Dynamic Theory of Picosecond Optical Pulse Shaping by Gain-Switched Semiconductor Laser Amplifiers

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Abstract—**A dynamic theory of semiconductor laser amplifiers is developed which takes into account the coherent time-dependent amplification of an incident optical pulse as well as the nonlinear dynamics of the semiconductor laser when driven by an unbiased injection current pulse. For suitable time delays between the optical and the electrical pulse, a strongly nonlinear self-induced shortening of the emitted laser pulse is predicted.**

Introduction

THERE is currently much interest in the generation and amplification of short optical pulses for the purposes of optical signal processing and optical communication systems. Unbiased gain switching by driving a semiconductor laser with an electrical injection current pulse of a few hundred picoseconds width and a current maximum of several times its CW threshold is a simple and reliable method of producing stable pulses of less than 15 ps full width at half maximum (FWHM) in the 0.8–1.3 μm wavelength range [1].

In this paper, it is shown theoretically that by operating an unbiased gain-switched semiconductor laser as a dynamic optical amplifier, one can obtain even shorter optical pulses. Experimentally, it has been demonstrated in a cross-correlation arrangement that a conventional gain-switched semiconductor laser acts as a high-speed optical gate which can be used to detect an optical test pulse from a second laser with a time resolution better than 10 ps when the integral optical output is measured by a slow, integrating photodiode as a function of the delay time between the optical and the electrical pulse [2]. A theoretical understanding of these phenomena requires modeling of the internal nonlinear dynamics of the semiconductor laser as well as the coherent amplification of the incident signal. Previous theories were mainly restricted to either of these two aspects. The incoherent laser dynamics is commonly described by rate equations for the photon and carrier densities inside the laser medium [3]. This approach can explain the strongly nonlinear relaxation oscillations emitted by unbiased gain-switched semiconductor lasers [4], [5], in very good agreement with experiments. Incoherent optical amplification is described within this framework by adding the optical pumping rate, i.e., the number of injected photons per unit time and unit volume, to the photon rate equation [6], [7]. Hereby, interference effects are ignored. The other approach is basically via a static active Fabry–Perot theory, assuming time independence of injection currents and incident optical signals [8]–[10].

In this paper, a new simple analytical model is developed which combines the nonlinear transient dynamics of the semiconductor laser driven by an injection current pulse with the time-dependent coherent amplification of an incident optical pulse.

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The Dynamic Model

The signal wave must be treated coherently since its transmission through the laser amplifier depends strongly upon its phase due to reflection at the end facets and interference between forward and backward propagating waves. It will be described by electric field amplitudes $E^+(z, t)$ and $E^-(z, t)$ traveling in the positive and negative z direction of the laser resonator, respectively. (E^+ and E^- are normalized to the dimensions of $(\text{length})^{-3/2}$, such that $|E^\pm|^2$ denote photon densities). All the remaining laser modes excluding the signal mode may be treated incoherently and described by a spatially averaged photon density $N(t)$; these determine the free-laser oscillation as a result of amplified spontaneous emission. It has been shown [9] that working with photon densities, i.e., intensities rather than field amplitudes, is equivalent to averaging the results over the photon energy between two cavity resonances. The use of axially averaged photon densities rather than z-dependent forward and backward propagating photon densities is a reasonable approximation for facet reflectivities larger than 20 percent [3].

The dynamics of the carriers can be described by a simple rate equation for the electron concentration n if the phase diffusion time of the electrons (given by the intraband collision time) is much shorter than the recombination time, i.e., if the electronic polarization [11], [12] relaxes faster than the electronic inversion and if lateral

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diffusion of carriers is ignored

$$\frac{dn}{dt} = \eta \frac{J(t)}{ed} - R_{sp}(n) - g(n)\left(N + \left|E^+(z, t)\right|^2 + \left|E^-(z, t)\right|^2\right). \quad (1)$$

Here $J(t)$ is the given time-dependent injection current density, η is the electron injection efficiency factor, d is the thickness of the laser-active region, $R_{sp}(n)$ is the spontaneous emission rate governed by monomolecular, bimolecular, or Auger-type recombination, and $g(n)$ is the modal gain function (assumed to be the same in all modes, including the signal mode, for technical simplicity, but not for principal reasons).

