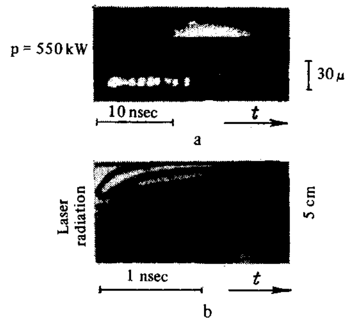

FIG. 1

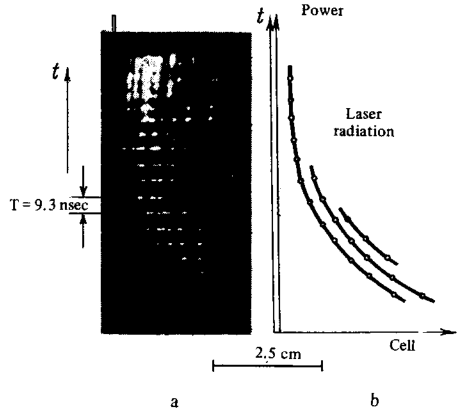

FIG. 2

ments, and the results could be explained only from the point of view of the moving foci.

In 1969, a detailed investigation was made[11] of the propagation of laser radiation in liquids. Single-mode radiation (one angular mode and one axial mode) passed through a cell with the investigated liquid. The radiation entering the cell had a plane phase front and approximately Gaussian transverse distribution. The process was registered with an electron-optical converter operating in the linear scan regime. Figure 1a shows a typical photograph of a time scan of the end face of the cell. This photograph shows clearly the successive passage of the foci of the multifocus structure through the end face (lower trace) when the laser-radiation power (upper trace) is varied. Figure 1b shows a scan of the scattered radiation obtained from the side of the cell. This photograph shows the motion of the individual foci in the direction towards the entrance face. Both the total number of foci and the maximum rate of their motion ($\sim 3 \times 10^9$ cm/sec) agree well with the theoretical calculated values.

The influence of the shape of the laser pulse on the character of the damage in glasses was investigated in[12]. In the case of a bell-shaped pulse, the damage was in the form of a long filament several microns in diameter; in the case of a rectangular pulse, the damage was produced in individual points. These results are likewise in good qualitative agreement with the concept of the multifocus structure. Similar conclusions are reached also by the authors of[13], who investigated the character of the damage in sapphire.

The multifocus-structure theory has predicted that in the case of ultrashort pulses, with duration $\tau < l/c$ (where l is the length of the cell), the focus splits after its production into front and rear focal points. The rear point becomes stationary relative to the medium after a certain time, while the front point moves with superluminal velocity. Loy and Shen[10] have confirmed experimentally that the front point moves with a superluminal velocity that reaches a value 2c.

We have investigated in detail the multifocus structure for the case of ultrashort pulses. Radiation from a neodymium laser operating in the axial-mode locking regime, passed through a cell with the investigated liquids (nitrobenzene or carbon disulfide). A typical photograph of the scan, obtained with an electron-optical converter, and taken from the side of the cell, is shown in Fig. 2a. It is seen from the photograph that the points where the rear foci stop actually exist. The intensity of scattering from these points is much higher, as it should be. The number of focal points increases with increasing beam power. These experiments are also in satisfactory agreement with the theory (Fig. 2b).

Mention should also be made of changes in the broadening of the spectrum of ultrashort pulses, which according to[14] can also be attributed from the point of view of the multifocus structure. Finally, only this theory can explain certain stimulated-Raman-scattering features connected with the generation of ultrashort pulses[10c].

Thus, summarizing all the results, it can be stated that the multifocus-structure concept has been experimentally confirmed.

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V. I. Talanov. Certain Problems in the Theory of Self-Focusing. The description of stationary self-focusing in the case of nonlinearity of the Kerr type is usually based on the parabolic equation[1]

$$2ikE'_z = \Delta_\perp E + k^2\varepsilon' |E|^2 E \tag{1}$$

for a slowly varying amplitude of the total field Ee^{-ikz}. A characteristic feature[2] of this description is the presence of foci with infinite field intensity at a beam power P higher than the critical self-focusing power P_{cr}. The field divergence at the foci can be naturally regarded as a consequence of the idealizations on which Eq. (1) is based. A very simple generalization of this equation, which leads to a limitation of the field, consists in taking the nonlinearity saturation into account by replacing $|E|^2$ by the function $|E_{sat}|^2 f(|E|^2/|E_{sat}|^2)$, such that $f(u) \to u$ at $u = |E|^2/|E_{sat}|^2 \ll 1$ and $(u) \to 1$ at $u \gg 1$.

