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If a monochromatic flux of phonons is excited in a crystal by an external source, the phonons will drag conduction electrons with them. As a result, an electric field will be produced in an open-circuited sample, and will cancel the effect of the phonon current in the sense that the total electric current through the sample will be zero. This phenomenon is the well known acousto-electric effect.

The cancellation of the charge current will, generally speaking, not be accompanied by cancellation of other fluxes (energy, momentum, etc.). The reason is that the perturbation of the electron distribution in momentum space, brought about by the external phonon flux, differs from that which produces the electric field (see formula (1)). Therefore the resultant force acting on each individual electron will depend on the energy of the latter, so that one part of the electrons will move predominantly in the direction of the phonon current, and the other part in the opposite direction. Since the average electron energy is generally speaking different in these two equal but opposite fluxes, a nonzero energy flux can be established in the open-circuited sample even under isothermal conditions. Accordingly, a temperature gradient will be produced under adiabatic conditions. In analogy with the acousto-electric effect, this phenomenon can be called the acousto-thermal effect.

Let us consider a non-piezoelectric crystal in which the carriers are characterized by an isotropic effective mass m and a relaxation time $\tau(\epsilon)$, both dependent on the energy. Assume that a hypersonic wave with wave vector \vec{q} , satisfying the inequality $l^{-1} \ll q \lesssim \sqrt{mkT/\hbar}$ (l = electron mean free path, T = crystal temperature). By regarding such a wave as a current of monochromatic phonons [1] and by calculating the integral of the collisions between the electrons and the phonon current [2] we obtain for the antisymmetrical part of the distribution function (taking into account the acousto-electric field E and the acousto-thermal temperature gradient ∇T)

$$\vec{f}_1(\epsilon) = -\frac{\hbar\tau(\epsilon)}{m} \frac{\partial f_0}{\partial \epsilon} \left\{ e\vec{E} - \left(\frac{\partial \xi}{\partial T} - \frac{\xi - \epsilon}{T} \right) \nabla T + \frac{3\pi\Xi^2 m^{1/2} q \vec{W}}{4\sqrt{2}\rho s^2 \epsilon^{3/2}} \vartheta(\epsilon - \epsilon_0) \right\} \quad (1)$$

Here ξ is the chemical potential of the electrons, ρ the crystal density, s the speed of sound, Ξ the deformation-potential constant, \vec{W} the flux density of sound energy, $\epsilon_0 = \hbar^2 q^2/8m$, $f_0(\epsilon)$ the electron equilibrium distribution function, $\vartheta(x) = 1$ when $x > 0$, and $\vartheta(x) = 0$ when $x < 0$.

Using (1), we can readily find the electric current density \vec{j} and the heat flux density \vec{Q} (the latter includes also the heat flux $\kappa_L \nabla T$ due to the lattice thermal conductivity κ_L). These flux densities are linear functions of \vec{E} and ∇T . Imposing the boundary conditions $\vec{j} = 0$ (open-circuited sample) and $\vec{Q} = 0$ (adiabatic mode), we obtain for the acousto-thermal temperature gradient

$$\nabla T = \frac{\alpha \vec{W}}{n s} T [1 + (\kappa_L / \kappa_e)]^{-1} \frac{\langle \langle \epsilon \tau \rangle \rangle \langle \tau \rangle - \langle \langle \tau \rangle \rangle \langle \epsilon \tau \rangle}{\langle \epsilon^2 \tau \rangle \langle \tau \rangle - \langle \epsilon \tau \rangle^2}, \quad (2)$$

where α is the electronic coefficient of sound absorption, and κ_e the electronic thermal conductivity,

$$\langle \dots \rangle = - \frac{(2m)^{3/2}}{3\pi^2 n^3} \int_0^\infty \langle \dots \rangle \frac{\partial f_0}{\partial \epsilon} \epsilon^{3/2} d\epsilon, \quad \langle \langle \dots \rangle \rangle = - \frac{1}{f_0 \epsilon_0} \int_{\epsilon_0}^\infty \frac{\partial f_0}{\partial \epsilon} d\epsilon,$$

and n is the electron density.

Putting $\tau(\epsilon) \sim \epsilon^\nu$ and considering a nondegenerate electron gas, we obtain for $\epsilon_0 \ll kT$

$$\nabla T = - \frac{\alpha \vec{W}}{k n s} [1 + (\kappa_L / \kappa_e)]^{-1} \frac{9\sqrt{\pi}}{8} \frac{\Gamma(1+\nu)}{\Gamma[(7/2)+\nu]}. \quad (3)$$

We see that in this case the heat rise is produced in that end of the sample in which the sound "enters" (since $\nu > -1$). In the other limiting case ($\epsilon_0 \gg kT$) we have

$$\nabla T = \frac{\alpha \vec{W}}{k n s} [1 + (\kappa_L / \kappa_e)]^{-1} (\epsilon_0 / kT)^{1+\nu} \frac{\Gamma(s/2)}{\Gamma[(7/2)+\nu]}. \quad (4)$$

Thus, the effect reverses sign when the frequency of the hypersound is raised.

Let us estimate the magnitude of the effect. Putting $\kappa_e \gtrsim \kappa_L$ we get for $\epsilon_0 \lesssim kT$, $m \sim 0.1 m_e$, $q \sim 10^6 \text{ cm}^{-1}$, $\Xi \sim 10 \text{ eV}$, and $T \sim 10^\circ \text{K}$

$$|\frac{\nabla T}{W}| \sim 10^4 \left(\frac{\text{deg}}{\text{cm}} \right) / \left(\frac{W}{\text{cm}^2} \right).$$

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- [1] L. E. Gurevich, Izv. AN SSSR ser. fiz. 21, 112 (1957), transl. Bull. Acad. Sci. Phys. Ser. p. 106.
- [2] E. M. Epshtein, FTT 8, 552 (1966), Soviet Phys. Solid State, in press.

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In formula (4) on p. 269 read $\Gamma(5/2)$ instead of $\Gamma(s/2)$.