Generation of magnons by carriers in magnetic semiconductors

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An analysis is made of the Cherenkov process of magnon generation by carriers in ferromagnetic and antiferromagnetic semiconductors. Attention is drawn to the different nature of the magnon emission processes in these two classes of substances. In ferromagnetic semiconductors a strong spin splitting of the electron energy bands in the molecular field is responsible for the threshold nature of the Cherenkov (three-particle) interaction of conduction electrons with magnons. Therefore, processes of this kind are activated by the heating of carriers in a strong static electric field. In antiferromagnetic semiconductors, in which there is no spin splitting of the energy bands in the absence of a magnetic field, three-particle electron-magnon collisions do not have a threshold and magnon emission is due to cold electrons.

1. INTRODUCTION

Basically new properties of magnetic semiconductors, which distinguish them from ordinary semiconductors and from magnetically ordered insulators, result from strong exchange interaction between carriers (electrons and holes) and the spin-ordered subsystem (s-d exchange interaction[1]). Therefore, the type and nature of the spin system may have a significant influence on the nature of the s-d interaction and, consequently, on the transport properties of magnetic semiconductors. We shall analyze the Cherenkov generation of magnons by carriers in ferromagnetic and antiferromagnetic semiconductors.[2] In the case of ferromagnetic semiconductors the threshold processes of the Cherenkov interaction of electrons[1)] with magnons are activated by the heating of an electron gas in a strong static electric field. In the case of antiferromagnetic semiconductors the process of Cherenkov generation of magnons can occur even when electrons are cold.

In wide-gap magnetic semiconductors the microscopic Hamiltonian of the s-d exchange interaction[1] can be simplified greatly at low temperatures. This is due to the fact that the exchange interaction is of the short-range nature with its radius of the order of the lattice constant a, and the de Broglie wavelength of electrons is long because the electron density is low ($\lambda_B \gg a$), so that the exchange integral contained in the s-d Hamiltonian may be replaced with the delta function. It is then found that the s-d interaction in typical semiconductors can be described within the framework of the molecular field model with the Hamiltonian

$$\mathscr{H}_{s-d} = -\sum_j A_j \boldsymbol{\sigma} \mathbf{M}_j, \tag{1}$$

where $j = 1,2,...$ is the index which labels the sublattices; A_j are the exchange constants, $\mathbf{M}_j$ is the macroscopic magnetization of the sublattices; $\boldsymbol{\sigma}$ are the Pauli spin matrices.

Since we shall be concerned mainly with magnons having small quasimomenta ($aq/\hbar \ll 1$), it follows that a macroscopic description of the spin system by means of the magnetization vectors $\mathbf{M}_j(r)$ is also valid.

We can show that in magnetic semiconductors with electron densities in the range $n < 10^{19}$ cm^{-3} we can ignore the influence of the s-d interaction on the spin system, so that the magnitude and the equilibrium conditions of the magnetic moments of the sublattices in the semiconductor subjected to a magnetic field are identical to those in a magnetic insulator. This is due to the low density of carriers compared with the d or f electrons, whose density is of the order of one per atom: $n_d \propto a^{-3}$. Therefore, the additional terms describing the influence of the s-d interaction on the spin system contain a small factor na^3. For example, such a factor is included in the correction to the dispersion law of magnons because of the s-d exchange in ferromagnetic semiconductors.[3,4]

Heavily doped magnetic semiconductors constitute a topic by itself. In this case we can expect a number of important and interesting effects which cannot be allowed for on the basis of the model proposed above. These effects include, for example, the change in the orientation of the magnetic moments of the sublattices as a result of doping,[5] the giant Zeeman effect,[4] and several others.

