Aberrations and Extreme Divergences of Continuous Laser Radiation in Defocusing Media

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The main regularities of thermal defocusing of high-intensity laser radiation are discussed. A radiation equation is derived and employed for analyzing the ray paths in an absorbing medium with allowance for spherical aberrations. Astigmatism of thermal lenses induced in the medium by an elliptic-cross-section beam is considered in the paraxial approximation. The extreme defocusing properties of optically transparent media are determined. The highest intensities attainable at various distances from the entrance into a nonlinear media are estimated. Results are presented of experiments on thermal defocussing of radiation from a one-watt argon laser which passed through a cell with ethyl alcohol. The experimental data on nonlinear beam divergence and field restriction are in accordance with the theoretical calculations.

1. INTRODUCTION

THE thermal self-action of continuous radiation in media with $dn/dT < 0$ leads to defocusing of laser beams. Most papers (see, e.g.,[1-6]) have been devoted to external defocusing, when a layer of nonlinear medium acts on the beam as a thin thermal lens. The initial stage of development of the internal defocusing (the beam begins to diverge inside the nonlinear medium) was considered by Smith[7] and by Livingston[8]. In this paper we present the results of an investigation of internal defocusing in the case of strong broadening of a laser beam inside a cell filled with alcohol.

A theoretical analysis of the laws governing internal defocusing was carried out in the case of an immobile medium. We obtained the ray equation for cylindrically-symmetrical and two-dimensional beams with arbitrary initial amplitude profiles. This makes it possible to calculate the nonlinear spherical aberrations and to reveal the limiting defocusing properties of optically transparent media. The nonlinear divergence produced by a thick thermal lens is determined essentially by the total beam power (it depends weakly on the initial radius of the beam). Another class of aberrations is astigmatism of thermal lenses induced in the medium by a beam with elliptical cross section, and is considered in the paraxial approximation. The self-defocusing leads to a limitation of the field on the beam axis[5]. Estimates of the limiting intensities are obtained for different distances from the entrance into the nonlinear medium.

The experiments on the observation of internal defocusing were carried out with a beam of continuous radiation, from an argon laser, passing through a cell with alcohol. Particular attention was paid to two circumstances, namely, the lowering of the threshold of the transition from external to internal defocusing and preventing convection. To this end, we chose, first, the optimal absorption and, second, used focused beams. The contraction of the beam into a narrow neck (after passage through the lens) by a factor of 10 decreased the convection and at the same time extended the power range within which internal defocusing can be observed. We investigated experimentally the dependence of the nonlinear divergence of the beam on its power and on the absorption of the alcohol (with variation of the fuchsine concentration), and also the effect of limiting the intensity. The results of the experiments agree with the theoretical calculations.

2. NONLINEAR GEOMETRICAL OPTICS

Stationary thermal defocusing of wave beams is described in the quasioptical approximation by the system of equations[4]

$$\frac{\varkappa}{r}\frac{\partial}{\partial r}\left(r\frac{\partial T}{\partial r}\right)+\frac{cn_0}{8\pi}\alpha A^2=0, \qquad (1)$$

$$2\frac{\partial s}{\partial z}+\left(\frac{\partial s}{\partial r}\right)^2=-\frac{2}{n_0}\frac{dn}{dT}T+\frac{1}{k_0^2A}\left(\frac{\partial^2 A}{\partial r^2}+\frac{1}{r}\frac{\partial A}{\partial r}\right), \qquad (2)$$

$$\frac{\partial A^2}{\partial z}+\alpha A^2+\frac{1}{r}\frac{\partial}{\partial r}\left(rA^2\frac{\partial s}{\partial r}\right)=0. \qquad (3)$$

Here T is the change of the temperature of the medium under the influence of the beam, $n = n_0 - T dn/dT$ is the refractive index, κ is the coefficient of thermal conductivity, α is the absorption coefficient, $E = (1/2)A\exp \times [i(\omega t - k_0 z - k_0 s)] + c.c.$ is the electric field of the wave, k_0 is the wave number, and s is the increment to the eikonal of the plane wave. In (1)–(3), no account is taken of the diffusion of the temperature and of the complex amplitude along the beam axis (the Z axis). As seen from (1)–(3), the thermal self-action leads to nonlinear refraction of the rays. The defocusing properties of the absorbing medium then depend on the type of intensity modulation of the propagating beam. In the general case, self-induced thermal lenses are subject to aberration. In thin lenses (phase correctors), calculation of the spherical aberration entails no special difficulty[3,4,6]. At the same time, in the case of strong defocusing of the beams, the aberrations can develop in the nonlinear medium itself. The unique nature of the thermal self-action makes it possible to construct a nonlinear geometrical optics (within the framework of the quasioptical approximation), and thus analyze the nonlinear spherical aberrations for beam with arbitrary amplitude profiles at the entry into the medium.

