Generation of high-frequency magnons in a ferromagnetic semiconductor

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It is shown that an intense beam of high-frequency almost-monochromatic magnons can be generated in a ferromagnetic semiconductor at attainable levels of pumping of electrons with down spin into the conduction band.

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Parallel and perpendicular pumping methods with which to generate high-frequency magnons with wave vectors $q < 10^6$ cm^{-1} are widely known. [1] It is shown in the present paper that in a ferromagnetic semiconductor, at attainable spin-down electrons pumping into the conduction band, it is possible to generate an intense beam of high-frequency monochromatic magnons that are also narrowly directed under certain conditions.

In a ferromagnetic semiconductor, the conduction band is split into subbands with spins $\uparrow$ and $\downarrow$. We consider the relaxation of nonequilibrium electrons thrown into the conduction band. Inasmuch as the degree of ionicity in ferromagnetic semiconductors is large, an electron with kinetic energy ϵ_p much higher than the energy of the optical phonon is in most cases more likely to emit an optical phonon than to emit or absorb a magnon. For example, for EuO with dielectric constants $\kappa_0 = 4.6$ and $\chi_\infty = 26$, the frequency of the longitudinal optical phonon is $\omega_l = 1.3\times 10^{14}$ sec^{-1}, [2] the exchange gap in the electron spectrum is $\Delta = 2Is = 0.6$ eV, [3] and the ratio of the relaxation frequencies of the electrons on the optical phonons and on the magnons is of the order of 10^2. Thus, electrons thrown into the $\downarrow$ subband emit n optical phonons ($n = E(\epsilon_p/\epsilon_l)$, where E is the integer part) and end up in a state with kinetic energy $\omega < \omega_l$. They then go over to the subband $\uparrow$, emitting one magnon each, after which, emitting optical phonons, they collapse to the bottom of the band. It follows from the energy and momentum conservation law that the momenta of the emitted magnons q lie in the interval $q_1 \le q \le q_2$, where $q_{1,2} = p_0 \pm p$, $p_0 = (2m\Delta)^{1/2}$, m and p are the mass and the momentum of the electron in the $\downarrow$ band and $p \ll p_0$. It follows from the foregoing that number of emitted magnons is equal to the number of electrons thrown into the upper subband, and at a sufficient pump intensity it can reach an appreciable value, greatly exceeding the thermal background. In the stationary state, the magnon distribution function N_q in the interval $q_1 \le q \le q_2$ is determined from the condition that the number of generated magnons be equal to the number of those relaxing via various magnon—magnon processes:

$$N_{\mathbf{q}} = \left(N^{\circ}_{\mathbf{q}} + \frac{\Gamma_{me}}{\Gamma_m}\right)\left(1 - \frac{\Gamma_{me}}{\Gamma_m}\right)^{-1}, \qquad q_1 \leqslant q \leqslant q_2 \tag{1}$$

where N_q^0 is the equilibrium distribution function of the magnons, Γ_m is the magnon relaxation frequency, Γ_{me} is the magnon generation frequency as determined by the relation

$$\Gamma_{me}(q) = 2^{1/2}\pi s \frac{l^2 \nu m^{1/2}}{\omega^{1/2} q} \gamma_{e,m}^{-1}(\sqrt{2m\omega}), \tag{2}$$

and ν is the electron pump intensity per unit cell. The electron relaxation frequency is

$$\gamma_{em}(p) = \frac{s}{2\pi} \frac{l^2 m a}{p} \frac{T}{\Theta_c} Z . \tag{3}$$

Here T is the temperature, $Z = \int_{x_1}^{x_2} (N(x) + 1)\, dx$, and $x_{1,2} = \omega_{\mathbf{q}1,\mathbf{q}2} T^{-1}$. Since Γ_{me} is expressed in terms of Z, it follows that (1) is an integral equation with respect to N_q. If $T > T_0 = \Theta_c (ap_0)^2$, then the generated magnons are subthermal. Putting $\Gamma_m(q) \sim q^a$ with $\alpha > 0$, we obtain the following expression for Z (the magnon spectrum is $\omega_{\mathbf{q}} = \Theta_c (aq)^2$, a is the interatomic distance, and Θ_c is a quantity on the order of the Curie temperature

$$Z = \frac{2}{a+1} \ln \left[\frac{\left(\frac{q_2}{p_0}\right)^{a+1} Z - \frac{\nu}{\nu_0}}{\left(\frac{q_1}{p_0}\right)^{a+1} Z - \frac{\nu}{\nu_0}} \right] , \tag{4}$$

where

$$\nu_0 = \frac{1}{4\pi} \left(\frac{p_0}{k_0}\right) \left(\frac{T}{\Theta_c}\right) \Gamma_m(p_0), \quad k_0 = \pi/a . \tag{5}$$