Since $\mathbf{M}_j = \mathbf{M}_{0j} + \mathbf{m}_j$, where $\mathbf{M}_{0j}$ is the equilibrium magnetization and $\mathbf{m}_j$ is an alternating correction, the molecular field acting on the electron spin consists—in accordance with Eq. (1)—of static and alternating components. The static component of the molecular field lifts the spin degeneracy of carriers and splits the conduction band (as well as the valence band) into two subbands with different spin orientations. According to Eq. (1), in the case of ferromagnetic semiconductors with one magnetic sublattice this splitting is fairly large ($\Delta = 2AM_0$) and for typical materials it amounts to several tenths of an electron volt. However, in the case of antiferromagnetic semiconductors, in view of the full identity of the sublattices ($A_1 = A_2$), the constant component of the molecular field in weak magnetic fields is either identically equal to zero for substances of the easy axis type [$\mathscr{H}^0_{s\text{-}d} = -A_1\boldsymbol{\sigma}(\mathbf{M}_{10} + \mathbf{M}_{20}) \equiv 0$ if $H < H_{EA}$, where H_{EA} is the characteristic spin flopping field] or it is fairly small for substances of the easy plane type ($\mathscr{H}^0_{s\text{-}d} = -2A_1\sigma_z M_0 \cos\theta$, where $\cos\theta = H/H_E \ll 1$, and H_E is the exchange field). The case of strong magnetic fields is not of special interest, since it gives rise to a strong noncollinearity of the sublattice moments and, consequently, to a large spin splitting of the carrier energy and is essentially analogous to the case of ferromagnetic semiconductors.

The alternating component of the s-d Hamiltonian of

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Eq. (1) describes the interaction of electrons with magnons. Therefore, a very important circumstance is that the alternating component of the molecular field in antiferromagnetic semiconductors differs from the constant component, because it never vanishes and reaches a considerable value. Physically, this is due to the fact that in the course of oscillations the magnetic moments of the sublattices deviate from their equilibrium antiparallel orientations and this gives rise to an alternating ferromagnetic moment which is proportional to H_{EA}/H_e (Refs. 6 and 7), where $H_{EA} = (H_E H_A)^{1/2}$ and H_A is the anisotropy field.

2. AMPLITUDES OF THE INTERACTION OF ELECTRONS WITH MAGNONS

In writing down the *s-d* Hamiltonian in the second quantization representation it is necessary to expand, in terms of the eigenfunctions of free electrons, the electron wave functions in the expression for the average Hamiltonian (1) and to express the components of the alternating magnetic moment in terms of the magnon creation and annihilation operators. The ability to expand in terms of the proper wave functions of free electrons arises from two circumstances. Firstly, in the case of wide-gap semiconductors the *s-d* interaction of Eq. (1) can be regarded as a small and slowly varying perturbation in terms of the crystal fields and, therefore, it is possible to use the effective mass method. Secondly, if the spin-orbit interaction is ignored (it is assumed to be small compared with the exchange interaction), the electron wave function can be represented by a product of its orbital and spin components. The Fourier components of the operator of an alternating magnetic moment $\mathbf{m}_j(\mathbf{q})$, contained in the Hamiltonian, can be expressed in terms of the Holstein-Primakoff operators $b_j^+(\mathbf{q})$ and $b_j(\mathbf{q})$ and the latter can be obtained using the Bogolyubov transformation in terms of the magnon creation $\xi_j^+(\mathbf{q})$ and annihilation $\xi_j(\mathbf{q})$ operators. The relevant procedure is well-known (see, for example, Ref. 8) and will not be given here.

The *s-d* Hamiltonian subjected to second quantization and describing the electron-magnon collisions has the following form in the lowest order of perturbation theory.

1) In the case of ferromagnetic semiconductors this Hamiltonian is[1,9]

$$\mathscr{H}_{s-d} = \sum_{\mathbf{p},\mathbf{q}} \Psi_f c_{\mathbf{p}\uparrow}{}^+ c_{\mathbf{p}+\mathbf{q}\downarrow} \xi_{\mathbf{q}}{}^+ + \text{H.c.}, \qquad \Psi_f = -A_f (2\mu M_0/V)^{1/2}, \tag{2}$$

where $c_{p\sigma}^+$ and $c_{p\sigma}$ are the electron creation and annihilation operators for electrons with a momentum p and a spin $\sigma = \uparrow, \downarrow$; V is the volume; $\mu = g\hbar$; g is the gyromagnetic ratio; $\hbar$ is the Planck constant.

