3}	2.0
Master oscillator wavelength	λ_0 /nm	850
Width of the pump region	W/μm	6.0
Amplifier length	L/cm	1.0
Effective beam size along the layers of the heterostructure	W_0/μm	6.0
Effective beam size across the layers	$d \approx d_a \Gamma^{-1}$ / μm	0.8
Saturation intensity	I_s/W cm^{-2}	2.33×10^5
Saturation power	$P_s = WdI$/mW	11.18
Transparency current	J_{tr}/A	0.154
Proportionality factor between the normalised gain and pump current	β/mA^{-1}	0.26

$$\frac{\mathrm{d}P}{\mathrm{d}J} \approx \Gamma_{\mathrm{w}} \frac{\hbar\omega_0}{e} \frac{1-\exp(-\alpha L)}{\alpha L}. \tag{10}$$

It can be seen in the inset in Fig. 3 that, despite its approximate nature, Eqn (10) ensures good agreement with the numerically calculated amplifier efficiency.

As to the dynamic characteristics of the amplifier, note that, at any frequency Ω, they are determined by the transformation of the relative amplitudes of the side components, $V_{+1}(z)$ and $V_{-1}(z)$. According to theory [4], in our case (at a constant pump current of the amplifier) the coordinate dependence of these parameters takes the form

$$\begin{bmatrix} V_{+1}(z) \\ V^*_{-1}(z) \end{bmatrix} = K(u(z),\Omega)\, C_1 \begin{bmatrix} \mathrm{e}^{-\mathrm{i}\psi} \\ \mathrm{e}^{\mathrm{i}\psi} \end{bmatrix} + C_2 \begin{bmatrix} 1 \\ -1 \end{bmatrix}. \tag{11}$$

Here $K(u(z),\Omega) \equiv [\Phi(u(z),\Omega)]^{-1}$; $\tan\psi = R$; and $\Phi(u(z),\Omega)$ is a complex-valued function:

$$\Phi(u,\Omega) = \exp\left\{\frac{g+1}{g+1-\mathrm{i}\Omega\tau} \ln\left[\frac{(1+u-\mathrm{i}\Omega\tau)(1-u_0/g)}{(1-u/g)(1+u_0-\mathrm{i}\Omega\tau)}\right]\right\}. \tag{12}$$

It is easy to see that $\Phi(u_0,\Omega) = 1$ at the amplifier input. Therefore, the constants C_1 and C_2 will be expressed through the relative amplitudes of the side components at the input as follows:

$$C_1 = \sqrt{1+R^2}\,\frac{V_{+1}(0)+V^*_{-1}(0)}{2},$$
$$C_2 = \frac{V_{+1}(0)-V^*_{-1}(0)}{2} + \mathrm{i}R\frac{V_{+1}(0)+V^*_{-1}(0)}{2}. \tag{13}$$

Equations (11)–(13) and the known dependence $u(L)$ allow one to calculate the transformation of a harmonically modulated input signal described by $E_0(z = 0)$, $V_{+1}(z = 0)$ and $V_{-1}(z = 0)$ into an output signal, which can be described by $E_0(z = L)$, $V_{+1}(z = L)$ and $V_{-1}(z = L)$.

It is easy to note from (11) and (13) that, when a purely phase-modulated signal ($C_1 = 0$) is fed to the input of the amplifier, only a phase-modulated signal, with the same modulation depth, will emerge at its output. At the same time, if the input signal has amplitude or mixed ($C_1 \neq 0$) modulation, the output signal will always have mixed modulation. Thus, the amplifier will in general significantly change the nature and parameters of optical beam modulation.

