Frequency tunable generation of giant pulses in spectrally inhomogeneous media

V. B. Markov, T. Ya. Marusiĭ, M. S. Soskin, and A. I. Khizhnyak
Physics Institute, Ukrainian Academy of Sciences
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Generation of giant pulses by spectrally inhomogeneous media is investigated theoretically and experimentally. An expression for the energy emitted in a giant pulse is obtained from the kinetic equations. It is shown that the emission energy increases on retuning of the generation frequency onto the wing of the luminescence line. The theoretical results qualitatively explain the experimental results for an Nd^{3+} glass laser.

The influence of spectral inhomogeneity of the active medium of the laser, due to inhomogeneous broadening of the working luminescence line, on the characteristics of stimulated emission has been investigated to date only for the regime of stationary generation (see, e.g.,[1]). On the other hand, problems of nonstationary generation, and particularly generation of individual giant pulses by such a medium, have not been investigated theoretically, although it is understandable that the parameters of isolated pulses should depend on the character of the spectral broadening of the luminescence line of the working transition. Thus, the energy radiated in a single pulse depends on the ratio of the homogeneous (δ) and inhomogeneous (Δ) broadenings, for when the inhomogeneity parameter $q = \Delta/\delta$ is large the energy radiated is only a small fraction of the stored energy, on the order of $1/q$. This explains the need for knowing the connection between the parameters of such a pulse and the structure of the luminescence line.

The present paper is devoted to a theoretical and experimental investigation of the characteristics of the giant pulse with tunable frequency, when the active medium has an inhomogeneously broadened working-transition line. The theoretical study is based on the kinetic equations, and the experiments were performed on a neodymium-glass laser Q-switched with a passive shutter.

1. THEORETICAL INVESTIGATIONS OF GIANT-PULSE PARAMETERS

It is well known that the process of generation of a giant pulse can be broken up into several stages[2]. In the present paper we are interested in the last stage, the emission of the giant pulse. The field produced during the preceding stages does not depend on the character of the active-medium luminescence-line broadening; its value was determined earlier[3].

The equations describing the last stage become much simpler if it is assumed that there is no mode locking and the migration of the energy over the luminescence lines during the stage of interest to us is neglected[4,5]. In addition, we consider two limiting cases: the lifetime of the active centers at the lower working level $T_L = 0$ or $T_L = \infty$. Finally, we recognize that in the case of solid-state lasers at room temperature the intermode interval is much smaller than δ, so that we can introduce the spectral density $m(\nu, \tau)$ of the number of quanta at the frequency ν. Starting from these assumptions, we obtain the following system of kinetic equations:

$$\begin{aligned} dm(\nu',\tau)/d\tau &= \left\{\mu\int_{-\infty}^{+\infty}[n_2(\bar\nu,\tau)g_{21}(\bar\nu,\nu') - n_1(\bar\nu,\tau)g_{12}(\bar\nu,\nu')]d\bar\nu - \varphi(\nu')\right\}m(\nu',\tau),\\ d[n_2(\bar\nu,\tau)-n_1(\bar\nu,\tau)]/d\tau &= -d\int_{-\infty}^{+\infty}[n_2(\bar\nu,\tau)g_{21}(\bar\nu,\nu') - n_1(\bar\nu,\tau)g_{12}(\bar\nu,\nu')]m(\nu',\tau)d\nu', \end{aligned} \quad (1)$$

where

$$m(\nu',\tau)=BT_r{}^0M(\nu',\tau),\quad n_i(\bar\nu,\tau)=BT_r{}^0N_i(\bar\nu,\tau),$$
$$g_{ij}(\bar\nu,\nu)=B^{-1}B_{ij}(\bar\nu,\nu),\quad \varphi(\nu)=T_r{}^0T_r{}^{-1}(\nu),$$
$$\tau=\frac{t}{T_r{}^0},\quad \mu=\frac{V_a}{V},\quad B=\int_{-\infty}^{+\infty}B_{21}(\bar\nu,\nu)d\bar\nu,$$
$$\alpha=\begin{cases}1, & \text{if } T_L=0,\ n_1(\bar\nu,\tau)=0\\ 2, & \text{if } T_L=\infty,\end{cases}$$

$N_i(\bar\nu, \tau)$ is the spectral density of the active centers at the level i in a unit volume (1 and 2 denote the lower and upper working levels, respectively); $T_r(\nu)$ is the lifetime of the photons of frequency ν in the resonator when the shutter is bleached, and T_r^0 is the same at the frequency with maximum Q; $B_{ij}(\bar\nu, \nu')$ is the Einstein spectral coefficient of an induced transition from the level i to the level j ($B_{ij}(\bar\nu, \nu')$ includes a coefficient that takes into account the temperature distribution of the population over the Stark components of the level i); V and V_a are the field-occupied volumes of the resonator and of the active medium.