2) The corresponding expression for antiferromagnetic semiconductors of the easy axis type is [10]

$$\mathscr{H}_{s-d} = \sum_{\mathbf{p},\mathbf{q}} \Psi_{af}(\mathbf{q}) c_{\mathbf{p}+\mathbf{q}\uparrow}^+ c_{\mathbf{p}\downarrow} (\xi_{1-\mathbf{q}}^+ + \xi_{2\mathbf{q}}) + \text{H.c.}, \tag{3}$$

where

$$\Psi_{af}(\mathbf{q}) \approx -A_{af}\left(\frac{2\mu M_0}{V}\frac{\hbar\omega_{\mathbf{q}}(0)}{\mu H_E}\right)^{1/2}, \qquad H<H_{EA}, \tag{3a}$$

$$\hbar\omega_{\mathbf{q}}(0) = (\Theta_N{}^2 a^2 q^2/\hbar^2 + \mu^2 H_{EA}^2)^{1/2}$$

is the dispersion law of magnons in the absence of a magnetic field and Θ_N is a parameter of the order of the Néel temperature.[6,7]

3) For antiferromagnetic semiconductors of the easy plane type, we have

$$\mathscr{H}_{s-d} = \sum_{\mathbf{p},\mathbf{q}} \{\Psi_{af}'(\mathbf{q}) c_{\mathbf{p}+\mathbf{q}\uparrow}^+ c_{\mathbf{p}\downarrow} \xi_{1\mathbf{q}} + \Psi_{af}''(\mathbf{q}) c_{\mathbf{p}+\mathbf{q}\sigma}^+ c_{\mathbf{p}\sigma} \xi_{2\mathbf{q}}\} + \text{H.c.} \tag{4}$$

In general, the expressions for the amplitudes of a transition in the case of easy plane semiconductors are fairly cumbersome. Therefore, we shall only give the expressions in the limiting case of zero field:

$$\Psi_{af}'(\mathbf{q}) = -A_{af}\left(\frac{2\mu M_0}{V}\frac{\hbar\omega_{1\mathbf{q}}(0)}{\mu H_E}\right)^{1/2},$$
$$\Psi_{af}''(\mathbf{q}) = -iA_{af}\left(\frac{2\mu M_0}{V}\frac{\hbar\omega_{2\mathbf{q}}(0)}{\mu H_E}\right)^{1/2}, \tag{4a}$$

where $\hbar\omega_{1\mathbf{q}}(0)$ is identical with the dispersion law of magnons in the easy axis case and $\hbar\omega_{2\mathbf{q}}(0) = \Theta_N aq/\hbar$ is of acoustic nature if we ignore the anisotropy in the basal plane.

We shall now consider only the case of a weak magnetic field, because otherwise a field-induced noncollinear magnetic structure creates a large spin splitting of the electron spectrum and we are then essentially dealing with the case of a ferromagnet.

It follows from Eqs. (2)–(4) that the electron-magnon collisions involve the emission or absorption of a magnon by an electron and as a result of these collisions the electron spin may be reversed or may remain unchanged, as determined by the laws of conservation of the energy, momentum, and magnetic moment, since the exchange processes (in contrast to the relativistic processes) conserve the total moment of the colliding quasiparticles. It is known[11] that the magnon spin is equal to unity and that the eigenvalues of the operator of the spin projection s_z expressed in units of $\hbar$ are -1, 0, and 1, respectively. Therefore, different branches of the magnon spectrum which correspond to circular precession of localized spins with the right- and left-rotation, respectively, should have different values of the spin projections. It is also known[7] that acoustic magnon oscillations in materials of the easy plane type occur without excitation of the total ferromagnetic moment of the sublattices, so that quite naturally in the process of interaction of electrons with such magnons the electron spin remains unchanged.

In ferromagnetic semiconductors the large spin splitting of the conduction band means that the strongest three-particle electron-magnon collisions associated with the Cherenkov emission (or absorption) of a magnon by an electron are in fact normally forbidden, because of the large threshold energy Δ. Therefore, only four-particle processes of electron scattering by magnons have been considered earlier, although they describe the electron-magnon interaction in the second order of perturbation theory.[12] Activation of three-particle collisions result from heating of conduction electrons in a static strong electric field with the aim of increasing the temperature to the value of the gap Δ (Ref. 13).

In antiferromagnetic semiconductors, which do not have a constant component of the molecular field and, therefore, do not exhibit spin splitting of the electron energy bands, three-particle electron-magnon collisions are threshold-free and consequently the Cherenkov emission of magnons can be expected also in the case of cold electrons.

