

NONLINEAR OPTICAL PHENOMENA

PACS numbers: 42.25.Fx; 42.25.Ja
DOI: 10.1070/QE2002v032n05ABEH002209

Self-diffraction of light upon optical poling of glass

M.K. Balakirev, L.I. Vostrikova, V.A. Smirnov

***Abstract*. The self-diffraction of light appearing upon optical poling of bulk glass samples is studied. During poling, a refractive-index grating is accumulated in the medium, on which one of the beams or both beams diffract efficiently. A theoretical expression is obtained for the amplitude of diffracted radiation. The results of the experimental study of this phenomenon in oxide glasses are in agreement with the theory, which explains the formation of a spatially periodic electric field in the medium by the coherent photogalvanic effect. The self-diffraction of light can be efficiently used for studying the physical properties of purely optical poling of media.**

Keywords: *glass, optical poling, self-diffraction, coherent photogalvanic effect.*

1. Introduction

Illumination of glass by mutually coherent two-frequency laser radiation of the first and second harmonics results in the accumulation of the reversible long-lived static polarisation in the glass. This effect is assigned to the formation of a spatially periodic electric field in the medium under the action of multifrequency coherent light, leading to a change in the polarisation of the medium [1]. As a result of this effect, which is called purely optical poling (OP) [2], the glass loses its symmetry and acquires the properties of a uniaxial crystal. In the region of interaction of light beams, the modulation of the refractive index appears (the Δn anisotropic grating) [3], which is responsible for birefringence, and the second-order polarisability (the $\chi^{(2)}$ grating) [1] responsible for the appearance of three-wave interactions [4, 5], which were earlier forbidden by the glass symmetry.

Optical poling was often studied by investigating the SHG on an induced $\chi^{(2)}$ grating [5, 6]. This phenomenon can be also studied from diffraction of light by the inhomogeneity of the refractive index induced in glass. The sensitivity of the latter method is higher because the SHG intensity is proportional to the square of a static field, whereas the intensity of diffracted radiation is proportional to the fourth power of this field. The first observations of light scattering from the Δn grating induced in glass are reported in papers [3, 7].

In this paper, we studied the self-diffraction of light from accumulated refractive-index gratings in bulk glasses. We obtained the theoretical expression for the amplitude of diffracted radiation and studied the properties of this radiation. The self-diffraction of light was experimentally investigated in some oxide glasses. The model of this phenomenon presented in the paper can be used for studying the physical picture of the OP of media.

2. Theory of self-diffraction of light upon OP of glasses

Consider the self-diffraction of the fundamental-harmonic radiation upon the OP of glass by crossed Gaussian beams. In a classical variant, the OP of a sample is performed by the two-frequency mutually coherent radiation of the first ($\boldsymbol{E}_1 = \boldsymbol{e}_1 E_1(\boldsymbol{r})\exp[\mathrm{i}(\boldsymbol{k}_1\boldsymbol{r} - \omega t + \psi_1)]$) and second ($\boldsymbol{E}_2 = \boldsymbol{e}_2 E_2(\boldsymbol{r})\exp[\mathrm{i}(\boldsymbol{k}_2\boldsymbol{r} - 2\omega t + \psi_2)]$) harmonics of a laser. It is assumed that in this case, the coherent photogalvanic (CP) current

$$\boldsymbol{j}(\boldsymbol{r}) = [\sigma_1\boldsymbol{e}_1(\boldsymbol{e}_1\boldsymbol{e}_2) + \sigma_2\boldsymbol{e}_2]E_1^2(\boldsymbol{r})E_2(\boldsymbol{r})\cos(\Delta\boldsymbol{k}\boldsymbol{r} + \Delta\psi), \quad (1)$$

appears in the region of the interaction of light beams [1, 8, 9], where $\Delta\boldsymbol{k} = 2\boldsymbol{k}_1 - \boldsymbol{k}_2$; $\Delta\psi = 2\psi_1 - \psi_2$; and σ_1, σ_2 are the CP constants; and $\sigma_1 = 2\sigma_2$ for isotropic media of the most general symmetry class $\infty\infty m$ [9].

