the first term of the expansion in p:

$$q(p) \approx \frac{1}{15} a_1 p^2 + \ldots; \quad a_1 = \int_0^\infty \Gamma(x) x \, dx \approx 1. \tag{30}$$

Using the formulas (29) and (30), we can find the absorption coefficients of the regular monochromatic plane waves at $\omega = \omega_0$ in a plasma with fluctuations of the electron density. Calculation shows that at the resonance frequency, the longitudinal waves are damped according to the law $e^{-r/l}$ and the electromagnetic waves (at $k_0 l \sigma_\varepsilon^{1/2} \ll 1$), are damped like $e^{\gamma r}$, where

$$\gamma = (2\pi)^{-1/4} k_0 \sigma_\varepsilon^{1/2}. \tag{31}$$

It is useful to give the results, which are easily obtained from (29), for $\varepsilon^l_{eff}(\omega_0, k=0)$ in the case of the Cauchy probability density distribution

$$W(\xi) = \frac{1}{\pi} \frac{\sigma_\varepsilon}{\sigma_\varepsilon^2 + \xi^2}. \tag{32}$$

The fact is that an exact value of the quasielectrostatic ε^l_{eef} for this distribution has been obtained in the work of Abramovich.[10] At $k=0$ we have, from (29),

$$\varepsilon^l_{eff}(\omega_0, k=0) = 1/\alpha_0. \tag{33}$$

Finding α_0 with the help of (32), we get

$$\varepsilon^l_{eff}(\omega_0, k=0) = i\sigma_\varepsilon \tag{34}$$

—the value found by Abramovich.[10]

The author expresses his gratitude to V. I. Tatarskiĭ for significant remarks.

[1] W. F. Brown, Jr., J. Chem. Phys. 23, 1514 (1955).
[2] L. D. Landau and E. M. Lifshitz, Elektrodinamika sploshnykh sred (Electrodynamics of continuous media) Fizmatgiz, 1957, Sec. 9 [Pergamon, 1959].
[3] V. M. Finkel'berg, Zh. Tekhn. Fiz. 34, 509 (1964) [Sov. Phys. Tech. Phys. 9, 396 (1964)].
[4] I. Webman and J. Jortnner, Phys. Rev. B16, 2593 (1977).
[5] L. P. Kadanoff *et al.*, Rev. Mod. Phys. 39, 395 (1967).
[6] I. M. Lifshitz, M. I. Kaganov and V. M. Tsukernik, Uchenye zapiski KhGU (Kharkov State University) 2, 41 (1950).
[7] E. A. Kaner, Izv. vuzov. Radiofizika 2, 827 (1959).
[8] V. M. Finkel'berg, Zh. Eksp. Teor. Fiz. 46, 725 (1964) [Sov. Phys. JETP 19, 494 (1964)].
[9] Yu. A. Ryzhov, V. V. Tamoĭkin and V. I. Tatarskiĭ, Zh. Eksp. Teor. Fiz. 48, 656 (1965) [Sov. Phys. JETP 21, 433 (1965)].
[10] B. S. Abramovich, Zh. ksp. Teor. Fiz. 77, 87 (1979) [Sov. Phys. JETP 50, 44 (1979)].

Translated by R. T. Beyer

Neodymium laser with a mirror plasma-optical *Q*-switch

V. A. Batanov, D. A. Dement'ev, A. N. Malkov, A. M. Prokhorov, and V. B. Fedorov

P. N. Lebedev Physics Institute, Academy of Sciences of the USSR, Moscow

(Submitted 24 April 1979)
Zh. Eksp. Teor. Fiz. 77, 2186–2199 (December 1979)

The use of laser radiation self-action in a plasma to develop plasma-optical elements for laser resonators is considered. Experiments with a neodymium laser show that a resonator plasma can simultaneously act as a resonator mirror, a laser Q-switch, and a nonlinear element for mode locking during giant pulse generation. The generation kinetics of a neodymium laser is investigated in detail. The possibility in principle of controlling the generation regime and varying the radiation characteristics over a broad range is demonstrated. A mean radiation power $P \approx 1$–10 GW is attained in a train of nanosecond radiation spikes; the power density at the plasma mirror is estimated at $I \approx 10^{13}$–10^{14} W/cm^2. There are indications that self-focusing of the radiation occurs at the plasma mirror. One manifestation of this is the selection of transverse modes in the laser resonator.

PACS numbers: 42.55.Px, 42.60.Da

1. INTERRELATED GENERATION OF HIGH-POWER LASER RADIATION AND OF A REFLECTING RESONATOR PLASMA

A new physical effect, the generation of high-power laser pulses in a resonator with plasma mirror and interrelated with high-temperature heating of the plasma mirror was reported earlier.[1] In the reported experiments the laser plasma produced in a gas or solid target produced in the focus of a lens performs two functions. First, it is the mirror of the resonator (resonator plasma) and produces positive feedback for the development of the lasing. The strong nonlinear increase of the plasma reflection in the course of the growth of the lasing amplitude ensures rapid switching-on of the Q of a resonator with such a plasma mirror and generation of giant laser pulses. In other words, during the course of generation the dense reflecting plasma serves as a mirror laser Q-switch. Second, the plasma mirror of the resonator serves also as a target for the laser radiation, whose absorption produces in the target a high temperature and a near-criti-

 0038-5646/79/121049-08$02.40

cal plasma density for the laser-generation wave.

The preceding experiments[1] were performed with a single-channel neodymium-glass laser with aperture 4.5 cm, with the ordinary electro-optic Q-switch replaced by a laser shutter based the plasma produced by optical breakdown of atmospheric air or on plasma from graphite or polyethylene $(CH_2)_n$ targets placed in the focus of an $F = 50$ cm lens. This laser generate giant pulses of wavelength $\lambda = 1.06$ μm and amplitude 10^2–10^3 times larger than that of the initiating radiation, as well as hard x-rays from the plasma mirror from all the employed targets and neutrons if a deuterated polyethylene target $(CD_2)_n$ was used.

