Coherent effects in generation of ultrashort light pulses in a semiconductor injection laser

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Theoretical and experimental investigations were made of the generation of ultrashort pulses in a Q-switched semiconductor laser. Experiments demonstrated splitting of the pulse envelope into a series of subpulses, as well as stimulated Raman scattering of a train of picosecond pulses by the active medium. The observed effects can be attributed to the coherence of the interaction of a light pulse with the active medium when the intraband relaxation is suppressed by the strong electromagnetic field. A theoretical analysis of the operation of the laser was made using the Maxwell–Bloch equations and allowing for the field dependence of the polarization relaxation time, for the inhomogeneity of the gain profile, and for the possibility of stimulated Raman scattering in the active medium of the laser.

1. INTRODUCTION

Generation of coherent light radiation by a two-level medium in its purest form could be regarded as due to an ensemble of particles with infinite longitudinal and transverse relaxation times T_1 and T_2, and with a particle polarization vector amplitude P varying in accordance with the law

$$P \propto \sin \frac{\mu}{\hbar} \int_{-\infty}^{t} E(t)\,dt,$$

where μ and E are the dipole moment of the transition and the amplitude of the electric field of the optical wave. Depending on the angle (pulse "area")

$$\psi = \frac{\mu}{\hbar} \int_{-\infty}^{t} E\,dt$$

particles either absorb or emit photons in resonance with the transition frequency; a medium of this kind can emit photons even when the population is not inverted. Until the appearance of our earlier paper,[1] there have been no experimental realization of coherent emission from a medium consisting of optical-range oscillators. This has been due to the obvious difficulties in satisfying simultaneously the conditions for pulsed emission of photons from particles and the requirement that the pulse duration τ should be short compared with the polarization relaxation time. For example, in the case of semiconductors the transverse relaxation time T_2 is 10^{-13} s at 300 K.

In the Bloch equations, used widely in studies of coherent effects in the interaction of light with matter, the times T_1 and T_2 are assumed to be independent of the optical field: $T_1(E) = T_1(0)$ and $T_2(E) = T_2(0)$. Then, since $T_2 < T_1$, the conditions for coherent emission from a medium reduce to the requirement $\tau < T_2(0)$. It should be pointed out that T_2 is usually due to several physical mechanisms. In the case of semiconductors, for example, the transverse relaxation time T_2 is governed by electron–electron, electron–phonon, electron–impurity, etc. collisions and each of them has its own correlation time τ_c. In a strong electromagnetic field when the Rabi frequency $\Omega = \mu E/\hbar$ becomes comparable with the characteristic frequency τ_c^{-1} of perturbations in one of the relaxation channels, the channel in question becomes suppressed.[2-4] The physics of the effect is related to the circumstance that the field photons also participate in each elementary relaxational collision. Consequently, the time T_2 becomes dependent on E: as E is increased, the value of $T_2(E)$ rises. Even if the interaction coherence condition $\tau < T_2(E) = T_2(0)$ is not obeyed in a weak field, it may be satisfied in strong fields.

We shall give the results of theoretical and experimental investigations of coherent emission from the active medium of an injection laser in the case when the condition $\tau < T_2(0)$ is not satisfied (in fact, conversely, we have $T_2(0) < \tau$). The use of a new method for Q-switching the laser cavity[5] enabled us to reach radiation power densities of the order of gigawatts per square centimeter in the active region for pulses shorter than 5 ps and a value of $\Omega T_2(0)$ of the order of unity. The parameters of the generated pulses were sufficient for strong stimulated Raman scattering (STRS) by the active medium of the laser.[6] We could assume that the Raman scattering occurred also in the case of coherent transformation of light, when relaxation of the phase of the Raman-active modes was suppressed in GaAs: the phase relaxation time T_2 for STRS was found to be in the picosecond range. In the concluding section we shall discuss the possibility of using coherent STRS in generation of femtosecond pulses in an injection laser.

