Instability of electron-hole semiconductor plasma in a strong magnetic field following interband absorption of light

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We investigate the macroscopic instabilities of a semiconductor electron-hole plasma in a quantizing magnetic field, obtained as a result of monochromatic optical pumping. It is shown that the character of the instability depends on the ratio of the carrier lifetime τ_R to the energy relaxation time τ_ϵ. At $\tau_R < \tau_\epsilon$ the instability is due to the absolute negative conductivity of one of the types of carrier across the magnetic field. In the absence of an external electric field the instability can set in if the ambipolar diffusion coefficient of the carriers is negative. If the ambipolar diffusion coefficient is positive, the instability threshold in an external electric field is determined by the value of this coefficient. If $\tau_R > \tau_\epsilon$, then instability due to the decrease of the rate of carrier generation and to the increase of the electric field can set in under conditions when magneto-optical oscillations of the absorption coefficient are observed in crossed electric and magnetic fields.

1. The present paper is devoted to a study of the instabilities of an electron-hole plasma of semiconductors under conditions of interband absorption of light in strong magnetic fields, when the carrier spectrum is quasidiscrete, i.e., when the conditions

$$\omega_c, \quad \omega_c m/M \gg \nu \tag{1}$$

are satisfied, where ω_c is the cyclon frequency of the electrons, m and M are the effective masses of the electrons and holes, and ν is the characteristic frequency of the collisions with the impurities and the phonons. The possibility of such instabilities is due to features of the density of states when relations (1) are satisfied, and to the specifics of the energy distribution of the carriers.

The energy distribution of the carriers generated in interband transitions depends essentially on the ratio of the carrier lifetime τ_R to the energy relaxation time τ_ϵ. For electrons in a number of semiconductors, for example in InSb, it is possible to have both $\tau_R > \tau_\epsilon$ and $\tau_R < \tau_\epsilon$, depending on the doping and on the temperature [1,2]. At $\tau_R < \tau_\epsilon$ the energy distribution of the photoelectrons can deviate strongly from equilibrium. The last condition is usually impossible to satisfy for holes, and therefore the energy distribution of the holes remains close to quasiequilibrium.

Under monochromatic illumination, the energy distribution of the photoelectrons, if their lifetime is short in comparison with the energy relaxation time, is close to a δ-function. In strong magnetic fields, this photoelectron distribution leads to an absolute negative conductivity (ANC) of the photoelectrons transversely to the magnetic field [2-4].

However, a monoenergetic (δ-like) electron distribution is generally speaking unstable [5]. The reason why such a distribution can be unstable is that the system of monoenergetic electrons constitutes an aggregate of opposing electron streams. The presence of such streams, if their intensity is high enough, can lead to two-stream instability. This instability greatly broadens the electron distribution function and this in turn can cause the ANC to vanish in a transverse electric field.

If the concentration of the strongly non-equilibrium photoelectrons is high enough, so that their plasma frequency exceeds the plasma frequency of the thermalized carriers (in particular, holes), the monoenergetic distribution of the photoelectrons is stable if

$$\omega_p^2 < \max\{\nu\nu_R, \nu_R^2\}, \tag{2}$$

where ω_p is the plasma frequency of the photoelectrons and $\nu_R = \nu_R^{-1}$. The stabilization conditions become much weaker at high concentration of the quasiequilibrium carriers. The criterion of the stability assumes in this case the form

$$\omega_{pT}(n/n_T)^{1/2} < \nu, \tag{3}$$

where ω_{pT} is the plasma frequency of the thermalized carriers, n_T is their density, and n is the concentration of the photoelectrons.

We assume that the photoelectron distribution is monoenergetic at $\tau_R < \tau_\epsilon$, assuming by the same token that the stabilization conditions are satisfied. In this case, as already noted, the transverse conductivity connected with the photoelectrons is negative.

If the absolute value of the contribution of the photoelectrons to the conductivity exceeds the contribution of the holes, then the homogeneous state of the electron-hole plasma turns out to be unstable because of the negative total conductivity [6]. However, a homogeneous carrier distribution may turn out to be unstable also if the contribution of the photoelectrons to the conductivity is relatively small, so that the total conductivity of the electron-hole plasma is positive. In the absence of an external electric field, this instability is possible if the coefficient of ambipolar diffusion is negative [7]. The homogeneous distribution is unstable in this case against quasi-neutral perturbations of the electron and hole density. The reason for this instability lies in the possibility of enhancement of the density fluctuations as a result of the negative ambipolar diffusion of the carriers.