Apart from the standard rotating wave and slowly varying amplitude approximations [13], our main assumption is now that the electron concentration $n(t)$ varies little during the single-pass time τ of the signal, which is typically on the order of 4 ps, e.g., for a cavity length of $L = 300\ \mu$m and an optical group index of $n_g = 4$. This approximation seems appropriate since the electron lifetime is several orders of magnitude longer than the photon lifetime, and hence the electrons generally respond more slowly than the photons. Note that we still allow for changes in $n(t)$ on a slower time scale from one single-pass to the next one. We can then approximate the z-dependent signal intensity in (1) by the axially averaged traveling signal photon density

$$\overline{S}(t) = \frac{1}{L}\int_0^L \left\{\left|E^+\left(z, t - \frac{L - z}{v_g}\right)\right|^2 + \left|E^-\left(z, t + \frac{L - z}{v_g}\right)\right|^2\right\} dz, \quad (2)$$

with similar arguments as given by Adams *et al.* [9] in the case of time-independent amplification of CW signals. We obtain the following coupled equations, which form the basis of our model:

$$\frac{dn}{dt} = \eta \frac{J(t)}{ed} - R_{sp}(n) - g(n)(N + \overline{S}) \quad (3)$$

$$\frac{dN}{dt} = (\Gamma g(n) - \kappa)N + \beta R_{sp}(n) \quad (4)$$

$$\frac{\partial E^\pm}{\partial t} \pm v_g \frac{\partial E^\pm}{\partial z} = \frac{1}{2}(\Gamma g(n) - \alpha)E^\pm - ikv_g E^\pm. \quad (5)$$

Here Γ is the optical confinement factor, β is the spontaneous emission factor, $v_g \equiv c/n_g$ is the group velocity where c is the vacuum velocity of light and n_g is the group index, α is the optical loss constant for absorption and scattering in the active and cladding regions, κ is the total inverse photon lifetime including absorption, scattering, and mirror losses, and k is the wave vector of the signal carrier wave in the semiconductor. The left-hand side of (5) is the total time derivative as seen by an observer moving with the traveling wave [8], and the factor 1/2 on the right-hand side of (5) accounts for the net *amplitude* gain [8], [9]. Equation (5) has to be supplemented by boundary conditions for the crystal facets at $z = 0$ and $z = L$:

$$E^+(0, t) = t_1 E_{in}(t) + r_1 E^-(0, t) \quad (6a)$$

$$E^-(L, t) = r_2 E^+(L, t) \quad (6b)$$

$$E_{out}(t) = t_2 E^+(L, t). \quad (6c)$$

Here t_1, t_2, and r_1, r_2 are the amplitude transmission and reflection coefficients of the two facets, respectively; they are related to the reflectivities R_1, R_2 by

$$R_i = r_i^2 = 1 - t_i^2 \qquad (i = 1, 2). \quad (7)$$

$E_{in}(t)$ and $E_{out}(t)$ are the incident and the outgoing signal field amplitudes at the facets normalized to the dimensions of (length)$^{-3/2}$.

Under our basic approximation of averaged $n(t)$, (5) can be integrated from $z = 0$ to $z = L$ for the forward or backward propagating signal wave, yielding the single-pass intensity gain

$$G_s(t) \equiv \left|\frac{E^+(L, t)}{E^+(0, t - \tau)}\right|^2 = \left|\frac{E^-(0, t)}{E^-(L, t - \tau)}\right|^2 = \exp\left\{\left[\Gamma g(n(t)) - \alpha\right]L/v_g\right\} \quad (8)$$

where $\tau = L/v_g$ is the single-pass time. The total outgoing signal amplitude $E_{out}(t)$ at a given time t is obtained by summing over all forward and backward propagating waves with appropriate time delays and phase factors, using (6):

$$E_{out}(t) = t_1 t_2 e^{-ikL}\left\{\sum_{m=0}^{\infty} E_{in}(t - (2m + 1)\tau) \times (r_1 r_2)^m \prod_{l=0}^{2m} \left[G_s(t - l\tau)\right]^{1/2} \exp(-2mikL)\right\} \quad (9)$$

The corresponding output intensity (normalized to the dimension of a photon density) is

$$\left|E_{out}(t)\right|^2 = t_1^2 t_2^2 G_s(t)\, \Sigma(t) \quad (10)$$

where

$$\Sigma(t) \equiv \sum_{m=0}^{\infty} \sum_{m'=m+1}^{\infty} E_{in}^*(t - (2m + 1)\tau) \cdot E_{in}(t - (2m' + 1)\tau)(r_1 r_2)^{m+m'} \cdot \left[\prod_{l=1}^{2m} G_s(t - l\tau) \prod_{l'=1}^{2m'} G_s(t - l'\tau)\right]^{1/2} \cdot 2\cos\left[2(m' - m)kL\right] + \sum_{m=0}^{\infty} \left|E_{in}(t - (2m + 1)\tau)\right|^2 (r_1 r_2)^{2m} \cdot \prod_{l=1}^{2m} G_s(t - l\tau) \quad (11)$$

contains all delay and interference terms.