2. DYNAMICS OF COHERENT EMISSION FROM THE ACTIVE MEDIUM OF A LASER (THEORY)

We shall adopt a laser model described by the following system of the Maxwell–Bloch equations:

$$\frac{dE}{dt} = \dot{E} = -\frac{1}{\tau_{\text{ph}}} E + \frac{2\pi\omega}{\beta^2} \sum_k N_k P_{2k}, \tag{1}$$

$$\dot{P}_{2k} = -P_{2k}/T_2(E) + (\mu_k^2/\hbar) n_k E - (\omega-\omega_k) P_{1k}, \tag{2}$$

$$\dot{P}_{1k} = -P_{1k}/T_2(E) + (\omega-\omega_k) P_{2k}, \tag{3}$$

$$\dot{n}_k = -E P_{2k}/\hbar. \tag{4}$$

Here, τ_{ph} is the photon lifetime in the cavity; β is the refractive index of the medium; ω is the frequency of the generated radiation; P_{2k} and P_{1k} are the amplitudes of the active and reactive (to the field) components of the polarization vector of an electron with the quasimomentum k, frequency ω_k, and a dipole moment $\mu_k = \mu$; N_k is the number of such k electrons. The relaxation time T_1 is assumed to be infinite and it does not enter in the above equations. The Q factor of the cavity is switched on at a moment $t = 0$, when the difference between the populations in the bands $n_k(t)$ amounts to n_0 (if at $t = 0$ all the electrons are in the conduction band, then $n_k = +1$). In the course of generation of a pulse the quantity n_k is generally an oscillatory variable-sign function. The presence of an absorber, which Q-switches the laser cavity, governs the value of n_0 and has a strong influence on the envelope of an ultrashort pulse. The efficiency of the operation of this absorber is governed by a quantity η, which is a parameter representing an excess of the population inversion above the threshold value when the absorber is bleached at the moment when the Q factor is switched on:

$$\eta = \frac{2\pi}{\hbar\beta^2}\omega\mu^2\tau_{\mathrm{ph}}\left[T_2(0)\sum_k \frac{N_k n_0}{1+[T_2(0)(\omega-\omega_k)]^2}\right]. \qquad (5)$$

[Equation (5) is derived from Eqs. (1)–(4) with zero left-hand sides of the equations, as was done in Ref. 7.]

Our laser has the remarkable property that we can achieve $\eta \gg 1$, whereas other injection lasers are characterized by $\eta - 1 \ll 1$.

We shall first analyze the results obtained by solving Eqs. (1)–(4). If in Eqs. (2)–(3) we substitute $\dot{P}_{1,2k} = 0$ (which corresponds to the case of noncoherent emission from the investigated medium) and assume that T_2 is independent of E, we find that over periods of $(10\text{–}20)\tau_{\mathrm{ph}}$ the laser radiation is emitted in the form of single pulses for any $\eta > 1$. In the case of a homogeneously broadened line ($\omega_k = \omega$) this is the familiar result obtained by solving the laser rate equations.[8] Conversely, when the derivatives of P_{1k} and P_{2k} cannot be regarded as zero, effects typical of coherent emission from the active medium are observed. In particular, if η is large, then during the pulse duration τ the amplitude E of the field executes a whole series of oscillations passing through zero; at the end of a pulse we have $n_k = -n_0$. The energy initially stored in the active medium of the laser is transferred in succession from the medium to the field and back again. The laser pulse profile resembles the profile of a light pulse propagating in a resonantly amplifying relaxation-free medium[9-12]: the field consists of a sequence of subpulses of about $\pm 2\pi$ area which added up gives a total area approximately equal to π. Such a pulse extracts all the energy stored in the active medium of the laser, and not half of it as in the noncoherent case. Figure 1 shows typical time dependences of the amplitude of the field pulse in the case of homogeneously and inhomogeneously broadened gain profiles, obtained by numerical solution of the system of equations (1)–(4) allowing for the fact that $\eta \gg 1$. In this case the value of T_2 tends to infinity when E exceeds $E_0 = \hbar/5\mu\tau_{\mathrm{ph}}$. The increase in η and the reduction in E_0 do not alter qualitatively the situation: the pulse duration remains practically the same as before, and only the frequency and peak-to-peak amplitude of the oscillations increase.