The separation of charges by the CP current results in the formation of an 'imbedded' electrostatic field in the medium. In the general case, the spatial distribution of the field is complicated because the envelope of the current grating (1) has a shape determined by the intersection of the Gaussian beams, and the CP current contains components that are parallel ($j_{\parallel}$) and perpendicular ($j_{\perp}$) to the grating planes. However, the current $j_{\perp}$ causes the accumulation of charges directly on the grating planes, unlike the current $j_{\parallel}$, which causes the accumulation of charges at the boundary of the overlap region of the interacting beams. Therefore, under the condition $1/(\Delta k a) \ll 1$ (where a is the size of the overlap region), the current $j_{\perp}$ makes the main contribution to the formation of the periodic grating, and the field produced in the medium can be approximated by the expression

M.K. Balakirev, L.I. Vostrikova, V.A. Smirnov Institute of Semiconductor Physics, Siberian Branch, Russian Academy of Sciences, prosp. akad. Lavrent'eva 13, 630090 Novosibirsk, Russia;
e-mail: vostrik@isp.nsc.ru

Received 19 November 2001; revision received 20 March 2002
Kvantovaya Elektronika **32** (5) 416 – 420 (2002)
Translated by M.N. Sapozhnikov

$$\boldsymbol{E} = \boldsymbol{e}_q \frac{j_\perp}{\sigma}, \quad (2)$$

where $\boldsymbol{e}_q$ is the unit vector perpendicular to the grating planes and σ is the conductivity. Expression (2) is valid in most cases of the OP of samples because the relation $1/(\Delta k a) \ll 1$ is violated only when radiation is rather tightly focused into a spot of diameter of the order of several micrometres for small angles between the interacting beams (for example, in the case of OP in single-mode glass fibres [2, 5]).

The field induced in the medium changes the optical properties of the latter. Note that here we deal only with the reversible variations in the optical properties of the medium, which are caused by the spatial redistribution of the charge density, and do not consider the questions related to the rearrangement of the glass structure. The modulation of the refractive index appearing upon poling of the medium is

$$\Delta\hat{n} = \delta\hat{n} F(\boldsymbol{r}) \cos(\boldsymbol{qr} + 2\Delta\psi), \quad (3)$$

where $\boldsymbol{q} = 2\Delta\boldsymbol{k}$ is the wave vector of the refractive-index grating; $F(\boldsymbol{r})$ is the grating envelope; and the form of the tensor $\delta\hat{n}$ depends on the symmetry of the medium used. In glasses (isotropic media), the components of the tensor $\delta\hat{n}$ have the form [10]

$$\delta n_{ij} = \frac{2\pi\chi^{(3)}}{n} E_0^2 \delta_{ij} + \frac{4\pi\chi^{(3)}}{n} \delta_{ik}\delta_{jl} E_{0k} E_{0l}, \quad (4)$$

where n is the refractive index of the glass and E_0 is the field amplitude. One can see that the glass acquires anisotropy, which is typical of a uniaxial crystal with the optical axis directed along the photoinduced electric field. Two Δn gratings are formed from which diffraction of light occurs: the ordinary grating Δn_o and the extraordinary grating Δn_e with the reciprocal grating vector $\boldsymbol{q} = 2\Delta\boldsymbol{k}$. If the Δn grating is oriented at the Bragg angle to the incident radiation, the light is scattered in the direction determined by the condition of synchronism (Bragg condition), and the diffraction efficiency increases by many orders of magnitude. It is obvious that the conditions can be selected at which the first- and second-harmonic light beams involved in the grating formation will experience efficient diffraction (self-diffraction). Note that the polarisation dependence exists, and the extraordinary beam diffracts more efficiently than the ordinary beam.

Consider the self-diffraction of the fundamental radiation from the refractive-index grating (3). Let us choose the following interaction geometry (Fig. 1). We assume that the grating is produced by two crossing Gaussian beams of the fundamental ω and second-harmonic 2ω radiation. The incident beams are located on the x, y plane, the beam ω propagating along y axis while the beam 2ω propagating at an angle of α to the first beam. The origin of coordinates $(x, y, z) = (0, 0, 0)$ is located at the point of intersection of beam-waist centres, and w_1 and w_2 are the radii of the beams in the waist $[2(w_2)^2 \approx (w_1)^2]$. Because the interaction region is small, we can neglect the divergence of radiation over the grating size, and, assuming that the angle α of intersection of the beams is small, we set $\cos\alpha \approx 1$.

