NONLINEAR AMPLIFICATION OF LIGHT PULSES

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The passage of a powerful light pulse through an optical quantum amplifier operating under saturation conditions is investigated theoretically and experimentally. It is shown that during nonlinear amplification the pulse peak moves with a velocity which considerably exceeds the velocity of light, the duration of the pulse remaining practically unchanged. In order to decrease the pulse duration during nonlinear amplification, the slope of the incoming pulse front should be increased. This is realized in the experiments by cutting off of the leading front with the aid of an additional shutter. The duration of the pulse with such a front was made as low as 4.7 nsec during nonlinear amplification.

1. INTRODUCTION

THE desire to obtain light pulses with maximum energy and power leads unavoidably to use of optical quantum amplifiers operating in the nonlinear mode connected with the saturation effect.[1] It had been expected[2-5] that nonlinear amplification would result in further increase in power, due to the preferred amplification of the leading front and reduction in the duration of the input pulse. Experimentally, however, no appreciable reduction in the pulse duration during amplification has been attained to date. Experiments reported in [6] show that nonlinear amplification of the pulse of a Q-switched laser results not in the shortening of the duration of the leading front, but in a displacement of the pulse as a whole, with the maximum of the amplified pulse appreciably leading the maximum of the pulse causing the amplification, so that the rate of propagation of the maximum pulse greatly exceeds the speed of light. The propagation of the maximum of the pulse in the amplified medium, with a velocity exceeding that of light, does not violate the causality principle, since such a propagation is at the expense of amplification of photons emitted by the driving laser at earlier instants of time. The motion of the pulse can continue only until the shutter of the driving laser is turned on. A shortening of the pulse duration should be observed simultaneously with the approach of the pulse near the instant of shutter switching.

In the present investigation we used an additional shutter (Kerr cell) to reduce the pulse duration in nonlinear amplification; this shutter increased the slope of the leading front of the amplified pulse. We were able to observe in this case an appreciable reduction of the pulse duration.

II. THEORY OF NONLINEAR AMPLIFICATION

The theory of nonlinear amplification of a light pulse was considered theoretically by several workers.[4, 5, 7-14] It has been shown that the duration of the pulse propagating in a medium with population inversion can be reduced; in a medium with radiation losses, the pulse energy becomes practically constant after passing through a definite distance; the minimum pulse duration obtainable with nonlinear amplification is of the order of the reciprocal of the width of the radiation transition. We consider below the singularities connected with propagation of a pulse with a velocity exceeding that of light.

1. Fundamental equations

We consider a medium of two-level "atoms" with inverse level population. Although we have in mind the case of impurity ions in a crystal, the results obtained are valid for any quantum system with homogeneous broadening of levels. To describe the medium we use Boltzmann's equation for the density matrix $\hat{\rho}$ with longitudinal and transverse relaxation, and to describe the field **E** we use Maxwell's equations in a medium having also nonresonant radiation losses:

$$\frac{\partial^2 \mathbf{E}}{\partial t^2} + c^2 \operatorname{rot} \operatorname{rot} \mathbf{E} + \gamma c \frac{\partial \mathbf{E}}{\partial t} = -4\pi N_i \frac{\partial^2}{\partial t^2} \operatorname{Sp}(\boldsymbol{\mu}\hat{\rho}), \quad (1)^*$$

*rot ≡ curl.

$$i\hbar\frac{\partial\hat{\rho}}{\partial t} = [\mathscr{H}_0 - \mu E, \hat{\rho}] - i\hbar\Gamma\hat{\rho}, \qquad (2)$$

where c is the velocity of light in the medium without ions, γ the coefficient of radiation loss per unit length in the medium, N_i the ion density, $\mathscr{H}_0$ the unperturbed Hamiltonian of the ion, and μ the operator of the electric dipole moment of the ion. The term $\Gamma\hat{\rho}$ describes phenomenologically the relaxation of the density matrix element: the diagonal elements are $(\Gamma\hat{\rho})_{nn} = (\rho_{nn} - \rho^0_{nn})/T_1$, the off-diagonal elements are $(\Gamma\hat{\rho})_{mn} = \rho_{mn}/T_2$ (m, n = 1, 2), where T_1 is the time of longitudinal relaxation, equal to the lifetime of the ion in the excited state, and T_2 is the time of transverse relaxation, determined essentially by the ion-phonon interaction that leads to homogeneous broadening of the levels. In crystals in which population inversion is attained we have $T_1 \gg T_2$.