In the case of a homogeneously broadened gain profile and high values of η the profile of a subpulse in the vicinity of a field maximum can be described analytically. If we assume that a subpulse appears at the leading edge of a radiation pulse and if we ignore the terms E/τ_{ph} and P_{2k}/T_2 in Eqs.

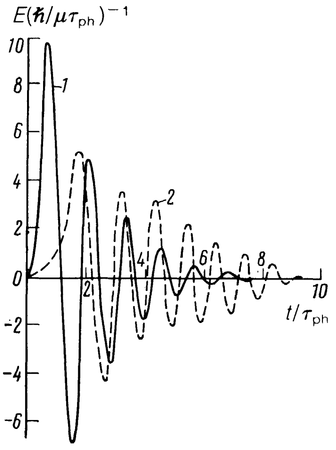

FIG. 1. Calculated envelopes of the field pulses in the case of a coherent interaction in an injection laser with a homogeneous (1) and an inhomogeneous (2) gain profile when $T_2 \to \infty$ if $E > E_0$, $T_2(0) = 0.03\tau_{\mathrm{ph}}$, $T_2(E>4) \gg 1$, $\eta = 6$.

(1)–(4), i.e., if we assume that the subpulse is short compared with the photon lifetime in the cavity and that the process of phase relaxation of the active medium is suppressed, we find (see also Ref. 7) that

$$E(t) = 2\hbar/\mu\tau_c\,\mathrm{ch}[(t-t_c)/\tau_c],\quad \tau_c = [T_2(0)\tau_{\mathrm{ph}}\eta^{-1}]^{1/2}, \qquad (6)$$

where t_c is the moment at which the field intensity in the subpulse reaches the maximum. At $t = t_c$ it follows from Eqs. (5) and (6) that the field energy density is

$$E^2/8\pi = \hbar\omega n_0 N \left(N = \sum_k N_k\right)$$

and the whole energy stored initially in the laser is now in the form of an electromagnetic field. The characteristic time τ_c is inversely proportional to the square root of the number of the inverted particles:

$$\tau_c = [\hbar/2\pi\omega\mu^2 n_0 N]^{1/2}.$$

The number of oscillations of the field during the pulse duration τ increases with η as $\eta^{1/2}$.

If the rise of T_2 on increase in the field intensity is limited to the finite value $T_2(\infty)$, the pattern of emission of an ultrashort pulse changes. Figure 2 shows the dependence of the pulse $E(t)$ on the field amplitude when the polarization relaxation time is

$$T_2(E) = T_2(0) + [T_2(\infty) - T_2(0)][1-\exp(-|E|/E_0)]. \qquad (7)$$

We carried out calculations assuming for the function $T_2(E)$ the following parameters which are closest to the parameters of an injection laser: $T_2(0) = 10^{-13}$ s, $\tau_{\mathrm{ph}} = 3\times 10^{-12}$ s, $E_0 = \hbar/3\mu T_2(0)$, $\mu = 2\times 10^{-17}$ cgs esu, $\eta = 10$, $T_2(\infty) = 3\times 10^{-12}$ s.

We considered the cases of homogeneous and inhomogeneous line broadening. Allowance for the finite value of $T_2(\infty)$ reduces considerably the number of oscillations of the field (cf. Fig. 1). It is worth noting that the interaction of the field with the active medium of the laser is coherent also when $T_2(\infty) < \tau_{\mathrm{ph}} \sim \tau$. It should also be mentioned that the $T_2(E)$ dependence selected by us in the form of Eq. (7) is not decisive in the effects under consideration and the interac-

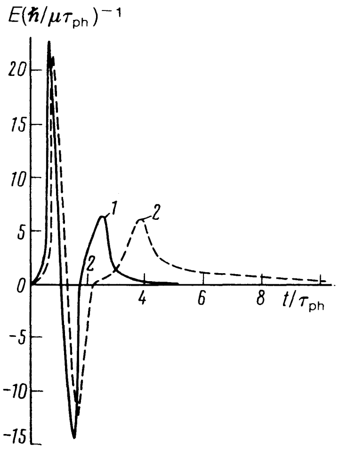

FIG. 2. Calculated envelopes of the field pulses for a finite time T_2 and a homogeneous (1) or an inhomogeneous (2) gain profile. Here, $T_2(E)$ is given by Eq. (7), $T_2(0) = 0.03\tau_{\rm ph}$, $T_2(\infty) = \tau_{\rm ph}$, $E_0 = 10$, and $n_0 = 0.05$.