Let us introduce the vectors $\boldsymbol{k}_\perp = \{k_x, k_z\}$ and $\boldsymbol{r}_\perp = \{x, z\}$ directed perpendicular to the propagation of light. As a result, the grating envelope (3) can be approximated by the expression

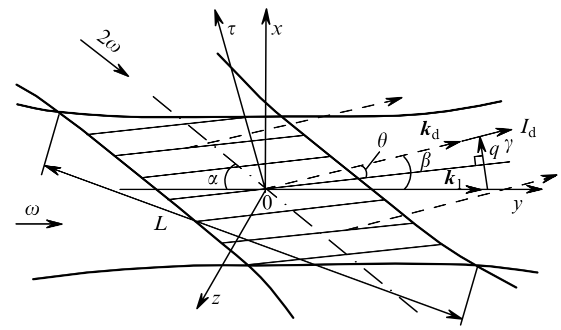

Figure 1. Formation of the Δn grating and self-diffraction of light.

$$F(\boldsymbol{r}) = \exp\left[-\frac{2}{w_1^2}\left(2\boldsymbol{r}_\perp^2 + 2xy\sin\alpha + y^2\sin^2\alpha\right)\right]. \quad (5)$$

We assume that the fundamental radiation

$$\boldsymbol{E}_{\rm in}(\boldsymbol{r}) = \boldsymbol{A}\left(-\frac{\boldsymbol{r}_\perp^2}{2w_1^2}\right)\exp({\rm i}k_1 y) \quad (6)$$

propagating along the y axis is incident on grating (3) and diffracts from it.

In the given field approximation, the wave field describing diffraction in the frequency representation has the form

$$k_1^2\boldsymbol{E} + \Delta\boldsymbol{E} = -\frac{2\left(\Delta\hat{n}^+ + \Delta\hat{n}^-\right)k_1^2}{n_1}\boldsymbol{E}_{\rm in}, \quad (7)$$

where $\Delta\hat{n}^\pm = \delta\hat{n}F(\boldsymbol{r})\exp[\pm{\rm i}(\boldsymbol{qr} + 2\Delta\psi)]$. One can see from (7) that two waves can exist in the system: $\boldsymbol{E} = \boldsymbol{E}_{\rm d}^+ + \boldsymbol{E}_{\rm d}^-$, where $\boldsymbol{E}_{\rm d}^+$ and $\boldsymbol{E}_{\rm d}^-$ are the diffracted waves with the wave vectors $\boldsymbol{k}^+ = \boldsymbol{k}_1 + \boldsymbol{q}$ and $\boldsymbol{k}^- = \boldsymbol{k}_1 - \boldsymbol{q}$; and $\boldsymbol{E}_{\rm d}^+(\boldsymbol{q}) = \boldsymbol{E}_{\rm d}^-(-\boldsymbol{q})$.

Let us derive the expression for $\boldsymbol{E}_{\rm d}^+$ (we omit the sign '+' below). The expression for the wave $\boldsymbol{E}_{\rm d}^-$ can be obtained by the replacement $\boldsymbol{q} \to -\boldsymbol{q}$ in the final expression for $\boldsymbol{E}_{\rm d}^+$. We seek the solution of equation (7) by the Fourier transform method. As a result, the expression for the amplitude of the scattered field has the form

$$\boldsymbol{E}_{\rm d}(\boldsymbol{r}) = \frac{k_1^2}{4\pi^3 n_1}\int_{-\infty}^{+\infty}\frac{\Delta\hat{n}(\boldsymbol{r}')\boldsymbol{E}_{\rm in}(\boldsymbol{r}')}{k_y^2 + k_\perp^2 - k_1^2}\exp[{\rm i}k_y(y - y')]$$
$$\times\exp[{\rm i}\boldsymbol{k}_\perp(\boldsymbol{r}_\perp - \boldsymbol{r}'_\perp)]{\rm d}k_y{\rm d}\boldsymbol{k}_\perp{\rm d}\boldsymbol{r}'. \quad (8)$$

Because $q/k_1 \ll 1$, upon scattering the fundamental radiation weakly deviates from its initial direction of propagation (i.e., $k_\perp^2 \ll k_1^2$). Taking this into account, we perform integration over the region of the complex variable k_y. Because the waves with $\boldsymbol{k}_\perp \approx \boldsymbol{q}_\perp$ make the main contribution to the scattered field, we perform the change of variables $\boldsymbol{s}_\perp = \boldsymbol{k}_\perp - \boldsymbol{q}_\perp$ and pass to the integration over ${\rm d}\boldsymbol{s}_\perp$. Because the value of $\boldsymbol{s}_\perp$ is small, the higher-order terms in $\boldsymbol{s}_\perp$ in exponentials of the integrand can be neglected.

