Wave-front reversal on a semiconductor surface during plasma reflection

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The reversal of a wave front on the surface of a semiconductor is investigated theoretically and experimentally. The inhomogeneous reflectance was recorded by using plasma reflection. The experiments were performed on surfaces of Ge, GeAs, InSb, and Si at wavelengths 1.06 and 0.69 μm. The energy of the reversed wave was several per cent of that of the incident. The experimental results are in reasonable agreement with the theory.

Wave-front reversal by a surface (WFR-S)[1] can be realized on the interface of two media if the amplitude reflection coefficient r of the interface depends on the local intensity J of the incident radiation: $r = r(J)$. Let the incident radiation consist of a reversible signal wave E_s and of a plane reference wave E_0: the surface reflection coefficient will then be modulated in accord with the relation

$$r=r_0+\frac{\partial r}{\partial J}(E_0E_s^*+E_0^*E_s),$$

i.e., a reflective dynamic hologram will be recorded, Fig. 1. When the reference wave is oriented normal to the surface, the conjugate image $E_0(\partial r/\partial J)E_0E_s^*$ reconstructed by the reference wave in real time will propagate counter to the signal and form an inverted wave $E_{\rm rev} = [(\partial r/\partial J)E_0^2]E_s^*$ with efficiency $|E_{\rm rev}/E_s|^2 = |(\partial r/\partial J)J_0|^2$, where $J_0 = |E_0|^2$ is the intensity of the reference wave.

WFR-S has been realized with various mechanisms in which the radiation modulates the reflection coefficient of the nonlinear surface.[2–8] An advantage of this method is the possibility of working with opaque nonlinear media, meaning in a wide spectral range, and with radiation of various nature.[5] What is still urgent, however, is a search for surfaces with a sufficiently strong $r(J)$ dependence and with reversible rapidly relaxing nonlinearity. Promise is offered from this viewpoint by the use of dynamic holograms on the surfaces of a number of semiconductors.[6,7,9–11]

We note that a thermal mechanism is resorted to in Refs. 6 and 9–11 to explain phenomenologically how holograms are written on a semiconductor surface. Very simple estimates show that to produce surface holograms with the experimentally recorded diffractive efficiencies $\sim 10^{-3}$ it is necessary to have a holographic grating with relative variations on the order of 10%, from line to line, of the reflection coefficient, meaning also of the dielectric constant of the substance in the subsurface layer. Neither thermal expansion nor equilibrium thermal overpopulation of the conduction band can ensure a reversible optical-density change of this scale in the semiconductor on such a scale within a time $\sim 10^{-8}$ sec with the pulses used in these studies. In addition, the qualitative form of the experimentally obtained dependences of the grating efficiencies on a number of parameters[6,11] does not agree with the predictions of the thermal model.[6,9–11]

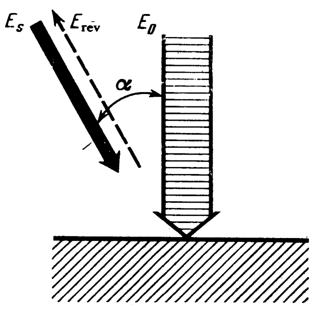

FIG. 1. Scheme of the method of WFR by a surface. Interference picture of plane reference wave E_0 normal to the surface and of the reversed signal E_s is recorded in the form of a reflection-coefficient grating. One of the diffraction maxima of the grating yields in real time the reversed wave $E_{\rm rev} \propto E_s^*$ when read by the reference wave.

It was shown in Ref. 12 that direct photoexcitation of an electron-hole plasma in the subsurface layer of a semiconductor with band gap E_g that is less than the quantum energy $\hbar\omega$ of the incident laser radiation makes the main contribution to the change of the surface reflection coefficient. It is therefore reasonable to propose the following mechanism of recording dynamic surface holograms. When the subsurface plasma is excited by radiation with intensity distribution that is not uniform along the surface, the plasma density is likewise nonuniform and duplicates the interference pattern of the writing field. The plasma alters the dielectric constant of the surface layer, meaning also the phase and amplitude of the reflected wave. As a result, the surface reflection coefficient acquires a spatial modulation corresponding to the intensity profile, i.e., a reflective surface hologram is recorded. We present first some results of model calculations of this mechanism, and then discuss in their light the experimental results.

