Temporal Development of Giant Pulses in Passively Q-switched Laser Oscillators

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The development of giant pulses in a passively Q-switched laser oscillator has been investigated in the case of a high power, multimode system. An experimental technique is described which allows the study of the pulse development over more than two decades of laser power. For powers $\lesssim 0.1$ maximum power an exponential increase of light flux is found. The time constant of this exponential part is calculated, and the relevant parameters are discussed. Calculated and measured time constants are in satisfactory agreement. In addition, it is shown that there is a connexion between the risetime of the pulse and the output energy, power, and divergence of the laser beam. As an example, the optimalization of a laser system for microplasma production is considered.

In a number of papers [1–5] the dynamics of passively Q-switched giant pulse lasers have been discussed. The main subject of these theoretical and experimental investigations is the determination of pulse characteristics such as pulse width, risetime from 0.1 to 0.9 maximum power, output energy, and peak output power. In this paper, data on the temporal development of the light pulse, starting from power levels less than 0.01 maximum laser power, are presented for the first time. The knowledge of the rise of the pulse over such a wide range of power is very useful in various applications of Q-switch lasers, e. g. for experiments on nonlinear optics and laser-produced plasmas.

I. Theory

The temporal development of a giant pulse after switching is governed by the well-known rate equations for the light flux and the population inversion in the active medium (in our case ruby). If optical pumping and spontaneous emission are neglected, these equations may be written

$$\begin{aligned} \mathrm{d}\varphi/\mathrm{d}t &= \varphi(\sigma\, n_{\mathrm{i}}-\gamma)\; c\, l/L\,, \\ \mathrm{d}n_{\mathrm{i}}/\mathrm{d}t &= -2\,\sigma\, n_{\mathrm{i}}\,\varphi/h\,\nu\,. \end{aligned} \tag{1}$$

φ is the light flux, c the velocity of light in vacuum, γ the total loss per unit length in the resonator, and L the optical length of the resonator. The population inversion n_{i}, the absorption cross-section σ, the photon energy $h\nu$, and the length l, are characteristic parameters of the active medium. The set of Eq. (1), which describes the formation of a photon avalanche is strictly valid only for a single resonator mode or a number of totally decoupled modes. If there is some relation between the phases of longitudinal modes, temporal fluctuations in the light pulse will occur. In our case very smooth pulses are found. For this reason the model of the photon avalanche seems applicable though there is a great number of modes present in our high power laser. The following assumption has entered Eq. (1): Active material, non-active material, and losses are distributed uniformly in the resonator. This assumption is made for more ease in the calculation. It is well justified in all cases treated in this paper. An especially simple solution of the rate equations is obtained if we assume in addition that i. there are no losses which depend on time or light intensity, and ii. the population inversion in the active medium varies neither with position nor with time. The first of these assumptions will be abandoned later on, the second should be regarded as a definition of the interval in which our analysis holds. Now Eq. (1) is greatly simplified. It describes an exponential increase of light flux:

$$\mathrm{d}\varphi/\mathrm{d}t = \varphi\, c\, \alpha^*\, l/L\,. \tag{2}$$

Here α^*, the "maximum effective gain constant", has been substituted for the expression $\sigma\, n_{\mathrm{i}}-\gamma$ in Eq. (1). Now we assume that the Q-switch bleaches from an initial transmission T_{i} to the final transmission $T_{\mathrm{f}}=1$, as soon as the gain in the system overcomes the losses. Then α^* may be expressed in terms of T_{i} in the following way:

$$\alpha^* = -(1/l)\cdot \ln T_{\mathrm{i}}\,. \tag{3}$$

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[2] R. W. HELLWARTH, Advances in Quantum Electronics, Columbia University Press, 1961, p. 334.
L. M. FRANTZ, Appl. Optics **3**, 417 [1964].
[3] A. SZABO and R. A. STEIN, J. Appl. Phys. **36**, 1562 [1965].
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[5] H. WEBER, Z. Angew. Phys. **27**, 1 [1969].