In the representation in which $\mathscr{H}_0$ is diagonal, the equation (2) for the density matrix can be reduced to equations for polarization $\mathbf{P} = N_i \mathrm{Sp}(\mu\hat{\rho})$ and the density of the inverse population $N = N_i(\rho_{22} - \rho_{11})$:

$$\frac{\partial^2 \mathbf{P}}{\partial t^2} + \frac{2}{T_2}\frac{\partial \mathbf{P}}{\partial t} + \omega_0^2 \mathbf{P} = -N\mathbf{E}\frac{\omega_0}{\hbar}|\mu_{12}|^2, \qquad (3)$$

$$\frac{\partial N}{\partial t} + \frac{1}{T_1}(N - N_0) = \mathbf{E}\frac{\partial \mathbf{P}}{\partial t}\frac{2}{\hbar\omega_0}, \qquad (4)$$

where $N_0 = N_i(\rho^0_{22} - \rho^0_{11})$ is the inverse-population density in the absence of the field, $\omega_0 = \omega_{21}$, and small terms have been omitted under the assumption that $\omega_0 \gg 1/T_2$. For the case $T_1 = T_2$ these equations were obtained by Faĭn[15] and by Oraevskiĭ.[16]

The field equation (1) and the material equations (3) and (4) together with the initial condition completely describe the resonance interaction of a light pulse propagating in a medium with population inversion.

In luminescent crystals usually $\sigma_0 N_0, \gamma \ll k$, and $T_2 \gg 1/\omega_0$, where σ_0 is the cross section, defined below, for the radiative transitions between levels, and k is the wave vector. We can therefore go over to "slow" variables $\mathscr{E}$, $\mathscr{P}$, φ, ψ:[1)]

$$E = {}^1\!/_2\mathscr{E}(t, \mathbf{r}) \exp\{i[\varphi(t, \mathbf{r}) + \omega t - \mathbf{kr}]\} + \text{c. c.},$$
$$P = {}^1\!/_2\mathscr{P}(t, \mathbf{r}) \exp\{i[\psi(t, \mathbf{r}) + \omega t - \mathbf{kr}]\} + \text{c. c.} \qquad (5)$$

Then the equation for the field and the material equations reduce to the following system of five equations:

$$\frac{\partial\mathscr{E}}{\partial t} + c\frac{\partial\mathscr{E}}{\partial x} + \frac{\gamma}{2}c\mathscr{E} = 2\pi\omega\mathscr{P}\sin(\psi - \varphi),$$
$$\mathscr{E}\left(\frac{\partial\varphi}{\partial t} + c\frac{\partial\varphi}{\partial x}\right) = -2\pi\omega\mathscr{P}\cos(\psi - \varphi),$$
$$\frac{\partial\mathscr{P}}{\partial t} + \frac{1}{T_2}\mathscr{P} = \frac{\sigma_0 c}{4\pi\omega T_2}N\mathscr{E}\sin(\psi - \varphi),$$
$$\frac{\partial\psi}{\partial t}\mathscr{P} + (\omega - \omega_0)\mathscr{P} = \frac{\sigma_0 c}{4\pi\omega T_2}N\mathscr{E}\cos(\psi - \varphi),$$
$$\frac{\partial N}{\partial t} + \frac{1}{T_1}(N - N_0) = -\frac{1}{\hbar}\mathscr{P}\mathscr{E}\sin(\psi - \varphi), \qquad (6)$$

where we introduce the symbol for the cross section of the radiative transition between levels at the central frequency ω_0:

$$\sigma_0 = 4\pi T_2\omega_0|\mu_{12}|^2/\hbar c. \qquad (7)$$

Before we investigate the equations in (6), let us determine the conditions under which they go over into the usual velocity equations for radiation transfer. We integrate the polarization equation (6_3) under the initial condition $\mathscr{P}(-\infty) = 0$ and substitute the resultant expression in the field equation (6_1) and in the equation for the inverse population density (6_5):

