

speed of sound. This gives two discrete values of the threshold velocity, $v_{n,1}$ and $v_{n,2}$, corresponding to transverse and longitudinal phonons; these reflect correctly the positions of the above-described oscillations on the J and V axes. The calculated length of the channel is ~ 90 Å, and for the period of the oscillations in the first series we obtain the value $\Delta v_1 \approx v_{1a}/L = 0.45$ mV, where $a = 5$ Å is the Pb lattice constant. Both values are in reasonable agreement with experiment.

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SPECTRAL BROADENING IN SELF-FOCUSING OF SINGLE ULTRASHORT LIGHT PULSES IN GLASSES

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When the radiation from a mode-locked laser is self-focused in a transparent isotropic medium, one usually observes at the output intense scattered radiation with a broad spectrum in the Stokes and anti-Stokes regions [1 - 5]. The frequency-angle diagram of this scattering has characteristic "whiskers" corresponding to conditions of four-field synchronism [1, 2], and a central lobe attributed presumably to phase modulation [1, 4, 5] or to other factors [2]. We present here results of an experimental study of the previously observed [3] frequency structure of the central lobe following pumping by a group of 1 - 3 ultrashort pulses (USP) with decreasing amplitude. The period of the observed structure agrees in order of magnitude with the results of theoretical estimates based on the model of the four-field parametric interaction of amplitude-modulated waves with wave and frequency detuning.

The pump used was a mode-locked neodymium-glass laser. A burn-out mirror separated 1 - 3 pulses with a total energy $\sim 6 \times 10^{-4}$ J from the complete train of USP.

The filtered infrared radiation was focused with lenses having focal lengths from 20 to 60 cm into glass samples 0.5 - 45 cm long. The power density at the focus of the lens, at an estimated USP duration 5 psec, was 10^{11} W/cm². The scattering was observed only in samples longer than 2 cm. This is apparently due to the action of the self-focusing, which increases the power density to the previously-obtained threshold value 10^{12} W/cm² [1, 2]. There was no additional scattering-intensity gain in longer samples. The anti-Stokes part of the frequency-angle distribution of the scattering was registered with an ISP-51 spectrograph. Typical spectrograms are shown in Figs. a and b, together with a mercury reference spectrum (Fig. c). At a slight excess above the threshold power, the scattering spectrum, with the exception of isolated cases, always had a periodic structure of the type shown in Fig. a (fused quartz 5 cm long). When the input power was increased, this structure became more



Typical observed broadening spectra with a fused-quartz sample 5 cm long: a) threshold pumping; b) twofold excess above threshold; c) spectrum of mercury with lines of neodymium laser (1.06 μ) and He-Ne laser (0.63 μ). The ordinates are the angles.

complicated by a superposition of several periods and an additional large-scale structure (Fig. b, twofold excess over threshold). Under identical conditions, the period of the structure was not reproducible from pulse to pulse, and could be different in different sections (table). The values of the periods and their non-reproducibility do not make it possible to attribute the observed structure to molecular scattering mechanisms (for SRS the shift is 600 cm^{-1} in fused quartz and 400 cm^{-1} in glass [6]). On the other hand, in the experiments performed in samples of fused quartz and K-8 glass, the periods of the structure were approximately the same in the two cases and did not coincide with the values indicated for SRS.

Measured periods of the frequency structure with fused quartz 5 cm long, a 30-cm focusing lens, and threshold power density

Spectrogram No.	wavelength, nm	structure period, cm^{-1}
23 a	700	91
	500	92
25b	700	64
30a	850	63
33a	800	77
51a	700	58
	600	45
	540	30
32a	1060	60

Let us consider, in the given-pump-field approximation, the collinear (along the z axis) interaction of four waves $E_i(z, t) = A_i(z, t)\exp[i(\omega_i t - k_i z)] + \text{c.c.}$ ($i = 0, 2$) with the synchronism conditions approximately satisfied: $2\omega_0 = \omega_1 + \omega_2 + \Delta\omega$ and $2k_0 = k_1 + k_2 + \Delta k$. For plane waves, collinear interaction without mismatch will have little effect, owing to dispersion. In the presence of mismatch, however, we can attempt to find the solution in the form of waves with AM (or FM) modulation. We shall use the equations for the slow amplitudes of waves of the type given in [7], adding to the right-hand sides the corresponding reactive terms [8, 9]:

$$\frac{\partial A_{1,2}}{\partial z} + \nu_{1,2} \frac{\partial A_{1,2}}{\partial \eta} - \frac{i}{2} (k_{1,2})''_{\omega\omega} \frac{\partial^2 A_{1,2}}{\partial \eta^2} = i a_{1,2} \times$$

$$\times [A_0^2 A_{2,1}^* e^{i(\Delta\omega\eta - \Delta\tilde{k}z)} + 4A_0^2 A_{1,2}] . \quad (1)$$