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It should be noted that α^* is independent of constant resonator losses, as scattering or neutral density filters, which appear in the threshold condition for laser action [see Eq. (9)]. This is easily understood because the available inversion (equal to α^*/σ) is given by the ratio of transmissions before and after switching; the threshold inversion, however, by the transmission itself. For this reason only time-dependent losses (see the discussion below) will influence the gain constant. In calculations which involve absolute values of inversions, constant losses will appear (this is the case with threshold relations and the calculation of the output energy; see Sect. IV).

Returning to Eq. (2), the risetime is defined as the time constant of the exponential increase of light flux:

$$\tau = L/(l\, c\, \alpha^*) . \tag{4}$$

This risetime represents a lower limit since time-dependent losses have been neglected as well as the decrease of population inversion with increasing light intensity. A more accurate prediction of the risetime requires the consideration of two kinds of losses which both reduce the effective gain constant: a) The "residual absorption" of the switching dye, and b) diffraction losses.

To a): It has been shown [6] that the usual dye switches, which are only considered here, do not bleach to 100% transmission at high laser powers. The remaining absorption per pass V_a may be considered in Eq. (2) by subtracting from α^* a loss constant

$$\gamma_a = -(1/l) \cdot \ln(1 - V_a) . \tag{5}$$

To b): The pure diffraction loses are negligible for laser oscillators of high Fresnel number as is the case in our experimental system. The following loss mechanism, however, has to be considered: When the inversion is not distributed uniformly over the ruby cross-section, laser action will start in a region of maximum inversion. From there it spreads gradually because of diffraction [7]. In this way regions of lower inversion and, correspondingly, lower effective gain constant participate in the formation of the giant pulse. For this reason not the maximum gain constant, given by Eq. (3), but an average value has to be used in calculating the risetime. One may account for this loss-mechanism in Eq. (2), in a similar way as for the residual absorption, by intrdoucing the loss constant γ_d i. e. the difference between maximum and average gain constant. Obviously γ_d will slightly increase with time as the lasing region spreads by diffraction. This effect is partly cancelled by the increasing transmission of the switching dye, and the net deviation from an exponential increase of laser power is expected to be small. The modified rate equation now reads

$$d\varphi/dt = \varphi\, c(\alpha^* - \gamma_a - \gamma_d)\; l/L \tag{3 a}$$

and the corresponding time constant is given by

$$\tau = L/[l\, c(\alpha^* - \gamma_a - \gamma_d)] . \tag{4 a}$$

The exponential part of the light pulse is limited, on the low power side, by the rapid bleaching of the switching dye before the residual absorption is reached. On the high power side, the exponential rise of the pulse terminates when the inversion begins to decrease. Both effects have been treated by various authors, e. g. by SZABO and STEIN [3]. An estimate, following their analysis, gives an exponential increase extending over about two decades of laser power for medium inversions (initial transmission $T_i \approx 0.4$ to 0.5). The exponential part gets shorter for low inversions ($T_i \gtrsim 0.6$) and vanishes at $T_i \approx 0.7$ in agreement with experiments (see also Sect. III).

II. Experimental Determination of Losses

The laser oscillator employed in the subsequent experiments is shown in Fig. 1. The system is built compact in order to obtain a favourable ratio L/l for short risetimes [see Eq. (4 a)]. The resonator consists of two totally reflecting prisms. Only a small portion of the light flux (0.175 per pass) is deflected from the resonator by a glass plate. The small deflection allows the use of low initial dye transmissions with resulting high gain constants α^*. The length of the ruby is 15 cm, its diameter 1.0 cm, and the optical length of the resonator 43 cm.

The residual absorption of the Q-switch inside this laser oscillator is determind experimentally. The two photocells (TRG 105 B) are calibrated with the dye cell in position 2 (see Fig. 1). When the switching cell is in position 1, the photocell PH 2 measures a pulse which has been attenuated by two passages through the dye cell. The signals from both photocells are displayed on a Tektronix 519

[6] C. R. GIULIANO and L. D. HESS, IEEE J. Quant. El. **QE-3**, 358 [1967].