There is special interest in one possible type of mixed modulation, when a monochromatic beam is modulated in a diode amplifier/modulator [4]. According to previous work [4], $V_{+1}(0)$ and $V_{-1}(0)$ then have the form

$$\begin{bmatrix} V_{1+}(0) \\ V^*_{-1}(0) \end{bmatrix} = V(\Omega,\gamma)\begin{bmatrix} \mathrm{e}^{-\mathrm{i}\psi_\mathrm{m}} \\ \mathrm{e}^{\mathrm{i}\psi_\mathrm{m}} \end{bmatrix}, \qquad \tan\psi_\mathrm{m} = R_\mathrm{m}, \tag{14}$$

where $V(\Omega,\gamma)$ is a common complex-valued factor, which depends on the relative amplitude of the modulating current (γ), modulation frequency (Ω) and parameters of the amplifier/modulator, and R_m is the amplitude–phase coupling factor for the gain medium of the amplifier/modulator. Comparison of Eqns (11) and (13) with (14) leads us to conclude that, if $\psi = \psi_\mathrm{m}$, i.e. if the amplitude–phase coupling factor of the medium of the amplifier/modulator, R_m, is identical to that of the medium of the power amplifier, R, we have $C_2 = 0$ and, hence, the nature of the signal modulation is not influenced by signal amplification. Only the depth of such modulation can change. In the general case ($R_\mathrm{m} \neq R$), both the depth and nature of modulation change. From (11), (13) and (14), one can readily obtain the ratio of the intensities of the side components of the output beam to the intensities of the input beams:

$$\left|\frac{V_{\pm 1}(L)}{V_{\pm 1}(0)}\right|^2 = \{|K(u(L),\Omega)|^2(1+R^2) + (R_\mathrm{m}-R)^2 + 2\,|K(u(L),\Omega)|$$
$$\times (R_\mathrm{m}-R)\sqrt{1+R^2}\sin(\psi_\mathrm{m} \mp \delta(\Omega))\}\,(1+R_\mathrm{m}^2)^{-1}, \tag{15}$$

where $\delta(\Omega) = \arg[K(u(L),\Omega)]$.

According to (11) and (13), the linear relation between $V_{+1}(z = 0)$ and $V_{-1}(z = 0)$ at the amplifier input and $V_{+1}(z = L)$ and $V_{-1}(z = L)$ at the output makes it possible to find, within our model, a transformation that characterises the change of an arbitrarily modulated signal in the amplifier. Indeed, representing the field intensity $\mathcal{E}$ at $z = 0$ by analogy with (2) in the form

$$\mathcal{E}(0,t) = \frac{1}{2}\{E_0(0)\exp(-\mathrm{i}\omega_0 t)$$
$$\times[1 + a(t) + b(t)] + \mathrm{c.c.}\}, \tag{16}$$

where $a(t)$ is a real function and $b(t)$ is an imaginary function [$a(t) = a^*(t)$, $b(t) = -b^*(t)$], both slowly varying with time, which represent amplitude and phase modulations, it is easy to find the field intensity at the amplifier output, $\mathcal{E}(z = L, t)$, in general form. To this end, $a(t)$ and $b(t)$ should be represented as a Fourier integral in terms of frequency Ω, and then the Fourier components $a(\Omega)$ and $b(\Omega)$ should be transformed according to (11). As shown in Appendix 2, the result has the form

$$\mathcal{E}(z = L, t) = \frac{1}{2}\{E_0(L)\exp[\mathrm{i}(k_0 L - \omega_0 t)]$$
$$\times[1 + A(t) + B(t)] + \mathrm{c.c.}\}, \tag{17}$$

where $E_0(L)$ is defined in (4), and $A(t)$ and $B(t)$ are the real and imaginary parts of the modulation amplitude at the amplifier output. The real function $A(t)$, corresponding to amplitude modulation of the output signal, can be expressed through the input signal $a(t)$ as

$$A(t) = \int_{-\infty}^{\infty} a(\Omega)K[u(L),\Omega]\exp(-\mathrm{i}\Omega t')\mathrm{d}\Omega, \tag{18}$$

where $a(\Omega)$ is the Fourier transform of the function $a(t')$, corresponding to amplitude modulation of the input signal (16), at time $t' = t - L/v_\mathrm{gr}$ ($v_\mathrm{gr} = c/n^*$). The time variable t' differs from t by the time delay L/v_gr, needed for a wave packet with a group velocity v_gr to pass through an amplifier of length L. According to (18), the complex-valued function $K(\Omega) \equiv K(u(L),\Omega)$, defined by (11) and (12), acts as a frequency filter for the modulation signal when an optical beam passes through the amplifier and is amplified.