$$\frac{\partial\mathscr{E}}{\partial t} + c\frac{\partial\mathscr{E}}{\partial x} + \frac{\gamma}{2}c\mathscr{E} = \frac{\sigma_0 c}{2T_2}\sin(\psi - \varphi)\int_{-\infty}^{t} N(t')\mathscr{E}(t')$$
$$\times \sin(\psi - \varphi)\exp\left(\frac{t' - t}{T_2}\right)dt',$$
$$\frac{\partial N}{\partial t} + \frac{1}{T_1}(N - N_0) = -\frac{\sigma_0 c}{4\pi\hbar\omega_0 T_2}\mathscr{E}\sin(\psi - \varphi)$$
$$\times \int_{-\infty}^{t} N(t')\mathscr{E}(t')\sin(\psi - \varphi)\exp\left(\frac{t' - t}{T_2}\right)dt'. \qquad (8)$$

If the changes of the field, of the inverse population, and of the phase difference of the field are negligibly small within a time T_2, then the integral in (8) reduces to $T_2 N(t)\mathscr{E}(t)\sin(\psi - \varphi)$ and the factor $\sin^2(\psi - \varphi)$ is equal to $T_2^{-2}/[(\omega - \omega_0)^2 + T_2^{-2}]$, that is, it describes a Lorentz emission line shape. Then $\sigma_0\sin^2(\psi - \varphi) = \sigma(\omega)$ is the cross section for the radiative transition at the frequency ω. Going over to the radiation flux density

$$I = \frac{1}{\hbar\omega_0}\frac{c}{8\pi}\mathscr{E}^2 \ [\text{photon/sec-cm}^2]$$

we can rewrite (8), subject to the foregoing conditions, in the usual form

$$\frac{\partial I}{\partial t} + c\frac{\partial I}{\partial x} + \gamma cI = \sigma(\omega)cIN,$$
$$\frac{\partial N}{\partial t} + \frac{1}{T_1}(N - N_0) = -2\sigma(\omega)IN. \qquad (9)$$

Thus, the assumption that the change in all the

[1)]For simplicity we confine ourselves to the one-dimensional case of propagation of linearly polarized light.

variables within the time T_2 is slow is essential when going over from the rigorous equations (6) to the velocity equations (9). In the analysis of processes that evolve during a time comparable with the transverse relaxation time T_2, it is essential to use (6).

2. Dynamics of a Light Pulse Under Nonlinear Amplification

We now consider the interaction of a radiation pulse of duration much longer than the transverse-relaxation time T_2. In this case we can use (see Sec. 1) the velocity equations (9), which reduce to a nonlinear partial differential equation:

$$\frac{\partial I}{\partial t}+c\frac{\partial I}{\partial x}=cI\left[\sigma_0 N_0\exp\left(-2\sigma_0\int_{-\infty}^{t}I(t',x)\,dt'\right)-\gamma\right], \tag{10}$$

where we neglect the spontaneous decay of the inverse population, since the pulse duration is in practice always $\tau_p \ll T_1$.

The nonstationary solution of (10) with $\gamma \neq 0$ cannot be obtained analytically. In the case of zero radiation losses ($\gamma = 0$), an analytic solution can be obtained[4, 5, 10, 11] but the solution is excessively idealized, since, as will be shown below, it is just the radiation losses which determine essentially the evolution of the pulse. However, it is possible to draw the principal conclusions without considering nonstationary solutions of (10).

The propagation equation (10) has a stationary solution of the form $I(t - x/v)$ ($v \neq c$). A stationary solution corresponds to a pulse that develops after a sufficiently long propagation in the medium and moves with velocity v; this solution is determined by

$$\dot{I}\left(\frac{1}{c}-\frac{1}{v}\right)=I\left[\sigma_0 N_0\exp\left(-2\sigma_0\int_{-\infty}^{\tau}I(\tau')\,d\tau'\right)-\gamma\right], \tag{11}$$

where $\tau = t - x/v$. Going over to the function

$$R(\tau)=\int_{-\infty}^{\tau}I(\tau')\,d\tau'$$

and integrating under the initial condition $\dot{R}(-\infty) = 0$, we reduce (11) to the form

$$\left(\frac{1}{c}-\frac{1}{v}\right)\dot{R}=\frac{N_0}{2}\left(1-e^{-2\sigma_0 R}\right)-\gamma R. \tag{12}$$

As can be seen from (12) a stationary solution with physical meaning (the total energy of the pulse $R(+\infty)$ satisfies the condition $0 < R(+\infty) < +\infty$), exists and is stable under the conditions

$$\sigma_0 N_0 > \gamma > 0, \quad v > c. \tag{13}$$

The first condition is trivial: the initial amplification of the medium $\sigma_0 N_0$ is larger than the nonzero losses γ. The second condition signifies that the propagation velocity of the developing stationary state of the pulse exceeds the velocity of light in the medium c. The physical meaning of this consists in the following.