We shall characterize an elementary ray $r = r(z)$ by its inclination to the beam axis

$$\theta = \partial s / \partial r \tag{4}$$

and by the power contained in a ray tube of radius r,

$$P_r = \frac{cn_0}{4} \int_0^r A^2 r dr. \tag{5}$$

Changing over to geometrical optics (we let $k_0 \to \infty$ in (2)), we obtain from (1)–(3) with allowance for (4) and (5) the system

$$\frac{\partial \theta}{\partial z} + \theta \frac{\partial \theta}{\partial r} = \frac{\alpha P_r}{2n_0 r \pi \varkappa} \frac{dn}{dT}, \tag{6}$$

$$\frac{\partial P_r}{\partial z} + \theta \frac{\partial P_r}{\partial r} + \alpha P_r = 0. \tag{7}$$

We use the method of characteristics to obtain the ray equation

$$\frac{d^2 f}{dz^2} = \frac{\alpha e^{-\alpha z}}{R_T f}, \tag{8}$$

where $f(z) = r(z)/r_0$ is the relative distance from the ray to the beam axis, r_0 is the initial coordinate of the ray, the quantity

$$R_T = \frac{2n_0 \varkappa}{I_{av}(r_0)\, dn/dT} \tag{9}$$

characterizes the strength of the nonlinear refraction, $I_{av}(r_0) = P_0(r_0)/\pi r_0^2$ is the average field intensity in a ray tube with initial power $P_0(r_0) = P_{r=r_0}$.

Equation (8) describes the course of the light-beam rays at an arbitrary initial divergence $\theta_0(r_0) = dr/dz|_{z=0}$ with arbitrary initial intensity profile $I_0(r_0)$. In a linear medium, $dn/dT = 0$, the rays are straight lines

$$f = 1 + \theta_0 z / r_0. \tag{10}$$

In a defocusing medium, the ray trajectories become bent, and different beam rays experience different nonlinear refractions. Only in a beam with a homogeneous intensity distribution over the cross section ($I_{av} \equiv I_0$) is the refraction strength R_T (9) constant and are there no nonlinear spherical aberrations (there will undoubtedly be spherical aberrations connected with the distortion of the initial wave front). The aberration picture produced for other beams is that of thermal self-defocusing.

3. NONLINEAR SPHERICAL ABERRATIONS

The region of effective nonlinear refraction can be limited, as seen from (8), by the finite length of the medium $0 < z < l$, by the optical thickness of the transparency layer $0 < z \lesssim \alpha^{-1}$, or by the region of expansion of the ray tube. In analyzing the action of thermal lenses, it is advisable to subdivide them into two groups (thin and thick), with respect to the character of the limitation of the action of the refraction strengths.

Thin thermal lenses. Assume that the beam barely broadens in the nonlinear medium, and that its wave front becomes curved. Such an absorbing layer operates as a phase corrector. Assuming in the right-hand side of (8) $f = 1$, we obtain the law whereby a thin thermal lens of thickness l changes the slope of the ray:

$$\theta = \theta_0(r_0) + r_0(1 - e^{-\alpha l}) / R_T(r_0). \tag{11}$$

We see that the focal distance of the thermal lens depends in the general case on the initial coordinate of the ray

$$F_T = R_T(r_0)(1 - e^{-\alpha l})^{-1}. \tag{12}$$

In a Gaussian beam, the average intensity decreases with increasing radius of the ray tube (a is the half-width of the beam):

$$I_{av} = I_0[1 - \exp(-2r_0^2 / a^2)](a^2 / 2r_0^2).$$

The largest angular deflection is experienced by a beam with coordinate $r_0 = 0.8a$. In the absorptionless approximation, formulas analogous to (11) and (12) were obtained earlier for a weakly absorbing medium[1] and were then generalized to include the case of arbitrary attenuation[2].