[7] N. G. BASOV, Sov. Phys. Dokl. Akad. Nauk **10**, 311 [1965].

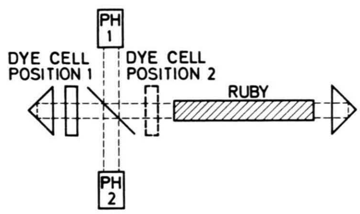

Fig. 1. Laser oscillator used for experiments, and arrangement for residual absorption measurements.

oscilloscope. Their height is compared at various moments of the giant pulse development. In this way, the residual absorption as a function of light power is readily obtained in a single shot.

Results for solutions of cryptocyanine in methanol are presented in Fig. 2. It is seen that the residual absorption per pass V_a is not constant but slightly decreasing with increasing laser power. The residual absorption also decreases with increasing initial transmission of the dye solution. The loss constants γ_a corresponding to the measured values of V_a are easily calculated from Eq. (5).

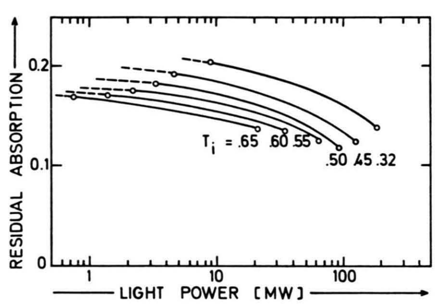

Fig. 2. Residual absorption V_a of cryptocyanine solution vs. power inside laser oscillator for various initial transmissions T_i. Oscilloscope traces have been evaluated for powers between ca. 0.04 maximum power and maximum power, as is indicated by open circles.

It remains to determine the loss constant γ_d. For this purpose, the variation of the gain constant as a function of position within the ruby cross-section has to be measured. Bearing in mind a cylindrically symmetric distribution, the gain constant may be written as a function of the radius r or, more conveniently, as a function of the circular area F which lies inside this radius. The area F is normalized to the total cross-section of the ruby crystal. The "loss" $\alpha^* - \alpha(F)$, i. e. the deviation from maximum gain constant at some position F within the ruby cross-section, contributes to the pulse development with a weight given by the local light intensity $I(F)$. Then γ_d is found by integrating $[\alpha^* - \alpha(F)] \cdot I(F)$ from zero to F_0, the total area which shows laser action:

$$\gamma_d = [1/F_0 I(0)] \int_0^{F_0} [\alpha^* - \alpha(F)] \cdot I(F) \, dF. \quad (6)$$

The gain constant as a function of position was found from near-field photographs of normal laser action by inserting defined neutral density filters into the resonator. The experimental results for our oscillator are well approximated by

$$\alpha(F) = \alpha^* - 0.064\, F \text{ cm}^{-1}. \quad (7)$$

The difference between maximum and local gain constant is proportional to the area F which has α^* in its center and $\alpha(F)$ on its edge.

The intensity distribution during the exponential rise of the giant pulse has not been measured. A sufficiently good fit to experiment is expected by the estimated distribution $I(F) = I(0) \cdot (1 - F/2\,F_0)$. Inserting this intensity distribution and Eq. (7) into Eq. (6), γ_d turns out to increase linearly with F_0, the area exhibiting laser action:

$$\gamma_d = 0.021\, F_0 \text{ cm}^{-1}.$$

An attempt has been made to calculate F_0, following the analysis of LEONTOVICH and VEDUTA [8]. In this analysis the divergence of the laser beam is ascribed to transverse modes of high order *. F_0 then is found from divergence measurements treating the ruby inhomogeneities as a lens, in our case with an estimated focal length of 150 cm. Numerical values of the experimentally determined divergence (half width Θ during the exponential rise of the giant pulse), the caculated area F_0, and the corresponding loss constant γ_d, are compiled in Table 1 for three typical initial dye transmissions. For comparison, the losses by residual absorption γ_a in the cryptocyanine solution are also given. They have been evaluated from the measurements presented in Fig. 2; the residual absorption at 0.05 maximum power was taken to be characteristic for the exponential part of the pulse. A comparison of the loss constants γ_a and γ_d shows that the residual absorption clearly dominates the loss term γ_d within the considered regime of initial dye transmissions. From

[8] A. M. LEONTOVICH and A. P. VEDUTA, Sov. Phys. JETP **19**, 51 [1964].