The development of such a laser with a plasma-mirror shutter has uncovered interesting prospects for the use of a laser plasma in the technique of high-power pulse lasers. These prospects are connected with the possibility of producing extremely simple and effective large-aperture Q-switched lasers.

The results of Ref. 1 were extended in Ref. 2 to include CO_2 lasers. A single-channel electro-ionization CO_2 laser with a solid-target plasma mirror was used to obtain and investigate lasing with a time structure governed by the mode locking of the plasma-mirror resonator.[2-5] Similar generation by a CO_2 laser, but using plasma produced inside the resonator by optical gas breakdown, was obtained also in Ref. 6.

In the present study (see also Refs. 7 and 8) we have resolved the temporal structure, due to resonator mode locking, of a giant radiation pulse of a neodymium laser with a mirror plasma-optical shutter. We have shown by the same token that the resonator plasma can possess simultaneously the properties of a lser shutter and of a nonlinear element that locks the resonator modes.

The cited references constitute at present the entire literature on the use of backreflection of radiation in a laser plasma to generate high-power radiation and to simultaneously heat the plasma itself by the generated radiation. From among other publications, mention should be made of papers that report, for CO_2 and neodymium lasers, data on the positive-feedback mechanism connected with reflection of the radiation from the plasma into the laser aperture.[9-12] This mechanism appears in a number of reported experiments[9,11,12] in the form of additional lasing against the background of the fundamental lasing in the absence of the plasma. Realization of a CO_2 laser operating with an arc-discharge-plasma mirror is reported in Ref. 10. In contrast to the last-cited papers, this feedback mechanism was used in Ref. 1 to realize a new radiation self-action process, namely generation of laser pulses via backreflection of the light from the resonator laser plasma.

The present paper deals primarily with the laser aspect of the effect of interrelated generation of a powerful radiation pulse and of a reflecting resonator plasma, and its purpose is to investigate the time and energy characteristics of the generation processes in single-channel neodymium lasers with plasma-optical mirror shutters.

2. EXPERIMENTAL RESULTS

1. We investigated two four-stage single-channel generator setups (Fig. 1). Each stage consisted of a GOS-1000B illuminator with an active element of neodymium glass GLS-1 of 4.5 cm diameter and 68 cm length, with end faces cut at an angle 5° to the rod axis. The stage was pumped by four IFP-20 000 xenon lamps with discharge-gap length 56 cm. The maximum pump energy in the illuminator was 80 kJ within a time ~1 msec.[1)]

In the first setup (Fig. 1a)[1] the optical axis of the generator is set by a confocal telescopic system of mirrors with focal lengths $f_1 = 1$ and $f_2 = 2$ m, inside of which is located the first stage of the laser. The plasma mirror is produced on a target at the focus of an $F = 50$ cm lens. The targets are graphite in air. Data of the preceding experiments,[1] with atmospheric air as the target, are also used.

In the second setup (Fig. 1b), the first stage does not differ in any way from the remaining ones; the optical resonator has a simpler configuration and consists of a flat mirror, which determines the direction of the axis, and a mirror plasma shutter similar to that of the first system. The graphite target is displaced away from the laser behind the focus of the lens by 1–2 cm.

In the experiments we investigated the lasing processes. To monitor the lasing, the radiation was diverted by a plane-parallel glass plate located between the focusing lens and the fourth stage of the setup (Fig. 1). The radiation was recorded, in a millisecond time scale, with the aid of FÉU-28 photomultipliers and an S8-2 oscilloscope. Simultaneously with oscillogram of the amplitude we recorded the time variation of its integral. The absolute calibration of the integrated amplitude was against a PTÉK-L calorimeter that recorded the total radiation energy. The structure of the giant pulses was investigated at nanosecond resolution with an FK-15 photocell and a 6LOR-04 oscilloscope.

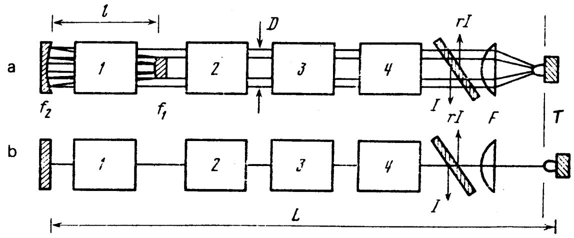

FIG. 1. Two optical setups of single-channel neodymium lasers with mirror plasma-optical shutters in the form of solid graphite targets in air: a) setup with confocal telescopic system of mirrors in the first stage b) setup with flat mirror: 1–4 stages of the setup, T) target, F) focusing lens, I and rI) beams extracted to measure the intensity of the radiation incident on the target and of the reflected radiation, L) optical length of the indicated setups, l) optical length of the telescopic stage, D) aperture of light beam.

2. The lasing regime of a neodymium laser with a plasma mirror on a graphite target, observed in the setup with the telescopic driving stage, are shown in Fig. 2. Reproducible generation of one and more giant pulses was obtained (Figs. 2a and 2b) during the time of the pump action. The giant pulses are preceded by smooth quasicontinuous lasing (Fig. 2a and 2b). Quasi-continuous lasing is observed also after the giant pulses for a time ~0.1 msec needed to restore the threshold level of the inversion in the active medium (Fig. 2a). Smooth lasing occurs also in the absence of giant pulses (Fig. 2c). The amplitude of the giant pulses exceeds by 10^2–10^3 times the level of the smooth lasing (Figs. 2a and 2b).