The imaginary function $B(t)$, which represents the phase modulation of the output optical beam, satisfies the relation (see Appendix 2)

$$B(t) = b(t') + \mathrm{i}R(a(t') - A(t)), \tag{19}$$

which reflects the fact that the phase-modulated component $b(t)$ at the amplifier input is transferred to the output without changes, in the form $b(t')$. This can also be seen from (11), where the spectral components V_{+1} and V_{-1}, proportional to the coefficient C_2, are also transferred from the input of the amplifier to its output without any changes. At the same time, according to (19) the output signal $B(t)$ has a component proportional to the amplitude modulation depth $a(t')$ at the input. Note in advance that, in cases of practical interest, where the optical beam power at the amplifier output is sufficiently high and there is strong gain saturation, as discussed below, we have $|K(\Omega)| \ll 1$ in a considerable frequency range, $\Omega \leqslant \Omega_\mathrm{b}$. We then have

$$B(t) = b(t') + \mathrm{i}Ra(t'), \tag{20}$$

Therefore, if an amplitude-modulated wave

$$\mathcal{E}^\mathrm{in}(t) = \frac{1}{2}\{E_0(0)\exp(-\mathrm{i}\omega_0 t)[1 + a(t)] + \mathrm{c.c.}\} \tag{21}$$

is directed to the amplifier input, at the output we obtain a phase-modulated amplified wave,

$$\mathcal{E}^\mathrm{out}(t) = \frac{1}{2}\{\sqrt{Q}\,E_0(0)\exp[\mathrm{i}(k_0 L - \omega_0 t + \varphi(L))]$$
$$\times[1 + \mathrm{i}Ra(t')] + \mathrm{c.c.}\}, \tag{22}$$

where $a(t')$ is a real function representing the amplitude modulation of the input signal at time t'. The quantity

$$Q = u(L)/u_0 \tag{23}$$

can be thought of as a power gain coefficient. It should be kept in mind of course that, because the amplifier operates deep in saturation, this quantity loses its usual meaning because it depends significantly on the input signal. In any case, however, relations (4), which define Q through u_0 and $u(L)$ in (23), remain valid.

It follows from analysis of (21) and (22) that amplitude modulation at the amplifier input transforms into phase modulation at the output. This is the well-known self-phase modulation of a wave packet when it passes through a nonlinear optical medium, in particular through the gain medium of a diode laser [15, 16]. In the case under consideration, nonlinearity is associated with gain saturation and the effect of inversion on the refractive index of the semiconductor.

According to (22), an important feature of the amplitude-to-phase modulation transformation is that it is possible in a wide spectral range. The reason for this is that the coefficient R, which characterises this transformation, is constant. R may vary significantly only in the optical range, so it is constant with high accuracy at frequencies $\omega_0 \pm \Omega$ when Ω lies in the microwave or rf range.

Thus, the universal nature of relations (21) and (22) is ensured by the condition $|K(\Omega)| \ll 1$, which is necessary for relation (20) to be fulfilled with sufficient accuracy.