* In our case, the experimentally determined values of the divergence correspond to modes of transverse orders between 8 and 35 for $T_i = 0.65$ and 0.33 respectively.

T_i	α^* $\cdot 10^{-2}\,cm^{-1}$	Θ mrad	F_0	γ_d $\cdot 10^{-2}\,cm^{-1}$	γ_a $\cdot 10^{-2}\,cm^{-1}$
0.33	7.52	4.7	0.29	0.61	1.53
0.49	4.80	3.5	0.16	0.33	1.37
0.65	2.86	2.2	0.06	0.13	1.25

Table 1. Experimental divergence Θ, calculated area with laser action F_0, and losses γ_d (calculated from divergence measurements) and γ_a (experimental) during the exponential part of giant pulses with different initial dye transmissions T_i. Corresponding values of the effective gain constant α^* are obtained from Eq. (3).

Table 1 it is also seen that at high initial transmissions the losses are comparable with the gain. In this case a large discrepancy between experimental risetimes and values obtained from Eq. (4) is expected. At lower initial dye transmissions the deviation should become smaller, since the gain constant increases much faster than the loss constants.

III. Risetime Measurements and Discussion

In the preceding section losses have been discussed. Now risetimes in the exponential part of a giant pulse can be calculated with the help of Eq. (4 a). For a comparison with measured risetimes an experimental technique was developed, which is capable of recording the rise of the pulse over a substantial interval in laser power. With one photocell the total pulse is registered, while a second photocell measures with larger gain the rise of the laser power far below the peak of the pulse. To avoid overloading of the second photocell, the arrangement depicted in Fig. 3 was used. The laser oscil-

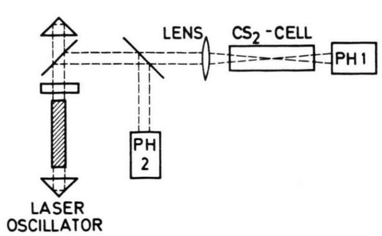

Fig. 3. Experimental setup for risetime measurements.

lator is the same as described above (see Fig. 1). Its output beam is focused into a cell filled with CS_2. When a certain threshold power is reached, the incident light is strongly reflected by stimulated Brillouin scattering (SBS), and the pulse transmitted through the CS_2-cell is cut off. With the rather large divergence of the laser beam in these experiments, a cut-off by a factor 3 to 10 was accomplished. The disturbing effect of the Brillouin light on the laser emission may be neglected because of the small amount of Brillouin radiation entering the resonator and the large distance between the CS_2-cell and the oscillator. The overall risetime of the detection system (0.3 nsec) is considerably smaller than the risetime of our giant pulses ($\gtrsim$ 1.5 nsec).

Oscilloscope traces of a typical giant pulse ($T_i = 0.54$) are shown in Fig. 4. The onset of SBS appears as a marked edge in the upper trace (pulse

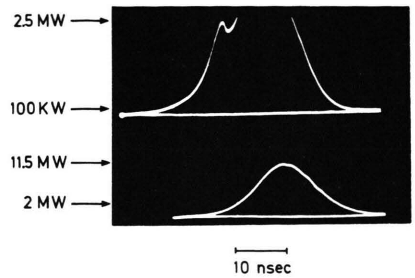

Fig. 4. Typical oscilloscope traces for $T_i = 0.54$; *upper* trace: Signal transmitted through CS_2; *lower* trace: Total laser pulse. Note the different ordinates for the two traces.

transmitted through the CS_2-cell). Fig. 5 shows the pulse development in a semilogarithmic plot as it has been constructed from the signals of the two

