Photostimulated acoustomagnetoelectric effect in semiconductors

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It is shown that the presence of a light wave lifts the hindrances imposed on the acoustomagnetoelectric (AME) effect by the degeneracy of the electron gas and by the energy and momentum conservation laws. In the case of impurity scattering of the electrons, such a photostimulated AME effect can be larger by one or two orders of magnitude than in scattering by acoustic phonons.

The acoustomagnetoelectric (AME) effect in monopolar conductors was predicted in[1] and observed experimentally in[2,3]. In this effect, an acoustic wave propagating in an electrically open-circuited sample in a direction transverse to a magnetic field H applied in the third (Hall) direction gives rise to an electric field

$$E_{\text{AME}} \sim \frac{\alpha W}{ens} \frac{\eta}{1+\eta^2} \quad . \tag{1}$$

where α is the coefficient of sound absorption by the electrons, W is the density of the acoustic flux, n is the electron density, s is the speed of sound, $\eta = \mu H/c$, and μ is the electron mobility. The AME effect is due to the dispersion of the electron-phonon interaction, i.e., to the fact that the acoustic flux, unlike the compensating electric field, acts in different manners on electrons with different energies. For this dispersion to appear, the electron relaxation time should depend on the energy. At $ql \gg 1$ and $q \gg p$ (q is the wave vector of the sound, l is the electron mean free path, $\bar{p}$ is the characteristic quasimomentum of the electron, we use a system of units with $\hbar = 1$), the AME effect is forbidden because of the turning off of the electron-phonon interaction as a result of the energy and momentum conservation laws; it vanishes also on going to a fully degenerate electron gas, for in this case only electrons with the Fermi energy ϵ_F participate in the absorption and emission of the phonons, and the dispersion of the electron-phonon interaction cannot come into play (we are considering semiconductors with spherical Fermi surfaces).

It is shown in this article that these two hindrances are lifted in the presence of a sufficiently intense optical wave whose frequency Ω satisfies the condition $\epsilon_F \ll \Omega \approx q^2/2m < E_g$ (E_g is the width of the forbidden band). Indeed, the electron-phonon interaction in the presence of an optical wave becomes possible as a result of the modification introduced into the conservation laws by the participation of the photons,[4] and the effective energy scatter of the electrons appears because the scattering of the electrons by the phonons in the presence of the optical waves is accompanied by photon absorption that leads to an appreciable increase of the electron energy in comparison with its initial one. Moreover, since this scatter greatly exceeds the thermal scatter, the value of E_{AME} for certain scattering mechanisms is much larger than the values that might be expected by starting from (1) and from the absorption of sound by electrons in the presence of the optical wave.[4]

To construct a quantitative theory of such a photostimulated AME we start from the kinetic equation[1,5]

$$e\mathbf{E}\frac{\partial f_{\mathbf{p}}}{\partial \mathbf{p}} + \frac{e}{mc}[\mathbf{p}\times\mathbf{H}]\frac{\partial f_{\mathbf{p}}}{\partial \mathbf{p}} = I_{\mathbf{p}}^{(1)}\{f_{\mathbf{p}}\} + I_{\mathbf{p}}^{(2)}\{f_{\mathbf{p}}\}, \tag{2}$$

where $\mathbf{E}$ is the constant electric field produced in the open-circuited sample by the acoustic wave, and $I_{\mathbf{p}}^{(1)}\{f_{\mathbf{p}}\}$ is the integral of the electron collisions with the scatterers (thermal phonons, impurities, etc.; the scattering is assume to be almost elastic),

$$I_{\mathbf{p}}^{(2)}\{f_{\mathbf{p}}\} = \frac{\pi\Lambda^2 W}{\rho s^3}\sum_{l=-\infty}^{\infty} J_l^2\left(\frac{e\mathbf{E}_0 q}{m\Omega^2}\right)\{(f_{\mathbf{p}+\mathbf{q}} - f_{\mathbf{p}})\delta(\epsilon_{\mathbf{p}+\mathbf{q}} - \epsilon_{\mathbf{p}} - \omega_{\mathbf{q}} - l\Omega) + (f_{\mathbf{p}-\mathbf{q}} - f_{\mathbf{p}})\delta(\epsilon_{\mathbf{p}-\mathbf{q}} - \epsilon_{\mathbf{p}} + \omega_{\mathbf{q}} - l\Omega)\} \tag{3}$$