The system (1) was obtained assuming a spatially-homogeneous distribution of the field in the resonator. This condition is well satisfied in a standing-wave resonator if a large number of modes, both longitudinal and transverse, are simultaneously emitted during the course of generation[6]. But in the case of generation with few modes, as is typical for Q-switching with a passive shutter and for a resonator with a diaphragm, the spatial inhomogeneity is maximal[7,8]. The small-scale longitudinal spatial inhomogeneity is eliminated by using a travelling-wave resonator, while the transverse inhomogeneity can be taken into account by using certain known results[8]. We shall dwell below on the question of taking the spatial inhomogeneity into account analytically, and will assume for the time being that the field is spatially homogeneous and use the system of equations (1).

It is clear that it is impossible to obtain an analytic solution of the system (1) and that numerical methods must be resorted to, but even this is different, by virtue of the integrodifferential character of the equations. It turns out, however, that under certain reasonable assumptions one can obtain a connection between the radiated energy and the peak power on the one hand and the initial inverted population and the spectroscopic parameters of the active medium on the other. Let us demonstrate this. If we assume that the spectral composition of the giant pulse remains unchanged during the time of its emission, as is apparently very nearly the case, i.e.,

$$m(\nu',\tau)=m(\tau)f(\nu'),\quad \int_{-\infty}^{+\infty} f(\nu')\,d\nu'=1, \tag{2}$$

then, using the system of equations (1), we can easily obtain an expression for the radiated energy[1)]

$$E=\int_0^{\tau_{fin}} m(\tau')\,d\tau'=\mu\int_{-\infty}^{+\infty}\{[n_2(\bar\nu,0)-n_1(\bar\nu,0)] -[n_2(\bar\nu,\tau_{fin})-n_1(\bar\nu,\tau_{fin})]\}\,d\bar\nu\Big/\alpha\int_{-\infty}^{+\infty} f(\nu')\varphi(\nu')\,d\nu'. \tag{3}$$

It follows from (3) that by determining the population difference that remains after the generation of the giant pulse, we can obtain the radiated energy.

To find the population difference remaining at the instant τ, we integrate the first equation of the system (1) with respect to ν' and the second with respect to $\bar\nu$, and divide the first equation by the second. Transforming the last term in the right-hand side of the resulting equation and using the second equation of the system (1), we obtain ultimately

$$m(\tau)-m(0)=\frac{\mu}{\alpha}\int_{-\infty}^{+\infty}\{[n_2(\bar\nu,0)-n_1(\bar\nu,0)]-[n_2(\bar\nu,\tau)-n_1(\bar\nu,\tau)]\}\,d\bar\nu + \frac{\int_{-\infty}^{+\infty}\varphi(\nu')f(\nu')\,d\nu'}{\int_{-\infty}^{+\infty}[g_{21}(\bar\nu,\nu')+g_{12}(\bar\nu,\nu')]f(\nu')\,d\nu'}\ln\frac{\int_{-\infty}^{+\infty}[n_2(\bar\nu,\tau)g_{21}(\bar\nu,\nu')-n_1(\bar\nu,\tau)g_{12}(\bar\nu,\nu')]f(\nu')\,d\nu'}{\int_{-\infty}^{+\infty}[n_2(\bar\nu,0)g_{21}(\bar\nu,\nu')-n_1(\bar\nu,0)g_{21}(\bar\nu,\nu')]f(\nu')\,d\nu'} \tag{4}$$

Using (4), we can determine two important giant-pulse parameters, the peak power and the energy. The power at the maximum is obtained from the condition $dm/d\tau = 0$, and to determine the energy we must put $m(\tau_{fin}) = m(0) = 0$ (see footnote 1), which leads to an expression for the remaining population difference.