A special role in the establishment of the lasing regime is played in these experiments by the first stage of the setup, which is in a confocal system of mirrors (Fig. 1a). If its pumping is lower than a certain threshold value ($U_1 = 0.20$ in the experiments represented in Fig. 2), then no giant pulses are produced even at maximum pumping of the succeeding three stages. At the same time, smooth quasicontinuous lasing can be observed, characterized by a relatively low pumping threshold of the stages of the setup. This regime, which sets in as a result of the positive feedback in the active medium + target system as a result of reflection of light from the surface of the target into the aperture of the active medium, is accompanied by a plasma flare and by a characteristic crater on the target. The observed pattern of the radiation (Fig. 2c) differs sharply from the spiked free lasing in the absence of a target, when the first stage operates as a driving generator and the succeeding ones act as amplifiers. The suppression of the spiked modulation and the smoothing of the radiation pulse in a system with a target is ensured by the drooping plot of the reflection r against the light intensity I at the target ($dr/dI < 0$), which is usually the case when target surfaces are damaged and ionized by radiation of relatively low intensity $I \sim 10^7$–10^{10} W/cm^2 (see, e.g., Ref. 13).

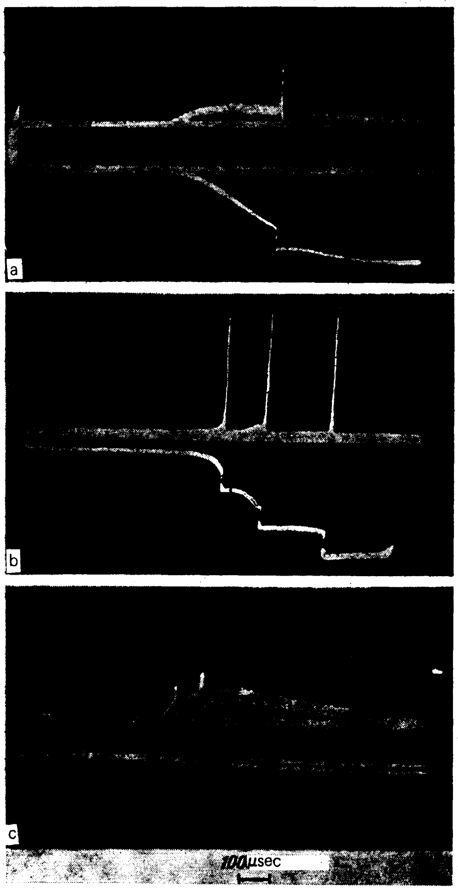

FIG. 2. Plasma-mirror lasing regimes (upper traces) and time dependence of the radiation energy (lower traces) in a neodymium laser with driving telescopic stage: a) single-pulse regime (pumps $U_1 = 0.20$, $U_2 = U_3 = U_4 = 0.36$, total radiation energy 350 J, single-pulse energy 77 J); b) regime with repeating giant pulses (pumps $U_1 = 0.25$ and $U_2 = U_3 = 0$, and $U_4 = 0.69$, total radiation energy 131 J, energy in the three gigantic pulses 76 J at 21, 17, and 28 J respectively in the first, second and third pulses); c) regime of smooth quasicontinuous lasing (the pump differs from case *a* in that U_1 is somewhat smaller, the radiation energy is close to the case *a*). Here and below, in all the millisecond scans, instrumental factors cause a lowering of the amplitude of the gigantic pulses compared with the quasi-continuous generation.

A sufficient condition for the appearance of giant pulses is that the pump threshold of the first stage be exceeded ($U_1 \gtrsim 0.20$) and that a certain minimum gain be produced by the succeeding stages (for example, $U_{2,3,4} \gtrsim 0.16$ or $U_2 \gtrsim 0.48$ and $U_{3,4} = 0$). The number of giant pulses generated during the pump time is determined only by the magnitude of the pump of the first telescopic stage and does not depend on the pumps of the three other stages. Thus, a small change in the low pump level of the telescopic stage (compared with those of the other three stages) of influences decisively the observed lasing picture (Fig. 2). This means that despite the lower pump, and hence the lower population inversion, than in the other stages, the first stage makes the principal contribution to the linear gain of the entire setup, because of the multiple passage of the radiation through the active medium in the telescopic resonator, meaning also the principal contribution towards the attainment of the thresholds of the observed lasing processes. We note that a practical consequence of this fact is the favorable possibility of controlling the number of generated giant pulses (Fig. 2) in this laser with mirror plasma-optical shutter.

3. The energy and temporal characteristics of the multipulse regime (Fig. 2a) are shown in Figs. 3 and 4. It is seen that an increase of the total lasing energy to a value $E \approx 500$ J as a result of pumping the second, third, and fourth stages (at a fixed pump of the first stage $U_1 \gtrsim 0.20$) does not increase the energy of the single pulse. In the interval 180 J < E < 460 J the energy of the giant pulse is approximately constant at $E_{\text{pulse}} \approx (63 \pm 15)$ J. The scatter is due to the insufficient reproducibility of the experimental conditions.