The behavior of the beam behind a thin lens obeys the laws of linear optics. The corresponding results are given in[3,4,6]. We note only that in accordance with (11) there can occur, in definite cases, compensation of the linear and nonlinear aberrations, so that $\theta(r_0) \sim r_0$.

The optical strength of a thermal lens increases in direct proportion to the intensity of the transmitted beam. At sufficiently high wave intensity, therefore, the beam should diverge to such an extent that the thin-lens approximation $F_T > \alpha^{-1}(1 - e^{-\alpha l})$ will no longer hold. This imposes limitations on the initial beam intensity

$$I_0 \leqslant I_{TT} = \frac{n_0 \varkappa \alpha}{(1 - e^{-\alpha l})^2\, dn/dT} \tag{13}$$

At higher intensities it is necessary to take into account the decrease of the strength of the nonlinear refraction as a result of the beam broadening.

Thick thermal lenses. In self-defocusing of high-intensity radiation one can neglect the dissipation of the wave in the effective-refraction zone. Putting $e^{-\alpha z} \approx 1$ in the ray equation (8) and integrating, we obtain the inclination of the ray $\theta = r_0 df/dz$ at the cross section z:

$$\theta^2(r_0, z) = \theta_0^2(r_0) + \theta_{nl}^2(r_0) \ln f(z), \tag{14}$$

where θ_{nl} characterizes the contribution of the thermal lens to the beam divergence:

$$\theta_{nl} = \left[\frac{\alpha P_0(r_0)\, dn/dT}{\pi n_0 \varkappa}\right]^{1/2}. \tag{15}$$

Integrating (14), we obtain the ray trajectories

$$z = 2F_{nl}(r_0) e^{-m^2} \int_m^{(\ln f + m^2)^{1/2}} e^{u^2} du, \tag{16}$$

where $m = \theta_0/\theta_{nl}$ and $F_{nl} = r_0/\theta_{nl}$ (see Fig. 2a, curve 1).

Let us examine the laws governing the defocusing of a Gaussian beam without initial divergence, $\theta_0 \equiv 0$. The aberration picture of the rays is shown in Fig. 1. At small beam broadenings ($1 < f \lesssim 2$) we obtain from (16) (see also[7,8])

$$f = \exp(z^2 / 4F_{nl}^2). \tag{17}$$

At distances $z \gg F_{nl}$, the beam broadens strongly, its intensity decreases, and the ray trajectories are almost straight lines:

$$f = (z / F_{nl}) \ln^{1/2}(z / F_{nl}). \tag{18}$$

Thus, the most effective defocusing of the beam is produced by a layer with thickness on the order of F_{nl}. In

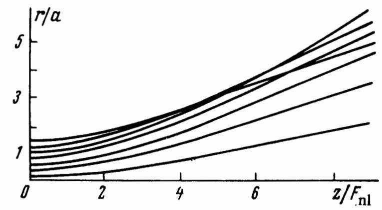

FIG. 1. Aberration picture of the course of the rays of a Gaussian beam in the case of internal defocusing.

essence, F_{nl} is the focal length of the thick lens, since the contribution made to the divergence by the succeeding regions of the medium is small. Indeed, whereas the divergence θ_{nl} is formed after the beam is broadened by a factor e, it is then doubled ($\theta = 2\theta_{nl}$) after the beam broadens by approximately 50 times (Fig. 2). The quantity θ_{nl} practically characterizes the limiting defocusing properties of the medium. It must be emphasized that the limiting divergence (15) depends on the total beam power: it increases monotonically with increasing distance between the ray and the beam axis ($r_0/a \to \infty$), whereas for a thin lens there exists a ray with maximum slope. This alters the defocusing picture, namely, in the case of internal defocusing the ring structure is much less pronounced (one external ring) than for a thin lens. This has been confirmed by the experiments (see Sec. 6).

The ray equation (8) was solved with a computer for the general case of arbitrary relations between the parameters α, l, and R_T. The numerical calculation has confirmed the possibility of introducing the thin- or thick-lens approximation in definite situations ($\alpha R_T \ll 1$ or $\alpha R_T \gg 1$). The ray trajectories in media with different absorptions are shown in Fig. 2.