Figures 4–7 illustrate to what extent and in what frequency ($\nu = \Omega/2\pi$) range these relations can be valid in real cases of practical interest. It can be seen from Fig. 4 that, in the frequency range $\Omega/2\pi \leqslant \Omega_b/2\pi \approx 10$ GHz, $|K(\Omega)|$ does not exceed 0.2. Thus, at frequencies below 10 GHz amplitude modulation converts into phase modulation with an accuracy far better than 20%. The conversion result is essentially independent of input optical beam power $P(0)$. It follows from (11) and (12) that, in the frequency range in question, $|K(\Omega)| \to 0$ with increasing output power ($\sim u$). This means that, when an optical signal is amplified, its amplitude modulation is effectively suppressed. Moreover, not only the relative modulation depth but also its absolute value decrease, i.e. the amplification of an optical beam in the amplifier under consideration is accompanied by some stabilisation of its output power. This is an expected result given the strong gain saturation in the problem under consideration.

In addition to the suppression of the amplitude-modulated signal, there is severe distortion of its shape. This is illustrated by the frequency dependence of $\delta n_{\text{gr}}^{\text{am}}$ in Fig. 5, which characterises the spectral dispersion in the amplifier:

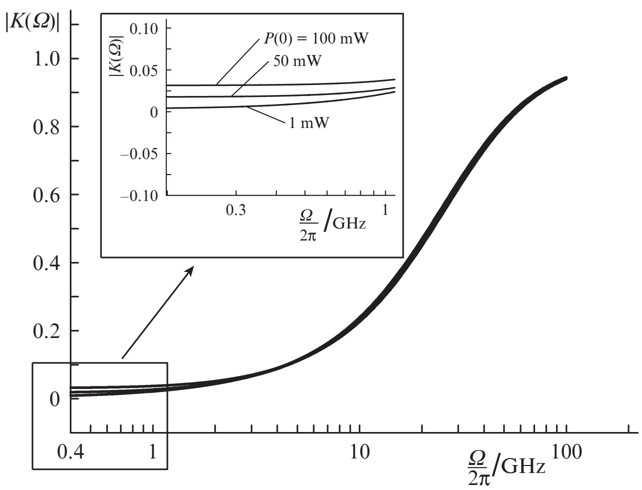

Figure 4. Magnitude of the transfer function of the amplifier, $|K(\Omega)|$, vs. frequency at different input signal powers $P(0)$ and a pump current $J = 2$ A.

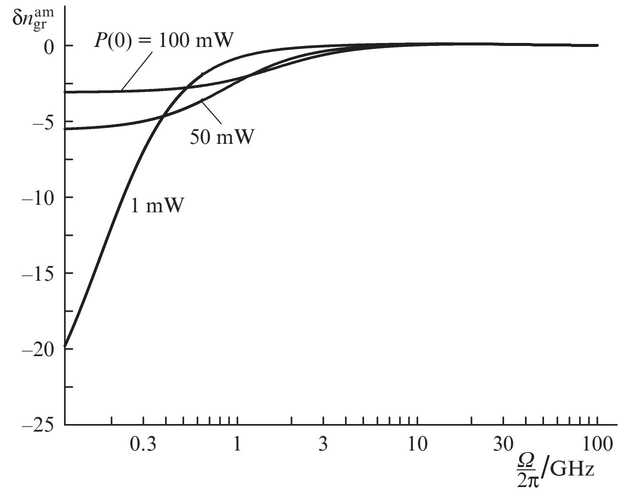

Figure 5. Frequency dependences of the additional dispersion term $\delta n_{\text{gr}}^{\text{am}}$ in the amplitude-modulated component of the signal at the amplifier output at different input signal powers $P(0)$ and a pump current $J = 2$ A.

$$\delta n_{\text{gr}}^{\text{am}} = -\frac{c}{L}\frac{\partial \arg(K(\Omega))}{\partial \Omega}. \tag{24}$$

If $\delta n_{\text{gr}}^{\text{am}}$ were constant in the frequency range of interest, $\Omega \leqslant \Omega_b$, and the condition of constant $|K(\Omega)|$ were fulfilled, it might be interpreted as an additional term in the group index n_{gr}. It is, however, seen from the curves in Fig. 5 that the change in the magnitude of $\delta n_{\text{gr}}^{\text{am}}$ across this frequency range exceeds n_{gr}: the heterostructures of interest for us typically have $n_{\text{gr}} \leqslant 4$. This points to distortion of the wave packet envelope such that $\delta n_{\text{gr}}^{\text{am}}$ loses the physical meaning of a change in the velocity of the wave packet.