By omitting calculations, we present the final expression for the amplitude of the diffracted wave on the right away from the inhomogeneity region

$$\boldsymbol{E}_{\rm d} = \frac{\boldsymbol{D}\pi\exp\left[-d_1\left(q_y + q_x^2/2k_1\right)^2\right]}{(f_1 f_2)^{1/2}\left[\left(1 + D_1^2\right)\left(1 + D_2^2\right)\right]^{1/4}}\exp({\rm i}k_1\gamma)\times$$

$$\times \exp\left[-\frac{z^2}{4f_1(1+D_1^2)}-\frac{\tau^2}{4f_2(1+D_2^2)}\right]$$

$$\times \exp\left[-\frac{\tau(\gamma-\tau q_x/k_1)(f_3/f_2)+f_3^2k_1^2}{4f_2(1+D_2^2)}\right]$$

$$\times \exp\left[\frac{\mathrm{i}z^2D_1}{4f_1(1+D_1^2)}+\frac{\mathrm{i}\tau^2D_2}{(1+D_2^2)}\right] \qquad (9)$$

$$\times \exp\left[-\mathrm{i}k_1\left(\frac{\tau f_3/(2f_2)+(\gamma-\tau q_x/k_1)f_3^2/(8f_2^2)}{1+D_2^2}\right)\right]$$

$$\times \exp\left[-\mathrm{i}\arctan\left(\frac{\left[(1+D_1^2)(1+D_2^2)\right]^{1/2}+D_1D_2-1}{D_1+D_2}\right)\right],$$

where

$$\boldsymbol{D}=\frac{(\delta\hat{n}\boldsymbol{A})\mathrm{i}k_1w_1^3\exp(\mathrm{i}2\Delta\psi)}{2(90\pi)^{1/2}n_1\sin\alpha};\quad d_1=\frac{9w_1^2}{40\sin^2\alpha};\quad d_2=\frac{4}{9}\sin\alpha;$$

$$D_1=\frac{y}{2k_1f_1};\quad D_2=\frac{y}{2k_1f_2};\quad f_1=\frac{w_1^2}{18}+\frac{d_1}{k_1}\left(q_y+\frac{q_x^2}{2k_1}\right); \qquad (10)$$

$$f_2=f_1+d_1\left(d_2+\frac{q_x}{k_1}\right)^2;\quad f_3=\frac{2d_1}{k_1}\left(d_2+\frac{q_x}{k_1}\right)\left(q_y+\frac{q_x^2}{2k_1}\right);$$

$$\tau\approx x\left(1-\frac{q_x2}{2k_1^2}\right)-y\frac{q_x}{k_1};\quad \gamma=y\left(1-\frac{q_x^2}{2k_1^2}-\frac{q_x^4}{8k_1^4}\right)+x\frac{q_x}{k_1}.$$

Note that the self-diffraction of the second-harmonic radiation involved in the grating formation can be described similarly. In this case, we obtain the expression for the amplitude of the diffracted second-harmonic radiation, which is similar to expression (9), with the corresponding replacements.

Let us analyse expression (9). One can see from (9) that the wave amplitude is maximal when

$$q_y=-\frac{q_x^2}{2k_1}. \qquad (11)$$

In this case, $f_3=0$, the wave propagates in the direction γ and has characteristic Gaussian shapes in the directions z and τ. The angle β at which the wave deflects from the initial direction of propagation y upon scattering can be determined from the expression $\sin\beta=q_x/k_1\approx q/k_1$ because $k_\mathrm{d}=k_1$ and $q\ll k_1$ (see Fig. 1). As a result, when the condition (11) is satisfied, we obtain the usual Bragg scattering from a phase diffraction grating, which is well known in acoustooptics [11] and is described by the expression

$$\sin\theta\approx\frac{q}{2k_1}, \qquad (12)$$

where $\theta=\beta/2$ is the Bragg angle.