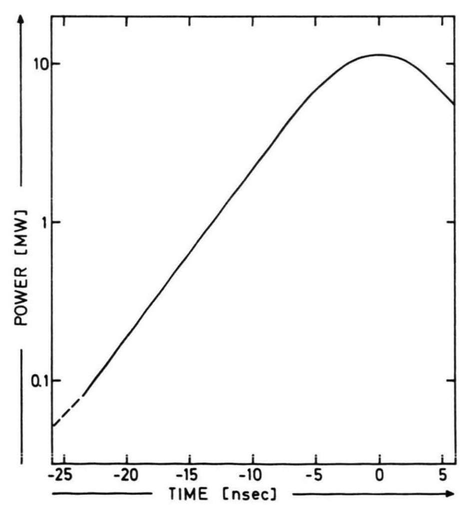

Fig. 5. Giant pulse of Fig. 4 in semilogarithmic plot.

photocells. The pulse shape turns out to be in good agreement with the considerations of the first section: For powers less than about 1/3 of maximum power, there is an exponential increase of light flux for at least two decades. This is typical for an initial dye transmission $T_i \approx 0.5$. With $T_i \approx 0.65$ the exponential part shrinks to about one decade and disappears for even higher transmissions. The pulse rise gets slower than exponential near maximum power, due to the decrease of inversion.

A series of measurements was made with varying values of T_i. Measured and calculated risetimes are given in Fig. 6 as functions of T_i and the reciprocal

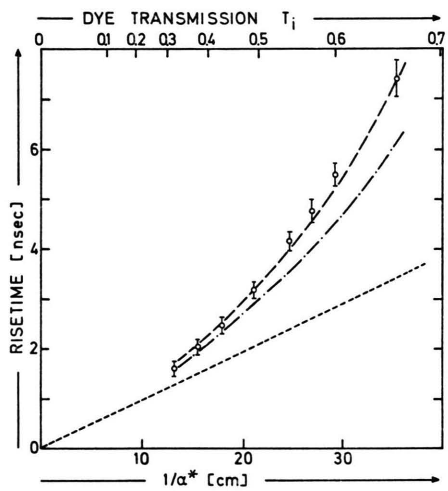

Fig. 6. Measured and calculated risetimes τ vs. reciprocal effective gain constant $1/\alpha^*$. ······ calculated from Eq. (4); —·—·— calculated from Eq. (4 a) considering residual absorption; — — — — calculated from Eq. (4 a) considering residual absorption and average gain constant for $F_0=0.15$.

effective gain constant $1/\alpha^*$, with α^* given by Eq. (3). Experimental points are shown as open circles. They were obtained with good reproducibility, as is seen from the error bars (maximum deviations from the average, found experimentally). The dotted, the dash-dot, and the dashed curves are calculated as follows: The dotted curve shows the linear dependence of the risetime τ on $1/\alpha^*$ as it is expected from Eq. (4). The agreement with the experimental points is rather poor, especially for large values of $1/\alpha^*$. It should be noted that the measured dependence of τ on $1/\alpha^*$ is stronger than linear. This behavior is expected from theory if the losses by residual absorption are taken into account. As an example the dash-dot curve has been calculated from Eq. (4 a) inserting the loss constant γ_a given by Eq. (5). Now the agreement of measured and calculated risetimes is quite satisfactory. It is further improved when γ_d, the difference between maximum and average gain constant, is taken into account in Eq. (4 a). A nearly perfect fit (dashed curve) to the experimental points is obtained, when γ_d is chosen to be $\gamma_d=3.2\cdot10^{-3}\ \mathrm{cm}^{-1}$, independent of the value of $1/\alpha^*$. This loss constant γ_d corresponds to a normalized area of laser action $F_0=0.15$. A glance at Table 1 shows that this value of γ_d is quite reasonable; it is a mean value of the experimental data presented there. Fig. 6 is in good agreement with the discussion of the preceding section; the losses γ_a and γ_d influence the risetime in a crucial manner except for low initial dye transmissions, where the measured points approach the dotted line.