The signal field distribution $E^{\pm}(z, t)$ inside the cavity can be obtained most easily by integration of (5) from z to L using the boundary conditions (6b), (6c), and (9):

$$E^{+}\left(z, t - \frac{L-z}{v_g}\right) = t_2^{-1} \exp\left\{-\tilde{g}(L-z) + ik(L-z)\right\} E_{\text{out}}(t)$$

$$E^{-}\left(z, t + \frac{L-z}{v_g}\right) = t_2^{-1} r_2 \exp\left\{\tilde{g}(L-z) - ik(L-z)\right\} E_{\text{out}}(t) \quad (12)$$

with

$$\tilde{g} \equiv (\Gamma g(n) - \alpha)/(2v_g).$$

The spatially averaged signal intensity (photon density) $\bar{S}(t)$ follows from (2), (12), and (10):

$$\bar{S}(t) = (G_s(t) - 1)(1 + r_2^2 G_s(t)) \cdot t_1^2 \Sigma(t)/\ln G_s(t). \quad (13)$$

Equations (9)–(13) are our main analytical results. In particular, (9) takes account of time-dependent multiple-beam interference. In the special case of a time-independent input signal E_{in} and a time-independent single-pass gain G_s, (9) reduces to the familiar result of an active Fabry–Perot cavity:

$$E_{\text{out}} = E_{\text{in}} \frac{t_1 t_2 G_s^{1/2}}{1 - r_1 r_2 G_s \exp(-2ikL)}, \quad (14)$$

which leads to the total amplifier intensity gain [7], [9]:

$$\left|\frac{E_{\text{out}}}{E_{\text{in}}}\right|^2 = \frac{(1-R_1)(1-R_2)G_s}{(1-\sqrt{R_1 R_2}G_s)^2 + 4\sqrt{R_1 R_2}G_s \sin^2(k-k_r)L} \quad (15)$$

where $k_r = \pi r/L$ ($r \in \mathrm{IN}$) corresponds to the cavity resonances. G_s is given by (8) where n is determined from (3) in the steady state.

Another limiting case, which occurs if the signal input pulse is short compared to the cavity round-trip time and G_s is time independent, is also contained in (9) and agrees with the results of Lenth [14]. The output signal then consists of a sequence of multiply reflected pulses which do not interfere with each other. The integral output intensity for $\int |E_{\text{in}}(t)|^2\, dt \equiv E_0$ is

$$\int_{-\infty}^{\infty} |E_{\text{out}}(t)|^2\, dt = E_0 (t_1 t_2)^2 G_s \sum_{m=0}^{\infty} (r_1 r_2 G_s)^{2m} = E_0 \frac{(1-R_1)(1-R_2)G_s}{1-(\sqrt{R_1 R_2}G_s)^2}. \quad (16)$$

Our result (9) generalizes (14) and (16) to a wide range of experimental conditions with both a time-dependent input signal and a time-dependent modal gain of the amplifier. In the general case, the infinite series (9) cannot be through $n(t)$ in (8) must be determined self-consistently from (3)–(5). To this purpose, the rate equations (3) and (4) with (13) must be solved numerically. They constitute a system of delay differential equations due to the occurrence of $n(t-\tau)$, $n(t-2\tau)$, $n(t-3\tau)$, $\cdots$ in $\bar{S}(t)$ by (11). Starting from an initial thermal equilibrium value $n(0)$ and $N(0)$ before the injection current pulse sets in, one can in principle integrate (3) and (4) for arbitrary injection current pulses $J(t)$ and incident signals $E_{\text{in}}(t)$.

The numerical effort can be greatly reduced by a further approximation which lets $\bar{S}(t)$ vary in discrete time steps T only. The rate equations (3) and (4) governing the internal laser dynamics are integrated with constant $\bar{S}(t)$ from the time t to $t+T$, yielding a new carrier density $n(t+T)$ and, by (8) and (13), a new signal photon density $\bar{S}(t+T)$, which is used in (3) and (4) during the next time step T. This procedure represents a reasonable approximation if, in addition to slowly varying $n(t)$, the incident signal $E_{\text{in}}(t)$ changes slowly during a time T.

Numerical Results

Without incident signal ($\bar{S} = 0$), (3) and (4) are the familiar semiconductor laser rate equations, which give rise to damped relaxation oscillations as a result of amplified spontaneous emission when driven with a suitable injection current pulse $J(t)$ [5]. In the following numerical calculations, $J(t)$ is modeled by an asymmetric Gaussian in order to allow for different rise and fall times t_r and t_f, respectively:

$$J(t) = \begin{cases} J_0 \exp\left\{-[(t-t_0)/t_r]^2\right\} & \text{for } t < t_0 \\ J_0 \exp\left\{-[(t-t_0)/t_f]^2\right\} & \text{for } t \geq t_0. \end{cases} \quad (17)$$

This closely approximates experimentally used current pulse shapes [1]. There exists a minimum value $J_0^{\min}$ of the peak current density J_0, considerably larger than the CW laser threshold $J_t = R_{\text{sp}}(n_t)ed/\eta$, $g(n_t) = \kappa/\Gamma$, below which no laser emission occurs. This effective dynamic threshold is pulse-shape dependent and is typically several times J_t for injection pulsewidths of a few hundred picoseconds. It is higher for shorter $t_r + t_f$. For sufficiently large $J_0 > J_0^{\min}$, more than one relaxation oscillation may be emitted. With increasing ratio $j_0 \equiv J_0/J_t$, the FWHM of the first relaxation oscillation decreases, the time difference between subsequent relaxation oscillations decreases, the peak optical power increases, and the delay time before the onset of the first oscillation decreases [5]. In the following, J_0 is optimized in the sense that one and only one relaxation oscillation is emitted.