On the basis of the results of[3] it is easy to see that the picture of the self-focusing of a beam with fixed profile and power $P > P_{cr}$ will go through the following successive phases when the saturation field E_{sat} changes from 0 to ∞: defocusing of the beam, complete trapping of the beam by the produced periodically-modulated dielectric waveguide, partial trapping of the beam by a highly irregular waveguide with rescattering of the power $P_f \sim P_{cr}$ in the vicinities of the first foci. Finally, in the limiting case $E_{sat} = \infty$ described by Eq. (1), the picture of stationary self-focusing can be illustratively represented in the form of energy streamlines that converge at individual foci and emerge from them in the form of infinitesimally thin radial scattering streams (they are infinitesimally thin because the longitudinal diffusion of the field amplitude is neglected). By changing over from the streams of finite width, it is easy to show that such radial streams do not influence the propagation of the peripheral radiation, which indeed makes up the picture of the successive foci with infinite intensity. The use of an implicit difference scheme[4] automatically excludes from consideration the scattered streams for any discretization of the coordinate z, and in this sense is equivalent to taking into account multiphoton absorption at the foci. Naturally, the self-focusing picture obtained thereby cannot go over continuously into the picture of self-focusing in a medium with saturating nonlinearity.

Multiphoton absorption can be taken into account also directly, by supplementing the right-hand side of (1) with the term $-ik\chi_m|E|^{2m}E$ (m = 1, 2, 3, ...). In this case the picture of the self-focusing changes with increasing coefficient χ_m, from a multifocus beam to a focal filament that is continuously supplied by a peripheral radiation flux.

The factors influencing the picture of the field in self-focusing include also stimulated scattering in the strong-field regions. In particular, at a large threshold increment of SRS, starting with the spontaneous-noise level, the scattered radiation can be localized near the axis of the beam of main radiation, causing the formation of self-focusing filaments.

In a theoretical analysis of the picture of nonstationary self-focusing, one observes the following effects: formation of converging filaments behind the traveling focus in the case of inertial nonlinearity $\Delta\epsilon \sim \int_{-\infty}^{t} |E|^2\,dt$, and production of large frequency shifts Ω upon "collapse" of the beam, in media having both inertial and non-inertial nonlinearity, as a result of the transverse compression of the self-focusing channel[5].

The latter effect can explain the experimentally observed broadening of the spectrum in self-focusing of ultrashort pulses. Thus, in the case of noninertial nonlinearity it follows from the approximate theory developed in[5] that the broadening of the spectrum to one side of the fundamental frequency at $z \approx z_{sf}$ and $P \gg P_{cr}$ should be of the order of

$$\Omega = [(2PP_{cr}^{-1} - 1)/4\,(PP_{cr}^{-1} - 1)]\,\Omega_{hom}\,a_0^2/a_f^2, \tag{2}$$

where $\Omega_{hom} = -2P_t'/P_{cr}(PP_{cr}^{-1} - 1)^{1/2}$ is the broadening of the spectrum in a homogeneous channel of radius a_0 over a self-focusing length z_{sf}. It it is assumed that the limitation of the radius of the focal spot a_f is due to the finite time τ of establishment of the nonlinearity, i.e., $a_f \sim (v_f\tau/k)^{1/2}$, where $v_f = (z_{sf})_t'$ is the velocity of the focal spot, then, taking into account the relation given for $z_{sf}(P/P_{cr})$ in (2) we can easily obtain the following estimate of the broadening of the spectrum:

$$\Omega\tau \sim (2PP_{cr}^{-1} - 1)/PP_{cr}^{-1}. \tag{3}$$

The right-hand side of this equation is determined in accordance with the aberration-free approximation[5] and has been corrected by the factor P_{cr}/P, with allowance for the fact that the power fed to the focus is $P_f < P_{cr} < P$. It follows from (3) that when $P/P_{cr} \gg 1$ the broadening of the spectrum saturates at the level $\Omega \sim 2/\tau$ determined only by the nonlinearity relaxation time τ. This time is of the order of $10^{-12}-10^{-11}$ sec for the orientational Kerr effect and of the order of 10^{-15} sec for the electronic effect. The small value of τ in the latter case explains the anomalously large broadenings of the spectrum, which overlap the entire visible band, in experiments on self-focusing of ultrashort pulses in glasses at the wavelength λ = 1.06 μ.[6]

If the medium contains two nonlinearity mechanisms (e.g., the orientational and electronic Kerr effect) with relaxation times $\tau_1 > \tau_2$, but $P_{cr1} < P_{cr2}$, then the broadening of the spectrum of pulses of duration $\tau_p > \tau_1 P_{cr1}/P_{cr2}$ is determined by the time τ_1 even when $P > P_{cr2}$. This is due to the fact that the low-threshold mechanism causes the beam to break up into regions with power $P \sim P_{cr1} \sim P_{cr2}$, inside of which the self-focusing due to the high-threshold mechanism is impossible.

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