Let us stop to discuss the energy characteristics of the giant pulse. Using (4) and (3), we obtain an equation for the radiated energy[2)]:

$$E=\frac{1}{\bar\varphi}\int_{-\infty}^{+\infty}\frac{n_2(\bar\nu,0)\bar g_{21}(\bar\nu)-n_1(\bar\nu,0)\bar g_{12}(\bar\nu)}{\bar g_{21}(\bar\nu)+\bar g_{12}(\bar\nu)}\{1-\exp[-E[\bar g_{21}(\bar\nu)+\bar g_{12}(\bar\nu)]]\}\,d\bar\nu, \tag{5}$$

where we have introduced, to simplify the notation,

$$\bar g_{ij}(\bar\nu)=\int_{-\infty}^{\infty} g_{ij}(\bar\nu,\nu')f(\nu')\,d\nu',\quad \bar\varphi=\int_{-\infty}^{+\infty}\varphi(\nu')f(\nu')\,d\nu'.$$

Expression (5) connects the radiated energy with the initial inverted population and with the spectroscopic parameters of the luminescence line.

We proceed now to an analysis of relation (5). To this end we determine the initial population difference from the first equation of (1)

$$\mu\int_{-\infty}^{+\infty}[n_2(\bar\nu,0)g_{21}(\bar\nu,\nu)-n_1(\bar\nu,0)g_{12}(\bar\nu,\nu)]\,d\bar\nu=\psi(\nu),$$

$$\psi(\nu)=T_r^0/T_r^*, \tag{6}$$

where T_r^* is the lifetime of photons with frequency ν in the resonator with the passive shutter not bleached, and ν is the frequency at which the condition (6) is first satisfied.

Expressions (5) and (6) simplify in the case when $\alpha = 2$, since in fact T_L is finite and we can therefore put $n_1(\bar\nu, 0) = 0$. In addition, if it is recognized that in the case of Q switching with a passive shutter the generation-spectrum width is very small, much narrower than δ and in most cases narrower than $\varphi(\nu)$, then we can simplify the calculation by putting henceforth $f(\nu') = \delta(\nu - \nu')$. Then (5) and (6) take the form

$$\frac{1}{\gamma}=y=\frac{1}{EJ(\nu)}\int_{-\infty}^{+\infty}\frac{\zeta(\bar\nu)g_{21}(\bar\nu,\nu)}{g_{21}(\bar\nu,\nu)+g_{12}(\bar\nu,\nu)}\{1-\exp[-E[g_{21}(\bar\nu,\nu)+g_{12}(\bar\nu,\nu)]]\}\,d\bar\nu,$$

$$\gamma=\frac{\psi(\nu)}{\varphi(\nu)},\quad n_2(\bar\nu,0)=n_2\zeta(\bar\nu)\quad\left(\int_{-\infty}^{+\infty}\zeta(\bar\nu)\,d\bar\nu=1\right); \tag{7}$$

$$J(\nu)=\int_{-\infty}^{+\infty}\zeta(\bar\nu)g_{21}(\bar\nu,\nu)\,d\bar\nu,$$

where $J(\nu)$ is the luminescence-line shape. We note that if the pump-pulse duration up to the instant of the start of the generation is much shorter than the time T_1 of spontaneous decay of the excited state of the active medium, then γ is the ratio of the threshold pump energy in the presence of a shutter to the pump energy in the absence of a shutter, if the latter is completely bleached.

Let us investigate the obtained relations for certain luminescence-line types.

Homogeneously broadened luminescence line

Let

$$n_{2,1}(\bar\nu,0)=n_{2,1}\delta(\bar\nu-\nu_0).$$

Expression (7) then takes the form

$$y=[1-\exp\{-EG(\nu,\nu_0)\}]/EG(\nu,\nu_0),\quad G(\nu,\nu_0)=g_{21}(\nu_0,\nu)+g_{12}(\nu_0,\nu). \tag{8}$$

It follows from (8) that at $\gamma = \mathrm{const}(\nu)$, the radiation energy will vary during the course of tuning of the generation frequency in accordance with

$$E(\nu)=E(\nu_0)G(\nu_0,\nu_0)/G(\nu_0,\nu),$$

i.e., it will increase when the frequency is tuned to the wing of the luminescence line. If the working levels are split, but the temperature equilibrium is not violated, then the radiated energy will increase on tuning to a less intense luminescence line. This effect can easily be attributed to the fact that the threshold condition contains the product $n_i g_{ij}(\nu_0, \nu)$, while the radiated energy is determined by the quantity n_i, i.e., when $g_{ij}(\nu, \nu_0)$ is decreased, the energy stored in the active medium increases and consequently the output energy increases.