On the other hand, if the coefficient of ambipolar diffusion is positive, then the instability is possible only in the presence of an external electric field. The instability threshold, as will be shown below, is determined in this case by the value of the ambipolar diffusion coefficient. This instability in an external field is analogous to the well known instability in a system with two types of carriers, one of which has negative differential conductivity [8]. However, unlike the latter, this instability has practically no threshold.

It should be noted that such a situation has, besides a lower threshold, also an upper threshold value of the electric field. The presence of the upper threshold is due to the fact that in strong electric fields the conductivity connected with the non-equilibrium photoelectrons becomes positive. This occurs when the effective energy relaxation time, which depends on the electric field, becomes comparable with the lifetime of the nonequilibrium carriers, as a result of which the monoenergetic distribution becomes much broader and the electrons populate the bottom of the Landau subbands.

The electric field at which the photoelectron mobility reverses sign can be estimated in the following manner: In an electric field, the effective energy relaxation time is

$$\tau_\epsilon^{\text{eff}} = \tau_\epsilon (\bar{E}/E)^2, \quad (4)$$

where $\bar{E}^2 = \frac{1}{2}(sH/c)^2(1 + \nu_a/\nu_i)^{-1}$, and ν_a are respectively the frequencies of the interaction with the impurities and with the acoustic phonons, and s is the speed of sound. Assuming that $\tau^{\text{eff}} \sim \tau_R$, we get for the critical field

$$E_{\text{cr}} \sim \bar{E}(\tau_\epsilon/\tau_R)^{1/2}. \quad (5)$$

In InSb we have at helium temperatures $\bar{E} \sim 0.1\text{–}1$ V/cm [9]. For $\tau_\epsilon \sim 10^{-7}\text{–}10^{-9}$ sec [10,11] and $\tau_R \sim 10^{-10}$ sec [2] we obtain $E_{cr} \sim 3\text{–}30$ V/cm, which agrees well with the experimental data [2].

Finally, in a transverse electric field, in the case of interband carrier generation, another type of instability, also due to negative differential conductivity of the electron-hole plasma, becomes possible. This instability takes place when the carrier lifetime is long in comparison with the relaxation time of their energy. This situation is realized in practice much more frequently. The photoelectron distribution function does not differ from Maxwellian in this case, at least qualitatively, and the conductivity via the photoelectrons is positive.

The physical nature of this instability consists in the following: It is well known [12] that the interband absorption coefficient of light in quantizing magnetic fields is an oscillating function of the frequency of the light and on the magnetic field intensity. These oscillations are due to the presence of singularities in the density of states when relations (1) are satisfied. The oscillations of the absorption coefficients lead to oscillations of different physical quantities, and primarily to oscillations of the carrier density and to ensuing oscillations of the photoconductivity. In crossed electric and magnetic fields the oscillation pattern shifts towards lower frequencies or lower magnetic fields [13]. In this case the rates of photoelectron and photohole generation become functions of the electric field. If the light frequency Ω is chosen such that $d\alpha/d\Omega < 0$ (α is the interband absorption coefficient), then the generation rate, and consequently also the carrier density, will decrease with increasing electric field. Under these conditions it is natural to expect the appearance of an instability analogous to recombination instability [14].

2. We proceed to a quantitative description of the instabilities considered above. As the initial model we consider an intrinsic semiconductor in crossed electric and magnetic fields. The extrema of the valence band and the conduction band of the semiconductor are located at one point of the Brillouin zone, so that the principal role in the photon absorption is played by direct interband transitions. We assume the condition (1) of quasi-discreteness of the carrier spectrum to be satisfied. For simplicity we assume that the dispersion for both bands is quadratic and isotropic. The latter assumption imposes a limitation on the electric field intensity. At not too large magnetic quantum numbers, this limitation takes the form [15–17]

$$\tfrac{1}{2}(m+M)c^2(E/H)^2 \ll \Delta_g, \quad (6)$$

where c is the speed of light and Δ_g is the width of the forbidden band. We also assume that the light intensity is not too large, so that the restructuring of the energy spectrum, due to frequent interband transitions, can be neglected (see, e.g., [18]).