In the stationary states the energy $\mathscr{E}_m = R(+\infty)$ of the pulse is determined from (12) under the condition $R(+\infty) = 0$, which leads to

$$1-\exp\left(-\frac{\mathscr{E}_m}{\mathscr{E}_s}\right)=\frac{\gamma}{\sigma_0 N_0}\frac{\mathscr{E}_m}{\mathscr{E}_s}, \tag{14}$$

where $\mathscr{E}_s = 1/2\sigma_0$ has the meaning of the saturation energy that reduces the inversion by a factor e. We shall henceforth confine ourselves to the case of practical interest $\gamma/\sigma_0 N_0 < 0.5$ (in practice $\gamma/\sigma_0 N_0 \ll 1$). In this case $\mathscr{E}_m \gg \mathscr{E}_s$, therefore the energy of the leading front of the pulse is perfectly sufficient for complete de-excitation of the active particles of the medium. As a result, the leading part of the pulse takes on the entire energy of the active particles, while the trailing part of the pulse propagates in the medium with loss. This mechanism ensures additional forward motion of the pulse maximum. Thus, owing to essentially nonlinear interaction between the radiation pulse and the inversely populated medium, the pulse becomes continuously deformed, thus ensuring its forward motion with effective velocity $v > c$.

The shape of the pulse in the stationary state can be obtained by approximately solving (12), expanding $\exp(-2\sigma_0 R)$ in a series up to R^3 inclusive. For the sake of brevity we do not present this solution, since all the principal parameters of the stationary state will be obtained below by a different method, and show for comparison in Fig. 1 a typical pulse shape in the stationary state, obtained by numerically solving (12).

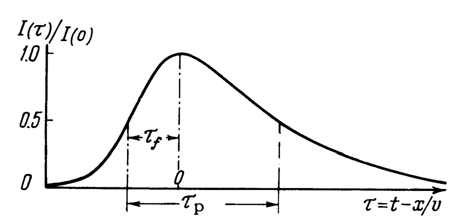

FIG. 1. Shape of a light pulse in the stationary state with a duration much longer than T_2 ($\gamma/\sigma_0 N_0 = 0.17$).

We note that one can speak of stationary states of a real pulse with leading front of finite length only in the case when the time necessary to establish a stationary state is much shorter than the total duration of the leading front. We shall show below that an exponential pulse approaches the stationary state within a time $\Delta t_s \approx (c\gamma)^{-1}$. The total length of the leading front of a Q-switched 1 laser pulse (delay time Δt_d) is in practice always much larger than $(c\gamma)^{-1}$.

Let us estimate the effective velocity of the pulse in the stationary state. Since the pulse as a whole moves with velocity v, it is sufficient to find, for example, the displacement velocity of the leading front. It follows from (11) that the leading front, propagating in the medium with practically constant inverse population N_0, grows exponentially:

$$I_f(t, x) \sim I_0 \exp(vt - x / c\tau_0).$$

In addition, the leading front satisfies the nonstationary equation (10), thus leading to the following expression for the effective pulse velocity:

$$v / c = 1 + c(\sigma_0 N_0 - \gamma)\tau_0. \qquad (15)$$

The slope of the leading front τ_0 is connected with the duration of the leading front of the pulse τ_f by the relation $\tau_f = f(\alpha)\tau_0$, where $\alpha = \gamma/\sigma_0 N_0$ is a dimensionless parameter. Examining the equations in (12) at the instants of time τ_{max}, $\tau_{max} - \tau_f$, and $\tau \to \infty$ we can obtain the function $f(\alpha)$. In the case when $\alpha > 10^{-2}$ we have

$$\frac{\tau_f}{\tau_0} = f(\alpha) \approx \frac{(1-\alpha)\ln(1/\alpha)}{(1-\alpha) - \alpha\ln(1/\alpha)}. \qquad (16)$$

