Picosecond relaxation processes in a semiconductor laser excited by a powerful ultrashort light pulse

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A method is proposed for describing the dynamics of laser emission and of an electron–hole plasma in a semiconductor laser with strongly heated nonequilibrium carriers. The method is based on perturbation theory in the small parameter τ_p/τ_T, where τ_p is the photon lifetime in the optical cavity and τ_T is the characteristic plasma-cooling time. Relaxation is investigated in a semiconductor laser excited by a pico- or subpicosecond light pulse. It is shown that the plasma cooling and the damping of the long-wave part of the laser spectrum are related so that the rates of these processes, as well as the rate of the stimulated recombination of the carriers, are determined by a single characteristic time τ_T. Analytic expressions are obtained for τ_T and for the laser intensity. It is shown that τ_T can be much longer than the time of carrier-energy relaxation by an optical phonon, a time that determines the rate of plasma cooling in the absence of lasing.

Interest in linear optical phenomena produced when intense optical fluxes interact with semiconductors, together with the goal of simultaneously increasing the operating speeds and powers of semiconductor devices, makes it necessary to study the relaxation times of an electron–hole plasma under conditions far from thermodynamic equilibrium.

The relaxation times of a plasma and of stimulated (laser) emission have been recently investigated by optical excitation of semiconductors and semiconductor lasers with powerful ultrashort light pulses. In this method the exciting light pulse is produced in the semiconductor by a hot (up to several thousand degrees) electron–hole plasma with carrier density 10^{17}–10^{19} cm^{-3}, which cools off subsequently mainly through generation of optical phonons. The relaxation times in semiconductors were determined from the change of the optical transparency and of the luminescence spectrum of the hot carriers.[1,2] It was observed that a dense plasma has a cooling rate several times larger than that calculated with the equation for the plasma energy loss in the Fröhlich interaction between carriers and optical phonons. The possible cause of the slower cooling rate was assumed to be screening of the electron–phonon interaction[3–5] and optical-phonon heating,[6,7] although to our knowledge no direct experimental data to confirm the decisive roles of these mechanisms are available.

The relaxation times in experiments with semiconductor lasers were determined by measuring the damping of the laser emission associated with the cooling of an electron-hole plasma excited by a strong subpicosecond light pulse.[8,9] In Ref. 9 was used a GaAs laser with an optical cavity having a short photon lifetime $\tau_p \approx 1$ ps and a cavity length $L = 200$ μm. Under these conditions, lasing occurred in a wide wavelength range. The measured damping time of the long-wave part of the lasing spectrum reached 10–20 ps, which exceeds τ_p by an order of magnitude. It was concluded in Ref. 8 that the duration of the damping of the laser light was governed by the plasma cooling rate. It might seem that so slow a plasma cooling in a semiconductor laser can also be attributed to screening of the carrier–phonon interaction or to heating of optical phonons. Another slowing-down mechanism, however, namely plasma cooling by predominant departure of "cold" carriers in stimulated emission,[1)] was pointed out in Ref. 8 in the course of a numerical simulation of the plasma dynamics and of the lasing. Note that this mechanism may be important also for semiconductors (i.e., in the absence of an optical cavity), since recent experiments on semiconductor excitation by ultrashort light pulses[11–13] have revealed an intense emission accompanying the cooling of the electron–hole plasma.

Three mechanisms have thus been proposed to explain the longer cooling time of the high-density plasma of a semiconductor and of a semiconductor laser. Two were investigated analytically,[4–7] while the stimulated-emission mechanism was studied only by numerical simulation, which is patently insufficient for comparison with other mechanisms and with experimental data.

Our purpose here is to develop an analytic method of describing the picosecond dynamics of an electron–hole plasma and laser emission when the nonequilibrium carriers are strongly heated. The method is based on the great differences between the time scales of different development stages of the stimulated processes. A numerical analysis is used to calculate the plasma cooling when stimulated emission sets in without an optical cavity.