3. FERROMAGNETIC SEMICONDUCTORS

The substances used at present exhibit a fairly high mobility as well as a fairly high electron density. An example is an n-type semiconductor $HgCr_2Se_4$, which at $T = 4.2$ K has the mobility $\mu_e \sim 500$–1200 cm^2·V^{-1}·sec^{-1}, an effective electron mass $m^* \approx 0.3m$ (m is the mass of a free electron), a free-electron density $n \approx 10^{17} - 5\times10^{18}$ cm^{-3}, and a spin splitting $\Delta = 0.2$ eV (Ref. 14).

Applying the formula for the frequency of electron-electron collisions[15]

$$\nu_{ee}(\varepsilon) = (2^{1/2}\pi n e^4 \Lambda/\varepsilon_0^2 m^{*1/2})\,\varepsilon(\mathbf{p})^{-3/2} \quad (5)$$

[e is the electron charge, $\varepsilon(p) = p^2/2m^*$ is the dispersion law of electrons, which is assumed to be isotropic and quadratic, ε_0 is the permittivity of the medium, and Λ is the Coulomb logarithm], we can readily show that, for example, in the case of $HgCr_2Se_4$ when the heating of electrons is strong [$\varepsilon(\mathbf{p}) \gtrsim \Delta$], we find that the frequency in question is $\nu_{ee}(\Delta) \lesssim 10^{12}$ sec^{-1}. We shall show below that in this case the frequency ν_{ee} is the highest compared with the frequencies of collisions of electrons with magnons ν_{es} and with phonons ν_{ep}. Consequently, we can solve the system of transport equations for the electron distribution functions in the spin subbands $f_\sigma(\mathbf{p})$ by the method of successive approximations employing a small parameter $\nu_{es}/\nu_{ee} \ll 1$, $\nu_{ep}/\nu_{ee} \ll 1$. The transport equation for magnons is not used, because the heating of magnons is ignored as unimportant in our case. We shall assume that the electric field $\mathbf{E}$ acting on electrons is constant and homogeneous. Then, in the zeroth approximation in terms of small parameters it follows from the vanishing of the electron-electron collision integral that the distribution function of electrons is a shifted Maxwellian function, i.e.,

$$f_\sigma^{(0)} = \exp[(\mu_\sigma - \varepsilon(\mathbf{p}) + \mathbf{pu})/\Theta],$$

where μ_σ is the chemical potential in each of the subbands, u is the drift velocity, and Θ is the electron temperature. This type of solution is associated with the circumstance that under the investigated conditions ($\varepsilon \gtrsim \Delta$, $n \lesssim 5\times10^{18}$ cm^{-3}) the electron gas is not degenerate and an allowance is made only for electron-electron collisions which do not alter the spin state of the particles. Spin-changing collisions are ignored, because they occur due to relativistic spin-orbit processes and are therefore $(v/c)^4$ times weaker than the Coulomb collisions.

The parameters μ_σ, Θ, and $\mathbf{u}$ are found from the condition of solubility of the transport equations in the first approximation, which corresponds physically to conservation of the number of particles, energy, and momentum in electron-electron collisions. Under steady-state conditions in the homogeneous case ($\partial/\partial t = 0, \partial/\partial \mathbf{r} = 0$) it follows from the conditions of particle balance in each of the spin subbands that the chemical potentials are equal: $\mu_\uparrow = \mu_\downarrow$. The energy and momentum balance equations yield the following system of equations

$$e\mathbf{uE} = -(\Theta - T)[\tilde\nu_{es}(\Theta) + \tilde\nu_{ep}(\Theta)],$$
$$e\mathbf{E} = -m^*\mathbf{u}[\nu_{es}(\Theta) + \nu_{ep}(\Theta)]. \quad (6)$$

Since collisions of electrons with magnons and phonons are quasielastic, the right-hand side of the system (6) includes frequencies associated with the transfer of momentum $\nu_{es}(\Theta)$ and $\nu_{ep}(\Theta)$, and of energy $\tilde\nu_{es}(\Theta)$, $\tilde\nu_{ep}(\Theta)$.