According to (20), the amplitude of the converted phase-modulated signal, $B(t)$, is a weak function of the gain (output power) of the amplifier. Possible distortions of the signal shape can be characterised by the parameter

$$\delta n_{\text{gr}}^{\text{ph}} = \frac{c}{L}\frac{\partial \arg(K(\Omega) - 1)}{\partial \Omega}, \tag{25}$$

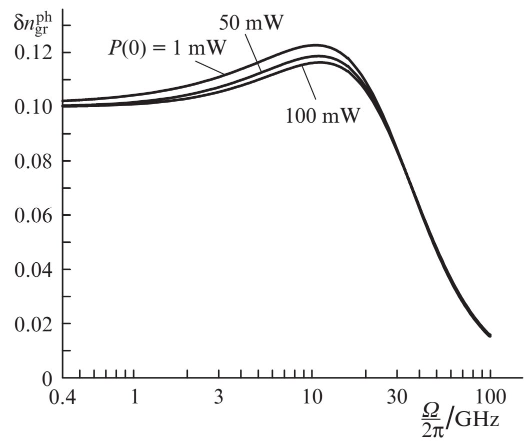

Figure 6. Frequency dependences of the additional term $\delta n_{\text{gr}}^{\text{ph}}$ in the group index of refraction of the phase-modulated component of the signal at the amplifier output at different input signal powers and a pump current $J = 2$ A.

which is shown in Fig. 6 as a function of frequency. It follows from the curves in Fig. 6 that, in this case, $\delta n_{\rm gr}^{\rm ph}$ can be viewed as an additional term in $n_{\rm gr}$ for the propagation of the wave packet of a phase-modulated signal. It is seen in Fig. 6, however, that the additional time delay due to $\delta n_{\rm gr}^{\rm ph}$ is substantially shorter than the duration of the packet, $1/\Omega_{\rm b}$. This means that, in analysing the dynamics of the propagation of a phase-modulated packet, the optical nonlinearity-induced dispersion can be neglected.

Figure 7 shows frequency dependences of $|K(\Omega)|$ at different pump currents of the amplifier. It is seen that the spectral band $\nu_{\rm b} = \Omega_{\rm b}/2\pi$ of the transformation and amplification of a phase-modulated signal at a pump current of e.g. 3 A may exceed 10 GHz. According to the data in Fig. 3, this corresponds to an output power of ~3 W. A favourable circumstance is that the spectral bandwidth increases with increasing pump current (output power).

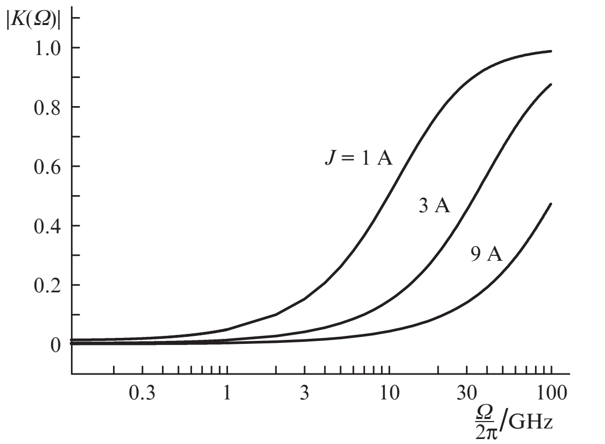

Figure 7. Magnitude of the transfer function of the amplifier, $|K(\Omega)|$, vs. frequency at different pump currents J, a signal power at the amplifier input $P(0) = 10$ mW and a background loss coefficient $\alpha = 0.5$ cm^{-1}.