Experiments were also carried out with other resonator geometries, other rubies, and different bleachable absorbers (vanadium phthalocyanine in nitrobenzene, 3-3′-diethylthiodicarbocyaniniodide in methanol, RG-8 filters, solid films of cryptocyanine). In all cases the risetimes follow with good accuracy an empirical relation of the form $\tau \propto (1/\alpha^*)^q$, with $q\approx1.6$. The deviation from the linear law of Eq. (4): $\tau \propto 1/\alpha^*$ is due to residual absorption losses and averaging of the gain constant.

In this connexion it should be pointed out that not only the risetime decreases with transmission T_i but also the pulse width at half maximum power. The fastest rising pulse, generated with the oscillator discussed here, had a risetime of 1.6 nsec and a half width of 8.3 nsec. With another oscillator, built in the same way, we obtained pulses with half widths of 6.5 nsec at $T_i=0.25$. On principle, there should be no difficulty to reduce the pulse width down to ca. 3 nsec by lowering the dye transmission.

Apart from the effective gain constant, the risetime will depend on the length of the resonator. It is expected from Eq. (4) that $\tau \propto L/l$. In Fig. 7 risetimes for different lengths of the optical resonator and for three different initial transmissions of the dye cell are plotted as functions of L/l. Experimental points were obtained as described above. They are shown as open circles, connected by solid lines. The dashed straight lines visualize the expected dependence $\tau \propto L/l$. The actual risetimes grow faster than linearly with L/l. In this experiment the residual absorption was not measured; a quantitative interpretation of our data was not possible.

According to Eqs. (3), (4 a) and Fig. 2 the initial transmission T_i and the ratio L/l are the major parameters which influence the magnitude of τ. Less important are the diameter of the resonator and the optical pumping arrangement which may affect γ_d.

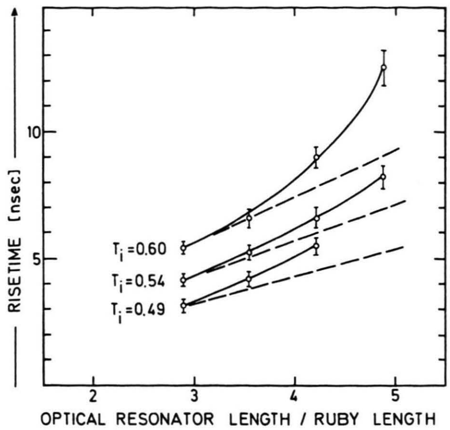

Fig. 7. Risetime τ vs. ratio of optical resonator length and ruby length L/l. ——— measured; – – – – calculated.

According to Eq. (4 a) no other parameter should influence the risetime. This fact was checked by varying different experimental parameters. For instance the pumping power was changed or constant losses (neutral density filters, ground glass plates) were inserted into the resonator without changing T_i. Within experimental error the measured risetimes remained unchanged. For this reason, and because of the good agreement of calculated and measured risetimes, Eq. (4 a) is expected to contain all laser parameters of relevance. We feel that Eq. (4 a) is a reliable basis for calculations of the risetime of a giant pulse laser.

IV. Concluding Remarks

It has been demonstrated that the risetime of a giant pulse is governed by the initial transmission T_i of the switching dye. It should be noted that a variation of T_i affects strongly the divergence, the output energy, and the power of the laser.

A summary of experimental results to this point is given in Table 2. It is seen that the output energy E_A, the output power P_A, and the divergence Θ

T_i	τ nsec (exp.)	E_A J (exp.)	(theor.)	P_A MW (exp.)	Θ mrad (exp.)
0.33	1.6	0.72	0.72	90	4.7
0.49	3.2	0.36	0.40	30	3.5
0.65	7.4	0.11	0.09	4.6	2.2

Table 2. Giant pulse characteristics for three different initial dye transmissions. E_A and P_A are total output energy and power, respectively (i. e. output to both sides of the oscillator shown in Fig. 1).

(half width), increase as the risetime is made shorter by reducing the initial transmission T_i.