Inhomogeneously broadened line (working levels not split and having equal degeneracy)

To establish the character of the behavior of E as a function of the inhomogeneous broadening, we put

$$n_2(\bar\nu,0)=\frac{n_0}{(\pi\Delta^2)^{1/2}}\exp\left\{-\frac{(\bar\nu-\nu_0)^2}{\Delta^2}\right\},\quad g_{21}(\bar\nu,\nu)=\frac{1}{\pi^{1/2}\delta}\exp\left\{-\frac{(\bar\nu-\nu)^2}{\delta^2}\right\}. \tag{9}$$

In this case expression (7) takes the form

$$y=\sum_{n=1}^{\infty}\frac{(-1)^{n-1}}{n!}\mathscr{E}^{n-1}\left(\frac{1+q^2}{1+nq^2}\right)^{1/2}\exp\left\{-\frac{(n-1)\theta^2}{nq^4+(n+1)q^2+1}\right\};\qquad(10)$$
$$q=\Delta/\delta,\quad \theta=(\nu_0-\nu)/\delta,\quad \mathscr{E}=2E/\pi^{1/2}\delta.$$

The series (10) was summed with a computer. By way of example, Fig. 1 shows the dependence of the radiated energy on the detuning θ for $\gamma = 2$. It is seen from the plots that the radiated energy depends little on θ already at q = 5, and it follows from (10) that at $q \gg 1$ the function y should not depend on q or θ. Calculations show that, starting with q = 10, y is practically independent of q and θ. Consequently, if we maintain Δ constant and increase q (i.e., we decrease the homogeneous width), then the radiated energy decreases like $E = \mathscr{E}\pi^{1/2}\delta/2$. This means that when an active medium having such a luminescence line is cooled, the radiated energy decreases in proportion to the decrease of δ, if γ is maintained constant, i.e., if the initial transmission of the passive shutter is not changed. If the homogeneous part of the luminescence line is described by a Lorentz curve

$$g_{21}(\bar{\nu},\nu)=\frac{\delta}{\pi(\delta^2+(\bar{\nu}-\nu)^2)}$$

then it follows from the calculations that the result remains qualitatively the same, although the quantitative changes are appreciable. An investigation of expression (7) shows that the more steeply the function $q_{21}(\bar{\nu}, \nu)$ falls off on detuning from the maximum, the smaller are the values of q at which the radiated energy ceases to depend on q and θ.

The foregoing examples demonstrate the influence of the detuning from the maximum of the luminescence line and of the parameter q on the characteristics of the giant pulse under conditions when the quantity γ that characterizes the excess over the free-generation threshold is constant. In the experiments, however, γ depends on the tuning frequency, by virtue of the frequency dependence of the initial shutter transmission, and, generally speaking, on resonator losses that are not connected with the passive shutter. It is therefore necessary to investigate the dependence of the radiated energy under conditions when γ is a function of the frequency.

From (7) we easily obtain a condition for finding the frequencies of the extrema of the radiation energy; this condition leads easily to the results obtained above, and states that at $d\gamma/d\nu = 0$ the radiated energy has a minimum at the frequency of the luminescence-line maximum. On the other hand, if $d\gamma/d\nu > 0$ ($d\gamma/d\nu < 0$), and $J(\nu)$ has a single maximum, then the radiated energy acquires a maximum at a frequency $\nu_M < \nu_0$ ($\nu_M > \nu_0$), and the minimum shifts towards the longwave (short-wave) side. As $d\gamma/d\nu \to 0$, we have $|\nu_M| \to \infty$, and at $|d\gamma/d\nu| \to \infty$ we have the limiting value of $|d\gamma/d\nu|$, above which the extrema merge together and the radiated energy becomes a monotonic function of the frequency, increasing together with γ. If the luminescence line has several maxima, then the number of extrema increases, but the picture remains qualitatively the same.

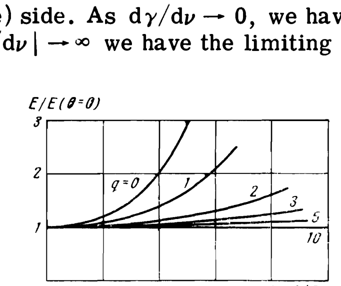

FIG. 1. Theoretical dependence of the giant-pulse energy on the generation frequency at different values of the inhomogeneous luminescence-line broadening parameter q. $\theta_{1/2}$ is the half-width of the luminescence line.