3. Discussion

The above relations (4), (5) and (16)–(19) allow one to model the key characteristics of a diode optical power amplifier having uniform parameters along its optical axis. At the same time, they offer the possibility of analysing more complex optical schemes, e.g. a module with an integrated multisection structure where each section may have its own guidance properties and pump intensity. To this end, these relations should be used recurrently, with the output signal of each section being the input signal of the next section.

Clearly, the present results are valid within the approximations made. Parameters necessary for calculations can be determined in independent experiments, as e.g. in Ref. [17]. We hope that the approximations made in the proposed model do not degrade the accuracy that is limited by the scatter in the real parameters of amplifiers, such as the differential gain σ, carrier lifetime τ and background loss α. Because of this, we believe that the approximations used in the proposed model do not affect the adequacy of the calculation results.

An important point in this study is that we have taken into account the optical nonlinearity due to gain saturation and to the effect of inversion level on the refractive index of the semiconductor. To this end, we have taken advantage of the feature of this nonlinearity that there is coupling only between harmonics symmetric with respect to the carrier frequency, without mixing with harmonics at other frequencies. As a result, the complex, nonlinear system of equations for a broadband signal breaks down into independent systems, each consisting of two linear equations for harmonics symmetric with respect to the carrier frequency. In these equations, nonlinearity is represented by a parameter. This type of optical nonlinearity was first studied in Ref. [18] and was analysed later in other reports (see e.g. Refs [19–24]). Runge et al. [25] used a similar approach in the most general form to investigate the optical amplification of a signal in a nonlinear semiconductor with various characteristics.

The low sensitivity of the output power of the amplifier to the input optical beam power makes it a convenient and efficient source of bright light. Basically, the brightness of such a source is only limited by the output power of single-transverse-mode lasers, which is related to catastrophic optical damage (COD) or disturbance of single-mode operation. In the current stage of technology development, these limitations are ~3 W [26], which can be thought of as the world's best results so far.

The 'rigid' connection between the phase of the output beam in the amplifier and that of the input beam and the possibility of controlling it create conditions for coherent summation of the power of a large number of amplifiers and, thus, for creating light sources with kilowatt output powers, a near diffraction-limited divergence and the possibility of controlling the spatial light beam direction. It is worth pointing out that such an emitting system differs fundamentally from a system of N optically coupled diode lasers (see e.g. Refs [27, 28]), where the optical coupling increases the number of degrees of freedom by N times. Because of this, coherent summation of beams in a system of coupled diode lasers occurs only when one of the N supermodes is excited, which is, as a rule, extremely difficult to achieve, and the result is unstable. In an emitter made up of amplifiers, optical coupling between them is not used at all, and the possibility of such coupling because of imperfections of the optical system can be minimised by using optical isolators.

In the model used here, we did not take into account spontaneous emission, assuming that the input power of the optical beam being amplified was well above the effective power of spontaneous emission. This power can be estimated as described by Bogatov [29]. Such estimates suggest that the ~1 mW level for the operation of an amplifier with the parameters indicated in Table 1 satisfies the above condition.

The spectral range where the amplifier suppresses amplitude modulation ($|K(\Omega)| \ll 1$) is limited by $\Omega_{\rm b}$. For Ω above this limit, $|K(\Omega)|$ rises, approaching unity. Physically, this is due to the weaker interaction of the side components with the carrier. The weakening is caused by the reduction in inversion oscillation amplitude when the frequency of such oscillations (beating) considerably exceeds the inverse characteristic response time τ^{-1} of inversion, i.e. $\Omega\tau \gg 1$. The role of the above nonlinearity mechanism in a travelling wave amplifier then decreases, becoming comparable to the role of the spectral dispersion of the material gain, which is neglected in this study because $\Omega \ll \Delta\omega$, where $\Delta\omega$ is the material gain linewidth. The material gain dispersion is essential in analysis of single-frequency diode laser operation. This is a separate problem, identical to that considered previously [21, 24].