So far, the spontaneous recombination rate and the modal gain have not been specified. While our general model allows for more realistic and sophisticated forms of $R_{\text{sp}}(n)$ [15] and $g(n)$ [16], [17], including the effects

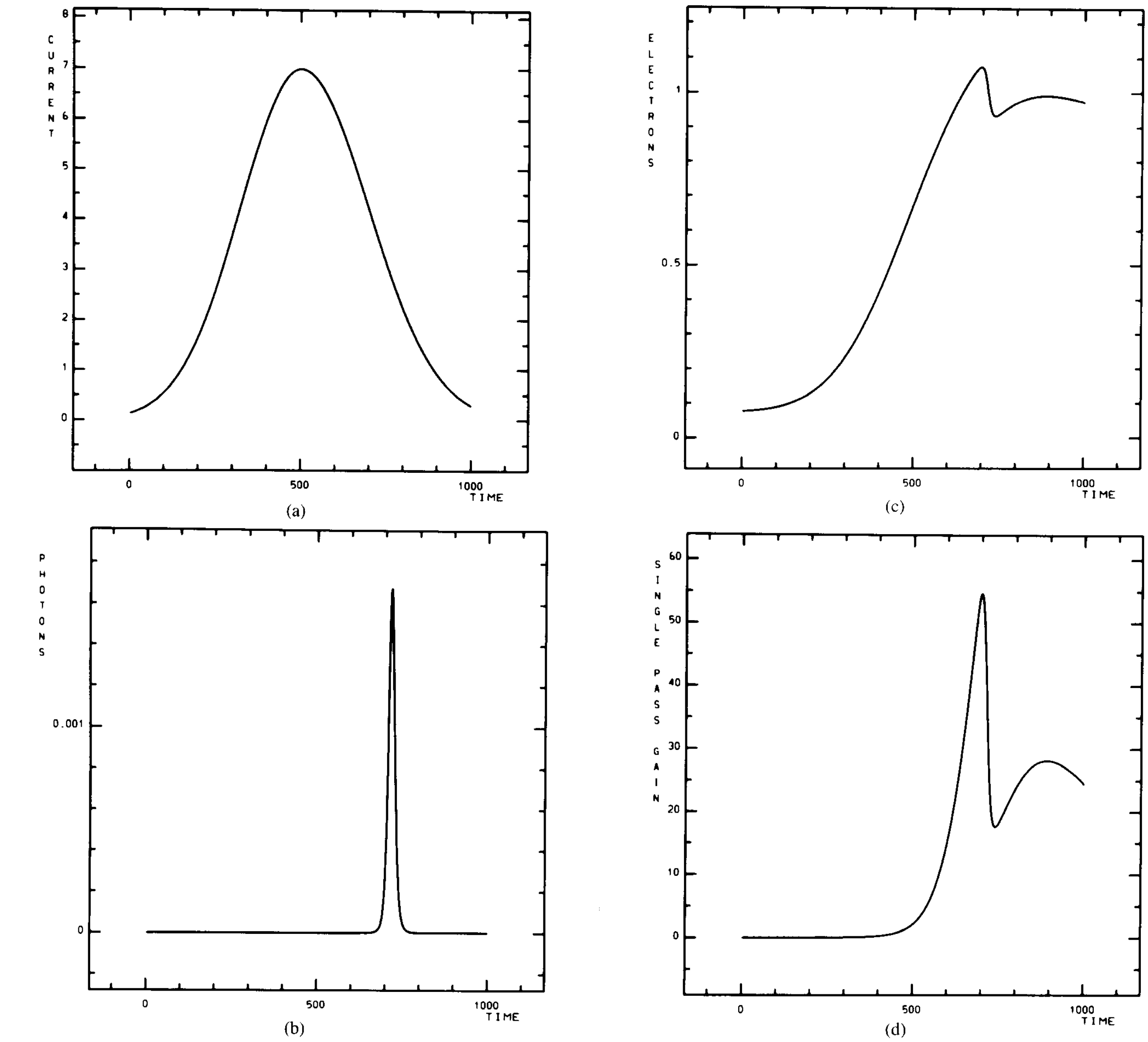

Fig. 1. Numerical solution of the laser rate equations without incident optical signal ($E_{in} = 0$). (a) The injection current density J/J_t plotted versus time in picoseconds. (b) The photon density N/n_t plotted versus time in picoseconds. (c) The electron density n/n_t plotted versus time in picoseconds. (d) The single-pass gain G_s plotted versus time in picoseconds. J_t and n_t are the respective CW threshold values of J and n. The material parameters are $B = 1.6 \times 10^{-10}$ cm^3/s, $g_0 = 4 \times 10^{-6}$ cm^3/s, $n_0 = 0.5 n_t$, $\kappa = 1$ ps^{-1}, $\Gamma = 0.2$, $\beta = 10^{-4}$, which gives $n_t = 2.5 \times 10^{18}$ cm^{-3}. The initial values are $N(0) = 0$, $n(0) \equiv N_D = 2 \times 10^{17}$ cm^{-3}.