Finally, we can find the total pulse duration τ_p in the stationary state. The total pulse energy is $\mathscr{E}_m \approx N_0/2\gamma$. From the condition that the pulse power be stationary at the maximum, it follows that

$$\sigma_0 N_0 \exp(-2\sigma_0 \mathscr{E}_f) = \gamma, \qquad \mathscr{E}_f = \int_{-\infty}^{\tau_{max}} I(\tau')\,d\tau',$$

where $\mathscr{E}_f$ is the energy of the leading front of the pulse. Approximately we have $\tau_f/\tau_p = \mathscr{E}_f/\mathscr{E}_m$ and consequently

$$\tau_f / \tau_p \approx \alpha \ln(1/\alpha). \qquad (17)$$

In order to trace fully the dynamics of the pulse and its approach to the stationary state, we have obtained the nonstationary solutions of (10) by numerical integration with an electronic computer. We present the results of the solution for a case corresponding to the propagation of a pulse in a ruby crystal with customary parameters. The initial pulse at the boundary of the medium has exponential form

$$I_0(t) = \frac{\mathscr{E}_0}{\tau_0} \frac{\exp(-t/\tau_0)}{[1 + \exp(-t/\tau_0)]^2},$$

which approximates well the leading front of the leading pulse of a powerful laser with instantaneous Q-switching. The real laser pulse has a somewhat asymmetrical form, owing to the long duration of the trailing fronts, but this is immaterial, since only the leading front determines the dynamics of the pulse. One of the typical solutions is shown in Fig. 2, which illustrates plots of I(t, x) in a coordinate system that moves with velocity c, after traversing a distance x in the medium with inverse population. It is clearly seen that as soon as the energy in the leading front of the pulse reaches the saturation energy $\mathscr{E}_s$, all the active particles are de-excited on the leading front and this leads to additional motion of the leading front of the pulse. Then, after traversing a distance $\sim v/c\gamma$, the energy of the pulse reaches the maximum value $\mathscr{E}_m$, and with further motion the pulse assumes a stationary form that propagates with velocity v. In addition, τ_f of the pulse in the stationary state is close to the duration of the leading front τ_f^0 of the initial pulse.

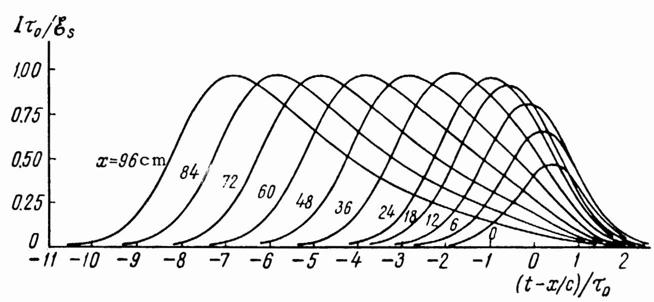

FIG. 2. Change in the shape of an exponential pulse with initial energy $\mathscr{E}_0 - \mathscr{E}$ propagating in a medium with inverse population ($\sigma_0 N_0$ = 0.2 cm^{-1}, γ = 0.03 cm^{-1}, x – distance traversed by the pulse in the medium).

Thus, the light pulse with exponential leading front of slope τ_0 propagating in a medium with inverse population, reaches after a time $\sim(\gamma c)^{-1}$ of nonlinear amplification a stationary state of duration $\tau_p \approx \tau_f[\alpha \ln(1/\alpha)]^{-1}$ moving with velocity $v = c[1 + c\tau_0(\sigma_0 N_0 - \gamma)]$. In this case the preferred amplification of the leading front does not lead to a reduction in the duration of the initial pulse.

3. Reduction of Pulse Duration

To reduce the pulse duration it is necessary to increase appreciably the slope of the leading front. To this end, one can pass the initial pulse through a shutter which opens at the instant of passage of the pulse maximum. In the case of zero initial transmission of the shutter, when total cutoff of the leading front takes place, the preferred ampli-