tion of a field pulse with the active medium of the laser is not altered qualitatively if we use other monotonic functions $T_2(E)$ satisfying the condition $T_2(0) \ll T_2(\infty)$. The parameters $T_2(0)$, $T_2(\infty)$, $\tau_{\rm ph}$, η in Eq. (7) represent in our opinion correctly the physical conditions in a Q-switched injection laser and the calculated pulse field profile describes approximately correctly the profile found experimentally.[1]

3. DYNAMICS OF COHERENT LASER EMISSION (EXPERIMENTS)

Our experiments were carried out on an AlGAs/GaAs injection heterolaser operating using modified Q-switching.[5] The active region of the laser along the cavity axis consisted of three sections: two of them were light-amplifying and one, enclosed between them, absorbed light nonlinearly. This central section was our saturable absorber (performing passive Q-switching of the cavity) and was characterized by an ultrashort recovery time (<100 ps). The main feature of our lasers, which distinguished them fundamentally from other lasers with passive Q-switching, was a strong electrical (and not only optical) coupling of the nonlinear absorber with the amplifying regions. The presence of a strong (10^3–10^4 V/cm) electric field in the semiconductor parts between the absorber and amplifier and the feasibility of free drift of carriers caused a large proportion of the excited electron–hole pairs in the absorber to leak to the amplifying parts within a time interval shorter than T_1 during the intervals between the pulses. The carrier drift made it possible to maintain a high value of the absorption coefficient of the bleachable absorber (Q-switch), avoid its bleaching by spontaneous radiation, and accumulate (before Q-switching) a population inversion in the amplifying region exceeding by a factor of 7–10 the threshold value for the bleached absorber. The calculated photon lifetime in the cavity with the Q switched on was 3–4 ps. Experiments yielded the following estimates: $\tau \leqslant 5 \pm 2$ ps, $\Omega > 10^{13}$ rad/s (at a GaAs power density $q > 4\times 10^8$ W/cm^2 and dipole moment was $\mu = 2\times 10$ cgs esu). Therefore, in contrast to other picosecond injection lasers, we were dealing here with the case of generation of a strong field in the form of picosecond pulses characterized by $\Omega\tau_c \sim 1$.

In our earlier study[1] we determined the profile of the intensity of picosecond pulses, confirming the splitting of a single pulse into three subpulses of a strong field. This splitting was observed only when pumping of the amplifying and absorbing parts of the laser was such that the value of η was maximal (it was estimated to be 10). The appearance of several pulses instead of one was unrelated to multiple reflection of the same pulse from the laser mirrors, because the difference between the intensity maxima varied with the rate of pumping and the round-trip time was fixed and amounted to 15 ± 1 ps. Unfortunately, there are at present practically no experimental methods for determination of the sign of the envelope of the field intensity of picosecond pulses. In view of the fact that the durations of the pulses and the separations between them were comparable with the time resolution of our "Agat" camera (2 ps), and also because of the method of imaging of an optical process used in this camera, the pulses might not be resolved completely (see Fig. 1 in Ref. 1).

4. STIMULATED RAMAN SCATTERING ASSOCIATED WITH GENERATION OF A PICOSECOND PULSE TRAIN (EXPERIMENTS)

The compound GaAs is a semiconductor which is Raman-active. For example, the scattering of laser radiation by infernal vibrations in GaAs was described in Ref. 13 (spontaneous Raman scattering occurred on plasmons and on longitudinal LO and transverse TO optical phonons). These vibrational modes of GaAs interact with one another, which results in their coupling. Coupled modes represent a mixture of a phonon and a plasmon, and the scattering occurs on plasmon–phonon modes. In our experiments it was found that stimulated Raman scattering occurred when a train of picosecond pulses with a repetition frequency in excess of 1 GHz was generated. Figure 3 shows the optical spectrum of pulses, beginning from the first pulse in a train. The intensity of the pulses was spread out along the horizontal (ω) axis by a 600-lines/mm diffraction grating, and along the vertical (t) axis by using the "Agat" image-converter camera. Clearly, the number of the Stokes components of the spectrum increased with time. The distance between the components was 220 ± 30 cm^{-1}. The energy of the Raman scattering components was comparable with the energy of the main pulse.