The result can be explained as being due to the weak dependence of the population inversion N_{inv} in the active medium on the pump at the instant when the single pulse is generated. This value of N_{inv} is formed by the free lasing which has a low threshold and precedes the gigantic pulse (Fig. 2a), and is given, for neodymium

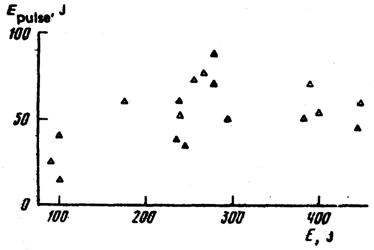

FIG. 3. Ratio of the energy of the single pulse E_{pulse} to the total radiation energy E in the laser setup with the telescopic stage.

glass, by the relation

$$N_{inv}/N = Q\tau/(1+W\tau),$$

where N is the density of the neodymium ions in the glass, τ is the lifetime of the neodymium ions on the upper level of the working transition, Q is the effective probability of the transition and describes the rate of pumping of the neodymium ions to the upper levels of the working transition, and W is the probability of the induced transition at the radiation wavelength and is approximately equal to

$$W = \frac{NQ\sigma}{\alpha}(1-e^{-\alpha x})$$

(α is the resonator loss coefficient at the lasing frequency and x is the length of the resonator with the active medium). At $W\tau > 1$, which corresponds in Fig. 3 to $E \gtrsim 200$ J, at a free lasing duration $t \sim \tau \sim 1$ msec, and at a radiative cross section of the working transition of the GLS-1 glass $\sigma = 2.5 \times 10^{-20}$ cm^2, the population inversion at the instant of the single-pulse generation

$$N_{inv}/N \approx Q/W \approx \alpha/\sigma(1-e^{-\alpha x})$$

is in fact, as stated above, independent of Q and is determined only by the parameters of the active medium in the resonator. This value yields also an estimate, corresponding to the experimental data of Fig. 3, of the radiation energy in the single pulse:

$$E_{pulse} \approx h\nu N_{inv} V \sim 100 \text{ J}.$$

At an active-medium volume $V \approx 3 \times 10^3$ cm^3, a loss $\alpha \approx 2 \times 10^{-3}$ cm^{-1} in the glass, a photon energy $h\nu = 2 \times 10^{-19}$ J, and under the condition $\alpha x > 1$.

Since the population of the upper working level of the active medium, which determines the energy of the giant pulse, is depleted by the quasicontinuous lasing, it makes sense to suppress the free energy so as to increase the contribution of the giant pulses to the total radiation energy. In fact, by placing into the resonator a solution of a standard dye (No. 3955) in nitrobenzene, with saturable absorption at an initial transmission 1.1% at the lasing wavelength, it is possible to obtain a giant pulse with energy 165–175 J at a total radiation energy of approximately 190 J (Fig. 4a). This experiment demonstrates the possibility of reducing the background and increasing the energy contrast in the single-pulse regime.

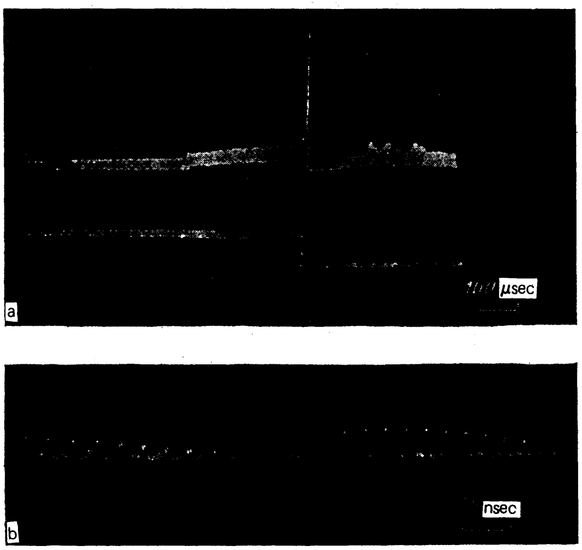

FIG. 4. Temporal characteristics of the single-pulse regime in the laser setup with telescopic stage. a) Time dependence of the amplitude (top) and of the energy (bottom) of the generation, with a saturating filter made of dye No. 3955, placed in the resonator between the first and second stages; pumps $U_1 = 0.20$ and $U_2 = U_3 = U_4 = 0.75$; $E = 190$ J, $E_{pulse} = 155$–175 J. b) Mode locking in resonator made up of the plasma-mirror shutter on a graphite target in air and the telescopic mirrors of the first stage of the setup, under pump conditions close to the experiment of the Fig. 2a.

4. An analysis of the time structure of the giant pulses shows the presence of mode locking in the resonator with the plasma mirror. A time scan of the single pulse, corresponding to the pump conditions close to the experimental data of Fig. 2a, is shown in Fig. 4b. The single pulse consists of two or three trains of nanosecond radiation peaks of duration $t_0 \leq 2$ nsec spaced $T_1 = 10$ nsec apart in the train. The interval between the maxima of the train envelope is $T_2 = 150 \pm 10$ nsec. These data attest to mode locking in the compound resonator made up of the telescopic resonator with optical length $l = cT_1/2 = 1.50$ m and the resonator with optical length $(L-1) + k \cdot 2l = cT_2/2 = 21$–$24$ m, made up of the plasma mirror on the target and one of the mirrors of the telescope. Here $L = 10$ m is the optical length of the setup, k is the number of double passes of the radiation in the telescopic resonator (see Fig. 1a). The last equation yields $k = 4$–5. Since the number of passes of a beam with aperture D and divergence φ in the telescopic resonator with magnification $m = f_2/f_1$ is limited by the condition

$$\varphi m^k = D/m^k f_1,$$

it follows that the value of k obtained from the oscillogram can be used to estimate the beam divergence φ in the resonator with the plasma mirror. At $f_1 = 100$ cm, $m = 2$, $k = 4$, and $D = 4.5$ cm we obtain $\varphi \approx 2 \times 10^{-4}$. We note that this value is close to the angle $\sim 2\lambda/A$ of the diffraction of the beam by the annular aperture $A \approx D(1 - 1/m)/2$ of the mirror telescope. Thus, the length $L + (2k-1)l$ indicated above is the real length of the resonator with the plasma mirror, corresponding to the time interval T_2, and the mirror system of the first stage produces a beam with a divergence $\varphi \approx 2 \times 10^{-4}$ close to the diffraction value.

Thus, in a neodymium laser the plasma on the target serves simultaneously as a mirror laser shutter and as the nonlinear mode-locking element. This conclusion is valid also for the case of generation of a laser based on reflection from the plasma of optical breakdown of atmospheric air (without a solid target).[1] The oscillograms obtained in this case are similar to those of Fig. 4b.