Calculated values of the output energy, which are also listed in Table 2, are of interest for two reasons: First, E_A is a parameter besides the risetime τ which may be calculated in a relatively simple way and with reasonable accuracy. Second, an analytic expression for E_A will be useful for optimizing the output [9], especially when risetime and divergence are involved. A problem of this kind will be considered below.

The calculation of the output energy has been the subject of several articles [1, 3]. Accurate results are obtained, when the losses by residual absorption, and the inversion ** distribution [in our case given by Eq. (7)] are considered. The expression for the output energy has the form

$$E_A = \left\{ \frac{\ln(1-A)}{\ln[(1-A)\,(1-V)]} \right\} \mathscr{V} h\,\nu\,\overline{(n_{ii} - n_{if})} \,. \tag{8}$$

A is the fraction of radiation deflected from the resonator per pass, V are all other losses per pass (scattering losses, neutral density filters, residual absorption), $\mathscr{V}$ is the ruby volume, ν the laser frequency. $(n_{ii} - n_{if})$ is the difference between initial and final inversion, averaged over the inversion distribution. Its calculation is not treated in detail since it follows closely the theory of Ref. [1]. Inserting our oscillator parameters ($A = 0.175$; $V = 0.78$) (l. c. [9a]) into Eq. (8), good agreement between calculated and measured output energies is found (see Table 2).

In Eqs. (4 a) and (8) expressions are given for two important characteristics of giant laser pulses: risetime and output energy. To conclude this section, an example is outlined, which demonstrates how these results may be applied to design special laser systems. The production of microplasmas by intense laser light is considered. To obtain energetic plasmas, high output power and energy are required [10]. In addition, the risetime of the generating pulse markedly influences the plasma temperature

[9] J. G. EDWARDS, Appl. Optics **6**, 1011 [1967].

** Following Ref. [1], the inversion n_i is used rather than the gain constant α. The connexion between these parameters is simply $\alpha = \sigma\, n_i$, where σ is the absorption cross-section of the laser line.

[9a] The large value of V is mainly due to a neutral density filter (transmission $T = 0.44$) which is placed inside the resonator to keep the laser power well below the damage threshold.

[10] See e. g.: H. OPOWER, W. KAISER, H. PUELL, and W. HEINICKE, Z. Naturforsch. **22 a**, 1392 [1967].

(l. c. [11–13]). Short risetimes are desired, suggesting the use of very short pulses [14] ($\approx 10^{-11}$ sec) and the shaping of usual giant pulses by external electro-optic switches [15, 16].

When a passively Q-switched giant pulse itself is employed for microplasma production, short rise-time and high output energy are simultaneously obtained at low initial transmission of the switching dye (see Table 2). The lowest possible initial transmission is determined by the maximum attainable inversion $n_i^{\max}$ and the resonator losses, as is seen from the threshold relation:

$$T_i [1 - A] [1 - (V - V_a)] \exp\{\sigma n_i^{\max} l\} \geqq 1. \quad (9)$$

The fraction A of light deflected from the oscillator is chosen in such a way that maximum energy is delivered. This optimum value of A is found from Eq. (8) depending on the initial transmisson T_i. Introducing the result into Eq. (9), the lowest possible T_i is obtained. For our oscillator (with the neutral density filter removed the total loss is $V = 0.5$) representative values are $A_{\text{opt}} \approx 0.6$ to 0.7 and $T_i \approx 0.25$, if a maximum inversion $n_i^{\max} = 8 \cdot 10^{18}$ cm^{-3} is assumed. The risetime of the light pulse is expected to be ca. 1 nsec according to Eq. (4), and the total output energy ca. 4 J from Eq. (8). A laser oscillator delivering giant pulses of this risetime and output energy is quite useful for microplasma production.

Frequently higher energies are demanded, and one or several amplifier sections are placed behind the oscillator. In this case, the divergence of the laser beam, which is growing as the initial transmission of the Q-switch is lowered, has to be thoroughly considered. A reduction of the divergence by mode-selecting elements may prove to be favourable for special applications.

The authors are indebted to Professor Dr. W. KAISER for many stimulating discussions.

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