In general, the expression for the frequencies of electron-magnon collisions is very cumbersome. Simple formulas can be obtained only in the case of strong heating of electrons when the electron temperature is considerably greater than the spin splitting ($\Delta/\Theta \ll 1$) and in particular greater than the lattice temperature. Then, apart from factors of the order of unity, we have

$$\nu_{es}(\Theta) \approx \frac{\Delta}{\hbar}\frac{\Delta}{\varepsilon_a}\left(\frac{\Theta}{\varepsilon_a}\right)^{1/2}, \quad \hbar\omega_{\mathbf{q}} \ll T,$$
$$\nu_{es}(\Theta) \approx \frac{\Delta}{\hbar}\left(\frac{\Delta}{\varepsilon_a}\right)^{1/2}\left(\frac{\Delta}{\Theta}\right)^{1/2}\frac{T}{\Theta_C}, \quad \hbar\omega_{\mathbf{q}} \gg T. \quad (7)$$

In both limiting cases we obtain

$$\tilde\nu_{es}(\Theta) \approx \frac{\Delta}{\hbar}\frac{\Delta}{\varepsilon_a}\left(\frac{\Theta}{\varepsilon_a}\right)^{1/2}\frac{m}{m_s}. \quad (8)$$

Here, $\hbar\omega_{\mathbf{q}}$ is the dispersion law of magnons; Θ_C is a parameter of the order of the Curie temperature; m_s is the effective mass of magnons; $\varepsilon_a = (1/2m)(h/a)^2$ is an energy of the order of the width of an electron band.

In the system (6) of the energy balance equations applicable to the case when $\hbar\omega_{\mathbf{q}} \gg T$ it is necessary to drop the term T compared with Θ. The expressions for the frequencies of electron-phonon collisions are found, for example, in the monograph of Bass and Gurevich[15] [see Eq. (5.15) and the equations that follow]. The frequencies $\nu_{ep}(\Theta)$ and $\tilde\nu_{ep}(\Theta)$ describe collisions of electrons with optical phonons, because at higher values of Θ such collisions are more probable than collisions of electrons with acoustic phonons.[15]

Using the formulas given in Ref. 15, we can show that

$$\frac{\tilde\nu_{ep}(\Theta)}{\tilde\nu_{es}(\Theta)} \sim \frac{m_s}{M}\left(\frac{\varepsilon_a}{\Theta}\right)^2\left(\frac{\varepsilon_{ef}}{\Delta}\right)^2,$$
$$\frac{\nu_{ep}(\Theta)}{\nu_{es}(\Theta)} \sim \frac{m^*}{M}\left(\frac{\varepsilon_a}{\Theta}\right)\left(\frac{\varepsilon_a}{\Theta_D}\right)\left(\frac{\varepsilon_{ef}}{\Delta}\right)^2, \quad (9)$$

where M is the mass of an ion; Θ_D is a parameter of the order of the Debye temperature; $\varepsilon_{ef} = \gamma Z e^2/a$ is a characteristic energy describing the electron-phonon interaction; Ze is the ion charge; γ is a dimensionless coefficient of the ion polarizability.[16]

Since $m^*/M \sim 10^{-5}$, $m_s/M \sim 10^{-4}$, and $\varepsilon_a/\Delta \sim 10$ it follows that the contribution of optical phonons to the processes of heating of the electron gas depend on the parameter γ. In the case of the substances for which γ is small the temperature of the electron gas is governed by the electron-magnon collisions. In the opposite case, we find from Eq. (6) that the heating is affected both by magnons and phonons, since the frequencies of collisions of electrons with quasiparticles of both types are of the same order of magnitude. The latter case is clearly realized in the compound $HgCr_2Se_4$.

According to Eq. (6), the heating field considered as a function of the electron temperature is described by the expression

$$E^2 \approx (\Theta m^*/e^2)(\tilde{\nu}_{ep}+\tilde{\nu}_{es})(\nu_{ep}+\nu_{es}). \quad (10)$$

In view of the threshold nature of the electron-magnon interaction the momenta of Cherenkov-emitted magnons are fairly large and satisfy the inequality $ql \gg 1$, where l is the mean free path of electrons [it follows from the laws of conservation that $q \propto p \propto (2m\Theta)^{1/2}$, $l = v\tau$, and $\tau = \mu_e m/e$]. In this limiting case the magnon emission coefficient is found by analogy with sound as the ratio of the power of the investigated magnons to the total energy stored in the magnon subsystem, and is of the form

$$\gamma_q^f = \frac{2\pi}{\hbar}\sum_p |\Psi_f|^2 [f_\downarrow^{(0)}(p+q) - f_\uparrow^{(0)}(p)]\delta(\mathscr{E}_{p+q}^\downarrow - \mathscr{E}_p^\uparrow - \hbar\omega_q), \quad (11)$$

where $\mathscr{E}_p^{\uparrow\downarrow} = \varepsilon(p) \mp \Delta/2$ is the dispersion law of electrons derived allowing for the spin splitting.

