Optical orientation of an electron–photon system under stimulated luminescence conditions

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Theoretical and experimental studies have been carried out of the behavior of the polarization of charge carriers and of recombination radiation (luminescence) photons at various excitation rates G of oriented particles. Although under optical orientation conditions the transition from spontaneous luminescence to lasing increased considerably the luminescence polarization, under steady-state lasing conditions the average electron spin s tended to zero. The dependence $s(G)$ had a maximum in the region of the transition from luminescence to lasing. Nevertheless, whereas in the spontaneous regime the spin relaxation time of carriers was short compared with their lifetime, the spin s was always small compared with the initial spin s_i of the generated particles. The nature of depolarization of the luminescence in a transverse magnetic field was determined. A comparison of the theoretical and experimental results for PbTe crystals showed that the proposed model provides a satisfactory description of the physical situation.

1. INTRODUCTION

Investigations of optical orientation are usually carried out at relatively low pumping rates when the excited luminescence is due to spontaneous transitions and does not react back on the electron system.[1] However, work has recently appeared on optical orientation under stimulated luminescence (lasing) conditions.[2–4] It is suggested in Refs. 2 and 3 that the reduction in the lifetime of excited carriers associated with induced transitions increases the steady-state polarization. (Under spontaneous luminescence conditions the average spin s of the oriented carriers increases as the ratio of the carrier lifetime τ is reduced to the spin relaxation time τ_s.)

It was found[2] that in the case of a PbTe crystal the degree of circular polarization ρ of the spontaneous luminescence was of the order of a fraction of a percent, whereas in the stimulated luminescence case excited by two-photon pumping with a CO_2 laser it amounted to 30–50%. Hence it was concluded that strong orientation of carriers was achieved. However, as we showed in a brief communication,[4] a strong nonlinear interaction between electron and photon spins results in redistribution of the spin between the photon and electron subsystems. Consequently, the average electron spin is anomalously low when the degree of polarization of the emitted luminescence is relatively high.

We report the results of experimental and theoretical investigations of optical orientation in the course of transition from the spontaneous luminescence to the stimulated emission regime. In Sec. 2 we describe an experiment carried out on PbTe crystals. In Sec. 3 we propose the simplest theoretical model which allows for spontaneous and stimulated optical transitions. Finally, in Sec. 4 we analyze the results obtained.

2. DESCRIPTION OF EXPERIMENTS

Favorable conditions for the investigation of the polarization of the electron–photon system under stimulated luminescence conditions are provided by two-photon pumping which makes it possible to orient carriers in the bulk of a sample. Our experiments were carried out on n-type PbTe samples with carrier densities $(4\text{–}7)\times 10^{17}$ cm^{-3}. Nonequilibrium carriers were excited by two-photon absorption of radiation from a Q switched CO_2 laser ($\lambda = 10.6$ μm)[1)] which generated pulses of 1–2 kW power, 300 nsec duration, and 250 Hz repetition frequency. The linearly polarized radiation of this laser was converted into circularly polarized radiation by a quarter-wave CdS plate and it was focused on samples of PbTe which were immersed directly in liquid helium inside a cryostat at a temperature of 1.8 K. The radius of the illuminated spot R was small compared with the thickness of the sample l. The recombination radiation, which was investigated in the transmission geometry, was separated from the high-power pump radiation by an LiF filter and detected with a Ge:Au photodetector. The luminescence polarization was analyzed with a $\lambda/4$ mica plate and a linear grating polarizer.

The luminescence depolarization curves were recorded in a transverse magnetic field (Hanle effect) modulating the polarization at a frequency of about 20 Hz by rotation of the plate.

Estimates deduced from the absorbed intensity and the luminescence spectrum showed that the density of excited carriers was less than 10^{15} cm^{-3} (which was considerably less than the equilibrium electron density) and the temperature of the electron–hole gas was approximately 25 K. In this case the electrons were degenerate with a quasi-Fermi level at ~ 10 meV and the holes had a Boltzmann distribution. The luminescence spectrum was located close to ε_g ($\hbar\omega_L \sim \varepsilon_g \approx 2\text{–}3$ meV). All this indicated that holes were oriented in n-type PbTe.