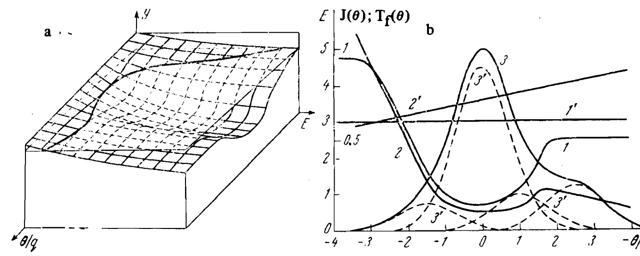

FIG. 2. The function y (θ, E) (a) and the corresponding dependence of the radiated energy (b): 1 – E (θ) ($T_f(\theta)$ – 1′), 2 – E (θ) ($T_f(\theta)$ – 2′), 3 and 3′–luminescence line and its components, T_f – initial transmission of the passive shutter.

Allowance for the splitting of the working levels makes the analysis of (5) very difficult. However, if the splitting is much larger than the luminescence line width of an individual center, the investigation becomes simpler. By way of example, Fig. 2a shows the function $y(\theta, E)$ for the case when each level splits into a pair of sublevels, the distance between which is the same for all centers. The function describing the distribution of the centers over the frequencies and the line shape of the individual center was assumed to be Gaussian (9), and q = 10. This has made it possible to use the results of the calculations for the second example. Figure 2b illustrates the shape of the luminescence line and its components. It also shows the dependence of the radiated energy for the cases $d\gamma/d\nu = 0$ (1) and $d\gamma/d\nu > 0$ (2). An investigation of expression (7) shows that when the working levels are split, the function y remains dependent on θ as $q \to \infty$, but unlike the case of the homogeneously broadened line the plot of the radiated energy against the generation frequency, during retuning at constant γ, does not take the form of an inverted luminescence line. Figure 2a shows the surface $1/\gamma$, the line of intersection of which with the surface y gives the values of the radiated energy as a function of the generation frequency. A maximum of the radiated energy on the wing of the luminescence line is clearly seen.

The foregoing analysis pertains to the case when the variation of $1/\gamma$ with frequency, due for example to the dependence of the initial transmission of the shutter on the frequency, is sufficiently smooth. On the other hand, in the case of a fast variation of $1/\gamma$ in comparison with y, for example, as a result of the increased losses, when the generation frequency during the course of tuning lags the tuning frequency strongly, the dependence of the radiated energy on the generation frequency can have a maximum at the frequency at which the luminescence line has a maximum, as will be seen from the experimental data.

The foregoing examples have demonstrated the influence of the character of the broadening and of the internal structure of the luminescence line on the energy of the giant pulse. The spectral composition of such a pulse is shaped during the first stages of development and therefore does not depend on the character of the luminescence-line broadening. However, generally speaking, the spectrum of the giant pulse can also experience changes during the last stage of generation, when an effective distortion of the gain contour takes place as a result of "hole burning" in the luminescence line. We did not take this mechanism into account, but we can state that it is of no importance if the width of the spectrum produced during the first stages is much smaller than the homogeneous width δ, since the width of the "hole" which is described under our assumptions by expression (4), is of the order of δ. The spectrum of the pulse can then be obtained in the same manner as in[3].

To find the temporal characteristics of a giant pulse, it is necessary to solve the system of equations (1), but the pulse duration (τ_p) can be estimated as the ratio of the energy radiated in the pulse to its peak power, the latter given by expression (4). In the general case, the relations obtained for the power at the maximum are cumbersome and will not be presented here. We cite directly the expression for τ_p in the case of a homogeneously broadened luminescence line:

$$\tau_p = \frac{1}{\varphi(\nu)} \frac{EG(\nu)}{\gamma - 1 - \ln\gamma}. \tag{11}$$

This leads to a result of practical importance, namely that the pulse duration remains constant when the frequency is tuned if $\gamma(\nu)$ and $\varphi(\nu)$ remain constant. This can easily be explained as being due to the fact that the pulse rise time is determined in the main by the quantity $\gamma(\nu)$, while its fall-off time is determined by the resonator loss $\varphi(\nu)$ and the pulse duration τ_p is equal to sum of the rise and fall-off times. This also explains the dependence of τ_p on the generation frequency if γ and φ depend on ν.