BASIC EQUATIONS

According to the experimental[14,15] and theoretical[16,17] results, an electron–hole plasma with carrier density $n \gtrsim 10^{17}$ cm^{-3} enters into a state of quasiequilibrium with a single particle within a time < 1 ps. We therefore use an approximation with a quasi-equilibrium distribution function of the nonequilibrium carriers. In addition, we confine ourselves to the parabolic model of the semiconductor band structure. The balance equations for the number of particles and for the plasma energy then take the form

$$\frac{dn}{dt} = \frac{\alpha(\Omega)}{\hbar\Omega} I + \int \frac{\alpha(\omega)}{\hbar\omega} B_\omega \, d\omega, \qquad n_e = n_h = n, \tag{1}$$

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$$\frac{dW}{dt}=(\hbar\Omega-E_g')\frac{\alpha(\Omega)}{\hbar\Omega}I+\alpha'\left(I+\int B_\omega d\omega\right)$$

$$+\int(\hbar\omega-E_g')\frac{\alpha(\omega)}{\hbar\omega}B_\omega d\omega-nJ,\qquad(2)$$

$$\alpha(\omega)=\alpha_0(\hbar\omega-E_g)^{1/2}(1-f_e-f_h),$$

where the integration is in the frequency interval. $E_g\ll\hbar\omega\ll\hbar\Omega$, E_g is the semiconductor band gap, $\hbar\Omega$ is the photon energy in the exciting pulse, $(\hbar\Omega>E_g^0)$, E_g^0 is the band gap of the unexcited semiconductor, $I(t)$ is the light intensity in the exciting pulse, $n_{e,h}$, is the electron or hole density, $E_g=E_g^0-\gamma n^{1/3}$, γ is the coefficient of band-gap renormalization due to the Coulomb interaction between the carriers

$$E_g'=E_g-n\frac{\partial E_g}{\partial n},$$

W is the plasma kinetic-energy density, α' is the coefficient of intraband absorption of light by the plasma (absoprtion by free carriers), $J=J_e+J_h$, where $J_{e,h}$ are the rates of plasma energy loss to the lattice per electron or per hole (see, e.g., Ref. 7),

$$f_{e,h}=\left[1+\exp\left(\frac{m_{h,e}}{m_e+m_h}\frac{\hbar\omega-E_g}{kT}-\eta_{e,h}\right)\right]^{-1}$$

are Fermi functions, $m_{e,h}$ are the effective masses of the electron and hole, respectively, T is the plasma temperature, $\eta_{e,h}=\mu_{e,h}/kT$, $\mu_{e,h}$ are the electron and hole quasilevels measured from the bottom of the conduction and valence bands, and B_ω is the spectral intensity of the stimulated (laser emission ($E_g\ll\hbar\omega\ll\hbar\Omega$) in the optical cavity. The equation for B_ω is

$$\frac{dB_\omega}{dt}=-\left[c\alpha(\omega)+\frac{1}{\tau_p}\right]B_\omega+g,\qquad(3)$$

where

$$\frac{1}{\tau_p}=c\left(\alpha'+\frac{1}{L\ln R}\right)$$

is the photon lifetime in an optical cavity of length L and a mirror reflection coefficient R ($L\gg\lambda$, where λ is the emission wavelength), c is the speed of light in the semiconductor, and g is the generation rate of the spontaneous recombination radiation. The quantity g comes into play only during the initial stage of the radiation development, i.e., it serves as a trigger. The first terms in the right-hand sides of Eqs. (1) and (2), which correspond to carrier and energy generation by the exciting pulse, have been written in the approximation $\alpha(\Omega)d\ll1$, where d is the semiconducting-layer thickness.

ANALYSIS OF EQUATIONS

Light amplification in a semiconductor sets in if $\alpha(\omega)<0$, i.e., when[18]

$$\mu_e+\mu_h>0.\qquad(4)$$

This connection between the carrier Fermi quasilevels is reached at an electron density $n>n_G(T)$, where

$$n_G(T)=N_C\mathcal{F}_{1/2}(\eta_G),\qquad(5)$$

and η_G is a number determined by the equation

$$N_C\mathcal{F}_{1/2}(\eta_G)=N_V\mathcal{F}_{1/2}(-\eta_G),$$

$N_{C,V}$ are the effective carrier densities in the conduction and valence bands, and $\mathcal{F}_j(\eta)$ is the Fermi integral. If the masses are equal ($m_e=m_h$) we have $\eta_G=0$, while for GaAs ($m_e=0.13m_h$) we have $\eta_G\approx2$.