We note that it is possible to obtain a single analytic formula for thin and thick lenses, analogous to (16), by approximating in (8) the exponential attenuation law by a slower one, $\sim(1+\alpha z)^{-2}$ (see[9]). In a converging beam, if $\alpha R_T = 4$ and $\theta_0 = \alpha r_0/2$, the trajectory approaches the beam axis monotonically:

$$f = e^{-\alpha z/2}. \tag{19}$$

4. ASTIGMATISM OF THERMAL LENSES

The nonlinear aberrations of thermal lenses induced in a medium by a laser beam are determined, as already mentioned, by the distribution of the intensity over the beam cross section. The aberrations produced in an axially-symmetrical field are spherical. Deviation from

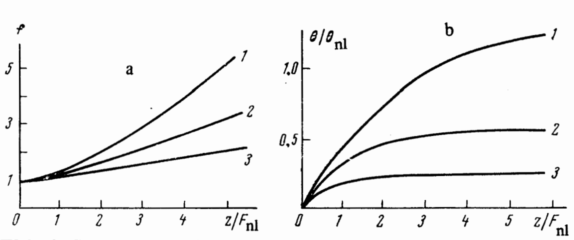

FIG. 2. Solutions of the ray equation, obtained with a computer for different absorptions of the medium: a—trajectory of a beam in a medium with absorptions αR_T equal to 0, 1, and 9; b—dependence of the nonlinear divergence of the paraxial ray tube on the traversed path in medium with the same absorptions, respectively, for curves 1, 2, and 3.

axial symmetry leads to other types of distortion of the wave or amplitude profile. In particular, defocusing of beams with elliptic cross sections (where the constant-intensity levels are ellipses) is accompanied by nonlinear astigmatism. Such an effect arises in self-focusing (of the Kerr[10,11] or thermal[12] type). We shall calculate the astigmatism of a defocusing medium in the paraxial approximation, taking into account the quadratic distortions of the wave front:

$$s = \frac{x^2}{2a_1}\frac{da_1}{dz} + \frac{y^2}{2a_2}\frac{da_2}{dz} + \varphi(z). \tag{20}$$

In this case the shape of the Gaussian beam is assumed constant:

$$A = E_0\left[\frac{a_{10}a_{20}}{a_1(z)a_2(z)}\right]\exp\left[-\frac{x^2}{a_1^2(z)} - \frac{y^2}{a_2^2(z)}\right]. \tag{21}$$

Here $a_1(z)$ and $a_2(z)$ are the principal radii of the elliptic profile in the cross section z; $a_{j0} = a_j(0)$ are the initial radii. Substituting the sought solutions (20)—(21) into the system (1)—(3), expressed in terms of the variables x, y, and z, we arrive at two ordinary equations (we assume $k_0 \to \infty$):

$$\frac{da_1}{dz^2} = \frac{2\alpha}{R_{T0}}\frac{a_{10}a_{20}e^{-\alpha z}}{a_1 + a_2}, \tag{22}$$

$$\frac{da_2}{dz^2} = \frac{2\alpha}{R_{T0}}\frac{a_{10}a_{20}e^{-\alpha z}}{a_1 + a_2}, \tag{23}$$

where the quantity R_{T_0} should be taken in accordance with formula on the beam axis (x, y → 0). The system (22) and (23) is subject to the conservation law

$$a_1 - a_2 = a_{10} - a_{20} + (\theta_{10} - \theta_{20})z, \tag{24}$$

which breaks up in a linear medium into two independent equations for the straight-line rays $a_j = a_{j0} + \theta_{j0}z$.

Adding (22) and (23), we arrive at a single equation

$$\frac{d^2 f_{av}}{dz^2} = \frac{\alpha e^{-\alpha z}}{R_{T0} f_{av}} \tag{25}$$

for the quantity

$$f_{av} = \frac{a_1(z) + a_2(z)}{2(a_{10}a_{20})^{1/2}}, \tag{26}$$

which is equal to the arithmetic mean of the radii, normalized to the geometric-mean radius of the initial cross section. Equation (26) coincides in form with the ray equation (8) considered above for axially-symmetrical beams.