In concluding Sec. 1, we call attention to the following circumstance. On the one hand, the system of kinetic equations which serves as the starting point for the derivation of the system (1) (see, e.g.,[1]) describes in formally identical fashion both the spectral and the spatial inhomogeneities, and on the other hand, the solution method described above does not, in principle, require introduction of the spectral density $M(\nu', \tau)$, i.e., the transition to integration with respect to ν' is not obligatory, and we can retain the summation over the modes. Consequently, in analogy with the procedure used above, we can take into account the spatial inhomogeneity together with the spectral inhomogeneity, but in this case the condition that the spatial structure of the field remain unchanged during the last stage (a condition similar to (2)) is much stronger and can be satisfied in most cases only for a laser operating in a single mode with respect to the transverse index[8]. On the other hand, if this condition is satisfied, then it is easy to obtain an expression for the radiation energy with the spatial inhomogeneity also taken into account. For lack of space, we shall not discuss this question, but note that the results presented here are not qualitatively altered, although the quantitative changes are appreciable.

This completes the discussion of the theoretical results; we now proceed to a description of the experimental results.

2. EXPERIMENTAL RESULTS

The experimental investigations were carried out with the aid of the setup shown in Fig. 3. The dispersion

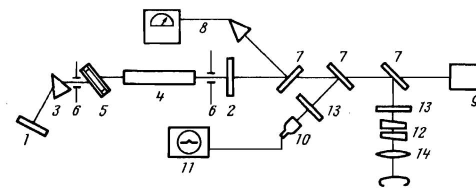

FIG. 3. Experimental setup: 1, 2 – resonator mirrors; 3 – dispersion prism; 4 – active element, 5 – passive-shutter cell, 6 – diaphragms, 7 – light-splitting plates, 8 – calorimeter, 9 – spectrograph, 10 – FÉK-09 photocell, 11 – oscilloscope, 12 – Fabry-Perot interferometer, 13 – attenuating filters, 14 – objective.

resonator, consisting of mirrors 1 and 2 and of a prism (or system of prisms) 3, contained the active element 4, which was placed in a double-ellipse illuminator with IFP-2000 lamps. When the angular dispersion of the prism used to narrow down the generation spectrum was low, the output mirror consisted of a stack of three glass plates, but when the angular dispersion corresponded to three prisms of STF-2 glass, the stack was replaced by a dielectric mirror. The Q switch, a solution of polymethyne dye in nitrobenzene, was poured into a cell 5 placed at the Brewster angle to the resonator axis. To increase the tuning range and to stabilize the transverse structure of the generation beam, two diaphragms 6 were used. Their diameter was chosen such as to obtain minimal divergence of the radiation. The active elements were rods of glass GLS-1, LGS-228-2, and LGS-250-3, with dimensions 10 × 130 mm and with end faces cut perpendicular to the axis and adjusted parallel to the mirror 2.

The radiation energies were measured with a thermocouple calorimeter and with a galvanometer 8. The frequency variation of the giant pulse was registered with a diffraction spectrograph 9, having a linear dispersion 13 Å/mm. To measure the temporal characteristics, we used FÉK-09 coaxial photocells 10 and an oscilloscope 11, of type S1-11 or 12-7. The width of the generation spectrum was monitored with the aid of Fabry-Perot interferometer 12.

A. Spectral Characteristics of Giant Pulse

The generation-wavelength tuning limits are determined by several parameters, namely, the luminescence line width of the active medium, the angular dispersion of the prisms employed, the spectral characteristics of the passive shutter, the optical quality of the active medium, and the capabilities of the pumping system. The obtained generation-frequency tuning characteristics are similar to the characteristics obtained in the free generation regime[9,10]. The reason is that the initial stage of development of the generation, during which the spectrum was formed, is similar for both regimes. The small difference between the generation frequencies at identical tuning in the two indicated regimes is due to the dependence of the initial transmission of the passive shutter on the frequency. By virtue of the indicated similarity, the requirements imposed on the active elements and on the resonator in order to obtain the necessary tuning characteristics remained the same as for the free-generation regime.