1. EFFECT OF SUBSURFACE PLASMA ON THE AMPLITUDE AND PHASE OF THE REFLECTED WAVE

If the incident-radiation photon $\hbar\omega$ exceeds noticeably the band gap E_g of the semiconductor, the light-absorption coefficient is large, $\alpha \sim 10^4$–10^5 cm^{-1}. We shall be interested in laser pulses with the following characteristic parameters: energy density $W \sim 0.2$ J/cm^2 and duration $\tau \sim 2\times 10^{-8}$ sec. The number of electron-hole pairs created in a subsurface layer of thickness α^{-1} corresponds to the number of absorbed phonons, equivalent in the visible band to approximately 10^{18} photoelectrons per unit surface area. With allowance for the diffusion of the electrons and holes from the surface, the excited zone is localized in a layer of thick-

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ness $\Delta z \sim (D\tau)^{1/2} \sim 10^{-4} - 10^{-3}$ cm, where D is the ambipolar diffusion coefficient; this corresponds to an excited pair density $N_e \sim 10^{21}$–10^{22} cm^{-3} (Ref. 12). The density of the subsurface plasma at the end of the pulse will apparently be determined by recombination processes that offset the direct photoexcitation.

The most substantial conclusion of this reasoning is a high value of the subsurface-electron density in the conduction band. For an effective mass of the plasma particles $m^* = 0.03 m_e$, the plasma frequency ω_p reaches the neodymium-laser emission frequency $\omega \approx 1.8 \times 10^{15}$ sec^{-1} at an electron-hole plasma density $N_e \approx 5 \times 10^{19}$ cm^{-3}. This means that in the indicated range of the parameters of the exciting beam the surface-reflection coefficient has a rather abrupt resonant dependence on the electron density, meaning also on the light intensity, near $\omega_p \approx \varepsilon_0{}^{1/2}\omega$—the so-called plasma resonance.[13–16]

We calculate now the efficiency of recording a holographic grating and of the WFR via the plasma-reflection mechanism. The dielectric constant ε of the excited semiconductor is made up of the contributions of the bound electrons and delocalized electrons and hole. We assume the $\varepsilon(N_e)$ dependence in the Drude-Lorentz model (see, e.g., Refs. 14 and 17) in the following form:

$$\varepsilon = \varepsilon_0 - g\left(\frac{\omega_p}{\omega}\right)^2 + i\frac{\nu_{\rm cr}}{\omega}\left(\frac{\omega_p}{\omega}\right)^2, \quad \omega_p = (4\pi e^2 N_e/\langle m^*\rangle)^{1/2}, \qquad (1)$$

where ω is the frequency of the exciting light, ω_p is the plasma frequency determined by the effective mass $1/m^* = (1/m_e^* + 1/m_h^*)$, averaged over the bands, where m_e^* and M_h^* are the effective masses of the conduction electrons and holes in the valence band and ε_0 is the dielectric constant of the unexcited sample. The factor $g \gtrsim 1$ takes into account the change of the contribution of the bound electron as their number is increased. The last term in (1) describes the light absorption by the free carriers, and $\nu_{\rm col}$ is the inelastic-collision frequency. In the optical frequency band $\nu_{\rm col}/\omega \ll 1$, nonetheless at high excitation energy $N_e \sim 10^{21}$ cm^{-3} this form of absorption can play a noticeable role, since $\nu_{\rm col}\omega_p^2/\omega^3 \sim 1$ (Refs. 16 and 18).