Figure 1 shows the experimentally determined values of the polarization ρ and of the integral intensity I of the luminescence obtained for different rates of excitation G of photocarriers. (In the case of two-photon pumping the rate of excitation G was proportional to the square of the intensity of the pump radiation: $G \propto J^2$.) Our experiments were carried out on the same sample and differed only in terms of the focusing of the pump radiation. The results presented in Fig. 1a were obtained for a smaller illumination spot than in the experiments represented by the results in Fig. 1b. Clearly, at

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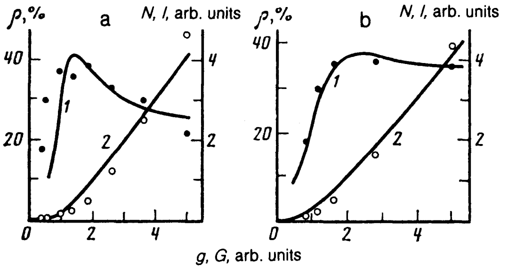

FIG. 1. Luminescence intensity and polarization versus the carrier excitation rate. The open circles represent the experimental values of the luminescence intensity I and the black dots represent the luminescence polarization ρ. Curves 1 and 2 give the theoretical dependences of the dimensionless variables $\rho(g)$ and $N(g)$, respectively. a) Experimental results for strong focusing compared with theoretical calculations calculated assuming the following parameters: $s_i = 0.09$, $\gamma = 0.025$, $t_s = 0.4$. b) Weak focusing ($R_b > R_a$), compared with theoretical calculations made on the assumption that $s_i = 0.12$, $\gamma = 0.1$, $t_s = 0.4$.

low pumping rates (spontaneous luminescence region) the differential quantum efficiency of the detected luminescence dI/dG was considerably less than in the lasing range. The luminescence polarization also increased strongly as a result of lasing. When the pump beam was focused more strongly, it was found that in the intermediate range of the values of G the polarization ρ had a clear peak, but if the illumination spot was larger such a maximum was practically undetectable. These differences were due to differences in the ratio of the radius of the illumination spot to the length of the active zone, as shown below.

Investigations of depolarization of the luminescence by a magnetic field H applied transverse to the direction of excitation (Hanle effect) showed that in the case of crystals with a broad lasing threshold the dependence $\rho(H)$ could be described by a Lorentzian profile with a half-width which increased with the pumping rate (Fig. 2). At high pumping rates the broadening of the Hanle curves was so large that the maximum field generated in our coils (~ 2 kOe) did not result in any significant depolarization. In the case of samples with an abrupt lasing threshold no depolarizing effect of the magnetic field was observed for any value of H. In the spontaneous luminescence range the value of ρ was negligible even for $H = 0$, whereas in the stable lasing range the influence of the applied magnetic field on ρ was negligible.

3. THEORETICAL MODEL

3.1. We shall use a system of rate (balance) equations[5] to describe optical orientation of carriers in the case of spontaneous and stimulated emission of photons.

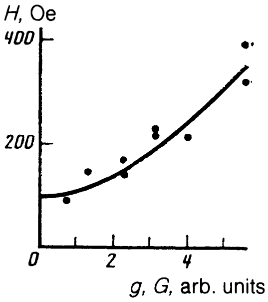

FIG. 2. Depolarization of the luminescence in a magnetic field (represented by the half-width of the Hanle curve as a function of the rate of excitation of carriers). The dots are the experimental values and the curve is a theoretical calculation made assuming that $s_i = 0.15$, $\gamma = 0.2$, $t_s = 0.1$, and that the g factor of carriers is 25.