Unless otherwise stipulated, we neglect thermal carriers. As to the energy distribution of the photoelectrons, we assume that it can be both monoenergetic and close to equilibrium, depending on the relation between the energy relaxation time and the recombination time. The holes, on the other hand, will always be assumed to be in quasi-equilibrium.

To describe processes whose characteristic frequencies are not too large[1)], if we confine ourselves to potential oscillations, we can use the continuity equations for the electrons and holes and the Poisson equation. These equations, assuming all the quantities to depend only on the coordinates transverse to the magnetic field, take the following form:

$$\frac{\partial n}{\partial t} - \frac{\partial}{\partial x_i}\left\{ n\mu_{ik}^{(n)}E_k + \frac{\partial}{\partial x_k}(D_{ik}^{(n)}n)\right\} = g(E) - \frac{n}{\tau_R}, \quad (7)$$

$$\frac{\partial p}{\partial t} + \frac{\partial}{\partial x_i}\left\{ p\mu_{ik}^{(p)}E_k - \frac{\partial}{\partial x_k}(D_{ik}^{(p)}p)\right\} = g(E) - \frac{p}{\tau_R}, \quad (8)$$

$$\frac{\partial E_i}{\partial x_i} = \frac{4\pi e}{\varkappa}(p-n). \quad (9)$$

Here n, $\mu_{ik}^{(n)}$, $D_{ik}^{(n)}$ and p, $\mu_{ik}^{(p)}$ $D_{ik}^{(p)}$ are respectively the concentrations, mobilities, and diffusion coefficients of the electrons and of the holes, $x_i = \{x, y\}$,

$$\mu_{xx}^{(n,p)} = \mu_{yy}^{(n,p)} = \mu_{n,p}, \quad \mu_{xy}^{(n,p)} = -\mu_{yx}^{(n,p)}, \quad D_{xx}^{(n,p)} = D_{yy}^{(n,p)} = D_{n,p}, \quad D_{xy}^{(n,p)} = -D_{yx}^{(n,p)}$$

and κ is the dielectric constant of the lattice. The mobilities and diffusion coefficients that enter in Eqs. (7) and (8) depend, generally speaking, on the electric field. We shall neglect this dependence, however, unless otherwise stipulated.

The term g(E) describes the generation of the carriers. If it is assumed that the carrier generation is due to transitions between zero-level subbands of the Landau valence band and conduction band, then

$$g(E) = \frac{N}{\tau_R}\rho\left(\frac{\Delta(E)}{\omega_c}\right), \quad (10)$$

where N is a quantity proportional to the radiation intensity, $\rho(z)$ is the dimensionless density of states in the zero-level Landau subband of the conducting band, $\Delta(E)/\omega_c = \delta + E^2/\tilde{E}^2$, and

$$\delta = \frac{\Omega}{\omega_c}\left(1 - \frac{\Delta_g}{\hbar\Omega}\right) - \frac{1}{2}\left(1 + \frac{m}{M}\right), \quad \tilde{E}^2 = \frac{2H^2\hbar\omega_c}{(m+M)c^2}.$$

In the absence of collision damping of the electron spectrum we have $\rho(z) \sim z^{-1/2}$. In the presence of this damping, the square-root singularity becomes somewhat smoother. Owing to the conditions (1), however, this singularity is quite strong, and the maximum absolute value of the ratio $(d\rho/dz)/\rho$ turns out to be of the order of $\omega_c/\nu \gg 1$.

We linearize Eqs. (7) and (8) relative to the perturbations of the homogeneous distribution. Then, taking Eq. (9) into account we obtain for perturbations of the type $\exp[i(kx - \omega t)]$ (the x axis is directed along the external electric field) the following dispersion equation

$$1=\frac{G/k-i\lambda_p}{\omega-ku_p+i(\nu_R+D_pk^2)}-\frac{G/k+i\lambda_n}{\omega+ku_n+i(\nu_R+D_nk^2)}. \quad (11)$$