With increasing dimensionless parameter $a = g(\omega_p/\omega)^2 = N_e N_{\rm cr}$ that describes the plasma density and is equal to the ratio of N_e to the density $N_{\rm cr}$ at which $\omega_p \approx \omega$, the refractive index n of the subsurface layer decreases and the absorption coefficient $\varkappa$ increases. The modulus of the dielectric constant $\varepsilon = (n + i\varkappa)^2$ goes through a minimum $|\varepsilon_{min}| \approx 2n_0\varkappa_0 + (\nu_{\rm col}/\omega)(\omega_p/\omega)^2$ at $a \approx n_0^2 - \varkappa_0^2$, where n_0 and $\varkappa_0$ are the refractive index and the absorption coefficient of the unexcited semiconductor. Their characteristic values are in the ranges $n_0 \approx 3$–5, $\varkappa_0 \approx 0.1$–1. At the same point, the values calculated with the formulas[17] $n = [(|\varepsilon| + {\rm Re}\varepsilon)/2]^{1/2}$ and $\varkappa = [(|\varepsilon| - {\rm Re}\varepsilon)/2]^{1/2}$ become approximately equal: $n \approx \varkappa \approx [n_0\varkappa_0 + 2\nu_{\rm col} a/\omega]^{1/2}$.

Since the plasma-density inhomogeneity scales are larger in our case than the incident-light wavelength, we calculate the amplitude coefficient r of the surface reflection using the known formula[19] $r = (n - 1 + i\varkappa)/(n + 1 + i\varkappa)$. The WRF efficiency is determined mainly by the derivative of the function $r(J)$ (Ref. 1). We obtain first the derivative of the reflection coefficient with respect to the parameter a. Neglecting the contribution made to the derivative by the weak dependence of the term $\propto \nu_{\rm col}/\omega$ we obtain

$$\frac{\partial r}{\partial a} \approx \frac{i\varkappa - n}{|\varepsilon|(n+1+i\varkappa)^2}. \qquad (2)$$

The presence of an imaginary part of $\partial r/\partial a$ indicates that the subsurface plasma alters not only the amplitude but also the phase of the reflected wave. By the same token, the holographic grating on the surface is of the amplitude-phase type.

2. CALCULATION OF THE WFR EFFICIENCY

Having relation (2) for the derivative $\partial r/\partial a$, we can obtain the sought dependence of the WFR efficiency $\eta = |E_{\rm rev}/E_s|^2$ on the incident-light intensity from the relation

$$\eta = \left|\frac{\partial r}{\partial J} J_0\right|^2 = \left|\frac{\partial r}{\partial a}\frac{\partial a}{\partial J} J_0\right|^2. \qquad (3)$$

This, however, calls for a determination of the dependence of the subsurface-plasma density on the local intensity of the incident radiation.

During the initial stage of the laser pulse, at $N_e \ll 10^{20}$ cm^{-3}, when the rate of the recombination processes is noticeably lower than the rate of photoelectron excitation, electron-hole pairs accumulate in the subsurface layer. During this stage, an important role is played by the diffusion of the photoelectrons. Therefore, on the one hand, the plasma density increases with time more slowly than linearly; when the pulse is instantaneously turned on, $N_e \sim 2J_0(t/\pi D)^{1/2}$ (Ref. 12). On the other hand, recording of the holographic lattice is nonlocal and becomes washed out after a time $\Delta t \sim \Lambda^2/(4\pi^2 D) \lesssim 10^{-9}$ sec, where Λ is the grating period. During the accumulation stage, furthermore, the efficiency of recording the dynamic grating is negligibly small because of the large distance from the plasma resonance: $a \lesssim 1$.

With increasing carrier density, the role of the recombination processes decreases. Near the plasma resonance of interest to us $a \sim n_0^2 \sim 10$–$25 \gg 1$, which corresponds to $N_e \sim 10^{21}$–10^{22} cm^{-3}, impact ionization is apparently the dominant process that balances the carrier excitation. It was proposed in Ref. 12 to calculate $N_e(J)$ from the simple electron balance equation in the conduction band, $\alpha J = \gamma N_e^3$, where α is the coefficient of interband absorption and γ is the impact recombination constant. At such high densities, however, the plasma can no longer be regarded as ideal and the impact recombination cannot be calculated from the three-particle-collision probability. Thus, in the Langevin theory for a dense plasma[20] the recombination probability is proportional to N_e^2. Moreover, in the recombination regime a noticeable role can be assumed by impact-ionization, by absorption saturation, and by heating of the electron gas via absorption of light by the free carriers. Unfortunately, physical processes in an electron-hole plasma of so high a density, which is furthermore acted upon by a high-power laser, remain practically uninvestigated. Their study is of independent interest,[21] but is outside the scope of the present paper. We assume therefore that in the recombination regime