simple functions

$$R_{sp}(n) = Bn^2 \tag{18}$$

$$g(n) = g_0 \cdot (n - n_0) \tag{19}$$

where n_0 denotes the transparency concentration and B, g_0 are constants. For typical material parameters and an injection current pulse with $j_0 = 7$, $t_r = 250$ ps, $t_f = 280$ ps, $t_0 = 500$ ps, the numerical results are shown in Fig. 1. A single relaxation oscillation is emitted [Fig. 1(b)]. Note that the single-pass gain $G_s(t)$, which is related to the exponential of $n(t)$ by (8), (19), represents an optical gate which very slowly regresses to its equilibrium value close to zero due to the slow decay of the carrier density to equilibrium after the optical pulse is emitted. There is amplification, i.e., $G_s(t) > 1$, even well below the CW threshold electron density $n_t = n_0 + \kappa/(\Gamma g_0)$, as long as

$$n > n_0 + \alpha/(\Gamma g_0). \tag{20}$$

Fig. 2 shows the numerical solution of (3), (4), (8), and (13) with the same material and excitation parameters for a rectangular incident optical signal of 100 ps width, an intensity three orders of magnitude smaller than that of

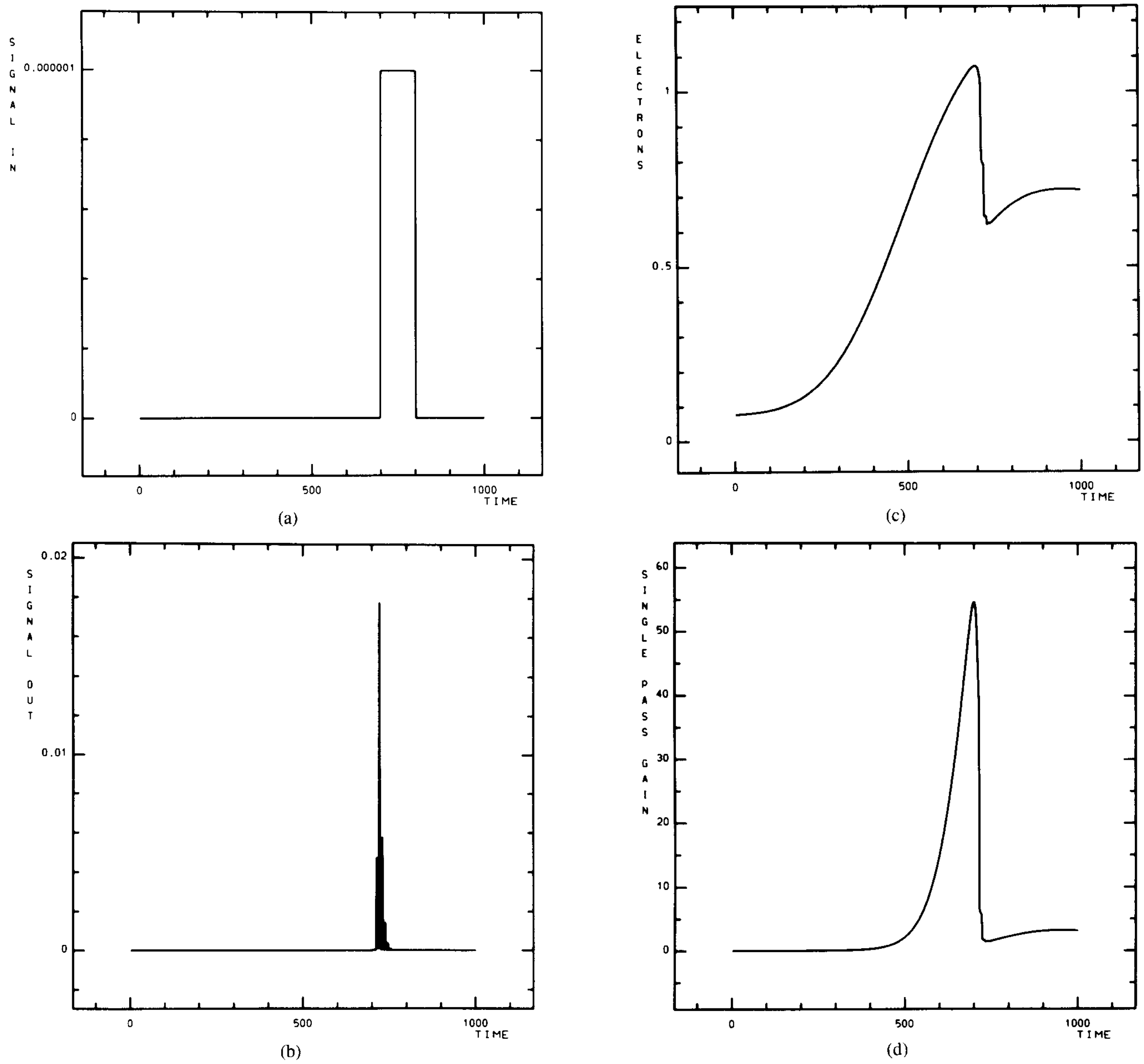

Fig. 2. Numerical solution of the amplifier dynamics. (a) The given incident signal photon density $|E_{\text{in}}|^2$ in units of n_t is plotted versus time in picoseconds. (b) The resulting emitted signal photon density $|E_{\text{out}}|^2/n_t$ plotted versus time in picoseconds. (c) The electron density n/n_t plotted versus time in picoseconds. (d) The single-pass gain G_s plotted versus time in picoseconds. The