is the integral of collisions with the acoustic phonons,[5,6] is the deformation-potential constant, ρ is the crystal density, $\mathbf{E}_0$ is the amplitude of the optical-wave field, and $J_l(x)$ is a Bessel function of real argument. It is assumed that $qR \gg 1$ (R is the Larmor radius), so that the magnetic field does not affect the integral of the collisions with the acoustic phonons. In the approximation linear in the light and sound intensities, we can replace f_p in (3) by a Fermi step, retaining only the terms with $l = \pm 1$ (the term with $l = 0$ drops out at $q \gg p_F$, since the argument of the corresponding δ-function cannot vanish), and we confine ourselves to the first term in the

expansion of the Bessel functions. To solve the obtained equation we use a method proposed in[7]. Multiplying (2) by $\mathbf{p}\delta(\epsilon-\epsilon_p)$ and summing over $\mathbf{p}$, we obtain for the current density

$$\mathbf{j}=\frac{e}{m}\int_0^\infty \mathbf{R}(\epsilon)d\epsilon, \tag{4}$$

where $\mathbf{R}(\epsilon)$ is defined by the equation

$$\frac{\mathbf{R}(\epsilon)}{\tau(\epsilon)}+\frac{e}{mc}[\mathbf{H}\times\mathbf{R}(\epsilon)]=\sum_{\mathbf{p}}\mathbf{p}\,\delta(\epsilon-\epsilon_p)l_{\mathbf{p}}^2\{f_{\mathbf{p}}\}+\frac{enE}{m}\delta(\epsilon-\epsilon_F), \tag{5}$$

and $\tau(\epsilon)$ is the electron relaxation time. Solving (5), substituting in (4), and assuming $\mathbf{j}=0$ (open-circuited sample), we obtain after eliminating the longitudinal component (along the acoustic flux) of the constant electric field the value of E_{AME}. The general expression is rather complicated and we present here only the result for $q=\sqrt{2m\Omega}\pm p_F$, $\mathbf{q}\parallel\mathbf{E}_0$, which corresponds to the boundary of the region of the electric-phonon interaction and to a maximum coefficient of absorption (amplification) of the sound by the electrons in the presence of the optical wave.[4] In this case we have

$$E_{\text{AME}}=\frac{a\mathcal{W}}{ens}\,\frac{\eta\gamma^\nu(\gamma^\nu-1)}{1+\eta^2\gamma^{2\nu}}, \tag{6}$$

where

$$a=\pm\frac{\Lambda^2m^2p_F}{2\pi\rho s}\frac{e^2E_0^2}{m\Omega^3} \tag{7}$$

is the corresponding value of the absorption coefficient, $\gamma=\Omega/\epsilon_F\gg1$, and $\nu=(\epsilon/\nu)(d\tau/d\epsilon)$. At $\eta=\gamma^{-\nu}$ the value of E_{AME} reaches its maximum

$$E_{\text{AME}}^{\max}=\frac{a\mathcal{W}}{ens}\,\frac{1}{2}(\gamma^\nu-1). \tag{8}$$

E_{AME} depends significantly on the scattering mechanism. Particular interest attaches to the case of impurity scattering ($\nu=3/2$), inasmuch as $E_{\text{AME}}^{(\max)}$ is reached in this case in weak magnetic fields ($\eta\ll1$) and exceeds by a factor $\gamma^{3/2}/2\gg1$ the value expected by substituting (7) in (1). The reason is that in the case of impurity scattering the electrons that absorb photons in the high-energy region are scattered more weakly than the remaining electrons and make the largest contribution to the kinetic effects. The sign of the effect also depends on the scattering mechanism.

The considered effect can be produced not only by the passage of a monochromatic acoustic wave, but also by propagation of thermal pulses.[8]

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