The calculations yield

$$\gamma_q^f(u) = \gamma_q^f(0)\left(1 - \frac{qu}{\hbar\omega_q}\right), \quad (12)$$

$$\gamma_q^f(0) \approx \omega_q \left(\frac{AM_0}{\Theta}\right)^{3/2} \left(\frac{AM_0 m}{q^2}\right)^{1/2} na^3 \exp\left[-\frac{(q^2+2m\Delta)^2}{8q^2 m\Theta}\right].$$

The strongest mechanism of damping of Cherenkov magnons is related to their collisions with thermal magnons and is described by the frequency[17]

$$\nu_{ss}^{(4)}(q) \approx \frac{\Theta_C}{\hbar}\left(\frac{T}{\Theta_C}\right)^2 \left(\frac{aq}{\hbar}\right)^4, \quad \hbar\omega_q \ll T. \quad (13)$$

For the compound $HgCr_2Se_4$ under investigation, we have $aq/\hbar \sim 10^{-1}$, $T/\Theta_C \sim 10^{-1}$, $\Theta \gtrsim \Delta$, $na^3 \sim 10^{-4}$, and the ratio of the magnon absorption and emission coefficients is

$$\frac{\nu_{ss}^{(4)}(q)}{\gamma_q^f(0)} \sim \left(\frac{aq}{\hbar}\right)^3 \left(\frac{T}{\Theta_C}\right)^{3/2} \left(\frac{\Theta}{\Delta}\right)^{3/2} \left(\frac{T}{\Delta}\right)^{1/2} \left(\frac{m_s}{m}\right)^{1/2} \frac{1}{na^3} \ll 1. \quad (14)$$

According to Eq. (10), the threshold field is $E(\Delta) \sim 10^3$ V/cm and in this case we have $\hbar\omega_q < qu$. Therefore, we can speak of generation of magnons by hot carriers in ferromagnetic semiconductors.

We shall now consider the physical aspects of magnon generation due to interactions of different types. It is known that in the case of the relativistic processes the number of magnons may not be conserved. Then true magnon generation takes place. On the other hand, the *s-d* interaction conserving the total number of magnons acts as a pump which distributes magnons between various parts of the magnon spectrum. For example, electrons absorb cold magnons which are in equilibrium with a magnetic lattice (this is the result of strong magnon-magnon collisions) and then after a certain time they emit hot magnons with a temperature Θ and these magnons are in equilibrium with the electron gas. This follows rigorously from the solution of the energy balance equation for each of the spin subbands in the case when $\partial/\partial t = \partial/\partial \mathbf{r} \equiv 0$.

It follows from the above estimates that if $\gamma \ll 1$, i.e., when the covalent type of binding in matter predominates strongly over the ionic binding, then magnons are created preferentially since the emission frequency of magnons by electrons is considerably higher than the emission frequency of phonons. In the case when $\gamma \sim 1$, we can expect simultaneous generation of both magnons and phonons. However, since these quasiparticles have different physical properties, the two processes can be easily distinguished experimentally.

The influence of a strong electric field on the power of electrical and magnetic noise in $HgCr_2Se_4$ at $T = 4.2$ K was investigated by Gal'dikas *et al.*[18] They observed a strong rise of temperature of magnetic radiation in fields $E \sim 500$–2500 V/cm (depending on the sample) and they attributed this effect to the generation of Cherenkov magnons by the drifting carriers.

4. ANTIFERROMAGNETIC SEMICONDUCTORS

Wide-gap antiferromagnetic semiconductors have parameters similar to those of ferromagnetic compounds. They are characterized by a high electron density, $n \sim 10^{17}$–10^{18} cm^{-3}, fairly high Néel temperatures $\Theta_N \sim 100$–300 K, and mobilities in the range $\mu_e \gtrsim 700$ cm$^2\cdot$V$^{-1}\cdot$sec^{-1}. One of such semiconductors is the compound MnTe (Ref. 4).