One of the reasons for the change in the number of the spectral components from pulse to pulse could be attributed to the following circumstance. In a time interval of $\Delta\tau \sim 10^{-9}$ s between the pulses, the dephasing of the Raman-active oscillators causes the polarization per unit volume of the medium to vanish, although the absolute value of the polarization vector of each of the oscillators does not change significantly in the time interval $\Delta\tau$. Therefore, the interaction of the next pulse with the Raman-active component of the active medium of the laser started with phase-locking of the oscillators, and then the Raman scattering occurred from the medium-polarization level established by the previous pulse. This was a cumulative effect and the transformation of the field integrated by the active medium became stronger the farther pulse was from the beginning of the train.

Another reason for the change in the number of the spectral components in Fig. 3 could be the transient nature

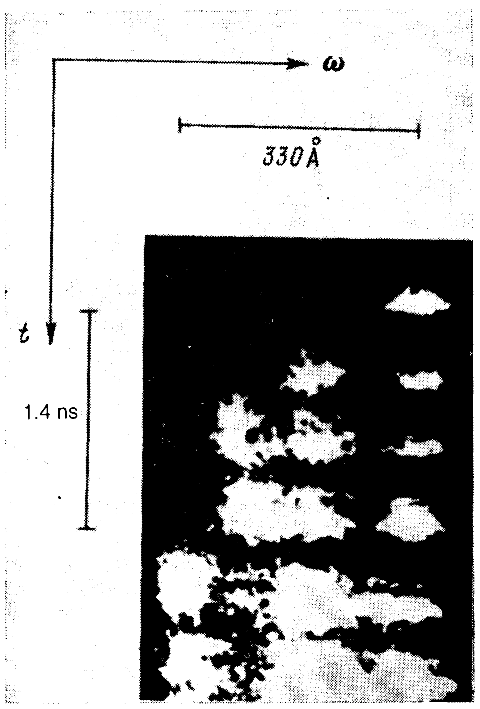

FIG. 3. Raman spectrum of a train of ultrashort pulses emitted by a Q-switched injection laser. Separation between the components 220 ± 30 cm^{-1}.

of the pulses used to pump the laser and also the excitation of collective oscillations with lifetimes in excess of 10^{-10}–10^{-9} s in a nonequilibrium electron–hole plasma. A satisfactory account of the physical mechanisms of spectral dynamics of a train of picosecond pulses is difficult and is at present largely lacking.

5. GENERATION OF FEMTOSECOND PULSES AS A RESULT OF COHERENT STIMULATED RAMAN SCATTERING IN THE ACTIVE MEDIUM OF AN INJECTION LASER

Since we can regard it as firmly established that the Raman scattering occurs in a train of ultrashort pulses, we shall consider the possibility of generation of femtosecond pulses by suppression of oscillations of $E(t)$ as a result of coherent STRS in the active medium of a laser.

If the pulse duration is such that the inequality $\tau > 2\pi/\Delta\omega$ is satisfied (here, $\Delta\omega$ is the difference between the frequencies of the main pulse and its first Stokes component), the STRS process is delayed because of amplification of the Raman components beginning from the spontaneous noise level.

The experiments described in the preceding section are characterized by $2\pi/\Delta\omega \sim 10^{-13}$ s. We shall assume that the condition $\tau < 2\pi/\Delta\omega$ is satisfied, which can be achieved for example by reducing the resonator length. In the case of STRS there is no longer any need for a "seed" field: the spectrum of a pulse contains initially a sufficient number of both Stokes and anti-Stokes components satisfying the STRS resonance condition. The STRS sets in simultaneously with the main field and, therefore, it should influence directly the dynamics of growth of a pulse. Since STRS occurs under the coherent interaction conditions, the absorption of light by the Stokes components is a cumulative effect. The absorption efficiency increases with time and this should reduce the intensity at the trailing edge of a pulse and, under certain conditions, it should shorten the pulse to a value of the order of the Rabi period.