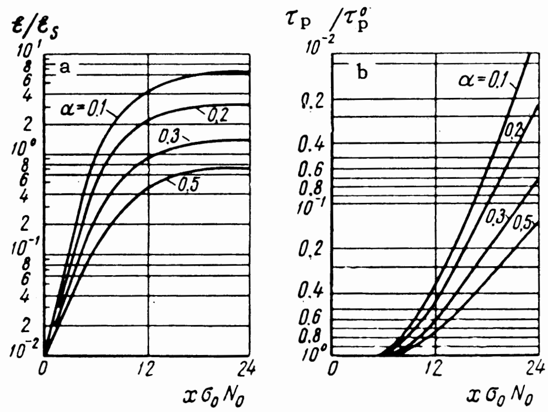

FIG. 3. Change in energy (a) and duration (b) of a pulse whose leading front is cut off by a shutter with zero transmission, propagating in a medium with inverse population ($E_0 = 10^{-2}\,\mathcal{E}_s$).

fication of the leading front leads to an effective reduction of the pulse duration. Figure 3 shows the change in duration and in the energy of the pulse as it propagates in a medium with inverse population. After the pulse energy becomes equal to the saturation energy, a sharp reduction in the pulse duration sets in. At nonzero initial transmission of the shutter, there is a definite limit on such a reduction.

The nonlinear absorption[17-19] of the leading front of the pulse as it propagates in a two-component medium (one component constitutes amplifying atoms and the other absorbing atoms) also leads to a suppression of the "runaway" of the leading front and makes possible an effective reduction of pulse duration. Figure 4 shows the change in the shape of an exponential pulse as it propagates in a two-component medium; this change was obtained by numerically integrating (10) and the equation for the number of absorbing atoms. It is clearly seen that the leading front of the pulse is continuously absorbed, bleaching the medium and ensuring effective amplification of the pulse maximum.

We note that there exists a limit beyond which the pulse duration cannot be reduced. Examination of the rigorous equations (6) shows that the shape of the pulse with minimum duration $\sim T_2$ is of the form ($\gamma/\sigma_0 N_0 < 0.5$):

$$I(\tau) = \frac{2(\sigma_0 N_0/\gamma - 1)^2 \exp(-t/\tau_0)}{\sigma_0 T_2 [1 + \exp(-t/\tau_0)]^2},$$

$$\tau_0 = \frac{T_2}{2}\left(\frac{\sigma_0 N_0}{\gamma} - 1\right)^{-1}. \qquad (18)$$

The pulse energy in the limiting stationary state $N_0(1/\gamma - 1/\sigma_0 N_0)$ is almost double the pulse energy

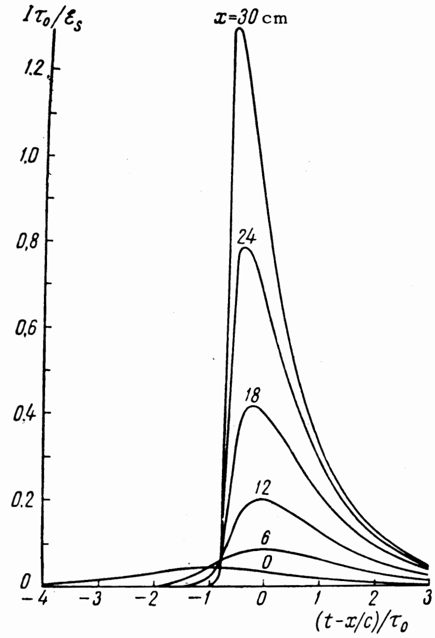

FIG. 4. Change in shape of exponential pulse propagating in a two-component medium when the initial pulse energy $\mathcal{E}_0$ is larger than the threshold ($\sigma_0^{(2)}/\sigma_0^{(1)} = 50$, $\sigma_0^{(2)}N_0^{(2)}/\sigma_0^{(1)}N_0^{(1)} = 5$, $\sigma_0^{(1)}N_0^{(1)} = 0.2$ cm^{-1}, where the index 1 pertains to the amplifying atoms, 2 pertains to the absorbing atoms, and $\gamma = 0.03$ cm^{-1}).

$\mathcal{E}_m$ at duration $\gg T_2$, for when the pulse duration is $\gg T_2$, the inverse population decreases to zero after the passage of the pulse, and in the limiting stationary state the inverse population varies in the following fashion:

$$N(\tau) = N_0 - \frac{2\gamma(\sigma_0 N_0/\gamma - 1)}{\sigma_0[1 + \exp(-t/\tau_0)]}, \qquad (19)$$

that is, the atoms become inverted after passage of the light pulse.