When antiferromagnetic semiconductors are subjected to a weak magnetic field ($\mu H \ll T$) we can ignore the spin splitting of carriers so that magnons may be emitted due to the Cherenkov mechanism by cold electrons which are in a degenerate state because of low temperatures and high densities. We can show that in the case of antiferromagnetic semiconductors we can expect also the quantum case when $ql/\hbar \gg 1$. In moderately high electric fields $\mathbf{E}$ at low temperatures, when the heating effects are unimportant and the electron distribution function can be approximated by a Fermi step with drift (which follows from the solution of the transport equation), the magnon emission coefficient is

$$\gamma_q^{af}(u) = \gamma_q^{af}(0)\left(1 - \frac{qu}{\hbar\omega_q}\right),$$
$$\gamma_q^{af}(0) = \frac{Vm^2}{\pi\hbar^3 q}\Theta\left(p_F - \frac{q}{2}\right)\sum_j |\Psi_{af}^j(q)|^2 \omega_{qj}. \quad (15)$$

Substituting in this formula the amplitudes of the transition from Eqs. (3a) and (4a), we obtain the following expressions for $\gamma_q^{af}(0)$, which are described by the same (within one order of magnitude) formula allowing for the nature of the branch of the magnon spectrum:

$$\gamma_{qj}^{af}(0) \approx \omega_{qj}\left(\frac{AM_0}{\varepsilon_a}\right)^{3/2}\left(\frac{AM_0 m}{q^2}\right)^{1/2}\left(\frac{\hbar\omega_{qj}}{\Theta_N}\right). \quad (16)$$

The most interesting mechanism of the damping of Cherenkov magnons in antiferromagnetic semiconductors is related to their interaction with thermal magnons, and it is described by the frequencies of four-magnon collisions:

1) in the case of easy axis materials,[7,19]

$$\nu_{ss}^{(4)}(q) \approx \frac{\Theta_N}{\hbar}\left(\frac{T}{\Theta_N}\right)^5, \quad \mu H_{EA} \ll T; \quad (17)$$

2) in the case of easy plane materials at temperatures in the range $\mu H_{EA} \ll T \ll \Theta_N$, the damping of magnons in both branches is of the same form[20]:

$$\nu_{ss}^{\prime(4)}(q) \approx \omega_{qj}\left(\frac{\hbar\omega_{qj}}{\Theta_N}\right)\left(\frac{T}{\Theta_N}\right)^3. \quad (18)$$

We can easily show that the ratio of the damping and emission coefficients of magnons for typical values of the parameters of materials ($\varepsilon_a/AM_0 \sim 10$, $T/\Theta_N \sim 10^{-1}$, $\hbar\omega_q/\Theta_N \sim 10^{-1}$) is

$$\frac{\nu_{ss}^{(4)}(\mathbf{q})}{\gamma_{\mathbf{q}j}^{af}(0)} \sim \left(\frac{\Theta_N}{\hbar\omega_\mathbf{q}}\right)\left(\frac{\varepsilon_a}{AM_0}\right)^2\left(\frac{T}{\Theta_N}\right)^5 \ll 1,$$

$$\frac{\nu_{ss}^{\prime(4)}(\mathbf{q})}{\gamma_{\mathbf{q}j}^{af}(0)} \sim \left(\frac{\hbar\omega_\mathbf{q}}{\Theta_N}\right)\left(\frac{\varepsilon_a}{AM_0}\right)^2\left(\frac{T}{\Theta_N}\right)^3 \ll 1. \qquad (19)$$

Therefore, in the case of antiferromagnetic semiconductors, in contrast to ferromagnetic materials, the generation of Cherenkov magnons is possible by cold electrons. Since the carrier heating is not necessary, characteristic electric fields in which magnon generation takes place are defined by the condition $E > v/\mu_e$ ($v = \Theta_N a/\hbar$ is the magnon velocity which has the numerical value $v \sim 10^5$ cm/sec) and they amount to several hundreds of volts per centimeter. It should be pointed out that if $\Theta_D > \Theta_N$, then generation of magnons belonging to the soft branch occurs in weaker electric fields compared with phonon generation.

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[1)] We shall consider the specific case of electrons, but all the results which will be derived will apply also to holes.

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