4. Conclusions

The present results demonstrate that an optical diode amplifier can serve as an output beam amplifier, ensuring multiwatt average output power at a near diffraction-limited beam divergence.

A distinctive feature of such an amplifier is that its output beam may have predominantly only phase modulation, independent of whether the input beam is amplitude- or phase-modulated. The gain band of a phase-modulated signal is only limited by the material gain bandwidth and may exceed several hundred gigahertz.

In the case of an amplitude-modulated signal, the amplifier serves in addition as an amplitude-to-phase modulation converter. In this case, however, its bandwidth is limited, but may exceed ~10 GHz. In this context, the most interesting application of such amplifiers is in systems that utilise bit phase shift keying (BPSK) modulation.

It seems likely that, in their application area, diode amplifier-based systems will be able to compete with fibre laser-based systems owing to their higher total efficiency and reliability and smaller dimensions. Moreover, a diode amplifier may serve as a building block of a high-power (kilowatt) laser system with a synthesised aperture of coherently summed optical beams and controlled angular position of the resultant optical beam.

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Appendix 1

We assume that the gain of the amplifier is deep in saturation and that the following inequalities are satisfied:

$$g \gg 1, \ \ln(u/u_0) \ll g\alpha L. \tag{A1.1}$$

Taking the logarithm of (5) and taking into account (A1.1), we obtain

$$g\alpha L = \ln\left(\frac{u}{u_0}\right) + (g+1)\ln\left(\frac{g-u_0}{g-u}\right) \approx g\ln\left(\frac{g-u_0}{g-u}\right). \tag{A1.2}$$

Thus,

$$\frac{g-u_0}{g-u} \approx \exp(\alpha L). \tag{A1.3}$$

Therefore,

$$u = u_0 \exp(-\alpha L) + g[1 - \exp(-\alpha L)], \tag{A1.4}$$

$$\frac{\mathrm{d}u}{\mathrm{d}g} = 1 - \exp(-\alpha L). \tag{A1.5}$$

It follows from (6)–(8) that

$$\frac{\mathrm{d}P}{\mathrm{d}u} = \frac{\hbar\omega_0 W d}{\sigma\tau}, \quad \frac{\mathrm{d}g}{\mathrm{d}J} = \frac{\Gamma\sigma\tau}{\alpha e W_0 d_{\mathrm{a}} L}, \tag{A1.6}$$

Finally, the derivative of the output power with respect to current has the form

$$\frac{\mathrm{d}P}{\mathrm{d}J} = \frac{\mathrm{d}P}{\mathrm{d}u}\frac{\mathrm{d}u}{\mathrm{d}g}\frac{\mathrm{d}g}{\mathrm{d}J} = \Gamma_{\mathrm{w}}\frac{\hbar\omega_0}{e}\frac{1-\exp(-\alpha L)}{\alpha L}.$$

Appendix 2

The field of an arbitrarily modulated signal propagating through an amplifier can be represented in the form

$$\mathcal{E}(z,t) = \frac{1}{2}\{E_0(z)\exp[\mathrm{i}(k_0 z - \omega_0 t)] \times [1 + \xi(z,t)] + \mathrm{c.c.}\}, \tag{A2.1}$$

where $E_0(z)$ and $\xi(z,t)$ are slowly varying functions of z. The complex-valued function $\xi(z,t)$, which determines the particular character of signal modulation, can be represented as the sum of a real and an imaginary function: $\xi(z,t) = a(z,t) + b(z,t)$, where

$$a(z,t) = \frac{1}{2}[\xi(z,t) + \xi(z,t)^*], \quad b(z,t) = \frac{1}{2}[\xi(z,t) - \xi(z,t)^*]. \tag{A2.2}$$