Some of the obtained tuning characteristics are given in Fig. 4a for a prism of TBF glass (angular dispersion $\partial\varphi/\partial\nu$ = 1.7 seconds of arc per cm). When

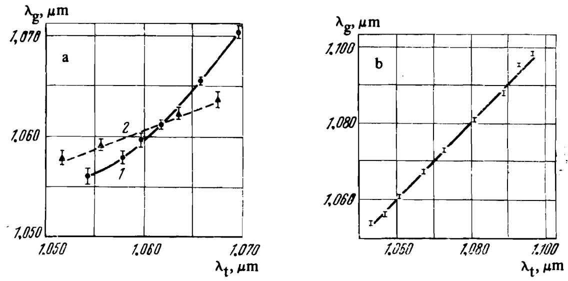

FIG. 4. Giant pulse generation frequency tuning characteristics: a – 1 – LGS-228-2, T_f = 0.51; 2 – LGS-228-2, T_f = 0.34; b – LGS-250-3, T_f = 0.60. The angular dispersion of the prism is $\partial\varphi/\partial\nu$ = 1.7 seconds of arc per cm. T_f – initial transmission of passive shutter at wavelength λ = 1.06 μm.

three prisms of STF-2 were used ($\partial\varphi/\partial\nu$ = 8.4 seconds of arc per cm), the generation frequency coincides with the tuning frequency, corresponding to the straight line in Fig. 4b. Just as in the free-generation regime, the tuning range increases when the prism dispersion is increased and is limited only by the capabilities of the pump and by the luminescence line width, while the generation frequency approaches the tuning frequency. In all cases, the width of the giant-pulse spectrum did not exceed 1 cm^{-1}, and this justifies the assumption made in the discussion following Eq. (1).

B. Temporal characteristics of giant pulse

Typical plots of the duration of the pulse as a function of the generation frequency are shown in Fig. 5a. It is clearly seen that pulse duration increases with increasing generation wavelength. This increase is due mainly to the increase of the initial transmission of the passive shutter (curves 2, 3, and 4 in Fig. 5a). The increase of duration in the case corresponding to curve 1 of Fig. 5a, on the other hand, is due to the increased loss during tuning, since the generation frequency lags the tuning frequency (Fig. 4a, curve 2). This also explains the growth of the pulse duration upon tuning to the short-wave side in the case of Fig. 5a, curve 4. It follows from the presented data that the last case is the most convenient from the point of view of constancy of the pulse duration, since we have τ_p = (40 ± 10) nsec in the greater part of the tuning region.

C. Energy characteristics

Figure 5b shows some of the obtained plots of the radiated pulse energy against the generation wavelength. It is clearly seen that they can be divided into three classes. The first includes the case when the maximum energy is radiated at the maximum of the luminescence line (Fig. 5b, curve 2). In this case it follows from Fig. 4a that the generation frequency lags the tuning frequency, i.e., the resonator is misaligned, and this leads to an increase of the loss in the empty resonator and consequently to a decrease of γ when the system is tuned away from the maximum of the luminescence line. The latter leads to a decrease of the radiated energy (see Sec. 1). In Sec. 1 we investigated the total radiated energy, while the energy emitted from the resonator was determined by the ratio of the useful and detrimental losses. In the present case, the useful losses increase, and this also leads to a decrease in the recorded energy. In addition, account must be taken of the influence of the inhomogeneity of the field in the transverse direction. It was shown in[8] that if the