We shall begin by estimating the conditions for the generation of such a pulse and we shall do this employing a classical description of the Raman scattering process. We shall represent the component of the permittivity of the active medium $\varepsilon(Q)$ responsible for STRS in the form $\varepsilon(Q) = Q\partial\varepsilon/\partial Q = \varepsilon'(0)Q$, where Q is the normal coordinate of a Raman-active mode. In this case the dipole moment per unit volume is $\varepsilon'(0)Q\mathscr{E}$ and Q varies in accordance with the law

$$m^*[\ddot{Q}+2\dot{Q}/T_2^*+(\Omega^*)^2Q]=-\alpha'(0)\mathscr{E}^2/2. \tag{8}$$

The left-hand side of the above expression describes an oscillator with an effective mass m^* and a frequency Ω^*, whereas the right-hand side is the force due to the polarizability $\alpha(Q)$ of the oscillators in a field of intensity $\mathscr{E}$. If we assume that $\tau \gg T_2^*, 2\pi/\Omega^*$, which allows us to substitute $\Omega^* = 1/T_2^* = 0$ in Eq. (8), we find that the energy absorbed per unit volume $N^*m^*Q^2/2$ is described by

$$\frac{N^*(\alpha')}{32m^*}\left[\int_{-\infty}^{t} E^2\,dt\right]^2.$$

We used here the relationship $\varepsilon' = \alpha' N^*$ to introduce the density N^* of the scattering centers, and the mean squared intensity $\overline{\mathscr{E}}^2 = E^2/2$ is expressed in terms of the square of its amplitude. If we use τ_c and $E_c = 2\hbar/\mu\tau_c$ from Eq. (6) in our estimates, we obtain

$$N^*(\alpha')^2E_c^4\tau_c^2/32m^*\approx Nn_0\hbar\omega,$$

which means that in a time τ_c of the order of the Rabi period the energy stored per unit volume of the active medium is absorbed by the oscillators. This condition can be rewritten in the form

$$\eta^*=\pi|\varepsilon'(0)\alpha'(0)|\hbar^2/\mu m^*\approx 1. \tag{9}$$

Therefore, when Eq. (9) is obeyed, the pulse duration is of the order of τ_c. We shall now consider the results obtained by numerical solution of the equations of the laser dynamics using a quantum description of STRS based on a two-level model. Equation (8) can be generalized on the basis of this formalism in the following self-evident manner: the right-hand side is multiplied by n^*, where n^* obeys

$$\dot{n}^*=-(\alpha'(0)/\hbar\Omega^*)\dot{Q}\mathscr{E}$$

and represents the difference between the populations of the excited and ground states. If $\tau \ll 2\pi/\Omega^*$, the material equations for a two-level medium can be solved for any field profile:

$$n^*=-\cos\psi^*,\qquad Q=(\hbar\Omega^*/2m^*)^{1/2}\int_{-\infty}^{t}\sin\psi^*\,dt,$$

where the "area" ψ^* under a pulse is given by

$$\psi^*=\alpha'(0)\,(2\hbar\Omega^*m^*)^{-1/2}\int_{-\infty}^{t}\mathscr{E}^2\,dt. \tag{10}$$

The lasing dynamics can now be described by a system

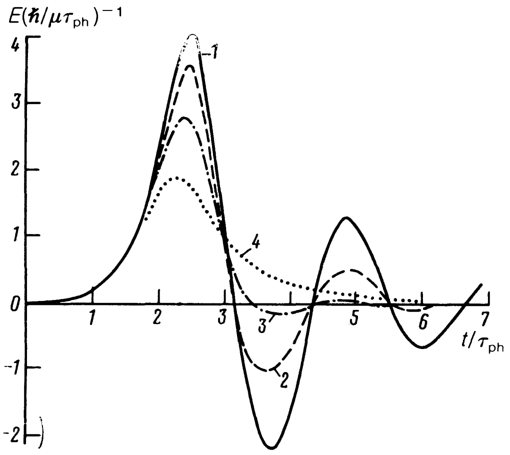

FIG. 4. Amplitude of pulse field in the case of stimulated Raman scattering in the laser oscillator medium for $\eta = 8$ and different values of η^*: 1) 0; 2) 0.0125; 3) 0.5; 4) 2.0.