In the case of defocusing of a beam that has no linear astigmatism, $\theta_{10} = \theta_{20}$, the difference between the radii (24) remains constant. At the same time, the average radius increases in accordance with (25). Thus, the more rapid defocusing along the minor diameter of the beam cross section causes its shape to become rounder. In the case of a strongly elongated cross section ($a_{20} \gg a_{10}$), as follows from (23), the minor radius of the beam a_1 increases during the initial stage of the defocusing, and the major radius remains practically unchanged, $a_2 \approx a_{20}$. For such a beam we can derive, in analogy with the axially-symmetrical problem, the following ray equation:

$$\frac{a_1}{a_{10}} = 1 + \frac{\theta_0(x_0)}{x_0}z + \frac{dn}{dT}\frac{P_1(x_0)}{n_0\varkappa\alpha a_{10}}(\alpha z - 1 + e^{-\alpha z}), \tag{27}$$

where $P_1(x_0) = \int_0^{x_0} I_0(x)dx$ is the running power of the ray tube. The behavior of the rays of a one-dimensional beam is analogous in many respects to the case of defocusing of a cylindrical beam. We note only that in accordance with (27) the nonlinear divergence of the rays increases monotonically with increasing distance from the rays to the axis. There will therefore be no aberration band pattern along the minor beam diameter.

5. LIMITATION ON FIELD INTENSITY

As the wave propagates, its intensity decreases, first, as a result of absorption and dissipation of the energy and second, as a result of the broadening of the cross section of the diverging beam. In a linear medium, the attenuation factor is a function of the parameters of the problem and is independent of the beam power. In a nonlinear defocusing medium, the beam divergence increases in proportion to the beam power (for a thin layer, the beam cross section is $S \sim P_0^2$, and for a thick one $S \sim P_0$). This circumstance imposes a limitation on the field intensity[5].

The level to which the field is limited depends on the observation point (the coordinate z) and is higher the closer we are to the entry into the medium. In the near zone ($z \leq R_d \equiv ka^2/2$) we can use for the calculation of the intensity on the axis of the cylindrical beam the geometrical-optics formula

$$I = I_0 e^{-\alpha z} f_0^{-2}, \tag{28}$$

where f_0 characterizes the broadening of the paraxial ray tube and is given by the solution of Eq. (8) as $r_0 \to 0$. In the far zone ($z \gg R_d$) diffraction effects play an important role. It is necessary here, generally speaking, to solve the complete system of Eqs. (1)–(3). The diffraction spreading of the beam can be taken into account relatively easily in the paraxial approximation (without allowance for spherical aberrations), in analogy with the procedure used in Sec. 4. Specifying the laser beam in the form (20) and (21), with $a_1 \equiv a_2$, we obtain the equation[1)]

$$\frac{d^2 f_a}{dz^2} = \frac{\alpha e^{-\alpha z}}{R_{T0} f_a} + \frac{1}{R_d^2 f_a^3}, \tag{29}$$

describing the relative change of the radius of a Gaussian beam, where $f_a = a(z)/a_0$ and $R_d = ka_0^2/2$ is the diffraction length of the beam. In the geometrical-optics approximation ($R_d \to \infty$), Eq. (29) goes over into Eq. (8) for paraxial rays. Equation (29) was solved with a computer for different values of the parameters αR_d and αR_{T0}. The character of the limitation of the field on the beam axis (28) is different for thin and thick thermal lenses.

Thick lens (internal defocusing). Typical plots of the field intensity against the input power in an optically transparent medium are shown in Fig. 3. We see that there exists an optimal input power, at which a maximum intensity is reached in a given cross section z. Besides the results of the numerical calculation, we can also derive theoretical formulas.