We shall assume that pumping at a rate G generates nonequilibrium carriers with an average spin s_i directed along the z axis in a semiconductor. These carriers recombine radiatively or nonradiatively (τ_{nr} is the nonradiative lifetime). The radiative transitions can be spontaneous (τ_r is the radiative lifetime) or stimulated, the latter representing the emission of photons in one of two lasing modes propagating within a small solid angle Ω near the z axis and differing in the sign of the circular polarization. The order-of-magnitude estimate is $\Omega \sim R^2/l^2$, where R is the radius of the illuminated spot excited by the laser beam and l is the length of the active region. Then, in the case of two-mode lasing,[5,6] the density of photocarriers n, the density of the spin of these photocarriers S, and the densities of the radiation in the modes with right- and left-handed circular polarizations I_+ and I_- are described by

$$\dot n = G - wn(I_+ + I_-) - 2n(\mathbf{Sk})(I_+ - I_-) - n/\tau,$$

$$\dot{\mathbf S} = G\mathbf{s}_i - w(I_+ + I_-)\mathbf S - \frac{u}{2} n(I_+ - I_-)\mathbf k + [\tilde\omega \mathbf S] - \mathbf S\left(\frac{1}{\tau} + \frac{1}{\tau_s}\right), \quad (1)$$

$$\dot I_+ = \left\{h\nu[wn + 2u(\mathbf{Sk})] - \frac{1}{T}\right\} I_+ + \frac{h\nu\gamma}{2\tau}\left[n + \frac{2u}{w}(\mathbf{Sk})\right],$$

$$\dot I_- = \left\{h\nu[wn - 2u(\mathbf{Sk})] - \frac{1}{T}\right\} I_- + \frac{h\nu\gamma}{2\tau}\left[n - \frac{2u}{w}(\mathbf{Sk})\right],$$

where $\mathbf k$ is a unit vector along the direction of lasing; T is the characteristic lifetime of a photon in a crystal; τ is the characteristic lifetime of a carrier in the case of nonradiative and spontaneous transitions ($\tau^{-1} = \tau_r^{-1} + \tau_{nr}^{-1}$); γ is a parameter showing what fraction of these transitions is used to excite laser modes: $\gamma \propto (R^2\tau)/(l^2\tau_r)$; τ_s is the spin relaxation time; w and u are coefficients describing respectively the average probability and the polarization dependence of the induced transitions. (If the spin pumping Gs_i is due to direct absorption of circularly polarized light, then $s_i = u/2w$. According to Ref. 7, in the case of PbTe, we have $u/w = 0.8$.) The term $[\boldsymbol\omega \cdot \mathbf S]$ in the second equation of the system (1) allows for the influence of a magnetic field, applied to the sample in order to study the Hanle effect, on the value of $\mathbf S$ ($\tilde\omega$ is the frequency of the Larmor precession of electron spins in this field), and $h\nu$ is the energy of the emitted photon.

If the pumping rate is so low that all the nonlinear terms can be ignored in the system of equations (1), then the system (1) reduces to ordinary equations for optical orientation under spontaneous luminescence conditions. The steady-state solution of this system is well known. For $\omega = 0$, $k_x = k_y = 0$, and $k_z = 1$, we have

$$n = G\tau, \quad S = Gs_i\tau\tau_s/(\tau+\tau_s), \quad I_\pm = h\nu\gamma GT[1 \pm s_i\tau_s/(\tau+\tau_s)]/2. \quad (2)$$

In the other limiting case (high pumping intensity and strong stimulated emission) the system of equations (1) is identical with the equations of Ref. 4 and has three steady-state solutions:

$$I_+ = h\nu GT, \quad I_- = 0, \quad n = (w - 2us_i)/[(w^2 - u^2)h\nu T],$$
$$S = (2ws_i - u)/[(w^2 - u^2)2h\nu T], \quad (3a)$$

$$I_\pm = h\nu GT(1 \pm 2s_i w/u)/2, \quad n = (wh\nu T)^{-1}, \quad S = 0, \quad (3b)$$

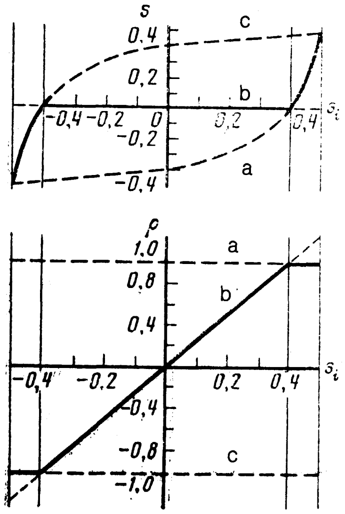

FIG. 3. Theoretically calculated average carrier spin s and degree of polarization of the luminescence under steady-state lasing conditions versus the initial value of the spin of the excited carriers s_i. Curves a, b, and c correspond to the states (3a), (3b), and (3c). The continuous curves correspond to stable states of the electron–photon system and the dashed curves represent unstable states. Numerical results are plotted for PbT ($u/w = 0.8$).