$\alpha J \approx \gamma N_e^p$, where p is an unknown quantity in the range $1 \lesssim p \lesssim 3$. We shall attempt to determine p from experimental results. Substituting the relation $N_e = (\alpha J/\gamma)^{1/p}$ in (3) we readily obtain

$$\eta = a_0^2 [p^2 |\varepsilon| (1+2n+|\varepsilon|)^2]^{-1}, \tag{4}$$

where $a_0 = g(\omega_p/\omega)^2$ corresponds to the plasma-density component homogeneous along the surface.

As expected, the dependence (4) of the reversal efficiency on the normalized electron density a is resonant. At $a \gtrsim 1$ the WFR efficiency is low: $\eta \approx a^2/p^2 n_0^6 \ll 1$. As the plasma resonance is approached, the efficiency increases abruptly, almost stepwise. The maximum value η_{max} is reached near the resonance point, where the real part of the dielectric constant vanishes, $a \approx n_0^2 - \varkappa_0^2$:

$$\eta_{max} \approx \frac{n_0}{8p^2\varkappa^3}\left[1 + \frac{1+2(n_0\varkappa_0)^{1/2}}{2n_0\varkappa}\right]^{-1}. \tag{5}$$

Here n_0 is the refractive index of the excited semiconductor, $\varkappa$ is the absorption coefficient near the resonance point with account taken of the contribution of the free carriers. In the Drude-Lonrentz model (1) we have $\varkappa = \varkappa_0 + \frac{1}{2} n_0 \nu_{col}/\omega$. For $p = 2$ and $\varkappa \times 1$ the maximum WFR efficiency is $\sim 10\%$. With further increase of the plasma density the efficiency of the grating recording decreases. It must be emphasized that the maximum efficiency and the width of the resonance itself depend strongly on the real absorption $\varkappa$ of the subsurface layer.

3. EXPERIMENTAL INVESTIGATION OF WFR EFFICIENCY

The experimental setup is shown in Fig. 2. A fraction ($\approx 3.4\%$) of single-mode emission of a ruby laser L operating in the single giant pulse regime ($\lambda = 0.6943$ μm, $\tau = 20 \times 10^{-9}$ sec) was diverted by wedge W_2 through diaphragm DP (1.5 mm diameter) to form the signal beam E_s. The radiation passing through the wedge W_2 was used as the reference wave E_0. The signal beam was directed by semitransparent mirror M_1 to the center of the reference-beam cross section on the surface of the employed semiconductor T. The reference wave was incident on the semiconductor surface practically normally. The beam-convergence angle was $\sim 2°$. To investigate the quality of the WFR, a phase plate PP was placed in the path of the signal beam. The ratio of the intensities of the reference and signal waves was maintained constant at $J_0/J_s \approx 60$. The coefficient η of reflection of the signal beam was determined from the readings of calibrated energy measuring devices CM_1 and CM_2. To compare the angle spectra of the signal and reversed beams, the intensity distribution at the focus of a spherical mirror M_2 ($f = 2300$ mm) was photographed through a stepped attenuator (SA) on photographic film F, followed by photometry. The time variation of the signal and reversed beams were monitored on an S7-10B high-speed oscilloscope in which an FK-19 photocell was used as the receiver. The WFR efficiency on the semiconductor surface was investigated by us also at neodymium-laser frequency ($\lambda = 1.06$ μm; $\tau \times 30 \times 10^{-9}$ sec). The experimental setup was similar. We investigated flat polished surfaces of InSb, GaAs, Ge, and Si.