Here

$$G=\frac{4\pi e}{\varkappa}\frac{dg(E)}{dE},\quad \lambda_{n,p}=\frac{4\pi eg(E)\tau_R}{\varkappa}\mu_{n,p},$$
$$u_{n,p}=\mu_{n,p}E$$

are the reciprocal Maxwellian relaxation times and drift velocities of the electrons and holes. We note that the dispersion equation (11) does contain explicitly the Hall parts of the mobilities. The reason is that we are confining ourselves to perturbations with wave vectors directed along the external electric field, i.e., directed across the drift of the plasma as a whole. In the more general case, when the projection of the wave vector k_y on the direction of the Hall drift is different from zero, it is necessary to replace $ku_{n,p}$ in (11) by $\mathbf{k}\cdot\mathbf{u}_{n,p}$, where the drift velocities $\mathbf{u}_n$ and $\mathbf{u}_p$ now depend both on the diagonal and on the Hall components of the mobility tensor. This, however, leads to a change in the real part of the frequency of the oscillations that build up. In addition, a term $G\cos\theta/k$ appears in (11) instead of G/k (θ is the angle between the wave propagation direction in the electric field). This dependence leads to higher thresholds of the considered instabilities with respect to the electric field for oblique waves ($\theta \neq 0$).

3. We consider first the situation in the absence of an external electric field. It follows from (10) that in this case $G = 0$. There is likewise no carrier drift, $u_n = u_p = 0$. The dispersion equation (11) has in this case the following roots:

$$\omega=-{}^1/_2 i[a\pm\sqrt{a^2-b}], \quad (12)$$

where

$$a=(2\nu_R+\lambda_n+\lambda_p)+(D_n+D_p)k^2,$$
$$b=4\{\nu_R(\nu_R+\lambda_n+\lambda_p)+[\nu_R(D_n+D_p)+(\lambda_nD_p+\lambda_pD_n)]k^2+D_nD_pk^4\}.$$

If the total conductivity of the electron-hole plasma is negative, then it follows from (12) that the homogeneous carrier distribution is unstable if

$$\lambda_n+\lambda_p<-\nu_R. \quad (13)$$

This is the known instability due to the ANC [6]. It has an aperiodic character ($\mathrm{Re}\,\omega = 0$). Notice should be taken of the stabilizing role of recombination, which suppresses the instability if the ANC has a sufficiently small absolute value.

More easy to realize in practice, however, is the situation in which the total conductivity of the system is positive, inasmuch as in strong magnetic fields, even at equal carrier densities, the mobility of the (heavy) holes has a larger absolute value than the mobility of the electrons.

In this case we obtain from (12) the following instability condition:

$$\lambda_nD_p+\lambda_pD_n<-\nu_R\left\{(D_n+D_p)+2\left[D_nD_p\left(1+\frac{\lambda_n+\lambda_p}{\nu_R}\right)\right]^{1/2}\right\}. \quad (14)$$

The increasing perturbations in this case are those whose length satisfies the condition

$$-\alpha-\sqrt{\alpha^2-\beta}<k^2<-\alpha+\sqrt{\alpha^2-\beta}, \quad (15)$$

where

$$\alpha=\frac{(\lambda_nD_p+\lambda_pD_n)+\nu_R(D_n+D_p)}{2D_nD_p},\quad \beta=\frac{\nu_R(\lambda_n+\lambda_p+\nu_R)}{D_nD_p}$$

Since the right-hand side of the condition (14) is negative, for the existence of instability it is necessary to have at least

$$\lambda_nD_p+\lambda_pD_n<0. \quad (16)$$

At large photoelectron and photohole concentrations ($\lambda_n + \lambda_p \gg \nu_R$), the condition (16) is in fact an exact instability condition, since the right-hand side of (14) is close to zero in this case.

The inequality (16) is none other than the condition that the ambipolar diffusion coefficient be negative

$$D_A=(\lambda_nD_p+\lambda_pD_n)/(\lambda_n+\lambda_p). \quad (17)$$

Thus, the physical nature of the considered instability lies in the possibility of negative ambipolar diffusion of the carriers into the region[2)] of quasineutral fluctuation of the carrier density. The quasineutrality of the stable perturbations can be easily verified directly from the dispersion equation (11). Indeed, neglecting in (11) the unity term in comparison with the right-hand part, we obtain

$$\omega=-i(\nu_R+D_Ak^2), \quad (18)$$

from which it follows that the unstable perturbations are those for which

$$k^2>-\nu_R/D_A, \quad (19)$$

which coincides with the lower limit of the instability region (15) at $\lambda_n + \lambda_p \gg \nu_R$.