III. EXPERIMENTAL INVESTIGATION

1. Measurement of the Rate of Propagation of the Pulse Maximum

The propagation of a powerful light pulse was investigated experimentally with the aid of a setup consisting of a Q-switched laser and an optical quantum amplifier. The experimental setup is shown in Fig. 5.

The light pulse was produced by the laser described in [3]. A small fraction of the laser light flux was diverted by means of glass plate 5 to a coaxial photocell 6. The greater part was fed to the input of an optical quantum amplifier consisting of two ruby crystals 7 of 24 cm length and 1.6 cm diameter each. To eliminate feedback, the end faces of the crystals were cut at the Brewster

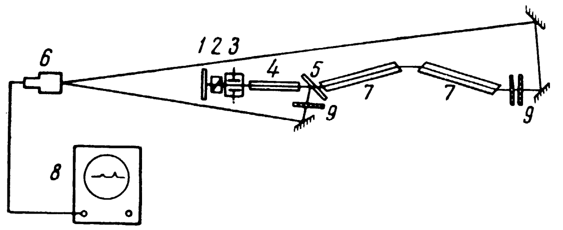

FIG. 5. Diagram of setup: 1 – laser mirror, 2 – polarizer, 3 – Kerr cell, 4 – laser ruby crystal, 5 – glass plate, 6 – coaxial photocell, 7 – amplifier ruby crystal, 8 – Sl-14 oscilloscope, 9 – neutral light filters. The input mirror of the laser is at the end of crystal 4.

angle. For a weak signal the overall amplification of the optical amplifier was approximately 50.

The light pulses before and after the amplifier were registered with the aid of a photocell 6, with the amplified pulse passing through a supplementary path of approximately 20 m and striking the photocell 56 nsec after the arrival of the input pulse.

The pulses from the photocell were registered with an S1-14 oscilloscope. The light fluxes were suitably attenuated by neutral filters 9 to ensure linear operation of the photocell and of the oscilloscope. Figure 6 shows a typical oscillogram of the input and output pulses for a weak input signal. The generator pulse was attenuated in this case by a factor 3×10^3 with neutral optical filters.

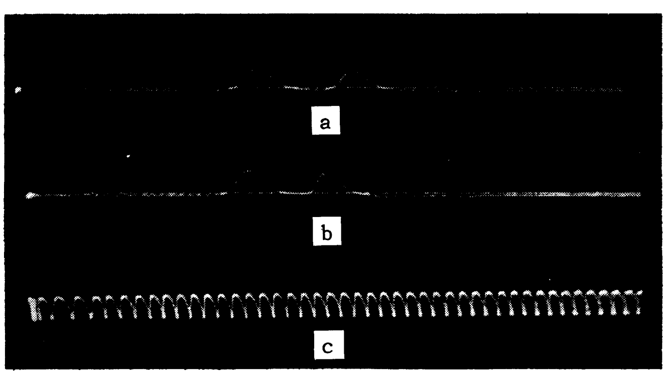

FIG. 6. Oscillograms of input (left) and amplified (right) light pulses: a – case of weak signal – linear amplification; b – case of strong signal – nonlinear amplification; c – calibration markers with period 8 nsec.

The parameters of the unattenuated input pulse were as follows: energy–1.3 J; duration at half-height–16 nsec; start of pulse determined by the instant of switching on of the shutter generator is 45 nsec away from the peak of the pulse; duration of the leading front (from the 0.5 level to the maximum)–8 nsec. Figure 6b shows an oscillogram of the input and output pulses with the signal unattenuated.

From a comparison of the oscillograms in Figs. 6a and b we can note that no appreciable shortening of the pulse takes place, but the output pulse in Fig. 6b turns out to be much closer to the input pulse than in Fig. 6a (the optical delay was not changed!). Reduction of the oscillograms yields for this shift a value $\Delta t_{exp} = 9 \pm 0.5$ nsec both for the maxima and for the half-heights of the leading fronts. In accordance with the ideas developed above, this effect is connected with the nonlinear amplification of the light pulse, which has an exponentially growing leading front.

The estimated pulse energy density reached the saturation energy $\mathscr{E}_s \approx 4\,J/cm^2$ at a distance $L = 10-15$ cm from the output end of the amplifier. (The pulse energy at the output of the amplifier reached 17 J.) The fact that the pulse leads by Δt means that it has moved $c\Delta t$ in space. The effective displacement velocity of the pulse maximum is determined from the relation $(v - c)L/c = c\Delta t$. We thus obtain in experiment $v/c = 10-15$, or a maximum displacement velocity 6–9 times larger than the speed of light in vacuum.