Because q depends on α, the condition (11) can be satisfied only for a certain angle between the crossed beams. The exact optimal angles between the crossing beams can be obtained by considering the interaction geometry. In this case, there are two possibilities. When $\eta=n_1/n_2<1$ (i.e., for the normal dispersion of the waves in the medium), the Bragg self-diffraction of the fundamental radiation is observed at the optimal angle between the beams $\alpha_\mathrm{m}=\arccos[(3\eta^2+2)/5\eta]$, while for the second-harmonic radiation, it is observed for $\alpha_\mathrm{m}=\arccos\eta$. Of more interest is the case $1<\eta<2$ (anomalous dispersion). The self-diffraction of fundamental and second-harmonic light beams for these values of η occurs simultaneously at the same optimal angle $\alpha_\mathrm{m}=\arccos[(\eta^2+2)/3\eta]$ between the beams forming the grating. Note that the angle between the crossing beams should be set sufficiently precisely in the experiment because the wave amplitude $[\sim\exp(-60k_1^2\times w_1^2\sin^2\delta\alpha)]$ sharply decreases upon deviations from the optimal angle ($\alpha_m\pm\delta\alpha$).

The diffracted wave diverges and has different curvatures of the phase front along the directions z and τ. The quantities D_1 and D_2 play the role of characteristic dimensionless diffraction lengths. Figure 2 shows the evolution of the beam radii $w_\tau=[2f_2(1+D_2^2)]^{1/2}$ and $w_z=[2f_1(1+D_1^2)]^{1/2}$ during the propagation of diffracted radiation along the directions τ and z, respectively, for the values of $w_1\approx$ 130 μm, $\alpha_\mathrm{m}\sim3.5°$, and $\lambda_1=1.08$ μm chosen by us in the experiment. Because the vector $\boldsymbol{q}$ lies in the plane (τ, γ), diffraction in this plane occurs over the entire length of the grating (see Fig. 1), and the effective width of the beam becomes the largest one (it exceeds the width of the incident beam by a factor of 2.4). In the plane (z,γ), diffraction occurs from the grating aperture. The grating aperture is determined by the convolution of the transverse distributions of the fundamental and second-harmonic beams, so that the width of the diffracted beam in this plane is minimal. The diffracted wave strongly diverges along the direction z. The beam, which was initially narrow along z, becomes symmetric at the point $y=y_0\approx12$ cm and has the same width along the directions τ and z, whereas for $y=20$ cm, the beam becomes rather strongly broadened along z.

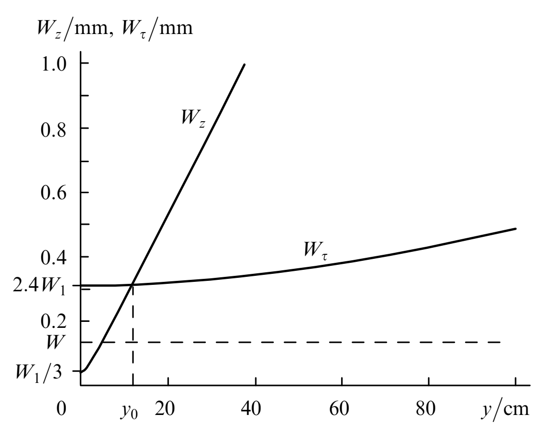

Figure 2. Evolution of the beam radii along directions z and τ during the propagation of diffracted radiation.

The diffraction efficiency, i.e., the integral ratio of the total energy flux of diffracted radiation

$$P_\mathrm{d}=\frac{c}{8\pi}\int\boldsymbol{E}_\mathrm{d}\boldsymbol{E}_\mathrm{d}^*\mathrm{d}\tau\mathrm{d}z$$

to the energy flux of the incident radiation

$$P_\mathrm{in}-\frac{c}{8\pi}\int\boldsymbol{E}_\mathrm{in}\boldsymbol{E}_\mathrm{in}^*\mathrm{d}x\mathrm{d}z$$

can be determined from the expression

$$\eta_{\rm d} = \frac{P_{\rm d}}{P_{\rm in}} = \frac{(\delta\hat{n}\boldsymbol{A})^2}{n_1^2 A^2}\,\frac{\sqrt{5}\pi k_1^2 w_1^2}{160\sin^2\alpha}\exp\left[-2d_1\left(q_y + \frac{q_x^2}{2k_1}\right)^2\right]. \quad (13)$$

Note finally that expression (9) was obtained for the wave $\boldsymbol{E}_{\rm d}^{+}$. The total solution contains two components: $\boldsymbol{E} = \boldsymbol{E}_{\rm d}^{+} + \boldsymbol{E}_{\rm d}^{-}$. To obtain $\boldsymbol{E}_{\rm d}^{-}$, as mentioned above, it is necessary to make the change of variables $\boldsymbol{q} \to -\boldsymbol{q}$. However, the phase matching condition cannot be satisfied for this wave in the case of normal dispersion in the medium. As a result, only one diffracted wave can propagate in the medium: the wave $\boldsymbol{E}_{\rm d}^{+}$ in the medium with normal dispersion and the wave $\boldsymbol{E}_{\rm d}^{-}$ in the medium with anomalous dispersion.