injection current pulse and the material parameters are the same as in Fig. 1; additionally, $\alpha = 0.15\ \text{ps}^{-1}$, $R_1 = R_2 = 0.33$, $L = 300\ \mu\text{m}$, $n_g = 4$, $k = 2\pi n_g/\lambda_0$ with $\lambda_0 = 0.8\ \mu\text{m}$ has been used. The corresponding single-pass time is $\tau = 4$ ps. $\overline{S}(t)$ has been varied in discrete time steps $T = 2$ ps.

the spontaneously emitted pulse of Fig. 1(b), and injected at about the time when the electron density is maximum. The emitted signal [Fig. 2(b)] is amplified by a factor of about 20 000 and drastically reduced in width well below 10 ps. This can be understood by noting the sharp reduction in n [Fig. 2(c)] after a few round trips of the incident signal. This is the result of strongly enhanced stimulated emission due to the amplified signal $\overline{S}(t)$, which in turn decreases the electron density by (3). Thus, the leading edge of the signal encounters a large amplification factor G_s, but immediately induces a sharp decrease of n and thus depletes the gain G_s, which leads to a shortening of the optical gate and the subsequent suppression of the signal field. This nonlinear negative feedback represents a simple and efficient mechanism of self-induced pulse shortening within the active laser cavity. It is reminiscent of active mode-locking [18], but differs essentially in that it generates single output pulses rather than a continuous pulse train. It requires only relatively broad optical and electrical pulses as input signal and driving injection cur-