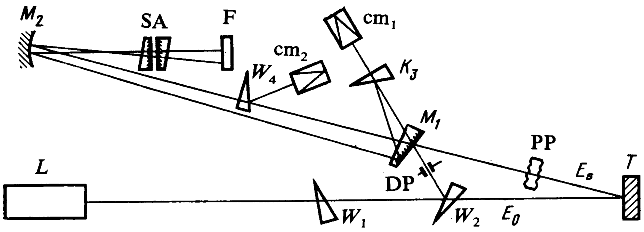

FIG. 2. Experimental setup.

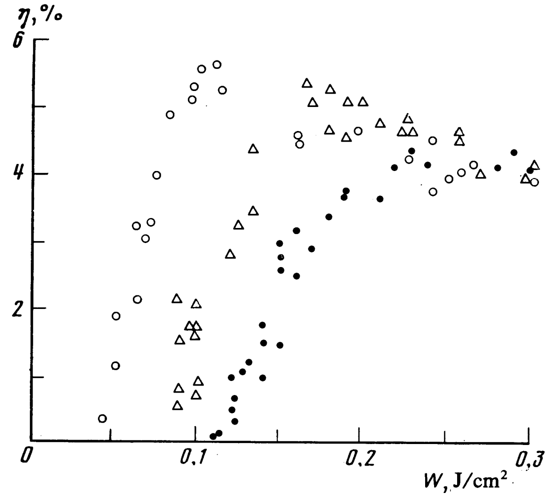

FIG. 3. Experimental dependences of the WFR efficiency η on the reference-wave energy density W when surfaces of InSb (○), GaAs (△) and Ge (●) are illuminated by a ruby laser; $\lambda = 0.6943$ μm.

Figures 3 and 4 show the experimental values of the energy coefficient $\eta(W)$ of reflection into a reversed wave of ruby and neodymium laser radiation, for surfaces of a number of semiconductors, as functions of the energy density W of the reference wave. The experimental curves have a clearly pronounced threshold due to the need to come close to

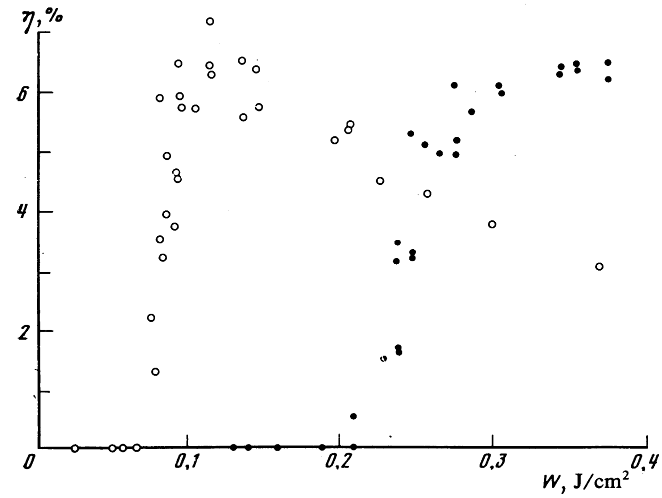

FIG. 4. Experimental dependences of the WFR efficiency η on the energy density W of the reference wave when a neodymium laser is used to irradiate the surfaces of InSb (○) and Ge (●); $\lambda = 1.06$ μm.

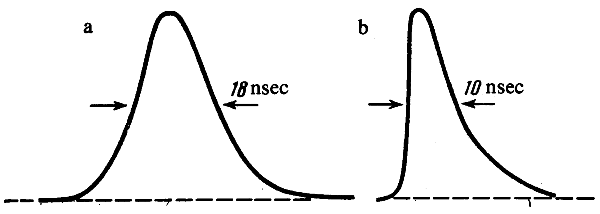

FIG. 5. Characteristic oscillograms of time dependence of the signal (a) and reversed (b) beams (twofold excess above threshold for the GaAs surface, and ruby-laser radiation).

plasma resonance $a \approx 10$–15 by exciting, during the accumulation stage, a sufficient density of photoelectrons that diverge by diffusion from the subsurface layer. This is confirmed also by the time dependence of the reversed pulse. Figure 5 shows typical pulse shapes of the signal and reversed pulses. The steepening of the leading front of the reversed pulse can be attributed to the existence of the accumulation stage. It must be noted that within the framework of the thermal model[9] the process has no threshold.