3. Experimental results and discussion

We have chosen for experiments the K8 and F4 oxide glasses from an available set of optical glasses, in which the diffraction of light was most efficient. The optical poling of samples was performed by fundamental and second-harmonic radiation from a 1.079-μm pulsed $\mathrm{Nd^{3+}:YAG}$ laser with a pulse energy of ~ 18 mJ, a pulse duration of ~ 15 ns, and a pulse repetition rate of 12.5 Hz. The second-harmonic conversion efficiency was ~ 10 %. The focused beams were crossed at an angle of α. The fundamental and second-harmonic light beams were linearly polarised, and their polarisation could be varied independently. The phase difference between these two fields could be also varied continuously.

The maximum peak intensity P_ω of the fundamental radiation in the focus was ~ 10^9 W cm^{-2} for the laser beam diameter of ~ 260 μm. The diffracted radiation was detected in the far-field zone with a photomultiplier and the data were processed with a PC. The threshold sensitivity of the detecting system was 1 μW pulse^{-1}. The refractive-index gratings written upon OP of glasses were stable and persisted for 4 hours in the absence of an external action. The process of grating recording was completely reversible, and no structural variations were observed in the samples under study. The kinetics of recording and relaxation of the Δn gratings is described in detail in paper [12].

Figure 3 presents the typical angular distributions of the intensity of fundamental radiation transmitted by a sample subjected to OP. The angular distributions were obtained by scanning with a slit of width ~ 60 μm with a step of ~ 10$'$ in the plane of crossing of the beams within the sector ±30° relative to the direction of propagation of the incident fundamental radiation.

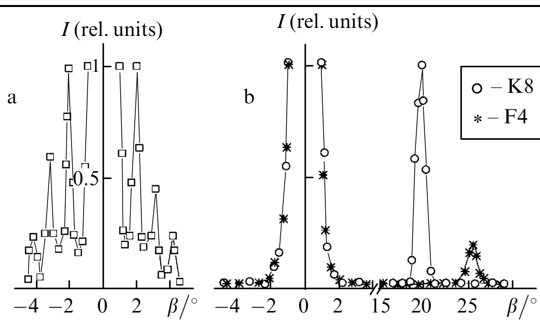

Figure 3. Angular distributions of the intensity of fundamental radiation (in the convergence plane of the beams) transmitted through a sample subjected to OP for $\alpha = 0$ (a) and $\alpha = \alpha_{\rm m}$ (b); the peak at $\beta = 0$ is the transmitted radiation, and the side peaks are diffracted radiation.

Figure 3a corresponds to the coaxial propagation of the beams. In this case, the planes of the produced grating are perpendicular to the direction of propagation of the beams, and the typical aperture diffraction by the inhomogeneity of the refractive index occurs. The curves in Fig. 3b were obtained for the angles between crossing beams equal to $\alpha = \alpha_{\rm m} = 3.19$ and 3.87°, which correspond to the Bragg diffraction of the fundamental radiation by produced Δn gratings in K8 and F4 glasses, respectively. One can see that in this case, diffraction of light occurs in the direction determined by the Bragg condition. In the case of Bragg diffraction, the intensity of the scattered beam was 300 times greater than in the case of 'aperture' diffraction.

The optimal angles between the crossing beams and Bragg angles obtained experimentally for K8 and F4 glasses were in agreement with the calculated values. According to expression (13), the maximum diffraction efficiency $\eta_{\rm d} \approx 10^{-3}$ achieved in our experiments corresponds to the relative change in the refractive index $\delta n/n_1 \approx 10^{-5}$. For oxide glass, $\chi^{(3)} \sim 10^{-18}$ cm^2 V^{-2} [13]. By substituting these values to expression (4), we obtain the estimate of the photoinduced electric field $E_0 = \left[n_1 \delta n\,(6\pi\chi^{(3)})^{-1}\right]^{1/2} \approx 10^6$ V cm^{-1}, which agrees, by the order of magnitude, with the estimates made in the studies of the photoinduced SHG upon the OP of oxide glasses [5, 6].