$$I_+=0,\quad I_-=h\nu GT,\quad n=(w+2us_i)/[(w^2-u^2)h\nu T],$$
$$S=(2ws_i+u)/[(w^2-u^2)2h\nu T]. \tag{3c}$$

Figure 3 shows how the polarization of electrons and phonons depends on the initial spin s_i. In the state (a) only a mode with right-handed circular polarization is generated and in this case the amplification is equal to the losses: $h\nu[wn + 2u(\mathbf{S}\cdot\mathbf{k})] = 1/T$. If

$$s_i > s_{cr} = u/2w, \tag{4}$$

then $S > 0$ holds and the gain experienced by the mode I_- is less than the losses, $h\nu[wn - 2u(\mathbf{S}\cdot\mathbf{k})] < 1/T$, and this mode is not excited. Solution (b) corresponds to $I_- < 0$ ($\rho > 100\%$), which is physically meaningless. However, if $|s_i| < s_{\rm cr}$ holds, then in the state (a) we have $S < 0$ and the gain of the mode I_- is greater than the losses so that a small fluctuation destroys this state. For similar reasons the state (c) is also unstable and, as pointed out in Ref. 4, the solution (b) is obtained. We can easily see that for this state the gain and the losses always balance one another. Finally, for $s_i < -s_{\rm cr}$, the system of electrons and phonons is in the state (c).

We shall now find the steady states for intermediate values of G when the investigated electron–photon system goes over from the spontaneous to the lasing regime. These states are given by the solutions of the following system of four nonlinear equations:

$$g-\eta N-2\xi\eta N(\mathbf{sk})\rho-\eta=0,\quad \eta N+\gamma\eta-N+2\xi\eta N(\mathbf{sk})\rho=0,$$
$$\eta N\rho+2\xi\eta N(\mathbf{sk})+2\xi\gamma\eta(\mathbf{sk})-N\rho=0, \tag{5}$$
$$g\mathbf{s}_i-\eta N\mathbf{s}-{}^1/_2\xi\eta N\rho\mathbf{k}+[\omega\mathbf{s}]\eta-\mathbf{s}\eta/t_s=0,$$

obtained from the system (1) by introducing the following dimensionless variables:

$$g=G\tau wT,\quad \eta=nwT,\quad N=(I_++I_-)w\tau/h\nu,\quad s=S/n,$$
$$\rho=(I_+-I_-)/(I_++I_-),\quad \xi=u/w,\quad \omega=\bar{\omega}\tau,\quad t_s=\tau_s/(\tau+\tau_s). \tag{6}$$

3.2. We shall first consider the situation in the absence of a magnetic field ($\omega = 0$). In the experiments the vector $\mathbf{s}_i$ is directed along $\mathbf{k}$, so the vectors $\mathbf{s}$ and $\mathbf{k}$ are always parallel to one another. We then have $\mathbf{s}\cdot\mathbf{k} = s$, and the system (5) reduces to four nonlinear algebraic equations for four unknowns η, N, ρ, and s. It is not possible to find the explicit dependences of these unknowns on the dimensionless pumping rate g and other parameters of the problem. However, we can derive analytic expressions for η, N, ρ, and g if we take the electron spin s as our argument.