The foregoing data lead to two important consequences. First, the number of radiation passes $2k = 8$ in the first telescopic stage, obtained essentially from experiment, makes quantitatively meaningful the statement made above in the part 2 of the present section, that the dominant contribution to the total linear gain of the fours stages, and by the same token to the attainment of the threshold of the entire setup, is made by the first stage. In fact, the linear gain of the second, third, and fourth stages at pumps $U_2 = U_3$, $U_4 = 0.36$ amounts to $\exp[(\beta - \alpha) \cdot 3l] \approx 27$, whereas the linear gain of the first telescopic stage at a smaller pump $U_1 = 0.20$ (see Fig. 2a) is much higher: $\exp[(\beta_1 - \alpha) \cdot 8l] \approx 100$. The pumped length of the stage is $l = 50$ cm, and the loss in the glass is $\alpha = 0.002$ cm^{-1}. The gain $\beta = 0.021$ cm^{-1} at the stage pumps $U_{2,3,4} = 0.36$ is determined from direct measurements that have yielded $\exp[(\beta - \alpha)l] = 3$. The value $\beta_1 = 0.012$ cm^{-1} at a pump $U_1 = 0.20$ was obtained by decreasing the measured value of β in the ratio $U_1/U_{2,3,4}$.

Second, the value $\varphi \approx 2 \times 10^{-4}$ enables us to estimate the dimension of the focusing spot on the target, $d = F\varphi \sim 10^{-2}$ cm, and accordingly the radiation intensity on the plasma mirror. Since the energy 70 J (Fig. 2a) goes to almost 30 peaks with approximate duration 2 nsec (Fig. 4b), the nanosecond-pulse power averaged over the entire series is not lower than 1 GW. Correspondingly, the intensity of the radiation on the target is of the order of $I \gtrsim 10^{13}$ W/cm^2. The obtained value of I is apparently in agreement with the previously attained[1] plasma parameters.

5. The results of the experiments with the second variant of the optical setup of the neodymium laser with plasma mirror (Fig. 1b) are shown in Fig. 5. Just as in the case of the first setup (Fig. 1a), free lasing on the target was obtained (Fig. 5a) as well as generation of giant pulses (Fig. 5b). In both setups, the radiation has a temporal fine structure due to the mode locking by the nonlinearity of the plasma mirror of the resonator (Figs. 5c and 5d). It is interesting that, in free lasing, stable mode locking of the resonator with the plasma mirror was observed also on the leading front of the smooth millisecond pulse (the initial wiggle on the oscillogram of Fig. 5a), as well as in other sections of the smooth pulse in its central part. The generation of the giant pulses contributes to the production, on the target, of a well developed damage crater in the focusing zone. The progressive development of the crater in a series of several successive experiments without change of target reveals an increase in the number of giant pulses during the pump time, from a single pulse at the start of the series to 5–7 pulses at the end.

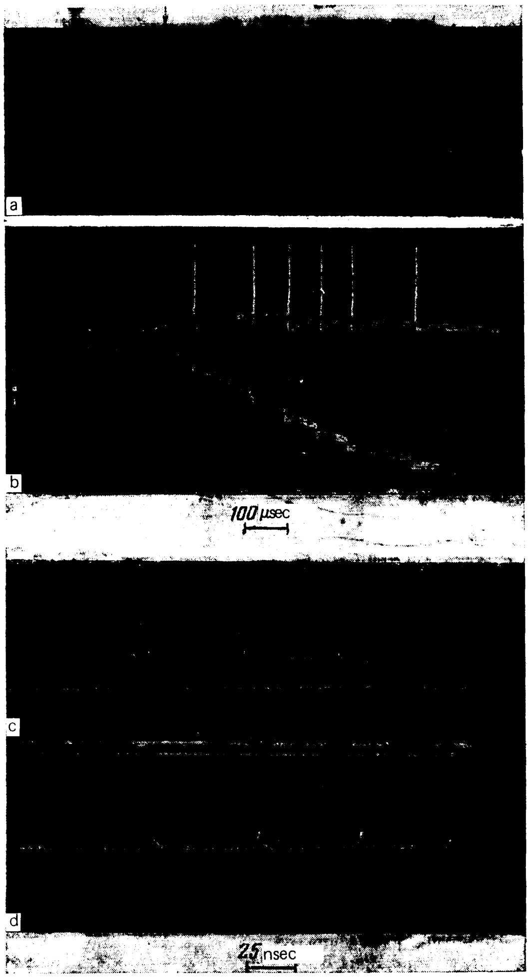

FIG. 5. Lasing regimes and temporal fine structure of the emission of a neodymium laser with a flat-mirror resonator and a plasma on a graphite target in air (Fig. 1b). a) Free generation regime on the target in the focal plane; radiation energy $E = 350$–400 J at pumps $U_1 = U_2 = U_3 = U_4 = 0.64$. b) Multiple Q switching, during the pump time, with the mirror plasma shutter on a target with a crater, placed 1.5 cm behind the focus; the total radiation energy and the pumping conditions are the same as in case a; the energy of the gigantic pulses, 190 J, consists of the first four pulses with 33 J each, two succeeding pulses of 25 J each, and one pulse of 8 J. (The lower oscillograms a and b show a time dependence of the radiation energy.) c) Mode locking of laser resonator in the free lasing regime; the scanning pertains to the instant of time marked by the arrow on the oscillogram a. d) Mode locking of resonator in gigantic radiation pulse (case b); the arrows show the pulses of the radiation reflected from the plasma on the target, delayed by 14 nsec relative to the incident radiation; 8% of the amplitude of the incident radiation is fed to the receiver.