Inside a weakly-absorbing nonlinear layer, the maximum intensity is reached, as follows from (28), under the condition $f = z\,df/dz$, having the clear meaning that

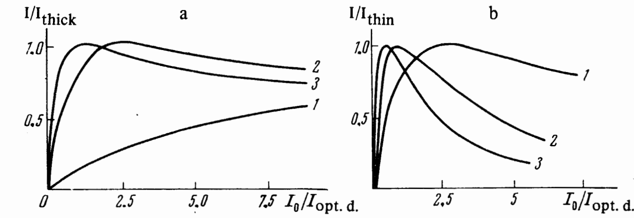

FIG. 3. Field intensity on the beam axis vs. input power: a—internal defocusing for cross sections z/R_d equal to 0.14, 1.14, and 3.0; b—external defocusing for z/R_d equal to 0.2, 1.0, and 3.0 respectively for curves 1, 2, and 3; $\alpha R_d = 10^{-3}$.

the focal length of the thick thermal lens be equal to the layer thickness z. The corresponding point can be found from Fig. 1 by drawing a tangent from the origin to the curve f(z). Regardless of the distance z, the maximum intensity is attained at the same beam broadening, f = 2.33. Such a defocusing sets in at $F_{nl} = 0.4z$, corresponding to an optimal initial intensity

$$I_{\text{opt. thick}} = 0.6\, I_{\text{lim}}, \quad I_{\text{lim}} = 2\, n_0 \varkappa / \alpha z^2 (dn/dT). \tag{30}$$

The limiting level is $I_{thick} = 0.11 I_{lim}$. In the region $I_0 \gg I_{lim}$, the intensity of the beam in a nonlinear medium decreases slowly with increasing input power (compare with the curves of Fig. 3):

$$I = I_{\text{lim}} \ln^{-1}(2 I_0 / I_{\text{lim}}). \tag{31}$$

In the far field ($z \gg R_d$), the limitation sets in at $F_{nl} \approx R_d$ or

$$I_{\text{opt. d}} = n_0 \varkappa (\alpha R_d\, dn/dT)^{-1},$$

i.e., at a constant input-power level. The field-limitation level is specified by the previous formula (see, Fig. 3, curve 3).

Thin lens (external defocusing). The cross section of a beam passing through a thin absorbing layer varies like

$$f^2 = (1 + z/R_{T0})^2 + z^2/R_d^2. \tag{32}$$

The maximum field

$$I_{\text{thin}} = \frac{n_0 \varkappa}{z(e^{\alpha l} - 1)\,[1 + (1 + z^2/R_d^2)^{1/2}]\, dn/dT} \tag{33}$$

is reached if the beam intensity at the entry to the medium is

$$I_{\text{opt. thin}} = \frac{2 n_0 \varkappa (1 + z^2/R_d^2)^{1/2}}{z[1 - e^{-\alpha l}]\, dn/dT}. \tag{34}$$

Unlike internal defocusing, the $I(P_0)$ dependence reveals a sharper maximum when the field is limited by a thin lens (see Fig. 3).

6. EXPERIMENTAL RESULTS

Experiments on thermal defocusing of continuous radiation were carried out with an argon laser operating in the regime of a single longitudinal (λ = 4880 Å) and the fundamental transverse mode. The Gaussian beam had a radius $a_0 = 0.45$ mm and a diffraction divergence $\theta_d \approx 1'$. The laser output power could be varied in the interval 1–300 mW, in which case the intensity in the beam ranged from 0.3 to 100 W/cm^2. The power was monitored with the aid of an F-18 photocell.

The laser beam passed through a cell (l = 10 cm)

[1)]We note that the average intensity of the ray tube with $r_0 = a$ is half as large as on the beam axis. Therefore, generally speaking, the strength of the nonlinear refraction in (29) is overestimated.

filled with ethyl alcohol ($n_0 = 1.36$, $dn/dT = 4.2 \times 10^{-4}$ deg^{-1}, and $\kappa = 1.5 \times 10^{-3}$ W/cm-deg). The cell was located at a distance 20 cm from the laser. In the alcohol, the transition from the thin to the thick lens occurs at an intensity (see (13))

$$I_{TT} = 4.85\alpha(1 - e^{-\alpha l})^{-2}. \tag{35}$$

The smallest threshold in (35) corresponds to an absorption $(\alpha l)_{min} = 1.25$. In our experiment $l = 10$ cm and $\alpha_{min} = 0.125$ cm^{-1}, $I_{TT} \approx 1$ W/cm^2. From these considerations, the absorption coefficient was chosen in the range $\alpha = 0.1$–0.5 cm^{-1} by varying the fuchsin concentration. It is obvious that even in the unfocused beam there was a sufficient power margin to observe internal defocusing. It is possible to lower the power threshold P_{TT} and by the same token to extend the range of observation of internal defocusing, by increasing the field intensity via contraction of the beam. Prefocusing of the laser beam onto the forward face of the cell by means of a lens with $F_L = 17$ cm ensured a lowering of the threshold power by a factor 100 in comparison with the diffraction beam.