Simple algebraic transformations of the first three equations in the system (5) yield expressions for N, g, and ρ in terms of η and s:

$$N=\gamma\eta/(1-\eta-z), \tag{7}$$

$$g=\eta[1+\gamma(\eta+z)/(1-\eta-z)], \tag{8}$$

$$\rho=2\xi s(1-z)/(1-\eta), \tag{9}$$

where

$$z=2\xi\eta(s\rho)=4\xi^2\eta s^2/(1-\eta+4\xi^2\eta s^2). \tag{10}$$

Substituting these expressions in the last of the equations in the system (5) and carrying out some involved algebraic transformations, we obtain the following relationship between the average spin and the carrier density:

$$[1-(2\xi s)^2][(1-\gamma)s_it_s-(1-\gamma t_s)s]\eta^2-\{[2-\gamma(1+(2\xi s)^2]s_it_s-[2-\gamma t_s(1+\xi^2)]s\}\eta+(s_it_s-s)=0. \tag{11}$$

For a fixed value of the dimensionless carrier density η this is a cubic equation for the spin s. If the value of s is fixed, then η is described simply by a quadratic equation which can readily be solved.

We can represent clearly the nature of the dependence $\eta(s)$ by solving Eq. (11) at several characteristic points: for example, if $s = 0$, then

$$\eta_1=1,\quad \eta_2=(1-\gamma)^{-1},$$

and for $s = s_it_s$, we have

$$\eta_1=0,\quad \eta_2=1-\xi^2t_s.$$

We can easily show that the state with $s = s_it_s$ and $\eta = 0$ and the adjoining state characterized by $0 < \eta \ll 1$ and $s \approx s_it_s$ correspond to the conventional spontaneous luminescence. The solution with $s = 0$ and $\eta = 1$ gives the second steady state of the electron–photon system in the stable lasing limit represented by Eq. (3b).

We shall now consider the existence of roots of Eq. (11). There are no solutions of this equation if the discriminant of Eq. (11)

$$D=y^2\left(1+\gamma t_s\frac{s_i-s}{s-s_it_s}\right)-\left(\frac{\gamma t_s}{4}\right)y(1-4ss_i)+\left(\frac{\gamma t_s}{4\xi}\right)^2[s_i-s(1+\xi^2-4\xi^2ss_i)]^2 \tag{12}$$

is negative. We then have $y = s(s - s_it_s)$.

We shall use the fact that in all cases of practical interest we can assume that $\gamma t_s \ll 1$. We shall show later that the average electron spin is also small ($s \ll 1$) and the most interesting values of s are those exceeding the electron spin in the spontaneous luminescence region (s_it_s). We then have $D < 0$

only if equation $D(y)=0$ has real roots. Ignoring all small terms, we have

$$D\approx y^2-2(\gamma t_s/8)y+(\gamma t_s/8)^2(2s_i/\xi)^2=0. \qquad (13)$$

We can see that apart from the spin given by Eq. (4), there is a second critical value of the spin of the newly created particles, $\tilde{s}\approx s_{cr}=\xi/2$, such that for $|s_i|>\tilde{s}$, we can solve Eq. (11) for any value of s, whereas for $|s_i|<\tilde{s}$ and some values of s, there are no real roots of Eq. (11). A more rigorous estimate gives

$$\tilde{s}=\xi/2+\gamma(1-\xi^2)/[(\xi^2+2\gamma/t_s)^{1/2}-\xi]. \qquad (14)$$

We can distinguish three characteristic ranges of the values of s_i. For $s_i>\tilde{s}$, then an increase in the pumping rate increases the electron spin monotonically from $s_0=s_it_s$ to

$$s_\infty=(2s_i-\xi)/[(1-2\xi s_i)2], \qquad (15)$$

and the luminescence polarization reaches 100%. This corresponds to the state (3a). As can be seen from Figs. 4a and 4b, the dependence of the electron density $\eta(g)$ is nonmonotonic for these values of s_i. If g is small, then $d\eta/dg>0$, whereas in the limit $g\to\infty$, we find that $d\eta/dg<0$.

In the range of intermediate values of the spin $s_{cr}<s_i<\tilde{s}$, an increase in the pumping rate increases monotonically the carrier density (Fig. 4c). In this case the average electron spin first increases, reaching a maximum

$$s_{max}=\frac{s_it_s}{2}\left[1+\left(1+\frac{\gamma}{2t_ss_i^2}\right)^{1/2}\right], \qquad (16)$$

and then falls to the value described by Eq. (15). In the limit $g\to\infty$, we again have the state (3a).