A giant pulse of duration 300–500 nsec (Fig. 5d) consists of 4–7 shorter peaks of radiation, with duration not more than 2.5 nsec, with a repetition time $T = 2L/c = 65$ nsec (±10%) where $L = 10 \pm 0.1$ m is the optical length of the resonator. Since the total energy 30 J (Fig. 5b) goes to 4 peaks (Fig. 5d), the average power on the target in a single-nanosecond peak is higher than in the case of the setup of Fig. 1a, and amounts to not less than 3 GW. Two circumstances point to a high intensity of the light on the plasma mirror in the setup with a flat mirror compared with the setup with the telescopic stage. These are, first, harder x radiation (a signal is registered with the same receiver at larger distances from the target in air) and, second, higher reflection from the plasma mirror.

Data on the reflection are shown in Fig. 5d. It is seen from the oscillogram that the reflection is maximal in the most powerful spikes, and a small difference between the spike amplitudes corresponds to a very large difference between the reflected signals. This means, first, that the function $r(I)$ is strongly nonlinear, and second, that $dr/dI > 0$ in the region of the action of the plasma-mirror shutter. The absolute value of the reflection in the most powerful nanosecond peaks of the radiation reaches, in the experiments with the graphite target in air, values $r_{max} \approx 0.003$ and 0.04 for the first and second variants of the optical setup of Fig. 1, and the corresponding electron density, calculated from the rejection coefficient assuming a sharp (in the radiation wavelength scale) plasma boundary reaches $n_e = 0.2n_{cr}$ and $n_e = 0.55n_{cr}$ respectively ($n_{cr} = 10^{21}$ cm^{-3}). It should be noted that the measured values of the reflection coefficient agree well with estimates of the lasing thresholds in which account is taken of the real lengths of the linear gain and of the losses in the active elements. Estimates yield for the setup of Fig. 1 a value $r \sim 10^{-3}$, and the setup of Fig. 1b a value $r \sim 10^{-2}$. The obtained values of the reflection of the plasma mirror, together with the observation of hard x rays synchronized with the laser pulses and duplicating the time pattern of the radiation (Figs. 4b and 5d), offer evidence that the plasma mirror is heated to a high temperature.

6. In conclusion, we present data on the dynamics of the development of the giant pulse. To this end we turn to Ref. 1, in which the optical shutter using the experiments was the reflecting plasma of optical breakdown in atmospheric air. During practically the entire time of the pump, the laser (with the setup of Fig. 1a) operates in the absence of coupling with the target and generates a sequence of microsecond radiation spikes. The first stage is in this case the laser, and the three succeeding stages are amplifiers. The coupling with the target arises only in one of the spikes, which causes optical breakdown of the air, whose plasma does in fact become the reflecting target. At that instant the setup begins to operate as a single laser with a mirror plasma shutter. The intensity of the radiation of the initiating spike lies in the interval $I \sim 10^{10}$–10^{11} W/cm^2, inasmuch as the energy in the spike is $E \sim 1$–10 J, the spike duration is $t \sim 10^{-6}$, and the beam divergence of the setup of Fig. 1a (Ref. 14) ensures a focusing down to $d \approx 10^{-2}$ cm. The optical-breakdown threshold of air at 1 atm when focused by a lens of $F = 50$ cm is close in these experiments to the value $I_{thr} \approx 0.5 \times 10^{10}$ W/cm^2.[15] It is of interest to note that in some experiments the air broke down during the pump time two or more times, but lasing on the plasma mirror and x-radiation were observed only in one of the breakdowns, and not always in the first.

It is seen from the oscillogram of the giant pulse in Fig. 6 that is preceded by a long (~1 μsec) precursor pulse that gives way to the short rapidly rising (within 0.2–0.3 μsec) main pulse. The absorption in the plasma produced by the precursor pulse rises. On the leading front of the main pulse, the signal radiation passing through the plasma is increased on account of the pulse-amplitude growth that occurs ahead of the growth of the absorption of the plasma mirror. However, the transmission of the plasma decreases soon past the maximum of the main pulse and then practically vanishes before this pulse terminates. The radiation single pulse has a temporal fine structure due to the mode locking of the laser resonator with the plasma mirror.

3. PHYSICAL MECHANISM OF GENERATION OF GIANT PULSES OF LASER RADIATION ON A PLASMA MIRROR

The transition from the free lasing regime on the target to the regime of giant pulses can apparently be characterized by a certain threshold value of the intensity of the radiation on the target, I^*, starting with which the smooth quasicontinuous lasing becomes unstable; as a result of the self-action processes on the plasma mirror, the reflection of the radiation increases, the Q factor changes abruptly, and the giant pulse is generated. To explain the existence of two lasing regimes it is necessary to invoke physical processes that produce an intensity dependence of the radiation with a minimum at the point $I = I^*$. The decrease of the reflection to the minimum in the region of low intensities $I < I^*$ can be due to the decrease of the Fresnel reflection from the surface of the target as a result of disintegration and ionization of the surface layer. At $I > I^*$ the process matching of the growth of the intensity of the beam and of the back reflection of the radiation into the laser aperture can be taken to be the increase,

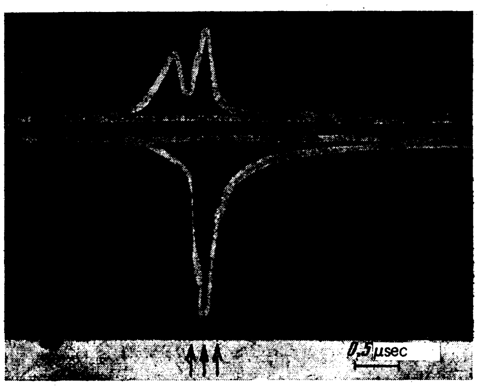

FIG. 6. Single-pulse generation by neodymium laser with a driving telescopic stage and a mirror optical shutter made up of atmospheric-air optical-breakdown plasma at stage pumps $U_{1,2,3,4} = 0.36$–0.49; envelopes of radiation incident on the focusing lens $F = 50$ cm with aperture 12 cm (below) and passing through the plasma into the aperture of a similar lens confocal with the first (top). The arrows mark the unresolved maxima connected with the mode locking of the resonator with the plasma mirror (see Fig. 4b).

due to the increased radiation amplitude, of the electron density of the plasma to a value close to the critical $n_{cr} = 10^{21}$ cm^{-3}, which can lead also to stimulated Mandel'shtam–Brillouin scattering.

