NEW OPTICAL AND ELECTRIC EFFECTS IN A STANDING LIGHT WAVE

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Excitation of a standing wave capable of ensuring impurity photoconductivity can cause a crystal to acquire fundamentally new optical and electric properties, which it does not possess in an ordinary traveling wave.

Assume that a standing monochromatic light wave, with a quantum energy smaller than the width of the forbidden band, but sufficient for the photoexcitation of the impurity centers present in the crystals, is produced in a semiconductor. In this case the wave field, which is zero at the nodes and is maximal at the antinodes, should lead to the photoexcitation of the free carriers (for examples, electrons and donors); this excitation is periodic along the beam direction (X axis) (Fig. 1), with a period equal to half the light wavelength λ in the crystal. If at the same time the screening length L is smaller than λ, a condition easily attained at a relatively low electron density[1)], then the resultant additional carrier density will be periodic along the X axis, which in turn leads to anisotropy of the photoconductivity in an initially isotropic semiconductor.

The periodic electron-gas density produces in turn a periodic modulation of the dielectric constant ε. For an additional probing light source (I_{pr}, Fig. 1), which is small enough not to disturb the spatial structure produced by the exciting light source (I_{exc}), such a modulation of the optical density will be quite similar to a diffraction grating that produces various types of diffraction. A similar effect is observed when ultrasonic waves pass through a substance (cf., e.g., [1]).

Finally, the inhomogeneous carrier distribution, which results from the equilibrium between the diffusion field and the electric field of attraction between the electrons and the donors, gives rise to a spatially periodic potential $\phi(x)$. Thus an additional potential

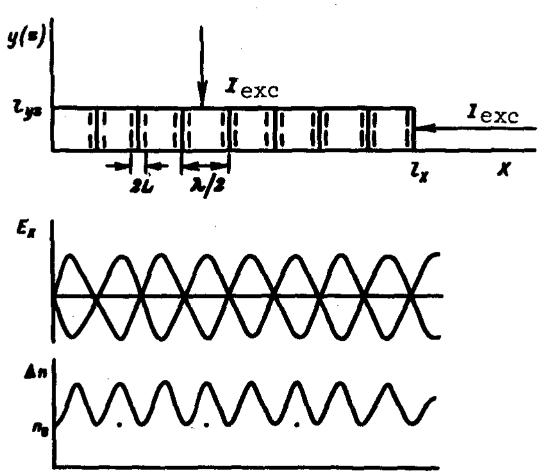

Fig. 1

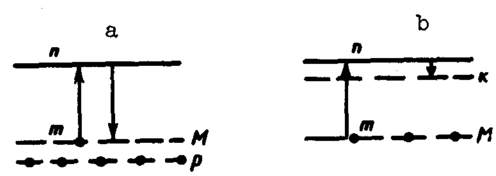

Fig. 2

Fig. 1. Experimental setup for the observation of the phenomena under consideration.

Fig. 2. Band schemes used in the calculations.

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1) These were precisely the considerations that let us to choose such a light-absorption scheme, since other mechanisms (say interband absorption) would produce a non-equilibrium electron distribution, determined by the diffusion length, which is always larger than λ.

substructure, with a period larger than the lattice potential, is superimposed on the crystal-lattice potential along the X axis. This should lead, in principle, to the appearance in the conduction band of a series of allowed and forbidden energy bands for the electron motion along the X axis, and by the same token change radically the electric and optical properties of the crystal. A similar model was discussed earlier in [2], but for a much more difficulty realizable variant, in which a hypersonic wave passed through the sample.

Let us obtain a quantitative estimate of the foregoing phenomena. Let the exciting light knock out electrons from the levels M (Fig. 2a); for simplicity, we assume these two levels to be little populated (owing to compensation by the P levels), i.e., $M \gg m$, where m is the electron density at the levels M, so that the inverse electron flux does not depend on the occupation of the levels M. A similar situation occurs (Fig. 2b) when the excited light tranfers the electrons from the levels M to the conduction band, and the inverse flux is determined by the levels K, with $K \gg n$. Such a scheme is realized in semiconductors of the CdS type [3]. The initial dark electron density (which is uniform throughout the sample) is assumed equal to n_0.