For nonequilibrium carriers photogenerated by short light pulses of duration Δt ($\Delta t\ll\tau_r\approx10^{-9}$ s, where τ_r is the spontaneous-recombination time) the stimulated-emission intensity can increase rapidly and strongly influence the plasma evolution, but only if a threshold condition stronger than (4) is met, namely,

$$-\alpha(\omega)>\frac{1}{c\tau_p},\qquad(6)$$

which is realized when $n>n_{th}$, where n_{th} ($n_{th}>n_G(T)$) is determined from the condition

$$-\alpha(\omega)=\frac{1}{c\tau_p}.$$

The quantity n_{th}, unlike $n_G(T)$, depends not only on the temperature but also on τ_p. The inequality (6) can be satisfied if the average light intensity $\langle I\rangle$ in the excitation pulse satisfies the condition

$$\frac{\alpha(\Omega)\langle I\rangle}{\hbar\Omega}\Delta t>n_{th}.\qquad(7)$$

The plasma thus evolves differently for $n<n_{th}$ and $n>n_{th}$. In the former case the condition (6) is not met and therefore the behavior of the plasma can be analyzed with the aid of the two equations (1) and (2), in which we put $B_\omega=0$. Such an analysis is reported in Refs. 16 and 17.

As soon as n exceeds n_{th} (the condition (6) is met), the stimulated (laser) emission begins to increase exponentially with an exponent

$$\xi=-\frac{\alpha(\omega)c\tau_p+1}{\tau_p}.$$

Below we consider the influence of this emission on the plasma evolution.

HIERARCHY OF THE EVOLUTION TIMES OF STIMULATED PROCESSES

Equations (1)–(3) can be investigated analytically because the stimulated processes in semiconductor lasers evolve over greatly differing time scales. The first scale, which determines the "ultrafast" processes, is characterized by the quantities τ_p and $1/k_mc$, where $k_m=\max[-\alpha(\omega)]$ is the maximum gain. The second time scale determines the "fast" changes and is governed by the characteristic plasma cooling time τ_T. For $1/k_mc\lesssim\tau_p\ll\tau_T$, the dynamics of the plasma and of the radiation are determined by two very different times τ_p and τ_T. In GaAs, for example, according to experimental data,[1] $\tau_T\sim5$–10 ps and $1/k_mc$ can be of the order of 0.1 ps.

To use the premise that different time scales are present, we express the density in the form

$$n=n_G(T)+\delta n. \tag{8}$$

The quantity $n_G(T)$ corresponds to a plasma density at which light is not yet amplified in the semiconductor, but any increase of the density at a constant temperature T gives rise to amplification. Clearly, the change of n_G is determined by the characteristic time τ_T of the temperature variation. The value of δn also depends on T, but can also undergo ultrafast changes.

ULTRAFAST EVOLUTION STAGE

In this stage, neglecting temperature variation, we get

$$\frac{\partial \delta n}{\partial t}=\frac{\alpha(\Omega)}{\hbar\Omega}I+\int\frac{\alpha(\omega)}{\hbar\omega}B_\omega\,d\omega, \tag{9}$$

$$T=\text{const},\quad n_e=n_h=n,$$

$$\frac{\partial B_\omega}{\partial t}=-\left[c\alpha(\omega)+\frac{1}{\tau_p}\right]B_\omega+g. \tag{10}$$

Equations (9) and (10) are similar to the rate equations describing the operation of a semiconductor laser. From the theory of such a laser[19] it is known that raising the lasing threshold first produces radiation having a broad spectrum, after which mode competition narrows down the emission line, and in the steady state the lasing is at only one frequency ω^*. The emission intensity takes then the form $B_\omega=B_s\delta(\omega-\omega^*)$, where $\delta(\omega)$ is the Dirac delta-function. The frequency ω^* is determined from the condition that the gain be a maximum

$$\frac{\partial k}{\partial\omega}(\omega^*)=0$$

and from the threshold condition

$$k_m=\frac{1}{c\tau_p}, \tag{11}$$

The steady-state value B_s of the radiation intensity is determined from (9) and (11):