An increase of the plasma density and of its corresponding reflections can be obtained also by cumulative compression of the plasma in the target crater. This process, however, seems to influence the physics of the lasing (see part 5 of Sec. 2) in a limited region of the experimental conditions. For example, experience with plasma initiation by a beam from an additional laser shows that giant pulses can be produced on a solid target with practically no crater formation. It is also possible to attribute to cumulative effects the results of experiments in air[1] when screening of half of the annular beam section did not change substantially the picture of the phenomenon. In addition, cumulative compression is probably difficult to reconcile with the relatively slow rate of development of the interrelated generation of the radiation and the growth of the reflection by the plasma.

In the discussion of the physics of the generation of giant pulses in lasers with plasma mirrors, it is necessary to point out a new process that can play an important role here. This is self-focusing of the generated radiation by the plasma mirror and simultaneous decrease of the transverse dimensions d of the plasma mirror itself. This decreasing-with-time self-concentration of the beam on a plasma mirror having an aperture is due to the fact that the monotonically increasing function $r(I)$ is transformed by the target into a monotonically decreasing time dependence $d(t)$ as a result of the presence of a maximum in the power distribution in the focus of the lens, and also of the rapid response of the dense hot plasma to changes in the radiation intensity. An important consequence of this process is the rapid acceleration of the rate of switching-on of the plasma mirror with increasing radiation at a limited growth of the lasing amplitude. Thus, the generation of a giant pulse can consist of two stages: first a relatively slow self-consistent stage that terminates by the turning of the plasma mirror, followed by a rapid state of development of the main large-amplitude pulse. This picture corresponds to the oscillogram of Fig. 6, which shows clearly a precursor pulse of duration ~1 μsec and an abrupt leading front of the main pulse of duration ~0.1 μsec.

Another important consequence of the results above is a new property of the plasma mirror, namely the self-selection of the transverse modes of the resonator and formation of an angle spectrum of the radiation by the plasma. The decrease of the aperture of the plasma mirror during the stage of the precursor pulse can occur in practice up to the diffraction limit $d \sim \lambda$, if processes in the heated plasma, distortions of the radiation (say, in the active medium of the laser) or additional diaphragms do not limit the value of d before this diffraction limit is reached. This consequence agrees with the fact reported in part 4 of Sec. 2, that the value $\varphi \approx 2 \times 10^{-4}$ rad for the setup with the telescopic stage agrees with the angle of diffraction by the aperture of the telescope. It can also be explained why hard x rays were observed in the setup with the flat rear mirror (part 5 of Sec. 2), since it points to the visibility, in this setup, of a mechanism of generation of a high-density light flux by the target. The minimum aperture of the plasma mirror is limited in this case by the distortion of the radiation of the active medium and is estimated at $\varphi \sim 10^{-4}$ rad.

To conclude this section, we examine the stage of the main pulse consisting of a series of nanosecond peaks that result from the resonator mode locking by the nonlinearity of the plasma mirror. The duration of the nanosecond pulse at a plasma expansion velocity $v \sim 10^7$ cm/sec (Ref. 1) is $t_0 \sim d/v \sim 1.0$ nsec, which agrees with the experimental data. It appears that the minimum duration of the short peak can be obtained by reaching the diffraction limit in the plasma-mirror aperture, and amounts to $t_{0\,min} \sim \lambda/v \sim 10^{-11}$–$10^{-12}$ sec.

Since the repetition period of the short radiation peak should not exceed the duration t of the giant pulse, the optical length L of the resonator is bounded by the conditions $t \gtrsim 2L/c$. At the observed value $t \approx 300$ nsec we obtain $L \lesssim 50$ m. At larger plasma-mirror resonator lengths the temporal fine structure of the giant pulse cannot be realized.

4. CONCLUSION

We have revealed in this paper a number of fundamental regularities of the processes that occur in laser systems with plasma reflections, and obtained new results on the use of laser plasma to control the generation by the laser itself. The plasma, which serves as a resonator mirror can at the same time incorporate the properties of an optical shutter for Q switching (with practically unlimited "beam endurance") and a nonlinear element for mode locking. Lasers with plasma elements are distinguished by their simplicity, the radiation parameters can be easily controlled by varying the optical setup, the pump level, or the focusing conditions.

A plasma-optical shutter makes it possible to synchronize and combine on the target beams from several laser channels. This was first demonstrated in Refs. 7 and 8 for nanosecond-pulse trains in experiments in which the setup consisted of 4 single-channel neodymium lasers with a common plasma mirror. The feasibility of replacing both mirrors of the resonator by plasma reflectors was also demonstrated. Both the millisecond-pulse regime and the Q-switching regime with mode locking were realized with a plasma resonator.

A number of characteristics of the radiation in the investigated setups and the parameters of the obtained laser plasma have come close to the record values obtained with high-power lasers even in the first experiments. The results demonstrate the promise offered by lasers of this type. At the same time, the experiments point to the need of improving certain parameters of the radiation, namely the contrast of the nanosecond pulses, the reproducibility of their characteristics, realization of maximum radiation energy, and development of a method for controlling the instant of Q switching in a resonator with a plasma mirror.