Assuming that the electric field of the light wave equals

$$E = E_0 \cos\left(\omega t - \frac{2\pi x}{\lambda}\right),$$

then the field in the standing wave, equal to

$$2E_0 \sin\omega t \cos\frac{2\pi x}{\lambda},$$

will determine the intensity of the incident light I_{exc}

$$I_{exc} = \frac{2\epsilon c E_0^2 \sin^2\omega t \cos^2\frac{2\pi x}{\lambda}}{8\pi N} = I_0 \cos^2\frac{2\pi x}{\lambda}, \tag{1}$$

(N is the refractive index), which turns out to be periodic with a period equal to $\lambda/2$. In the stationary state we obtain for the non-equilibrium electron increment

$$\frac{d\Delta n}{dt} = \beta\sigma_m m I_0 \cos^2\frac{2\pi x}{\lambda} - \Delta n\,\gamma_m (M - m) = 0, \tag{2}$$

where σ_m is the photon-capture cross section, β the quantum yield, and γ_m the cross section for the capture of an electron on the level M (for the case shown in Fig. 2b, the second term would contain the quantity $\gamma_k K_1$, which should not change the situation). From (2) we get

$$n(x) = n_0 + \frac{\beta\sigma_m m I_0}{\gamma'_m M}\cos^2\frac{2\pi x}{\lambda} = n_0 + \Delta n_0 \cos^2\frac{2\pi x}{\lambda}. \tag{3}$$

If we assume $\sigma_m \approx 10^{-18}$ cm^2, $\gamma_{m(k)} \approx 10^{-19}$ cm^3/sec, and $M \approx 10^{18}$ cm^{-3}, values typical to CdS, and $m/M \approx 0.1$, then it turns out that $\Delta n_0 \approx 10^{16}$ cm^{-3} even at a moderate exciting light intensity, $I_0 \approx 10^{17}$ qu/cm^2 sec. For the quasi-equilibrium state we obtain further from the Boltzmann distribution

$$\phi(x) = kT \ln\frac{n(x)}{n_0} \tag{4}$$

This means that at a readily realizable ratio $\Delta n_0/n_0$, equal to several orders of magnitude, the amplitude of the potential can reach several times kT.

In calculating the photoconductivity anisotropy it is necessary to take into account the fact that the density distribution n(x) was calculated by us neglecting the diffusion of the carriers to the light-wave nodes. We shall therefore assume henceforth that the electrons have a uniform distribution along the screening length $L = (\varepsilon kT/4\pi ne^2)^{1/2}$ near the nodes, with a concentration

$$n' \approx n_o + \Delta n_o \sin^2(2\pi L/\lambda) \approx n_o + \Delta n_o (2\pi)^2 (L/\lambda)^2.$$

In the approximation in which $\lambda \gg L$, $\Delta n_0 (2\pi L/\lambda)^2 \gg n_0$, $kT > \phi_0$, and $\ell_x = \ell_z = \ell_y$ (Fig. 1), an estimate yields

$$\sigma_x / \sigma_y = 2\pi^2 L/\lambda. \tag{5}$$

Let us consider now the diffraction of light. When the probing light is perpendicular to the exciting beam, the changes of the refractive index influence the phase of the wave of the probing light, so that the front of the wave emerging from the sample will be modulated at double the frequency of the exciting light, with the amplitude of the phase modulation (cf., e.g., [1]): $a = 2\pi \sqrt{\varepsilon}\, \ell_z/\lambda_{pr}$, where $\Delta\varepsilon$ is the change of the dielectric constant:

$$\Delta\epsilon = -\frac{4\pi ne^2}{m(\omega^2 - \omega_o^2)} + \frac{4\pi ne^2}{m\omega^2} \tag{6}$$

$h\omega_0$ is the energy of the donor electron ionization. The first term determines the change of ε due to the donor ionization, the second the change due to the interaction between the light and the free carriers. We see that when $\omega \gg \omega_0$ both terms cancel each other and in order to enhance the effect it is necessary to have ω not too far from ω_0.