The use of focused beams in the experiments turned out to be advantageous also for another reason. The point is that the theory developed above pertains to thermal self-action in immobile media. However, the laser beam may produce free convection in the liquid and thus distort the defocusing picture[2-4]. The convection effect depends strongly on the beam cross section and disappears in the case of a sufficient degree of focusing.

Let us stop to discuss this question in greater detail. The velocity of the vertical convection stream produced by horizontal propagation of a laser beam in a liquid is[3]

$$v_c \approx \beta g \alpha P_0 a^2 / 32\nu\kappa, \tag{36}$$

where β is the temperature coefficient of expansion, g the acceleration in the field of gravity, and ν the kinematic viscosity. The degree of the influence of the convection on the thermal defocusing depends on the relation between the thermal diffusion and the heat carried away by the convection stream, which is characterized by the Peclet number[13]

$$\gamma_c = a v_\kappa / 4\chi, \tag{37}$$

where χ is the coefficient of temperature conductivity. Substituting (36) in (37), we obtain the Peclet number of the convective stream in alcohol ($\beta = 1.1 \times 10^{-3}$ deg^{-1}, $\nu = 1.51 \times 10^{-2}$ cm^2/sec, $\chi = 7.2 \times 10^{-4}$ cm^2/sec):

$$\gamma_c \approx 5.2 \cdot 10^5\, \alpha P_0 a^3. \tag{38}$$

We see that the Peclet number depends strongly on the beam radius. The violation of the axial symmetry and the downward deflection of the laser beam occur under the condition $\gamma_c \gtrsim 1$. In an unfocused beam, $\gamma_c \approx 1$ and its cross section past the cell acquires the shape of a "half-moon." In this case the theory is valid after the temperature gradient is established, up to the instant of the onset of convection (on the order of 1 second). In a focused beam, $\gamma_c \approx 10^{-3}$, the convection exerts no influence whatever on the thermal defocusing, and we can use the results of theoretical calculations made for immobile media.

Typical pictures of the angle structure of the laser radiation after passing through a cell with alcohol are shown in Fig. 4 for two different beam-intensity levels: (a) $I_0 = 0.1$ W/cm^2 (external defocusing; photograph taken prior to the onset of convection); (b) $I_0 \approx 10$ W/cm^2 (internal defocusing, no convection takes place). We see that the fine structure of the spherical aberrations is less clearly pronounced in the case of a thick lens than in external defocusing.

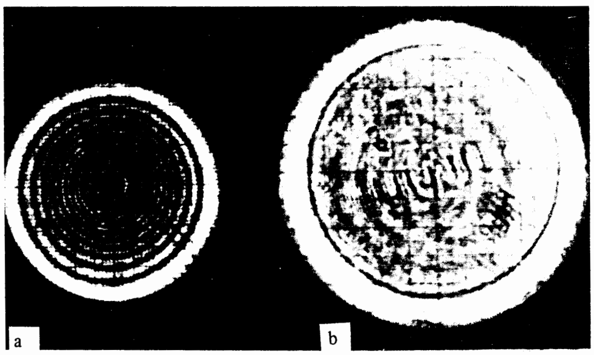

FIG. 4. Typical pictures of the angular structure of laser radiation passing through a cell with alcohol at two beam-intensity levels: a) $I_0 = 0.1$ W/cm^2 (external defocusing), b) $I_0 = 10$ W/cm^2 (internal defocusing).

The $\theta_{nl}(P_0)$ dependence for a thin thermal lens (11) was investigated in detail by a number of workers[1-5] and had the same character in our experiments. The main task of the present study was to investigate the thermal defocusing when $I_0 \gg I_{TT}$. Figure 5 shows the experimental plots of the beam divergence θ_{nl} (determined from the outer aberration ring) against the input power. The experiments were performed at different absorptions of the alcohol ($\alpha_1 = 0.10$ cm^{-1} and $\alpha_2 = 0.17$ cm^{-1}) with focused ($F_L = 13.5$ cm) and unfocused beams. The results are in satisfactory agreement with the theoretical ones. Under conditions of strongly developed internal defocusing ($I_0 \gg I_{TT}$), taking into account the limitation imposed on the thickness of the thick lens by the absorption length ($z \approx \alpha^{-1}$), the beam divergence according to (14), (15), and (18) is

$$\theta^2 = \frac{\alpha P_0 n_0}{2\pi\kappa}\frac{dn}{dT}\ln\left(\frac{I_0}{\alpha n_0 \kappa}\frac{dn}{dT}\right). \tag{39}$$