The condition (16) can be conveniently represented in the form

$$\mu_n/D_n+\mu_p/D_p<0. \quad (20)$$

Since the holes are thermalized, they satisfy the Einstein relation. On the other hand, the calculation of μ_n/D_n under the assumption that the photoelectrons populate only the lower Landau subband, yields

$$\frac{\mu_n}{D_n}=-\frac{e}{2\varepsilon_0}\zeta, \quad (21)$$

where ϵ_0 is the energy of the produced photoelectrons. The dimensionless factor ζ is determined by the broadening of the photoelectron distribution function. For a monoenergetic distribution we have $\zeta = 1$. Taking the Einstein relations for holes and expression (21) into account, we obtain from (20) the following instability criterion:

$$\varepsilon_0<{}^1/_2\zeta T. \quad (22)$$

Here T is the effective hole temperature (which generally speaking differs from the lattice temperature).

4. We now proceed to investigate the instabilities of a homogeneous carrier distribution in an external electric field. We first neglect the diffusion. Equating the real and imaginary parts of (11) separately to zero at the stability limit, we obtain

$$v_{ph}=\frac{\mathrm{Re}\,\omega}{k}=\frac{\nu_R(u_p-u_n)}{2\nu_R+\lambda_p+\lambda_n} \quad (23)$$

(v_{ph} is the phase velocity of the growing waves) and the following instability criterion:

$$\frac{G(u_n+u_p)+\nu_R(\lambda_n+\lambda_p+\nu_R)}{(\nu_R+\lambda_n)(\nu_R+\lambda_p)}<0. \quad (24)$$

It follows from (23) that at high carrier densities ($\lambda_n + \lambda_p \gg \nu_R$) the phase velocity of the growing waves is much smaller than the drift velocity of the carriers. At $\mu_n = \mu_p$ we have $v_{ph} = 0$, i.e., the perturbations increase aperiodically.

The instability criterion (24) is satisfied in two cases: a) either if

$$G(u_n+u_p)+\nu_R(\lambda_n+\lambda_p+\nu_R)>0,\quad (\nu_R+\lambda_n)(\nu_R+\lambda_p)<0 \tag{25}$$

b) or if

$$G(u_n+u_p)+\nu_R(\lambda_n+\lambda_p+\nu_R)<0,\quad (\nu_R+\lambda_n)(\nu_R+\lambda_p)>0. \tag{26}$$

The conditions (25) can be satisfied if the conductivity via the photoelectrons is negative ($\lambda_n < -\nu_R$). In this case, even if $G < 0$, instability is possible (if $\lambda_n + \lambda_p \gg \nu_R$) up to fields

$$E<E_{cr}^*=H[\hbar\nu/(m+M)c^2]^{1/2}. \tag{27}$$

An estimate of the critical field for InSb at $H \sim 10^{12}$ Oe, $M \sim 5\times 10^{-28}$ g (heavy holes), and $\nu \sim 10^{11}-10^{12}$ sec^{-1} yields $E_{cr}^* \sim 50-150$ V/cm. If this estimate is compared with the estimate of the critical field at which the dependence of the mobility of the photoelectrons on the electric field becomes significant (λ_n reverses sign) (see formula (5)), then it can be easily seen that the main cause of the stabilization of the instability in strong electric fields is the reversal of the sign of the photoelectron mobility, and not the decrease in the rate of carrier generation with increasing electric field. This is all the more correct when the rate of carrier generation increases with increasing field ($G > 0$).

It follows from (25) that the instability in question has no threshold, but this is true only if diffusion is neglected. If diffusion is treated in the dispersion equation (11) as a perturbation, then the following instability criterion results:

$$u_n+u_p>\left[-\frac{\nu_R(D_n+D_p)+(\lambda_n+\lambda_p)D_A}{(\nu_R+\lambda_n)(\nu_R+\lambda_p)}\right]^{1/2}(2\nu_R+\lambda_n+\lambda_p). \tag{28}$$

In the case when $D_A > 0$ and $\lambda_n + \lambda_p \gg \nu_R$ we obtain from (28) an expression for the threshold field

$$E_{\text{thr}}=\left[\frac{4\pi e g\tau_R}{\varkappa}\left|\frac{1}{\mu_n}+\frac{1}{\mu_p}\right|D_A\right]^{1/2}. \tag{29}$$

At $D_A < 0$, the instability has no threshold and constitutes, in the absence of an electric field, the already considered ambipolar instability.

As noted above, the considered instability in an external electric field is analogous to the instability in a system with two types of carrier, one of which has negative differential conductivity [8]. However, unlike the latter, the threshold of this instability is quite small. The instability sets in immediately as soon as the electric field reaches a value sufficient to produce an independent drift of electrons and holes.