When gratings were written by light beams polarised in the plane of their crossing and tested by light corresponding to the extraordinary beam in the grating, the efficiency of Bragg diffraction $\eta_{\rm d} = P_{\rm d}/P_{\rm in}$ was maximal and equal to 1.2×10^{-3} and 4×10^{-5} for K8 and F4 glasses, respectively. When gratings were written by light beams with orthogonal polarisations (under the condition that the CP current vector lies in the plane of crossing of the beams), the diffraction efficiency of the extraordinary beam was approximately a hundred times lower than the maximum diffraction efficiency. The diffraction efficiency of the ordinary beam was in both cases approximately ten times lower than that of the extraordinary beam. Finally, if the grating was written by light beams polarised perpendicular to the plane of their crossing, then, independently of the polarisation of the testing light, the diffraction efficiency was low (~ 10^{-4} of the maximum efficiency).

This is explained by the fact that in this case the CP current flows parallel to the grating planes, the charges are accumulated at the periphery of the interaction region of light beams, and no diffraction grating appears. The results of the polarisation studies are in agreement with the model of formation of the refractive-index grating in glass due to the CP effect and confirm the validity of expression (1) and the applicability of the symmetry ($\infty\infty m$) of isotropic centrally symmetric media for the description of glasses under study. The fact that the coherent current in glasses is determined by two CP constants [expression (1)] was confirmed in the study of polarisation properties of photo-induced SHG [14].

Experiments with the variation of the phase difference of light beams gave the following results. A change in the phase difference of the fields on passing from one writing cycle to another changes neither the intensity nor the diffraction pattern. However, if the phase difference is modulated during the grating writing with frequency ~ 1 Hz, the diffraction efficiency decreases to zero. When the phase difference is changed very slowly ($\sim 10^{-3}$ Hz), the intensity of diffracted light becomes periodically modulated. This agrees with the conclusions that follow from expressions (1) and (3). Indeed, the diffraction efficiency is the same for any constant phase differences $\Delta\psi$ because a change in $\Delta\psi$ results only in the phase shift of the diffraction grating in space. If the phase shift $\Delta\psi$ changes in time with a period that is much smaller than the time of accumulation of a separated charge, the diffraction grating does not appear.

Studies of the spatial distribution of the diffracted beam showed that a wave weakly diverges in the convergence plane of the beams and strongly diverges in the perpendicular direction. Figure 4 shows the experimental relative diffraction efficiencies $\xi_z = \eta_d(z_0)/\eta_d^m$ and $\xi_\tau = \eta_d(\tau_0)/\eta_d^m$ as functions of the widths z_0 and τ_0 of slits in front of the photomultiplier along these directions [$\eta_d(z_0)$ and $\eta_d(\tau_0)$ are the diffraction efficiencies for slit widths in front of the photomultiplier along the direction z and τ, respectively; η_d^m is the diffraction efficiency in the absence of a slit in front of the photomultiplier; the distance from the crossing centre of the beams to the photomultiplier is $y \approx 2.5$ cm]. Figure 4 also shows the relative diffraction efficiencies calculated by the expressions

$$\xi_z = \frac{1}{\left[2\pi f_1\left(1+D_1^2\right)\right]^{1/2}} \int_{-z_0/2}^{z_0/2} \exp\left[\frac{-z^2}{2f_1\left(1+D_1^2\right)}\right] dz\,,$$
$$\xi_\tau = \frac{1}{\left[2\pi f_2\left(1+D_2^2\right)\right]^{1/2}} \int_{-\tau_0/2}^{\tau_0/2} \exp\left[\frac{-\tau^2}{2f_2\left(1+D_2^2\right)}\right] d\tau\,, \qquad (14)$$

which were obtained using expression (9). The spatial distribution of diffracted radiation at the point $y \approx 2.5$ cm corresponding to the experimental data is shown in Fig. 5. One can see that the effective width of this distribution in the direction τ is larger than that in the direction z. The character of the beam distortion and the parameters of diffracted radiation are in agreement with the experimental data.