The theoretical value of the delay time can be estimated from formula (15) for the pulse velocity in the stationary state, since the velocity in nonlinear amplification is close to the stationary value even far from the stationary state. According to (15),

$$\Delta t_{theor} = \frac{v - c}{c}\frac{L}{c} = (\sigma_0 N_0 - \gamma) L \tau_0.$$

The slope of the input pulse is $\tau_0 \approx 4 \times 10^{-9}$ sec. Therefore $\Delta t_{theo} = 5-7$ nsec. Since the mechanism of displacement of the pulse maximum goes into operation somewhat before the pulse energy reaches $4\,J/cm^2$, the experimental lead should exceed this estimate. The experimental value obtained 9 ± 0.5 nsec does not contradict the representations developed above.

2. Shortening of Pulse Duration

As shown earlier (Sec. 3 of Ch. II), to reduce the duration of the leading front and of the entire pulse it is necessary to cut off the gently sloping part of the leading front of the initial pulse. This is done with an additional Kerr shutter placed past the generator and ahead of the diverting plate 5 (Fig. 5). The additional shutter was opened with a delay of 48–50 nsec relative to the generator shutter. The opening time of the Kerr shutter, between the 0.1 and 0.9 levels, was $\tau_{sw} = 8$ nsec. In the

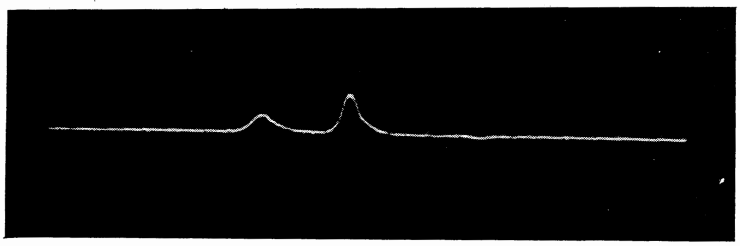

FIG. 7. Oscillograms of input (left) and amplified (right) light pulses following additional cutoff of the leading front of the generator pulse. Two-stage amplification.

closed state the shutter transmitted approximately 3% of the light ($\eta_1 = 0.03$), while in the open state it transmitted $\eta_2 = 50\%$. To compensate for the losses in the solution, an additional amplifier stage was added.

The delay between the instant of opening of the generator shutter and the appearance of the pulse maximum depends on population inversion of the crystal. It was therefore possible to cut off the pulse from any intensity level by varying the generator pump.

Figure 7 shows an oscillogram of the pulses ahead of the amplifier (but past the second shutter) and past the two-stage amplifier, when the nonlinear mode just begins to come into play. It is seen from the oscillogram that the leading front begins to shorten.

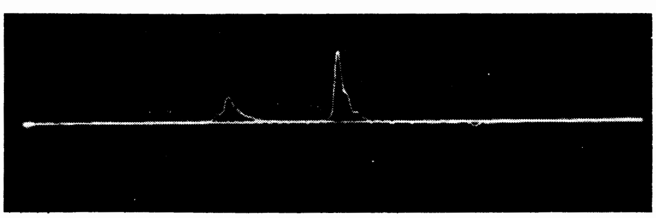

FIG. 8. Same as in Fig. 7. Three-stage amplification.

Figure 8 shows the pulses before and after amplification with three amplifier stages. In this case the pulse cutoff was from a higher level. The duration of the input pulse at half-height was $\tau_p^0 = 8.7 \pm 0.5$ nsec, and the duration of the amplified pulse was $\tau_p = 4.7 \pm 0.5$ nsec. The durations of the leading fronts from the 0.5 level to the maximum were $\tau_f^0 = 3.7 \pm 0.5$ nsec and $\tau_f = 1.9 \pm 0.5$ nsec for the input and output pulses respectively. A theoretical estimate for the possible reduction in the duration gives in our experiment a shortening by a factor of two, which agrees with the experimental value. In our experiment, the initial transmission of the additional shutter is $\eta_1 \neq 0$. This is what imposes a limit on the duration of the leading front of the amplified pulse.

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