Finally, in the range where $s_i\leqslant s_{cr}$ (Fig. 4d), which corresponds to the usual optical pumping, at high values of the pumping rate ($g\to\infty$) the system goes over to the state (3b). The maximum value of the spin is then

$$s_{max}=\frac{s_it_s}{2}\left\{1+\left\{1+\frac{\gamma}{2t_ss_i^2}\left[1-\left[1-\left(\frac{2s_i}{\xi}\right)^2\right]^{1/2}\right]\right\}^{1/2}\right\}. \qquad (17)$$

An analysis of the above relationship readily shows that if $\gamma/t_s\gg1$, then the maximum of the dependence $s(g)$ becomes sharper:

$$s_{max}\approx\frac{1}{2}\left\{\frac{\gamma t_s}{2}\left[1-\left[1-\left(\frac{2s_i}{\xi}\right)^2\right]^{1/2}\right]\right\}^{1/2}\gg s_it_s. \qquad (18)$$

However, if the spin relaxation time is short ($t_s\ll1$), then the maximum average spin is $s_{max}<t_s^{1/2}$ and it remains small for all the generation rates.

3.3. Although Eqs. (7)–(11) give the general solution of the problem, simpler and clearer relationships can be obtained in some limiting cases. For example, if the pumping rate is low so that

$$g<1 \quad \text{and} \quad [g-1/(1+z)]^2\gg\frac{4\gamma g}{1+z},$$

we have

$$N\approx\frac{\gamma g}{1-g(1+2\xi(s\rho))}, \qquad \eta=\frac{g-N}{1-\gamma}, \qquad (19)$$

whereas for excitation above the threshold, $g>1$ and $[g-1/(1+z)]^2\gg4\gamma g/(1+z)$, we find that

$$N\approx g-\frac{1}{1+2\xi(s\rho)},$$

$$\eta=\frac{1}{1-\gamma}\left[\frac{1}{1+2\xi(s\rho)}-\frac{\gamma g}{g(1+2\xi(s\rho))-1}\right]. \qquad (20)$$

It is clear from the above analysis that if $2s_i\leqslant\xi s\ll1$, the term $z=2\xi(s\rho)$ can usually be neglected. In this case the luminescence intensity and the carrier density are practically independent of their polarization. It is then convenient to express s and ρ in terms of η and N:

$$s=\frac{gs_it_s(1-\eta)}{\eta[(1-\eta)(Nt_s+1)+\xi^2\eta(N+\gamma)t_s]}, \qquad (21)$$

$$\rho=\frac{2\xi gs_it_s(N+\gamma)}{N[(1-\eta)(Nt_s+1)+\xi^2\eta(N+\gamma)t_s]}. \qquad (22)$$

Substituting here the expressions for η and N, we can for example find that for $\gamma\gg t_s$, the maximum photon polarization is

$$\rho_{max}\approx\frac{2s_i}{\xi}\left(1+\frac{\xi^2}{4}\frac{t_s}{\gamma}\right) \qquad (23)$$

and it is attained for

$$g\approx1+\frac{2}{\xi^2}\frac{\gamma}{t_s}. \qquad (24)$$

In the limit of very high values of g the polarization of light tends to $2s_i/\xi$. In the limit $g\to\infty$, we have $\rho\to2s_i/\xi$, so that at high values of the ratio γ/t_s the maximum of the dependence $\rho(g)$ is very flat.

When Eqs. (19)–(24) are used, we must remember that although z is $\ll1$ and has practically has no effect on the values of N and η, the formulas for s and ρ include the difference $1-\eta$ which is small in the lasing range. The correction