Here

$$\eta = t - \frac{z}{v_0}; \quad v_i = \frac{1}{(k_i)'_{\omega}} = \left(\frac{\partial k_i}{\partial \omega_i} \right)^{-1}; \quad \nu_i = (k_i)'_{\omega} - (k_0)'_{\omega};$$

$$i = 1, 2; \quad (k_i)'_{\omega\omega} = \frac{\partial^2 k_i}{\partial \omega^2}; \quad \Delta\tilde{k} = \Delta k - (k_0)'_{\omega} \Delta\omega; \quad a_i = \frac{2\pi\chi k_i}{\epsilon_i};$$

χ is the component of the nonlinear-susceptibility tensor. We seek the solution in the form

$$A_i = A_2 = A^0 \exp[(\beta + i\gamma)\eta - (\lambda + i\mu)z + i\phi]. \quad (2)$$

Substituting (2) in (1), we obtain

$$\gamma = -\frac{\Delta\omega}{2}; \quad \mu = -\frac{\Delta\tilde{k}}{2} (\Delta\tilde{k} = \Delta\tilde{k}(\Delta\omega)); \quad (3)$$

$$-\lambda + \nu_i \beta + 2(k_i)''_{\omega\omega} \beta\gamma = a_i A_0^2; \quad -\mu + \nu_i \gamma - (k_i)'_{\omega} (\beta^2 - \gamma^2) = 4a_i A_0^2. \quad (4)$$

It is seen from (4) that the gain of the modulated waves is maximal and equals $a_i A_0^2$ (the value for plane waves at synchronism) for waves with $\beta = 0$. This requirements follows also from the condition that the solution be bounded at infinity. Taking this into account, we obtain approximately two roots that give the periods of the amplitude modulation:

$$\gamma_1 \approx -\frac{(k_i)''_{\omega\omega}}{(k_i)'_{\omega}}, \quad \gamma_2 = -\frac{1}{(k_i)'_{\omega}} \left(\frac{\Delta k}{2} - 4a_i A_0^2 \right). \quad (5)$$

The first root does not satisfy the conditions that the amplitudes be slow, and is discarded. An estimate of the second root for the degenerate case when $\Delta k = 0$ yields a modulation frequency corresponding to shifts of 50 cm^{-1} at $\lambda_0 = 1.06 \mu$, $\chi = 10^{-13}$ cgs esu and a power density 10^{12} W/cm^2 . Estimates of Δk for fused quartz in the nondegenerate case [10] have shown that this dispersion part of the wave mismatch does not exceed 10 cm^{-1} .

Thus, in collinear four-field interaction, waves can be amplified in a wide frequency range, with an envelope modulated at a certain frequency that is determined by the dispersion and the power. For three-field generators, such processes were considered already in [7]. In our experiment we apparently observed the four-field analog of such a phenomenon, which, owing to the participation of the reactive terms, exhibits considerable differences.

It is seen from (5) that at high pump powers the modulation frequency will be determined entirely only by the field, and will be the same for all waves. The phases of the amplified waves are also rigorously connected with the pump phase [9], and it follows therefore that all the amplified frequency components will be synchronized, and their total envelope will constitute a train of

USP. In such a case, the observed total width of the spectrum will correspond to USP of approximate duration 10^{-14} sec. We arrive at analogous conclusions if we consider the consecutive collinear interaction of AM components, which is described by equations of the Mathieu type with a spectrum of multiple frequencies.

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METHOD OF INCREASING THE EMISSION SPECTRUM WIDTH IN A NEODYMIUM-GLASS LASER WITH A PASSIVE SHUTTER

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Broadening the spectrum in a mode-locked laser is one of the ways of increasing the peak power of the optical emission. Usually, however, to obtain a higher degree of mode locking and to decrease the influence of nonlinear effects in a neodymium laser, one operates at pumps close to threshold. In these cases the generation spectrum width is relatively small, $4 - 6 \text{ cm}^{-1}$, and the spectrum lies near the maximum of the gain contour. Since in the general case the width of the generation spectrum is determined by the Q of the different modes of the resonator and by the width and shape of the gain contour, one can expect, if the mode selection is eliminated, that variation of the shape of the gain contour will also permit variation of the generation spectrum width. We have attempted in this study to control the width of the spectrum in such a manner.

In our experiments, the gain contour was deformed by "burning out" the inverted population and producing a dip in the center of the inhomogeneously broadened luminescence lines of the neodymium ions, and also by migration of the excitation energy from some regions of the active rods to others. Both processes exert a strong influence on the spectral and temporal characteristics of a solid-state laser and determine, in particular, the kinetics of the free-generation spectra of a neodymium-glass laser [1].

To investigate the indicated spatial interaction of the different regions of the active medium, and its influence on the width of the generation spectrum, we used a passive-shutter laser with two generation channels formed in