Further, we represent the functions $\xi(z,t)$, $a(z,t)$ and $b(z,t)$ as Fourier integrals:

$$\xi(z,t) = \int_{-\infty}^{\infty} \zeta(z,\Omega)\exp[\mathrm{i}\Omega(n_{\mathrm{gr}}z/c - t)]\,\mathrm{d}\Omega = \int_0^{\infty}\{V_{+1}(z,\Omega)\exp[\mathrm{i}\Omega(n_{\mathrm{gr}}z/c - t)] + V_{-1}(z,\Omega)\exp[-\mathrm{i}\Omega(n_{\mathrm{gr}}z/c - t)]\}\,\mathrm{d}\Omega,$$
$$a(z,t) = \int_{-\infty}^{\infty}\tilde{a}(z,\Omega)\exp[\mathrm{i}\Omega(n_{\mathrm{gr}}z/c - t)]\,\mathrm{d}\Omega,$$
$$b(z,t) = \int_{-\infty}^{\infty}\tilde{b}(z,\Omega)\exp[\mathrm{i}\Omega(n_{\mathrm{gr}}z/c - t)]\,\mathrm{d}\Omega. \tag{A2.3}$$

It follows from the definition of the functions $a(z,t)$ and $b(z,t)$ (A2.2) that their spectral components satisfy the relations

$$\tilde{a}(z,\Omega) = \tilde{a}^*(z,-\Omega), \ \tilde{b}(z,\Omega) = -\tilde{b}^*(z,-\Omega). \tag{A2.4}$$

For $\Omega \geqslant 0$, we have

$$\tilde{a}(z,\Omega) = \frac{1}{2}[V_{+1}(z,\Omega) + V_{-1}^*(z,\Omega)], \quad \tilde{b}(z,\Omega) = \frac{1}{2}[V_{+1}(z,\Omega) - V_{-1}^*(z,\Omega)]. \tag{A2.5}$$

Since relations (11)–(13) are valid for all the pairs $V_{\pm1}(z,\Omega)$ in (A2.5), we obtain for $\tilde{a}(z,\Omega)$ and $\tilde{b}(z,\Omega)$:

$$\tilde{a}(z,\Omega) = K(z,\Omega)\tilde{a}(0,\Omega), \quad \tilde{b}(z,\Omega) = \tilde{b}(0,\Omega) + \mathrm{i}R[1 - K(z,\Omega)]\tilde{a}(0,\Omega). \tag{A2.6}$$

Then, according to (A2.3) we have

$$a(z,t) = \int_{-\infty}^{\infty}\tilde{a}(0,\Omega)K(u(z),\Omega)\exp(-\mathrm{i}\Omega t')\,\mathrm{d}\Omega, \quad b(z,t) = \int_{-\infty}^{\infty}\tilde{b}(0,\Omega)\exp(-\mathrm{i}\Omega t')\,\mathrm{d}\Omega \tag{A2.7}$$

$$+\int_{-\infty}^{\infty} \mathrm{i}R[1-K(z,\Omega)]\tilde{a}(0,\Omega)\exp(-\mathrm{i}\Omega t')\,\mathrm{d}\Omega,$$

or

$$b(z,t) = b(0,t') + \mathrm{i}R[a(0,t') - a(z,t)]. \tag{A2.8}$$

Here, $t' = t - z/v_{\rm gr}$ ($v_{\rm gr} = c/n^*$). Finally, denoting $a(0,t) \equiv a(t)$, $a(L,t) \equiv A(t)$, $a(\Omega) \equiv \tilde{a}(0,\Omega)$, $b(0,t) \equiv b(t)$ and $b(L,t) \equiv B(t)$ we obtain

$$A(t) = \int_{-\infty}^{\infty} a(\Omega)K(u(L),\Omega)\exp(-\mathrm{i}\Omega t')\,\mathrm{d}\Omega,$$

$$B(t) = b(t') + \mathrm{i}R[a(t') - A(t)].$$

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