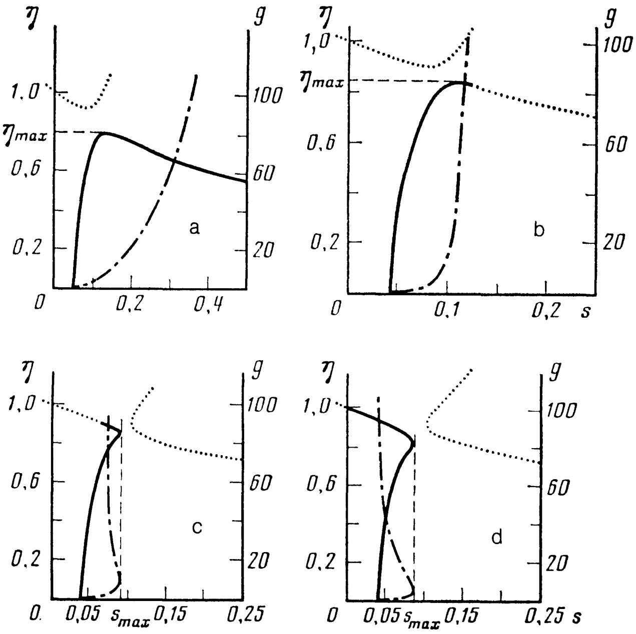

FIG. 4. Theoretically calculated photocarrier density η and rate of excitation g versus the average spin s for different values of the initial spin s_i: a) $s_i=0.5$; b) $s_i=0.44$; c) $s_i=0.42$; d) $s_i=0.4$. The continuous curves represent η in the states found experimentally and the dotted lines are the solutions of Eq. (11) which are physically meaningless. The chain curves are the dependences $g(s)$. The calculations were carried out for the parameters $\gamma=0.5$ and $t_s=0.1$.

to this difference can have relatively large values.

3.4. We shall now consider the behavior of the electron–photon system in the presence of a transverse magnetic field.

We shall consider the specific case when ω is directed along the y axis. The, the last vector equation in the system (5) reduces to two algebraic equations for the z and x projections of the carrier spin:

$$gs_i - N\eta s_z - {}^1/_2 \xi\eta N\rho - \omega\eta s_x - s_z\eta/t_s = 0, \tag{25}$$

$$-N\eta s_x + \omega s_z - s_x\eta/t_s = 0. \tag{26}$$

Solving Eq. (26) for s_x and substituting the resultant solution in Eq. (25), we obtain

$$gs_i - N\eta s_z - \frac{\xi}{2}\eta N\rho - \frac{s_z\eta}{t_s}\left[1 + \frac{\omega^2 t_s^2}{Nt_s+1}\right] = 0. \tag{27}$$

We can see that, as is usual in the Hanle effect, a transverse magnetic field effectively reduces the spin relaxation time:

$$t_s' = t_s[1 + \omega^2 t_s^2/(Nt_s+1)]^{-1}. \tag{28}$$

If the intensity of the incident light is low ($g \lesssim 1$), then $Nt_s \ll 1$ and the interaction between electrons and phonons does not affect the depolarization process. However, if the pumping rate is high ($g \gg 1$), we find from Eq. (20) that $Nt_s \approx gt_s$ and the effect of the magnetic field is weakened by a factor of $(gt_s + 1)$. Since we have $t_s \ll 1$, these two regions overlap and we can always use the approximate expression

$$t_s' = t_s[1 + \omega^2 t_s^2/(gt_s+1)]^{-1}. \tag{29}$$

If the parameters of the problem are such that we can employ approximate formulas from Sec. 3.3 (ignoring the dependence of η and N on s and ρ), it then follows from Eq. (22) that the $\rho(\omega)$ curve has a Lorentzian profile of half-width

$$\omega_{1/2} = \frac{1}{t_s}\left\{\frac{(Nt_s+1)(1-\eta) + \xi^2\eta(N+\gamma)t_s}{1-\eta}\right\}^{1/2} \tag{30}$$

At low pump intensities, as long as $(1-\eta) \gg \xi^2 Nt_s$, we have $\omega_{1/2} = t_s^{-1}$ whereas in the intermediate range when $(1-\eta) \ll \xi^2 Nt_s$ but $Nt_s \ll 1$, we find that $\omega_{1/2} = \xi(\gamma t_s)^{-1/2})(g-1)$ and this quantity depends linearly on the light intensity g. Finally, at high pumping rates ($g \to \infty$), when $Nt_s \gg 1$, we obtain $\omega_{1/2} = \xi g^{1/2}\gamma^{-1/2}$