When the reference-beam threshold energy density is exceeded, the reflection coefficient increases rapidly in a certain range of W and reaches saturation. The plots of $\eta(W)$ in Figs. 3 and 4 correspond to a reversible grating. With further increase of the reference-beam energy density, a residual diffraction grating is produced on the semiconductor surface and scatters the reference beam effectively, and with good quality, into an inverted wave when the signal beam is shut off. In this region, $\eta(W)$ increases rapidly because of the onset of destructive shaping of the grating surface.[2,8] On Ge, for example, residual gratings were produced under the conditions of our experiment at $W \gtrsim 0.3$ J/cm^2 for $\lambda = 0.6943$ μm and $W \gtrsim 0.4$ J/cm^2 at $\lambda = 1.06\,\mu$m.

4. COMPARISON OF THEORY AND EXPERIMENT

Inasmuch as in the literature there are no reliable data for the effective mass $\langle m' \rangle$, the ambipolar diffusion coefficient D, and the rates of recombination processes at high degrees of semiconductor excitation, a rigorous quantitative comparison of the theoretical and experimental results is quite difficult. Nevertheless, the theoretically obtained order-of-magnitude values $W_{\rm thr} \sim 0.2$ J/cm^2 of the thresholds and $\eta \sim 10\%$ of the maximum WFR efficiency agree well with the experimental results.

Noticeably more information is obtained by qualitative comparison of the behavior of $\eta(W)$ for different semiconductors and different exciting-radiation wavelengths.

We discuss first the threshold values. The lowest threshold at both wavelengths is observed for InSb. This substance has in the unexcited state the lowest effective mass and has therefore plasma resonance at the lowest excitation levels Ne. In addition, the carrier mobility in InSb is also lower, and the absorption coefficient higher, than in the remaining investigated semiconductors. As a result, the plasma-localization layer near the surface is a minimum, and this also makes it easier to reach the resonant density. For GaAs and all the more for Ge the effective masses and the ambipolar-diffusion coefficients are noticeably higher than for InSb. As seen from Figs. 3 and 4, the threshold for reaching plasma resonance is higher in GaAs than in InSb and higher in Ge than in GaAs. Besides InSb, GaAs, and Ge, several experimental points were obtained for an Si surface at $\lambda = 0.6943$ μm. The WFR efficiency for this material is very low. Thus, η did not exceed $(2–7)\times 10^{-2}\%$ at $W \sim 0.2$–0.3 J/cm^2. The reason is that the effective mass $m^* \approx 0.3 m_e$ for Si (Ref. 15) is smaller by an order of magnitude than for the remaining investigated materials. Therefore the threshold of the plasma resonance is not reached even at plasma densities $\sim 10^{22}$ cm^{-3} in Si. We note that Si is close in its thermophysical properties to the remaining investigated semiconductors, and in the framework of the thermal model[9] it should have reflection coefficients of the same order as they.

Observation of WFR by a GaAs surface at Nd-laser wavelength was reported in Ref. 6. This result contradicts the plasma-resonance model proposed in the present study, since the neodymium photon $\hbar\omega \approx 1.18$ eV is less than the band gap $E_g \approx 1.45$ eV of this material, and no subsurface plasma is produced in this case. In our experiments with an Nd laser, the results of Ref. 6 were not confirmed. A reversed wave was produced only if a GaAs sample with plane-parallel faces was used. In this case the reflection of the reference wave from the second face produces an opposing reference wave and the WFR takes place with four-wave displacement on bulk nonlinearity of the sample. No reversal of the signal wave was produced when the second face of the sample was dull-finished.