For pure ethyl alcohol

$$\theta^2\ (\text{deg}) = 0.2\alpha P_0 \ln(0.2 I_0/\alpha), \tag{40}$$

where α is given in cm^{-1}, P_0 in mW, and I_0 in W/cm^2. Formula (40) describes well a number of laws governing internal defocusing. For example, the ratio $(\theta_2/\theta_1)^2 = \alpha_2/\alpha_1 = 1.7$ has been confirmed experimentally. Likewise, Eq. (40) agrees with the increase of the beam divergence when the beam is focused with a lens of $F_L = 13.5$ cm (the beam intensity is increased 100 times). However, the absolute values of the divergence differ from the measured ones. This may be due to the fact that formulas (39) and (40) are asymptotic and overestimate the divergence, and also to factors unaccounted for by the theory (the joint manifestation of aberration and diffraction, longitudinal diffusion of the heat, etc.).

The second series of experiments was performed with the aim of observing the limitation of the intensity of radiation passing through a cell with alcohol at

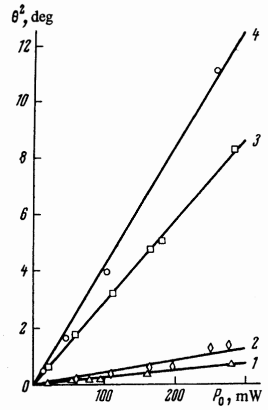

FIG. 5

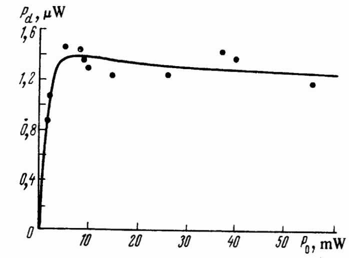

FIG. 6

FIG. 5. Experimental plots of the beam divergence θ_{nl} at different absorptions of the alcohol, $\alpha_1 = 0.10\ cm^{-1}$ and $\alpha_2 = 0.17\ cm^{-1}$, with focused beams (curves 3 and 4) and unfocused beams (1 and 2).

FIG. 6. Dependence of the diaphragmed beam power passing through a cell with alcohol ($\alpha = 0.46\ cm^{-1}$) on the initial power of an argonlaser beam.

$\alpha = 0.46\ cm^{-1}$. A diaphragm of diameter d = 0.35 mm was placed on the beam axis at a distance L = 65 cm from the cell. The transmitted radiation was registered with a photomultiplier. The dependence of the power passing through the diaphragm on the initial power of the argon-laser beam is shown in Fig. 6. We present theoretical estimates of the observed effect, using the thin-lens formulas (33) and (34). In accordance with the formulation of the experiment, we put $z = l + n_0 L$ (we take into account refraction at the exit window of the cell); in the near field $z \ll R_d$. Then the laser-beam intensity (34) corresponding to the maximum registered power is equal to $I_{opt.\,thin} = 0.52\ W/cm^2$ at a beam power $P_{opt} = 16$ mW. The maximum intensity on the axis at the point z (33) is $P_{thin} = 1.3$ mW, corresponding to a power flux $P_{d,\,thin} = 1.3\ \mu W$ through the diaphragm (experiment yielded respectively $P_{opt} \approx 10$ mW and $P_{d,\,thin} = 1.4\ \mu W$). The changeover from a thin to a thick lens occurred at $I_{TT} = 2.25\ W/cm^2$, $P_{TT} = 7.75$ mW, i.e., at the same power. Therefore, a changeover to a thick lens occurs in the case under consideration in the limitation region. After the transition to the thick lens, according to (30) and (31), we have $I_{lim} = 0.4\ mW/cm^2$ and $P_{d,\,thick} \approx 0.4\ \mu W$, which is somewhat lower than the observed result.

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