4. DISCUSSION

A comparison of the theoretical and experimental results is made in Figs. 1 and 2. The continuous curves represent the results of calculations made using the exact formulas (7)–(11) and (13). In calculating $\rho(g)$ and $N(g)$ we used the following parameters: $s_i = 0.09$, $\gamma = 0.025$, $t_s = 0.4$ (Fig. 1a); $s_i = 0.12$, $\gamma = 0.1$, $t_s = 0.4$ (Fig. 1b). Finally, the dependence of the half-width of the Hanle curve on the pumping rate was determined for $s_i = 0.15$, $\gamma = 0.2$, $t_s = 0.1$, and a g factor of carriers amounting to 25.

We can see that the theoretical model explains the main qualitative features of the behavior of the electron–photon system, i.e., the increase in the luminescence polarization and the half-width of the Hanle curve as a function of the pump power. In full agreement with the experimental results, it follows from the theoretical calculations that defocusing the pump beam (i.e., increasing the parameter γ) weakens the maximum in the dependence of the polarization luminescence on g [see Eq. (23)]. Even for $\gamma/t_s = 0.25$ (curve 2 in Fig. 1b) this maximum is practically undetectable. A satisfactory agreement between the theory and experiment is obtained for reasonable values of the parameters.

However, this comparison of the theoretical and experimental results also reveals certain quantitative discrepancies. The most obvious observation is that the initial spin s_i of the oriented electrons is in fact not calculated, but is a phenomenological parameter of the theory. In the optical orientation experiments it is indeed assumed that $s_i = s_{cr} = \xi/2$. This result follows directly from the selection rules for optical transitions. However, it is clear from Figs. 1 and 2 that a satisfactory agreement between the theory and experiment is obtained for $s_i \ll \xi/2$. Therefore, the spin of the carriers reaching energy levels participating in lasing is considerably less than the average spin of carriers at the moment of their creation. Since the luminescence photon energy $\hbar\omega_L \approx 190$ meV is considerably less than the energy of an electron–hole pair created by two-photon absorption $2\hbar\omega_p \approx 234$ meV, it is natural to assume that the oriented particles lose much of their polarization during the energy relaxation time. A more comprehensive theory of optical orientation in PbTe crystals should allow also for this process.

In spite of these discrepancies the phenomenological model proposed above describes quite fully the orientation characteristics of a system of electrons and photons under lasing conditions. Although the electron spin is not measured directly in these experiments, there is no doubt about the theoretical conclusion that for $t_s \ll 1$ the projection of the carrier spin along the direction of propagation of laser modes interacting with carriers always remains small compared with s_i.

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[1)] At this photon energy of the pump radiation the kinetic energy of the created particles (both electrons and holes) is about 20 meV.

[1] M. I. D'yakonov, I. A. Merkulov, and V. I. Perel', Izv. Akad. Nauk SSSR Ser. Fiz. **46**, 482 (1982); F. Meier and B. P. Zakharchenya (eds.), *Optical Orientation*, North-Holland, Amsterdam (1984) [Modern Problems in Condensed Matter Sciences, Vol. 8].

[2] A. M. Danishevskiĭ, Fiz. Tverd. Tela (Leningrad) **20**, 3150 (1978) [Sov. Phys. Solid State **20**, 1818 (1978)].

[3] A. M. Danishevskiĭ, Doctoral Thesis [in Russian], Leningrad (1986).

[4] B. P. Zakharchenya, M. S. Bresler, O. B. Gusev, and I. A. Merkulov, Dokl. Akad. Nauk SSSR **297**, 584 (1987) [Sov. Phys. Dokl. **32** (to be published)].

[5] V. M. Galitskiĭ and V. F. Elesin, *Resonant Interaction of Electromagnetic Fields with Semiconductors* [in Russian], Énergoatomizdat, Moscow (1986).

[6] W. E. Lamb Jr., Phys. Rev. **134**, A1429 (1964).

[7] E. L. Ivchenko, Fiz. Tverd. Tela (Leningrad) **14**, 3489 (1972) [Sov. Phys. Solid State **14**, 2942 (1973)].

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