We proceed now to discuss the absolute values of the WFR efficiency η. As aleady mentioned, when the threshold energy $W_{\rm thr}$ is exceeded the efficiency η increases quite rapidly and reaches a maximum, while for InSb it decreases slowly past the maximum, see Figs. 3 and 4. This behavior agrees with the resonant theoretical relation (4). For a quantitative comparison of the theoretical and experimental values of the maximum efficiencies $\eta_{\rm max}$ it is necessary to determine primarily the exponent p in the phenomenological relation $aJ = \gamma N_e^p$ of the subsurface plasma density N_e and the light intensity j in the recombination regime. To this end we turn to the experimental results.

We consider now the asymptotic behavior of the WFR efficiency as a function of $W \propto J_0$ in the region past the plasma resonance $a_0 \gg \varepsilon_0$. Here $n \approx 0, |\varepsilon| \approx a_0 - \varepsilon_0$, and for the efficiency we get

$$\eta \approx \frac{(a_0/p)^2}{(a_0-\varepsilon_0)^3} \approx \left(\frac{W}{W_m}\right)^{2/p} \frac{[(W/W_m)^{1/p}-1]^{-3}}{\varepsilon_0 p^2},$$

where W_m corresponds to the reference-wave energy density that ensures a maximum of the WFR reflection at $a_0 \approx \varepsilon_0$. By simply fitting it can be found that the decrease of the WFR efficiency of InSb past the maximum $\eta_{\rm max}$ is best described by this asymptotic formula with $p \approx 1$–1.5. At these values of p the experimentally attained values $\eta_{\rm max} \approx 4$–6% correspond in accord with Eq. (5) to excited-semiconductor absorption coefficients $\varkappa \approx 1.5$–2. This is in reasonable agreement with the experimental results recently obtained[18] for the absorption coefficient of a photoexcited subsurface layer. Thus, the dominant absorption mechanism in an excited subsurface layer is absorption by free carriers, since interband transitions yield much lower values of $\varkappa$ even if saturation of the interband absorption is not taken into account.

We must emphasize the similarity between the experimental $\eta(W)$ dependences and the practical equality of the maximum values η_{max} at different exciting-radiation wavelengths. This experimental result, obtained for Ge and InSb surfaces, agrees well with the conclusions of the theoretical model. With increasing wavelength λ, the plasma density corresponding to plasma resonance $a_0 \approx \varepsilon_0$ decreases $\propto \lambda^{-2}$. The cross section for absorption by free carriers, however, increases, $\sigma \propto \lambda^2$, so that the total absorption $\varkappa$ of the excited layer at resonance, which determines according to (5) the value of η_{max}, is practically independent of the wavelength of the exciting radiation.

Of all the experimental results of the plasma-resonance model, the only unexplained fact is the increase of the threshold of the process with increasing wavelength of the radiation. The physical cause of this might be sought in the fact that at equal intensity the longer-wavelength radiation heats the electron-hole plasma more strongly on account of absorption by free carriers. Therefore, despite the lower plasma density N_e needed to reach resonance at the Nd-laser wavelength it is more difficult to retain the electrons near the surface, since the diffusion coefficient increases with temperature: $D \propto T$. We hope to offer a more rigorous interpretation of the increase of the threshold after a detailed study of the properties of a subsurface plasma by powerful laser radiation.

5. MEASUREMENT OF THE REVERSAL QUALITY

Two causes of low quality of reversal in WFR-S are possible:

1) The very strong dependence of the WFR efficiency on the angle α between the reference and signal beams.[1,3] 2) The inhomogeneity of amplitude of the reference wave on the interaction area.[2,8] Figure 6 shows the experimental maximum WFR efficiencies η_{max} at different angles between the beams for a Ge surface and Nd-laser radiation. The shape of the $\eta(W)$ curves and the values of the threshold energy density are practically independent of α. So slow a decrease of the efficiency with increasing angle allows us to conclude that in the range $\alpha \lesssim 45°$ of interest to us the reflection-coefficient grating is recorded by a purely local mechanism. The cause of the decreased efficiency is apparently the decreased