$$B_s=\frac{\omega^*}{\Omega}\alpha(\Omega)c\tau_p I. \tag{12}$$

For the case

$$\frac{\mu_e+\mu_h}{kT}\ll 1, \tag{13}$$

from which it follows that $\delta n/n_G(T)\ll 1$, we get an analytic expression for ω^*:

$$\hbar\omega^*=E_g+(\mu_e+\mu_h)/3. \tag{14}$$

For $m_e=m_h$, we have analytically for the Fermi levels

$$\mu_e=\mu_h=\mu=\frac{3}{2}\left(\frac{kT}{2^{1/2}\alpha_0 c\tau_p}\right)^{2/3} \tag{15}$$

and for the nonequilibrium carrier densities

$$n=n_G(T)+\delta n,\qquad \delta n=\frac{\mu}{kT}n_G(T). \tag{16}$$

It follows from (15) that the inequality (13) holds if

$$1/3[\alpha_0(2kT)^{1/2}c\tau_p]^{2/3}\gg 1. \tag{17}$$

For the analytic results we confine ourselves hereafter to the approximation (13), in the framework of which $n_{\text{th}}\approx n_G$.

FAST EVOLUTION STATE

This state is connected with the change of the plasma electron–hole temperature. To obtain the equation that describes the temperature change in closed form, Eqs. (12) and (14)–(16) must be substituted in the energy-balance equation (2). First, however, we obtain more accurately (to first order in the small parameter τ_p/τ_T) the connection between the emission intensity and the plasma temperature:

$$\frac{d}{dt}n_G(T)=\frac{\alpha(\Omega)}{\hbar\Omega}I+\frac{\alpha(\omega^*)}{\hbar\omega^*}B, \tag{18}$$

It is implied in (18) that $B_\omega=B\delta(\omega-\omega^*)$. It follows from (11), (12), and (18) that

$$B=B_s-\hbar\omega^* c\tau_p\frac{d}{dt}n_G(T). \tag{19}$$

Substituting (19) in the energy balance equation and recognizing that $\delta n\ll n_G(T)$, we obtain the approximate equation

$$k\frac{dT}{dt}=\frac{2}{3}\frac{Q-n_G(T)J}{\beta n_G(T)}, \tag{20}$$

where

$$Q=\left[\frac{(\Omega-\omega^*)}{\Omega}\alpha(\Omega)+\alpha'+\alpha'\frac{\omega^*}{\Omega}\alpha(\Omega)c\tau_p\right]I \tag{21}$$

is the rate of plasma heating and

$$\beta=\frac{5}{2}\left[\frac{\mathcal{F}_{3/2}(\eta_G)}{\mathcal{F}_{1/2}(\eta_G)}+\frac{\mathcal{F}_{3/2}(-\eta_G)}{\mathcal{F}_{1/2}(-\eta_G)}\right]+\alpha' c\tau_p\frac{\hbar\omega^*}{kT}. \tag{22}$$

We now determine the characteristic times of electron and hole losses to optical phonons (see, e.g., Ref. 7):

$$\tau_e=\frac{W_e-W_e^L}{nJ_e},\qquad \tau_h=\frac{W-W_h^L}{nJ_h}, \tag{23}$$

where

$$W_e={}^3/_2kTN_C\mathcal{F}_{3/2}(\eta_e),\ W_h={}^3/_2kTN_V\mathcal{F}_{3/2}(\eta_h),$$

$W^L_{e,h}$ is the kinetic-energy density of a plasma of density n at a lattice temperature T_L. With allowance for the definitions (23) we can obtain, for $m_e=m_h$ and $\tau_e=\tau_h$,

$$k\frac{dT}{dt}=\frac{2}{3}\frac{Q}{\beta n_G(T)}-k\frac{T-T_L}{\tau_T}, \tag{24}$$

where

$$\beta\approx 5,5+\alpha' c\tau_p E_g/kT, \tag{25}$$

$\tau_T=(\beta/2.2)\,\tau_h$ is the characteristic plasma cooling time. The expression for β was obtained with allowance for $\hbar\omega^*\approx E_g$. An analytic expression for τ_T can also be obtained in the situation $m_e\ll m_h$, $\tau_e\gg\tau_h$, which obtains in GaAs. Recognizing that the hole gas is nondegenerate, we obtain an equation similar to (24), where the expression for Q coincides with (21), and

$$\beta=\frac{5}{2}\left[\frac{\mathcal{F}_{3/2}(\eta_G)}{\mathcal{F}_{1/2}(\eta_G)}+1\right]+\alpha' c\tau_p\frac{E_g}{kT},\qquad \tau_T=\beta\tau_h. \tag{26}$$

We note once more that expressions (20)–(26) are valid if $\tau_p/\tau_T \ll 1$.