At $Z \gg 1$ (the asymptotic behavior sets in in fact at $Z > 1$) and at $p \ll p_0$ we have

$$Z = \frac{\nu}{\nu_0} \left(\frac{p_0}{q_1}\right)^{a+1} \left\{ 1 + 2(a+1) \frac{p}{p_0} \exp\left[-\left(\frac{a+1}{2}\right) \left(\frac{\nu}{\nu_0}\right) \left(\frac{p_0}{q_1}\right)^{a+1} \right] \right\}, \tag{6}$$

i.e.,

$$N_{\mathbf{q}} = N_{\mathbf{q}}^{\circ} \left\{ 1 - \left(\frac{q_1}{q}\right)^{a+1} + 2(a+1) \left(\frac{p}{p_0}\right) \left(\frac{q_1}{q}\right)^{a+1} \exp\left[-\left(\frac{a+1}{2}\right) \left(\frac{\nu}{\nu_0}\right) \left(\frac{p_0}{q_1}\right)^{a+1} \right] \right\}^{-1} . \tag{7}$$

It follows from (7) that the number of magnons in an exponentially narrow momentum interval $\Delta q : q_1 \leq q \leq q_1 + \Delta q$,

$$\Delta q = 2(a+1)\, q_1 \left(\frac{p}{p_0}\right) \exp\left[-\left(\frac{a+1}{2}\right) \left(\frac{\nu}{\nu_0}\right) \left(\frac{p_0}{q_1}\right)^{a+1} \right] \tag{8}$$

increases exponentially rapidly with increasing pump intensity:

$$N_{\mathbf{q}} \sim N_{\mathbf{q}}^{\circ} \exp\left[\left(\frac{a+1}{2}\right) \left(\frac{\nu}{\nu_0}\right) \left(\frac{p_0}{q_1}\right) \right]^{a+1} . \tag{9}$$

At the same time, the number of magnons with momenta $q_1 + \Delta q < q \leq q_2$ is independent, with exponential accuracy, of the pump, and is equal to

$$N_{\mathbf{q}} = N_{\mathbf{q}}^{\circ} \left[1 - \left(\frac{q_1}{q}\right)^{a+1} \right]^{-1} . \tag{10}$$

Thus, if the pumping is strong enough, the magnon distribution function has a sharp peak at $q \approx q_1$. The described regime sets in at $Z > 1$, i.e., as seen from

(6), at a pump $\nu > \nu_0$. It is possible to consider in similar fashion the temperature region $T < T_0$, when the generated magnons are epithermal. We present results only for the case when the magnon damping Γ_m is determined predominantly by four-magnon exchange processes, i.e., $\Gamma_m(q) \sim q^3$. The exponential generation regime then sets in when

$$\nu > \nu_o' = \frac{\pi}{4}\left(\frac{p_o}{k_o}\right)^3 \Gamma_m(p_o) \; . \tag{11}$$

In this regime

$$N_{\mathbf{q}} = \frac{p_o}{8p} \exp\left[2\left(\frac{p_o}{q_1}\right)^6 \frac{\nu}{\nu_o'}\right], \quad q_1 \leqslant q < q_1 + \Delta q, \tag{12}$$

$$N_{\mathbf{q}} = \frac{q_1^4}{q^4 - q_1^4}\,, \qquad q_1 + \Delta q < q \leqslant q_2\,, \tag{13}$$

where

$$\Delta q = q_1 \frac{2p}{p_o} \exp\left[-2\left(\frac{p_o}{q_1}\right)^6 \frac{\nu}{\nu_o'}\right] \; . \tag{14}$$

In real crystals, the anisotropy of the magnon spectrum (due, for example, to dipole—dipole interaction) and the anisotropy of the electron spectrum cause, generally speaking, Γ_m, Γ_{me}, and γ_{em} to be dependent on the direction of the wave vector $\mathbf{q}$. As a result, the generation conditions turn out to be optimal for selected directions $\mathbf{q}$, and the generated magnons are concentrated in a narrow cone.

In conclusion, let us estimate the critical pump ν_0. Assuming $\Gamma_m(p_0) \approx 10^8$—$10^9$ sec^{-1}, $T/\Theta_c \approx 10^{-1}$, and $p_0/k_0 \approx 1/5$, we find that at $T > T_0$ the critical pump is $\nu_0 \approx 10^6$—10^7 sec^{-1}, and is perfectly feasible in experiment.

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