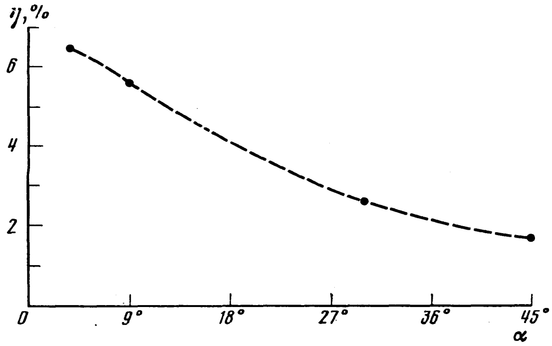

FIG. 6. Experimental dependence of the maximum values of the WFR efficiency η_{max} on the angle α between the recording beams.

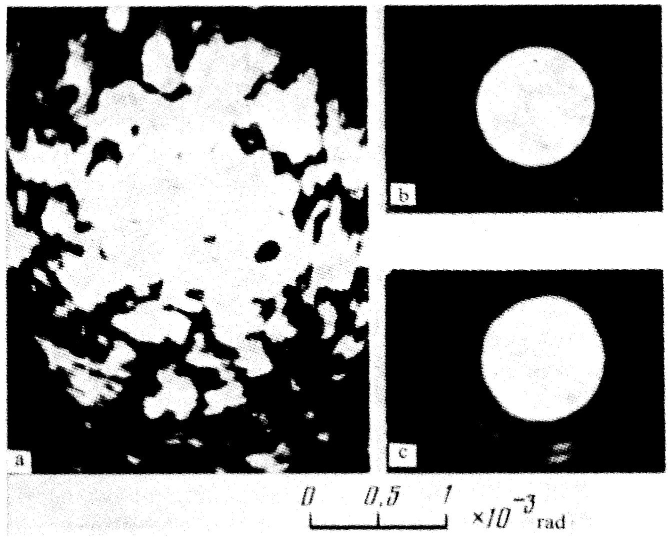

FIG. 7. Angular distribution of signal beam ahead of the phase plate (b), after its passage (a), and of the reversed beam rectified by the phase plate in the return pass (c).

transmission of the surface for oblique incidence of the signal beam and oblique for passage of the reversed beam. It is seen in addition that for the signal beams of practical interest, with divergence $\Delta\theta \lesssim 0.1$, the difference between the reflection efficiencies of different angular components does not exceed 10%, and with this efficiency the quality of the WFR can be regarded as ideal.

To improve the quality of the WFR, the beam-interaction band was localized in all experiments in the central part of the reference-wave spot, where the amplitude is constant. To investigate the WFR quality we used the standard method[22,23] of placing a phase plane in the path of the signal beam. Figure 7 shows the angular distribution of the signal beam ahead of the phase plate (b), after distortion by the phase plate (a), and the reversed beam rectified by the plate in its return pass (c). It can be seen that the reversed beam duplicates in the far zone not only the diameter and shape of the signal diffraction spot, but also fine details of its angular structure. Quantitative measurements of the fraction of the WFR (Ref. 24) were performed, using stepped attenuator SA (Fig. 2) by comparing the calorimetric and photometric measurements. At a high degree of homogeneity of the reference wave in the interaction region, the fraction of the WFR was high and did not drop, with allowance for the experimental error, below 85%.

6. DISCUSSION OF RESULTS

We have thus obtained in experiment WFR on surfaces of various semiconductors, with efficiency up to 7% and with good quality of the reversed beam. A theoretical model was proposed for recording surface holograms by a plasma reflection mechanism. The theoretical results are in reasonable agreement with experiment. On the basis of these results we can draw certain conclusions. First, WFR-S on semiconductors can find application in the long-wave region of the spectrum, where plasma resonance can be achieved with not too high densities of the excited photoelectrons. Second, this mechanism of surface nonlinearity is reversible and has low inertia. The grating relaxation time is determined mainly by a recombination time $\sim(\gamma N_e^2)^{-1}$ and amounts to $\sim 10^{-12}$ sec for $\gamma \approx 10^{-30}$ cm^6/sec, raising the hope that surface nonlinearities of semiconductors can be effectively used for WFR of broadband radiation.

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