It follows from (21) and (24) that the temperature, meaning also the plasma density and emission intensity, varies with the intensity of the excitation pulse. After the end of the excitation pulse the plasma temperature decreases, with a characteristic time τ_T, to the lattice temperature. The carrier density then drops to a value n_L given at $m_e = m_h$ by

$$n_L=n_G(T_L)+\delta n_L, \tag{27}$$

where

$$\delta n_L=n_G(T_L)\frac{\mu}{kT_L}=n_G(T_L)\frac{3}{[\alpha_0(2kT_L)^{1/2}c\tau_p]^{2/3}}.$$

The value of δn_L estimated from (27) is accurate to τ_p/τ_T.

It must be emphasized here that plasma cooling and laser-emission damping are interrelated processes. In fact, lowering the plasma temperature leads to an increase of the gain, meaning to a higher emission intensity. The latter, in turn, enhances the plasma heating (owing to intraband radiation absorption by the free carriers and to the faster rate of efflux of the "cold" carriers), thereby hindering the cooling and limiting the growth of the emission itself. This negative feedback (considered in Ref. 20) between the temperature and the emission intensity can make the characteristic time of the collective relaxation process much longer than the characteristic lifetime of the photon in the cavity and the energy-relaxation time of the semiconductor carriers.

SLOW EVOLUTION STAGE

After the end of the fast evolution stage of the electron–hole plasma is in a state with $T\approx T_L$ and $n\approx n_{th}(T_L)$. In this case $n>n_G(T_L)$ $(\mu_e+\mu_h>0)$, and therefore the stimulated processes continue. A numerical solution of Eqs. (1)–(3) and analytic estimates show that under these conditions the rate of stimulated carrier recombination and the emission damping are determined by times much longer than τ_T. We do not consider this stage of the evolution here, but we deem it important to note its existence, in view of reported experimental[8,13] observation of stimulated relaxation processes with times much longer than the plasma cooling time.

EXCITATION OF A SEMICONDUCTOR LASER BY POWERFUL SUBPICOSECOND LIGHT PULSES

Equations (1)–(3) are not valid for submicrosecond times, but it can be stated on the basis of experimental data[14,15] that the following processes occur during an excitation pulse: optical generation of electrons and holes, renormalization of the band gap, screening of excitons, generation of optical phonons by nonequilibrium carriers, and interactions between carriers. When a pulse is terminated in an electron–hole plasma of density $\sim 10^{17}$–10^{19} cm^{-3}, within less than a picosecond the carrier collisions establish one temperature common to the electrons and holes. From that instant on, the plasma evolution can be described with the aid of Eqs. (1)–(3).

If the initial density n_0 of a photoexcited plasma is insufficient to produce gain at a steady-state temperature, $n_0<n_G(T_0)$, the cooling takes place for some time without stimulated emission, and the plasma dynamics is determined by Eqs. (1) and (2) with $B_\omega\approx 0$ and $I=0$. In this time interval the carrier density remains practically constant, and the temperature decreases within a certain characteristic time. For a nondegenerate plasma with $m_e=m_h$ this time is found to equal τ_h, while for $m_e\ll m_h$ and $\tau_e\gg\tau_h$ it is equal to $2.5\tau_h$ (this last estimate was obtained under the assumption that only the holes are not degenerate).