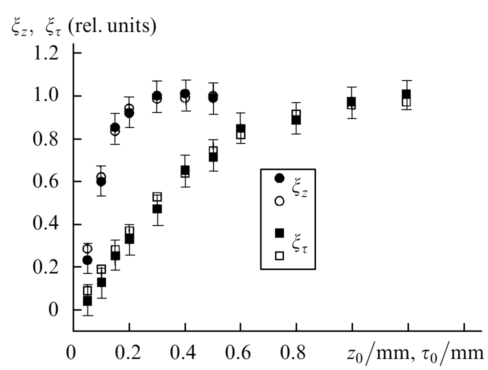

Figure 4. Relative efficiencies of light diffraction as functions of the width of slits in front of the photomultiplier (dark and light points are experiment and theory, respectively).

4. Conclusions

We studied theoretically and experimentally the self-diffraction of light upon the OP of glasses and obtained the expression for the amplitude of diffracted radiation and analysed its properties. The phase, polarisation, and angular dependences of diffraction in K8 and F4 oxide glasses, as well as the spatial distribution of diffracted light well agree with the theoretical dependences obtained from the model of OP of glasses caused by the coherent photogalvanic effect. The diffraction of light considered in the paper can be used for studying physical mechanisms of purely optical poling of media.

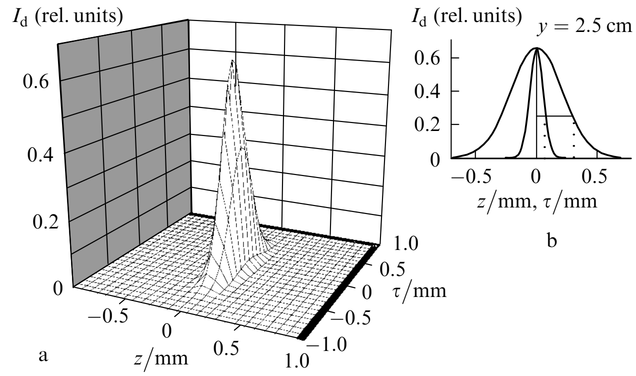

Figure 5. Spatial distribution of diffracted radiation at the point $y = 2$ cm (a) and its cross sections along directions τ and z at the beam centre (b).

Acknowledgements. This work was supported by the State Programme 'Optics. Laser Physics' of the Ministry of Science of the Russian Federation and by the Russian Foundation for Basic Research and Chinese State Foundation for Natural Sciences (Grant No. 99-02-39044).

References

1. Dianov E.M., Kazanskii P.G., Stepanov D.Yu. *Kvantovaya Elektron.*, **17**, 926 (1990) [*Sov. J. Quantum Electron.*, **20**, 849 (1990)].
2. Antonyuk B.P. *Opt. Commun.*, **174**, 427 (2000).
3. Balakirev M.K., Smirnov V.A. *Pis'ma Zh. Eksp. Teor. Fiz.*, **61**, 544 (1995).
4. Balakirev M.K., Smirnov V.A., Vostrikova L.I. *Opt. Commun.*, **178**, 181 (2000).
5. Dianov E.M., Starodubov D.S. *Kvantovaya Elektron.*, **22**, 419 (1995) [*Quantum Electron.*, **25**, 395 (1995)].
6. Antonyuk B.P., Antonyuk V.B. *Usp. Fiz. Nauk*, **171**, 61 (2001).
7. Balakirev M.K., Vostrikova L.I., Smirnov V.A., Entin M.V. *Pis'ma Zh. Eksp. Teor. Fiz.*, **63**, 166 (1996).
8. Baskin E.M., Entin M.V. *Pis'ma Zh. Eksp. Teor. Fiz.*, **48**, 601 (1988).
9. Sokolov V.O., Sulimov V.B. *Phys. Stat. Sol. (b)*, **187**, 189 (1995).
10. Landau L.D., Lifshits E.M. *Elektrodinamika sploshnykh sred* (Electrodynamics of Continuous Media) (Moscow: Nauka, 1992).
11. Tucker J.W., Rampton V.W. *Microwave Ultrasonics in Solid State Physics* (Amsterdam: NorthHolland, 1972; Moscow: Mir, 1975).
12. Balakirev M.K. *Phys. Vibrations*, **6**, 233 (1998).
13. Mizrahi V., Österberg U., Sipe J.E., Stegeman G.I. *Opt. Lett.*, **13**, 279 (1988).
14. Bolshtyansky M.A., Kapitzky Yu.E, Zel'dovich B.Ya., et. al. *Pure Appl. Opt.*, **1**, 289 (1992).