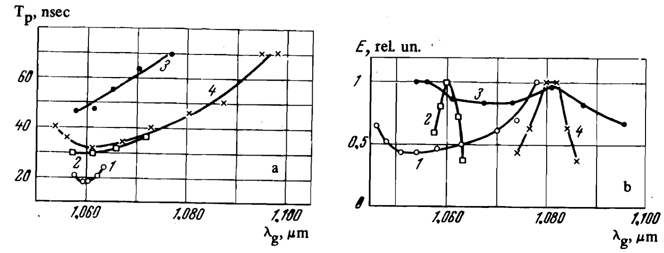

FIG. 5. a – Dependence of pulse width on generation wavelength: 1 – LGS-228-2, T_f = 0.34; 2 – LGS-228-2, T_f = 0.51; 3 – LGS-228-2, T_f = 0.60; 4 – LGS-250-3, T_f = 0.60. The angular dispersion of the prism is $\partial\varphi/\partial\nu$ = 1.7 seconds of arc per cm. b – Dependence of the radiation energy on the generation wavelength: 1 – GLS-1, T_f = 0.73, R = 45 %, $\partial\varphi/\partial\nu$ = 8.4; 2 – LGS-228-2, T_f = 0.34, R = 60%, $\partial\varphi/\partial\nu$ = 1.7; 3 – LGS-250-3, T_f = 0.60, R = 60%, $\partial\varphi/\partial\nu$ =1.7; 4 – GLS-1, T_f = 0.73, R = 88%, $\partial\varphi/\partial\nu$ = 8.4 ($\partial\varphi/\partial\nu$ is in seconds of arc per cm), R is the reflection coefficient of the output mirror. The second mirror had a reflection coefficient 99.8%. T_f is the initial transmission of the passive shutter at the wavelength λ = 1.06 μm.

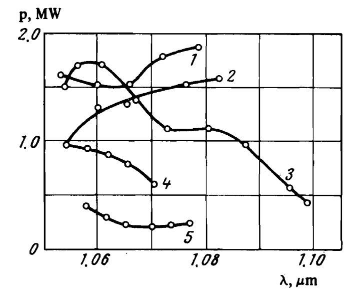

FIG. 6. Dependence of the pulse power on the generation frequency ($\partial\varphi/\partial\nu$ is in seconds of arc per cm): 1 – LGS-250-3, T_f = 0.56, $\partial\varphi/\partial\nu$ = 2.8; 2 – LGS-250-3, T_f = 0.60, $\partial\varphi/\partial\nu$ = 2.8; 3 – LGS-250-3, T_f = 0.50, $\partial\varphi/\partial\nu$ = 1.7; 4 – LGS-228-2, T_f = 0.51, $\partial\varphi/\partial\nu$ = 1.7; 5 – LGS-228-2, T_f = 0.62, $\partial\varphi/\partial\nu$ = 1.7.

resonator mirrors are tilted, the radiated energy decreases as a result of the distortion of the transverse field distribution.

The second class includes cases when the radiation energy has a maximum only at the edges of the tuning range (Fig. 5b, curve 1). In this case the limitation imposed on the tuning range by the capabilities of the pump did not make it possible to obtain either the maximum predicted in the first section of this paper or the subsequent fall-off of the radiated energy on the long-wave wing of the luminescence line. When the generation threshold was lowered by changing the mirror reflection coefficients, when the initial transmission of the passive shutter was increased, and when suitable active elements were selected, we were able to obtain the third class of dependences, in which the maximum occurs at a wavelength larger than 1.06 μm (curves 3 and 4 of Fig. 5b).

The cited examples can be explained qualitatively by means of the theoretical results obtained above. A quantitative comparison is difficult, since there are no exact data on the spectroscopic parameters of the luminescence lines in neodymium glasses.

Concluding the exposition of the experimental results, we present some of the dependences of the radiation power on the generation frequency. It follows from Fig. 6 that there exists several types of dependence of the radiation power on the frequency; nonetheless, most of them are characterized by only one type of dependence, namely a weak dependence of the radiation power on the generation frequency.

We note in conclusion that it is possible to increase appreciably the output energy and the power of the giant pulse by replacing the circular diaphragms used in the present paper by slit diaphragms, which ensure minimal divergence in the plane of the angular dispersion of the prism. This replacement may affect the dependence of the radiation energy on the generation frequency, since the transverse structure of the beam will no longer be fixed along the slit[7,8], while the characteristics of the frequency tuning remain unchanged.

Our investigation has made it possible to draw a sufficiently complete picture of the influence of the character of broadening and the structure of the luminescence line on the parameters of giant pulses with tunable frequency. On the basis of the derived relations, we can propose a new method for investigating the spectral characteristics of an active medium.

At the same time, the experimental results indicate that a neodymium-glass laser emitting giant pulses has a number of properties that are valuable for various applications, namely, a wide frequency tuning range, a narrow emission spectrum, and a power that varies little in the greater part of the tuning range. We note that the best results, in the sense of the giant-pulse parameters, were obtained with LGS-250-3 glass.

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1)Since we have in the case of the employed passive shutters $B_f >> B_a$ (where B_f and B_a are the Einstein coefficients for the shutter and the active medium), we can put $m(0) = 0$ in the integration.

2)For $\alpha = 1$ we must put in expressions (4) and (5): $n_1(\bar{\nu},\tau) = n_1(\bar{\nu}, 0) = 0$ and $g_{12}(\bar{\nu}, \nu') \equiv 0$.

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