As the plasma is cooled, the nonequilibrium-carrier distribution function initially becomes inverted $(\mu_e+\mu_h>0)$, and then lasing sets in. This occurs at a plasma temperature T_1 given by

$$n_0=n_{th}(T_1)\approx n_G(T_1),$$

whence

$$T_1=3{,}5\cdot 10^{-11}\frac{m}{m_e}\left[\frac{n_0}{\mathcal{F}_{1/2}(\eta_G)}\right]^{2/3}, \tag{28}$$

where m is the free electron mass. Thereafter the plasma cooling becomes qualitatively different. The ultrafast and fast evolution stages of the stimulated processes have already been discussed above. During the fast stage, which is described by Eq. (24) with $I=0$, the plasma temperature decreases to the lattice temperature, with a characteristic time τ_T. Let us estimate τ_T for a GaAs laser. According Ref. 21, $\alpha'=\sigma n$, where $\sigma=1.5\cdot 10^{-17}$ cm^2. Recognizing also that $n\approx n_G(T)$ and $\hbar\omega^*\approx E_g$ for the case considered in that paper, we get

$$\tau_T\approx(6{,}2+0{,}34\tau_pT^{1/2}E_g)\tau_h, \tag{29}$$

where τ_T, τ_p, and τ_h are in ps, T in K, and E_g in eV. For $\tau_p=1$ ps and two temperatures $T=300$ K and 500 K we easily find that τ_T equals $14.4\tau_h$ and $16.6\tau_h$, respectively. Dividing these values by the characteristic time $\approx 2.5\tau_h$ of the cooling in the absence of stimulated emission, we estimate that the cooling time is increased 6–7 times, respectively.

According to Ref. 7, at $T=500$ K and $T_L=300$ K we obtain $\tau_h\approx 0.22$ ps for GaAs if no allowance is made for the heating of the optical phonons, and $\tau_h\approx 0.8$ ps with this allowance, whence it follows that $\tau_T\approx 3.7$–13.2 ps. These values agree well with experiment.[8] Note that under the conditions of this experiment ($\tau_p=1.2$ ps, $T>300$ K) the inequalities (17) and $\tau_p\ll\tau_T$ for which the analytic equations were obtained are satisfied.

From (19) and (24) we obtain an expression for the intensity of the emission produced in the course of plasma cooling:

$$B={}^3/_2cE_g\frac{\tau_p}{\tau_T}n_G(T)\frac{T-T_L}{T}. \tag{30}$$

Using this estimate for τ_T, we obtain $B\approx 10^8$ W/cm^2 for $T_L=300$ K and $T-T_L=100$ K.

We obtained an analytic expression for τ_T by using a number of approximations and assumptions, particularly ultrafast mode competition in a semiconductor laser. To verify the analytic procedure we solved Eqs. (1)–(3) numerically. Simulation has shown that if τ_e and τ_h are constant the plasma cooling is indeed exponential for some time. The analytic value of τ_T and the argument of the exponential are equal with an accuracy 10–20% determined by the small parameter τ_p/τ_T.

SEMICONDUCTOR LASER EXCITATION BY STRONG PICOSECOND LIGHT PULSES

When a semiconductor is excited by a strong light pulse 10 ps long, a quasiequilibrium carrier distribution is established in the plasma as early as during the leading edge of the pulse. This permits Eqs. (1)–(3) to be used to investigate the dynamics of stimulated processes during the photogeneration of the electrons and holes. Before the threshold condition (7) is reached the stimulated emission is weak and its influence on the plasma dynamics can be disregarded. The ultrafast stage of the evolution of the stimulated processes sets in at $n > n_{\rm th}$. Subsequently the plasma dynamics is described by Eq. (24) under the assumption that the evolution is fast. It follows from (24) that the plasma is heated during the photogeneration. The reasons for the heating are 1) the optically excited carriers have an energy higher than the carriers that depart during the stimulated emission, and the energy difference goes to increase the plasma kinetic energy; 2) intraband absorption of the exciting light by the free carriers; 3) intraband absorption of the stimulated emission by free carriers. The above heating mechanisms are listed in the same order as the terms corresponding to them in Eq. (21).

Comparing the contribution of each mechanism we find, for example, that heating due to intraband absorption of stimulated emission is predominant if

$$\frac{\Omega-\omega^*}{\Omega}<\alpha' c\tau_p,\quad \alpha(\Omega)c\tau_p>1. \tag{31}$$

Thus, in the case of a GaAs laser, for which $\tau_p \approx 1$ ps and $\alpha' = 15\ \mathrm{cm}^{-1}$, we find that Eq. (31) is satisfied for $\hbar\Omega - E_g < 0.2$ eV and $\alpha(\Omega) > 10^2\ \mathrm{cm}^{-1}$.