Finally, as follows from (26), the homogeneous distribution of the carriers may turn out to be unstable also in the case when the conductivity of both types of carrier is positive ($\lambda_n > 0$, $\lambda_p > 0$), a situation realized in practice much more frequently. This is possible if the generation rate, and consequently also the carrier density, decreases with increasing electric field ($G < 0$). The instability criterion (26) then takes the form

$$\frac{d\ln g}{d\ln E}+1+\frac{\nu_R}{\lambda_n+\lambda_p}<0. \tag{30}$$

In the case of large carrier densities ($\lambda_n + \lambda_p \gg \nu_R$) it follows from (30) that such an instability is possible in electric fields $E > E_{cr}^*$, where E_{cr}^* is given by (27).

In the case of large carrier densities this instability, just as ambipolar instability, has a quasineutral character. Indeed, for quasineutral perturbations, the dispersion equation (11) has the following solution:

$$\omega=-i\nu_R\left[1+\frac{d\ln g}{d\ln E}\right],$$

from which it follows that if $\lambda_n + \lambda_p \gg \nu_R$ the condition for the buildup of the instability coincides with the criterion (30).

It should be noted that, owing to the rather high threshold of such a magnetooptical instability, the dependence of the mobilities and of the diffusion coefficients of the carriers on the electric field cannot be neglected when this instability is considered. Since the diffusion has practically no effect on the instability threshold (the criterion (30) was derived neglecting diffusion), it suffices to take into account the dependence of the mobility on the electric field. The mobility of the electrons and holes in the dispersion equation (11) should then be replaced by the differential mobilities

$$\mu_n^{\text{dif}}(E)=\partial u_n(E)/\partial E,\quad \mu_p^{\text{dif}}(E)=\partial u_p(E)/\partial E,$$

where $u_n(E)$ and $u_p(E)$ are the drift velocities of the electrons and holes.

In the case of carrier heating by the electric field, the dependence of the differential mobility on the electric field is determined by the mechanisms for the scattering of the momentum and energy of the carriers [19]. In particular, in the case of scattering by impurities and acoustic phonons, the differential carrier mobility is always positive in the classical limit ($\hbar\omega_c < T$, where T is the effective carrier temperature) and is negative in fields $E > \sqrt{2}\,\bar{E}$ ($\bar{E}$ is given by (4)) in the quantum limit ($\hbar\omega_c > T$) [20].

At $\mu_n^{\text{dif}}(E) > 0$ and $\mu_p^{\text{dif}}(E) > 0$ the analysis of the instability does not differ, at least qualitatively, from that given above. In particular, the instability criterion (30) takes in this case the form

$$\frac{dg}{dE}\frac{q(E)}{g}+1+\frac{\nu_R}{\lambda_n+\lambda_p}<0,$$

where

$$q(E)=\frac{u_n(E)+u_p(E)}{\mu_n^{\text{dif}}(E)+\mu_n^{\text{dif}}(E)}$$

The case when the differential mobility of one type of carrier is negative[3)] reduces to the previously considered instability of a system with negative photoelectron mobility in an external electric field. The instability threshold is determined in this case by the relation (if the differential electron mobility is negative)

$$g(E)\mu_n^{\text{dif}}(E)<-\frac{\varkappa}{4\pi e}\nu_R^2.$$

We note that the mechanisms that lead to the instability connected with the negative mobility or to the magnetooptical instability tend, as follows from the criteria (25) and (26), to stabilize each other.

In conclusion we note that the presence of thermal carriers (holes) exerts a stabilizing influence on the instabilities considered above. Thus, for example, the criterion of ambipolar instability (20) with allowance for the thermal carriers, takes the form

$$\mu_n/D_n<-\eta\mu_p/D_p,$$

where $\eta = 1 + n_T/g\tau_R$, and n_T is the density of the thermal holes. In the case of magnetooptical instability, the sensitivity threshold increases by a factor $\sqrt{\eta}$.

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[1)]It can be assumed that the expressions obtained below for the increments are valid up to the momentum-relaxation frequency.

[2)]In the case of negative fluctuations, the carriers will diffuse from the fluctuation region.

[3)]This situation is possible, for example, for carrier scattering by impurities and acoustic phonons if $\hbar\omega_c > T_n$ and $\hbar\omega_c m/M < T_p$, where T_n and T_p are the effective temperatures of the electrons and holes.

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