of equations (1)–(4) allowing for the fact that the right-hand side of the field-equation (1) is supplemented by

$$-2\pi|\varepsilon'(0)|E(\hbar\Omega^*/2m^*)^{1/2}\sin\psi^*,$$

and $\mathscr{E}^2$ in ψ^* is replaced by $E^2/2$. The solution of the system (1)–(4) now shows that mainly the Stokes components are excited during the initial stage of lasing if the total area under a pulse $\psi^*(\infty)$ is small and STRS involves absorption of the energy emitted by the active medium of the laser. Conversely, if $\psi^*(\infty)$ is large (i.e., if $\pi < \psi^*(\infty) < 2\pi$), the leading edge of a pulse excites the Stokes components and the trailing edge the anti-Stokes components, and the energy absorbed by the medium at the leading edge is restored by the anti-Stokes components to the field of the pulse. The results of a numerical solution of these equations showed that in the case of generation of pulses of minimum duration τ one should concentrate on the case of small values of $\psi^*(\infty)$ and large values of η. If ψ^* is small, so that $\sin\psi^*$ can be replaced with ψ^*, the influence of STRS on the pulse dynamics is characterized by one dimensionless parameter η^*. Figure 4 shows the dependence of the field amplitude $E(t)$ in a pulse as a function of η^*. At low values of η^*, it is found that $E(t)$ is an oscillatory variable-sign function. At high values of η^* the laser radiation becomes single-pulse and an increase in η^* causes the duration to increase slowly, whereas the amplitude falls. There is an optimal value η^*_{opt} when STRS selects only the first pulse of maximum amplitude and minimum duration from a train (curve 3 in Fig. 4). The numerical value of η^*_{opt} is approximately 0.5 and it is close to unity, i.e., it is close to the value obtained above in estimating the feasibility of shortening the pulses to durations of the order of the Rabi period. Therefore, when the condition (9) is satisfied, which can be achieved by the specific level of the doping of a semiconductor, we can use STRS to remove oscillations of the field and to obtain a single pulse of duration in the femtosecond range.

6. CONCLUSIONS

The present paper proposes a new approach to the dynamics of emission of ultrashort pulses by a Q-switched laser under conditions of generation of a strong field, based on a coherent interaction of an optical pulse with the active medium of the laser. In our opinion the proposed theoretical model, based on the Maxwell–Bloch equations and allowing for the dependence of the polarization relaxation time on the field amplitude, makes it possible to provide a qualitatively correct account of the observed behavior of lasing without any additional assumptions.

It follows from this model that at high values of the excess of the initial population inversion of the active medium above the threshold value, a field pulse splits into a series of 2π pulses. During about half the lasing time the energy stored in the active medium is in the form of an electromagnetic field: the pulse duration should then be $(3\text{–}4)\tau_{ph}$, where τ is the photon lifetime in the laser resonator, and in fact has this value (as confirmed experimentally). Since $\tau_{ph} \sim L/c$, it follows that a reduction in the cavity length L to 40–50 μm can be used to generate pulses of duration amounting to hundreds of femtoseconds. A further reduction in τ may be associated with the suppression of oscillations of the field amplitude and generation of a single pulse due to stimulated Raman scattering in the active medium. In this case the pulse duration should be of the order of the period of the Rabi oscillations of the field in the absence of Raman scattering and the value of τ should decrease as $\eta^{-1/2}$. We can therefore expect generation of pulses not exceeding tens of femtoseconds in duration. Moreover, the actual process of stimulated Raman scattering in a laser under coherent interaction conditions, when the spectral width of a pulse exceeds the frequency of Raman oscillators, is in itself of interest to physicists.

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