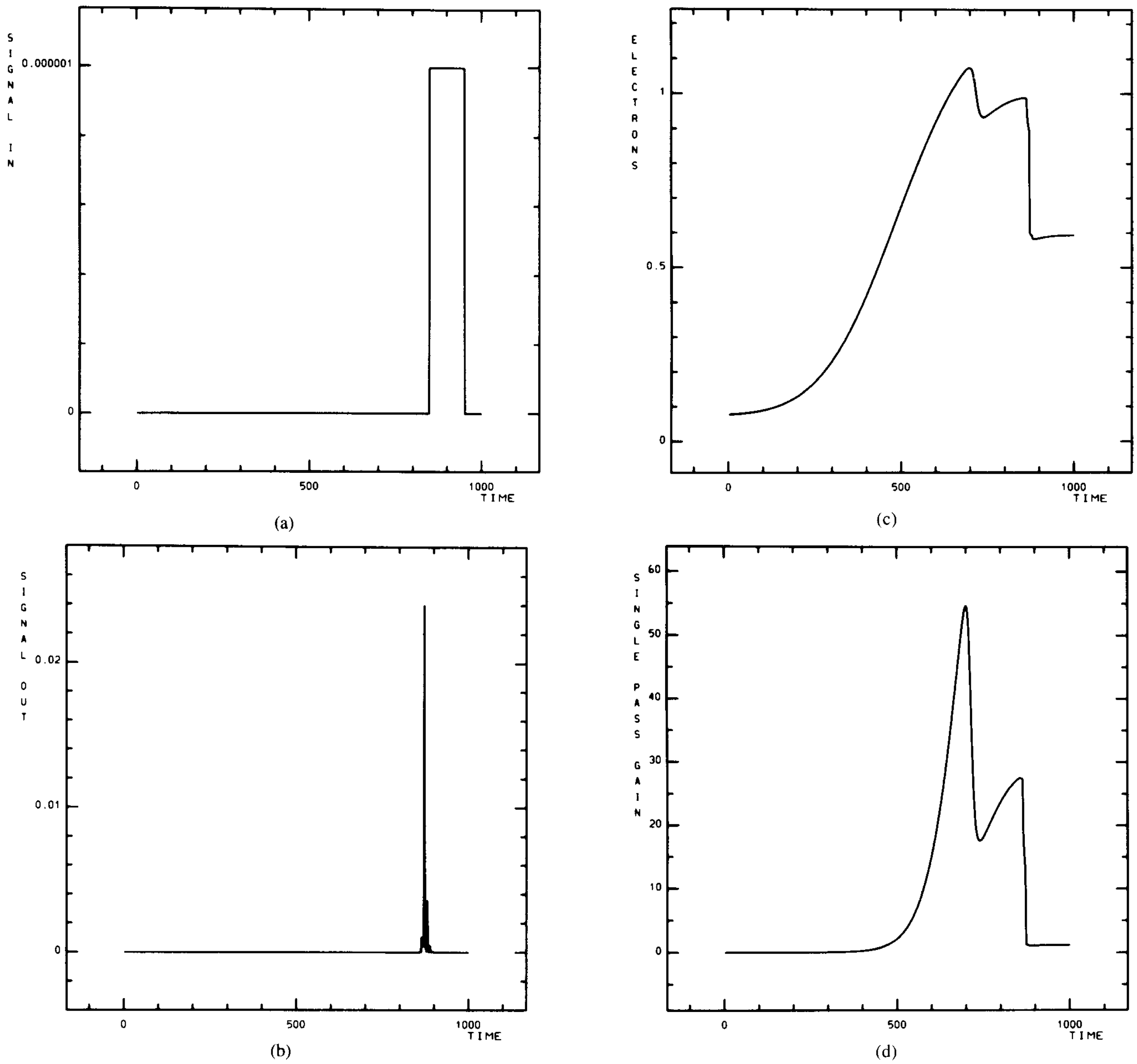

Fig. 3. Same as in Fig. 2, but with a different time delay between the injection current pulse $J(t)$ [Fig. 1(a)] and the incident optical signal $|E_{in}(t)|^2$ (a).

rent, respectively. Also, the resulting output pulse is relatively insensitive to the precise time delay τ_D between the optical and the electrical pulse as long as the optical signal coincides roughly with the broad maximum of $n(t)$. This has been checked numerically by varying τ_D. If the onset of the incident signal is shifted past the maximum of $n(t)$, the single-pass gain is smaller, but still greater than unity as long as (20) holds since $n(t)$ decays slowly. It takes longer for the signal wave to build up, and hence the self-induced gain depletion is not as pronounced, and the output signal is slightly broader (Fig. 3). If the input signal arrives well before the electron density has reached its maximum (Fig. 4), the single-pass gain G_s is too small to amplify the signal sufficiently fast in order to reduce G_s below unity. Rather, the remaining amplification is sufficient to initiate another buildup of the signal wave, and a second signal spike, and possibly subsequent ones, are emitted. Thus, a sequence of broad, damped relaxation oscillations is emitted [Fig. 4(b)].

Variation of the length and shape of the incident signal $E_{in}(t)$ yields little change in the emitted signal since only the leading edge of the incident pulse essentially contributes to the amplifier output $E_{out}(t)$. The material parameters that influence the FWHM of the output signal most heavily are the gain coefficient g_0 and the loss coefficient α. In particular, g_0 determines the time scale on which the signal field inside the amplifier builds up and on which the enhanced stimulated emission reduces $n(t)$.

Conclusions

Our proposed dynamic optical amplifier model predicts self-induced pulse shortening due to rapid gain depletion