It should be noted that similar experiments with light passing through a medium perturbed by an ultrasonic wave have shown that the diffraction pattern can be easily observed when $a \approx 1$ [1]. In our case, for light from a ruby laser ($\lambda_{pr} \approx 0.7\ \mu$) at $10 \geq \lambda_{exc}/\lambda_{pr} \geq 2$ we get $a \approx 1$ already for $\Delta n_0 \approx 10^{13} - 10^{14}\ cm^{-3}$, which can be readily attained experimentally. This form of diffraction becomes most clearly pronounced when $\lambda_{exc} \gg \lambda_{pr}$. When the frequencies of the exciting and probing light come closer together, the selective Bragg reflection becomes more perceptible; the angle of observation of this reflection is given by the well known relation

$$\sin\phi_B = \frac{k\lambda_{probe}}{\lambda_{exc}}, \text{ where k is an integer.}$$

Finally, let us consider the influence of the periodic potential $\phi(x)$ on the energy spectrum of the electrons. As is well known, the presence of a periodic potential leads to the appearance of allowed (ΔE_a) and forbidden (ΔE_f) energy bands determined by the magnitude of the potential ϕ_0 and by the period of its distribution (cf., e.g., [4]). Here $\Delta E_a \approx \hbar^2\pi^2/2m(\lambda/2)^2$ and $\Delta E_f \approx \phi_0$. It is clear that to observe such a picture experimentally it is necessary to have $\Delta E_f \approx \Delta E_a \geq kT$, i.e., it is necessary to increase the quantum energy and to choose materials with small effective mass. In CdS ($m \approx 0.2m_0$), when working with an argon laser ($h\omega = 2.4$ eV) we have $\Delta E_a \approx 3.5 \times 10^{-4}$ eV. In GaAs ($m = 0.065m_0$), when working with a laser of the same material ($h\omega \approx 1.4$ eV) we have $\Delta E_a \approx 4 \times 10^{-4}$ eV. As to the forbidden bands, $\Delta E_f > kT$ already when $\Delta n_0/n_0 \approx 10$. Thus, this effect can be registered at liquid-helium temperature.

In conclusion, the author thanks S. M. Ryvkin for interest in the work and N. I. Kramer and I. M. Fishman for useful discussions.

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VIOLATION OF P-PARITY IN INELASTIC SCATTERING REACTIONS WITH PHOTON EMISSION

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The existence of a weak interaction between nucleon causes the nuclear states to be no longer characterized by a definite parity. This results in a number of experimentally-observed effects such as the angular asymmetry of the photons emitted by polarized excited nuclei [1] and circular polarization of the photons emitted by unpolarized nuclei [2]. The weak interaction between nucleons is of the order of $G\mu^2 \approx 3 \times 10^{-7}$ ($G = 10^{-5}/m^2$ is the weak interaction constant, and μ and m are the pion and nucleon masses) relative to the strong interaction. As a result of a number of phenomena [3], the effects of weak nucleon interaction can be appreciably enhanced and reach a value on the order of 10^{-4}. One of the enhancement effects is dynamic enhancement, wherein [3] the amplitude a_{0n} of the admixture of the state n, with a parity opposite to that of the "principal" state 0 with a given spin and parity, can be greatly increased as a result of the proximity of these states, since $a_{0n} \sim (E_0 - E_n)^{-1}$. This effect may turn out to be significant at high excitation energies, where the level density is large and levels with the same spin values but with different parities may turn out to be very close to each other.

We consider in this paper circular polarization of photons from excited nuclei produced by inelastic particle scatterg.

Appreciable enhancement effects can be expected when the P-parity violation is connected with resonant scattering accompanied by formation of states of mixed parity. The enhancement effects may cause the weak nucleon interaction to reach a value on the order of 10^{-4}.

Let us spin of the incoming particle be j_1, the spin of the nucleus in the initial state j_2, the spin of the excited state j_3, the angular momentum of the emitted photon j_γ, and the spin of the nucleus produced after photon emission j_4. We expand the scattering amplitude in partial waves with definite total spin j, total orbital angular momentum L, and total angular momentum J:

$$< n' m_1' m_3 | S | n, m_1, m_2 > = \Sigma C^{jm}_{j_1 m_1; j_2 m_2} C^{JM}_{L\mu; jm} \times$$