It follows also from (24) that the temperature (meaning also the density) of the plasma varies with $I(t)$ but with a time lag determined by τ_T. The time variation of the stimulated-emission intensity is somewhat more complicated. That part of the radiation which corresponds to B_s in (19) follows the variation of $I(t)$, with a time lag determined by τ_p and $1/k_m c$. The second term in Eq. (19) corresponds to the radiation intensity, which varies with a time lag τ_T. After the end of the excitation pulse the plasma temperature and the emission intensity decrease with approximately the same characteristic time scale τ_T until the plasma temperature becomes equal to T_L.

We have simulated numerically the equations for a GaAs semiconductor laser with parameters $\tau_p = 1$ ps, $T_L = 300$ K, $\hbar\Omega - E_g^0 = 15$ meV, $\alpha' = 15\ \mathrm{cm}^{-1}$, $\tau_h = 0.3$ ps, a power $4\cdot 10^8$ W/cm, and a pulse length 25 ps at half-maximum. It is interesting that photogeneration of such relatively "cold" carriers ($\hbar\Omega - E_g^0 < kT_L$) heated the plasma to 600 K. The main contribution to the heating at these parameters is made by intraband absorption of the laser emission, since the inequality (31) is satisfied. Increasing $\hbar\Omega - E_g$ increases the relative contribution to the heating from the first term of (21), while at

$$(\hbar\Omega-E_g)/E_g=\alpha' c\tau_p \tag{32}$$

these contributions become comparable.

It is noteworthy that the reversible change of the plasma temperature in an optically excited semiconductor laser is accompanied by a reversible change of the carrier density. Under the conditions of the numerical computation, the density changed by more than a factor of 2 when the temperature changed from 300 to 600 K. The density change leads to a reversible optical bleaching at the excitation frequencies Ω and at other frequencies. In semiconductors, the effect of the reversible picosecond bleaching was observed experimentally[22] and discussed in a number of papers.[23–25] The reversible bleaching effect may turn out to be stronger in a semiconductor laser than in a semiconductor, other conditions being equal, since the laser-emission intensity can be many times larger than the light intensity in the excitation pulse or the stimulated-emission intensity in a semiconductor.

DISCUSSION

Let us examine which of the theoretical results obtained here for a semiconductor laser also hold for a semiconductor without a specially prepared optical cavity.

The main assumptions under which Eqs. (19)–(26) have been obtained are

$$\delta n \ll n_G(T),$$
$$\tau_p \ll \tau_T,$$
$$k_m = 1/c\tau_p.$$

Only the last condition, generally speaking, has a bearing on lasing. Reduction of the experimental results[24] has shown that the first two inequalities can be satisfied also in the absence of an optical cavity. Therefore, if we replace $c\tau_p$ in (20)–(22) and (24)–(26) by $\langle k_e\rangle^{-1}$, where $\langle k_e\rangle$ is the effective gain, these equations will correctly describe the plasma evolution in a semiconductor without an optical cavity. (The results indicative of the emission itself no longer hold here.)

CONCLUSION

We have proposed a method of describing the dynamics of laser emission and of a hot electron–hole plasma in a semiconductor laser in the case $\tau_p/\tau_T \ll 1$, where τ_p is the photon lifetime in the cavity and τ_T is the characteristic plasma cooling time. The method was used to investigate relaxation processes in a semiconductor laser excited by a strong pico- or subpicosecond light pulse. We have shown that the plasma cooling and the damping of the long-wave part of the laser-emission spectrum are interrelated, and the rates of these processes, as well as the rate of stimulated carrier recombination, are determined by a single characteristic time τ_T. An analytic expression was obtained for τ_T and the dependence of the lasing intensity on the plasma temperature was found. It was shown that τ_T can be much longer than the time of carrier-energy relaxation on optical phonons, a time that determines the rate of plasma cooling in the absence of laser radiation. The cause of the cooling time lag is the plasma heating during the emission on account of the preferred departure of the "cold" carriers and of the intraband absorption of the radiation. When the plasma temperature is raised, in turn, the gain decreases and with it the emission intensity.

As a result, relaxation with a single characteristic time is established in a self-consistent manner.

We have analyzed the effect of a reversible picosecond change of the optical properties of a semiconductor laser, an effect previously[22-24] studied for semiconductors without optical cavities.

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