[1)] The stage-pump energies U are indicated here and below as fractions of the maximum value. Thus, the $U_2 = 0.40$ means that the pump energy of the second stage is 32 kJ.

[1] V. A. Batanov, K. S. Gochelashvili, B. V. Ershov, A. N. Malkov, P. I. Kolisnichenko, A. M. Prokhorov, and V. B. Fedorov, Pis'ma Zh. Eksp. Teor. Fiz. 20, 411 (1974) [JETP Lett. 20, 185 (1974)].
[2] N. G. Basov, I. A. Bereshnoi, V. A. Boĭko, V. A. Danilychev, V. D. Evorykin, V. V. Ignat'ev, I. V. Kholin, and A. Yu. Chugunov, Pis'ma Zh. Tekh. Fiz. 1, 1105 (1975) [Sov. Tech. Phys. Lett. 1, 473 (1975)].
[3] V. A. Boĭko, V. A. Danilychev, V. D. Evorykin, T. G. Ivanova, I. V. Kholin, A. Yu. Chugunov, Kvantovaya Elektron. (Moscow) 4, 1307 (1977) [Sov. J. Quantum Electron. 7, 735 (1977)].
[4] N. G. Basov, V. A. Boiko, V. A. Danilychev, V. D. Evorykin, A. N. Lobanov, A. F. Suchkov, I. V. Kholin, and A. Yu. Chugunov, Kvantovaya Elektron. (Moscow) 4, 1761 (1977) [Sov. J. Quantum Electron. 7, 998 (1977)].
[5] N. G. Basov, V. A. Boiko, V. A. Danilychev, V. D. Evorykin, I. V. Kholin, and A. Yu. Chugunov, Kvantovaya Elektron. (Moscow) 4, 2268 (1977) [Sov. J. Quantum Electron. 7, 1300 (1977)].
[6] M. Piche and P. A. Belanger, Opt. Commun. 24, 158 (1978).
[7] V. A. Batanov, A. N. Malkov, A. M. Prokhorov, and V. B. Fedorov, FIAN Preprint, No. 78, 1978.
[8] V. A. Batanov, A. N. Malkov, A. M. Prokhorov, and V. B. Fedorov, Izv. Akad. Nauk SSSR Ser. Fiz. 42, 2499 (1978).
[9] K. Büchl, J. Appl. Phys. 43, 1032 (1972).
[10] N. V. Karlov, V. M. Kommisarov, G. P. Kuz'min, and A. M. Prokhorov, Pis'ma Zh. Eksp. Teor. Fiz. 16, 95 (1972) [JETP Lett. 16, 65 (1972)].
[11] P. Mulser, R. Sigel, and S. Witkowski, Phys. Rep. 6, 187 (1973).
[12] P. E. Dyer, S. A. Ramsden, J. A. Sayers, and M. A. Skipper, J. Phys. D: Appl. Phys. 9, 373 (1976).
[13] N. G. Basov, V. A. Boiko, O. N. Krokhin, O. G. Semenov, and G. V. Sklizkov, Zh. Tekh. Fiz. 38, 1973 (1968) [Sov. Phys. Tech. Phys. 13, 1581 (1968)].
[14] V. A. Batanov, I. A. Bufetov, S. B. Gusev, B. V. Ershov, P. I. Kolisnichenko, A. N. Malkov, Yu. P. Pimenov, and V. B. Fedorov, Kvantovaya Elektron. (Moscow) 1, 1544 (1974) [Sov. J. Quantum Electron. 4, 853 (1974)].
[15] B. V. Ershov, V. V. Savranskiĭ, and V. B. Fedorov, Abstracts, Seventh All-Union Conf. on Coherent and Nonlinear Optics, Moscow Univ. Press, 1974, p. 309.

Translated by J. G. Adashko

Resonant ionization of atoms under conditions of adiabatic level inversion

M. V. Fedorov

P. N. Lebedev Physics Institute, USSR Academy of Sciences

(Submitted 27 April 1979)
Zh. Eksp. Teor. Fiz. 77, 2200–2210 (December 1979)

Analytic expressions are obtained for the structure and the line shape of the dispersion dependence of the probability of the resonant polarization w on the radiation frequency ω under conditions of adiabatic level inversion due to the dynamic Stark effect in a field of variable amplitude. It is shown that under certain conditions the $w(\omega)$ dispersion curve is not characterized by a single parameter (width). A narrow and high principal maximum can appear against the background of the relatively broad $w(\omega)$ line. The parameters of the principal maximum and of the $w(\omega)$ as a whole are obtained. Conditions for the realization of the regime of adiabatic inversion of levels are investigated, particularly the conditions for the onset of the narrow principal maximum. The effect of spatial inhomogeneity of the field on the probability of resonant ionization under conditions of adiabatic level inversion is investigated.

PACS numbers: 32.70.Jz, 32.60. + i

1. Resonant multiphoton ionization of atoms has recently attracted considerable attention both experimentally[1-9] and theoretically.[10-18] In the experiment it is customary to investigate either the dependence of the ionization probability w on the intensity of the radiation at a given frequency ω, or the $w(\omega)$ dependence at a fixed field intensity. The experiments of the latter type are apparently most suitable for the determination of the physical nature of the resonant-ionization process, since they yield direct information on the width and position of the maximum and in general on the shape of the $w(\omega)$ dispersion curve. It is this which governs the formulation of the principal problems faced in the theory of resonant ionization: the determination of the shape of the dispersion curve and the elucidation of the physical mechanisms of its formation. From this point of view, an investigation of the excitation of the resonant level is insufficient for a complete solution of the problem of resonant ionization.

One of the principal parameters that characterize the $w(\omega)$ dispersion curve is its width Γ. Various physical

 0038-5646/79/121056-06$02.40