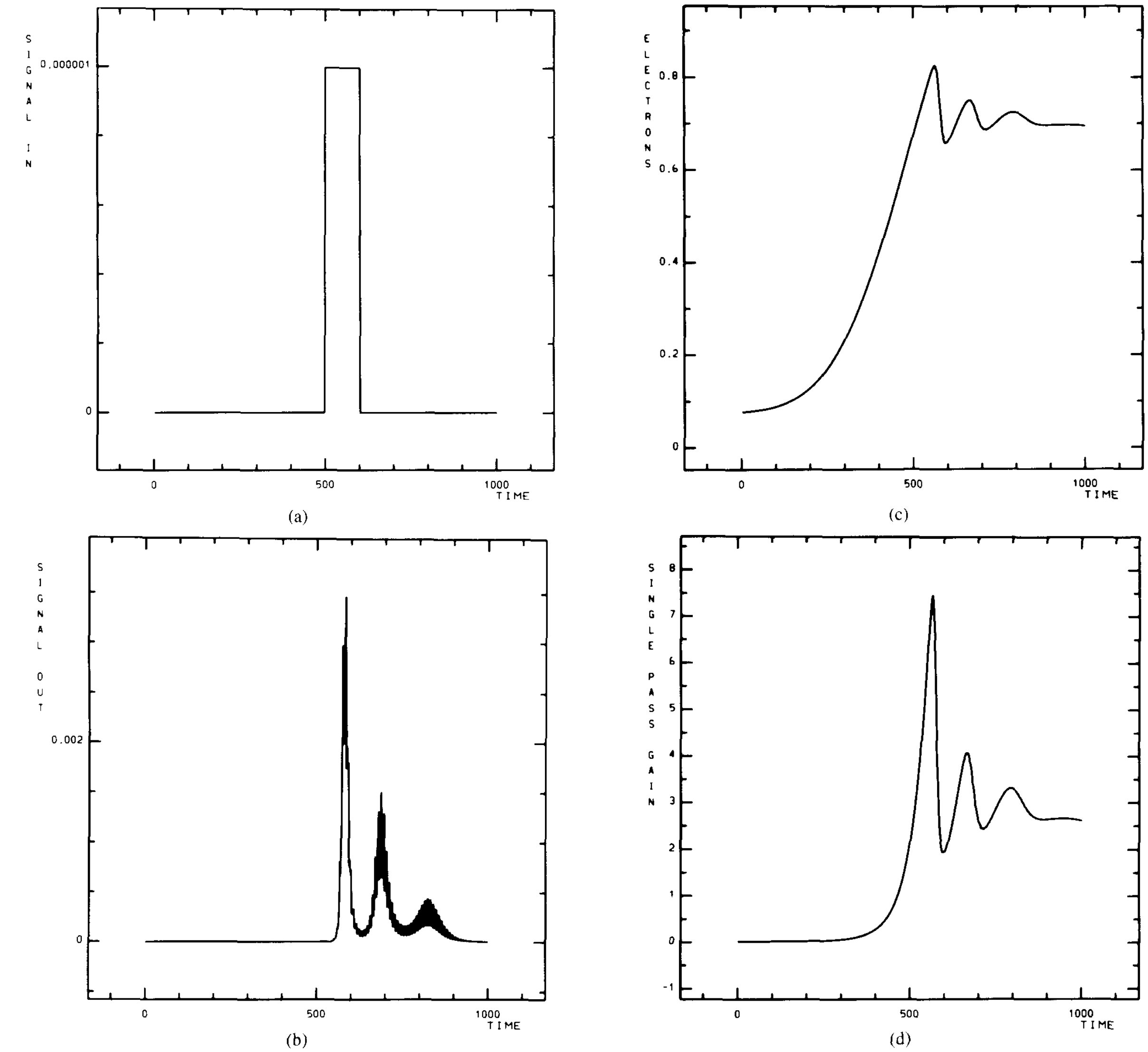

Fig. 4. Same as in Fig. 2, but with a different time delay between $J(t)$ and $|E_{in}(t)|^2$.

by the incident signal. This suggests a practical application as a very simple and inexpensive method of producing short single optical pulses. A possible experimental realization is sketched schematically in Fig. 5. It consists of an electrical pulse generator, which drives two semiconductor laser diodes l_1 and l_2 via a power divider and an electrical delay in such a way that they both emit single optical pulses under free-running conditions. The pulse emitted by l_1 is focused, possibly via an adjustable optical delay, onto l_2, which operates as a dynamic amplifier.

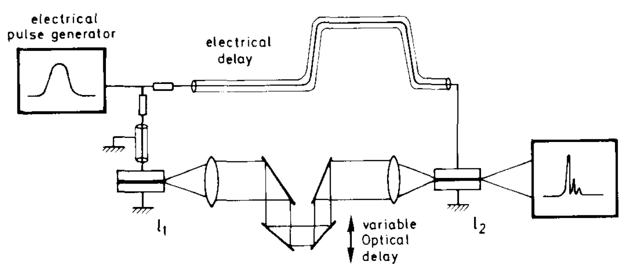

Fig. 5. Schematic experimental setup for generation of ultrashort optical pulses by self-induced pulse shaping. The pulse emitted by the laser l_1 is focused onto the laser l_2, which operates as an amplifier.

The effect of pulse shortening by gain depletion heavily relies upon the *transient* dynamics of the laser amplifier driven by a pulsed injection current since under steady-state conditions the electron concentration, and hence the gain, saturates at the CW laser threshold. An incoherent dynamic theory which treats the amplified signal on the basis of rate equations for the photon density with an external optical pumping rate [7] would also not appropriately take into account the gain depletion due to enhanced stimulated emission. Moreover, it would not explain the fine structure of the emitted optical pulse due to multiple reflections at the facets [cf. Figs. 2(b), 3(b), 4(b)]. Our theory can explain the observation [2], [15] that the dynamic gain decays much faster than given by the spontaneous recombination rate. Although our approximation of time-averaged n during a single pass tends to overestimate the self-induced pulse shortening, stable optical pulses well below 10 ps seem to be feasible.

A different application has been realized in cross-correlation experiments where the amplifier acts as an optical sampling gate, and the time integral of the optical output intensity $\int |E_{\text{out}}(t)|^2 \, dt$ is recorded as a function of the delay time between the injection current pulses of the two lasers [2]. In order to obtain a narrow optical gate, it is desirable that the incident optical signal have very little influence upon the internal dynamics of the amplifier; i.e., the single-pass gain is modified only slightly by the signal. This can be achieved by appropriate attenuation of the signal.

Possible future applications of the theory could include a feedback of the emitted optical pulse $E_{\text{out}}(t)$ onto the semiconductor laser amplifier itself through an optical fiber in an experimental setup similar to that of the soliton laser [19], [20]. In such a feedback loop, only one laser is needed.

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