Nonlocal phonon thermal conductivity under conditions of strong excitation

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Self-similar solutions, which describe the distribution of nonequilibrium phonon temperature due to laser excitation of electron-hole plasma, are obtained.

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A large number of nonequilibrium phonons with a frequency of the order of the Debye frequency ω_D are produced as a result of the generation of electron-hole plasma in a semiconductor by a laser beam during thermalization and recombination of the carriers in the absorption layer. The manner in which these phonons propagate in the bulk of a crystal is important in many physical phenomena.

The spectral evolution of the phonon distribution, which occurs because of three-phonon decay and fusion processes, is a characteristic feature of the propagation of phonons under these conditions. This evolution, which occurs simultaneously with the spatial diffusion due to scattering by the defects, affects the diffusion strongly, since the diffusion coefficient D depends on the frequency. On the other hand, the diffusion controls the phonon-occupation numbers n and therefore influences the three-phonon processes: if $n \ll 1$, there is a spontaneous decay; if $n \gg 1$, then the fusion will dominate; if, however, $n \approx 1$, then the decay and fusion will balance out and a Planck distribution will be established.

If the pumping is weak, the layer occupied by the phonons will expand faster during the diffusion than the number of phonons produced due to the decay. As a result, the occupation numbers decrease and the temperature cannot be stabilized.[1,2] In this letter we study another propagation regime which occurs during a strong pumping. The occupation numbers in this case reaches a value of ≈ 1 before the phonons leave the absorption layer, and the phonon temperature T_0 is established in the layer. If the temperature T_0 is sufficiently low and the transfer processes are unessential, then, in contrast with the standard thermal conductivity, the energy flux w will be determined by the cold phonons with $\omega \approx \tilde{\omega} \ll T$, rather than by the thermal phonons with $\omega \approx T$; the narrow point of the thermal conductivity, which is nonlocal, is the spectral energy transfer from the region $\omega \approx T$, in which it is "stored," to the region $\omega \approx \tilde{\omega}$ in which it is "transported" (Refs. 3 and 4).

The "initial" phonon temperature T_0 is determined from the balance $\epsilon(T_0) d = P$; here $\epsilon(T)$ is the thermal energy of the lattice in 1 cm^3, d is the absorption depth, and P is the energy absorbed in 1 cm^2 of the surface. The time for establishment of the temperature T_0 is $\tau_0 \equiv \tau(T_0)$, where $\tau(\omega)$ is the time of the three-phonon processes. During this time the phonons diffuse to a depth $l_0 \equiv l(t_0)$, where $l(\omega) \equiv [D(\omega)\tau(\omega)]^{1/2}$. The temperature is therefore established if $l_0 \ll d$, which gives $P \gtrsim 10^{-3}$ J/cm^2 for typical parameters for the semiconductors.

The equation, which determines the distribution of the phonon temperature, can be derived in the following way. Assuming that T is a known function of r and t, we shall determine the distribution of the cold phonons from the equation

$$\left[\frac{\partial}{\partial t} - D(\omega)\nabla^2\right] n = -\frac{1}{\hat{\tau}(T,\omega)}\left[n - \frac{T}{\omega}\right], \tag{1}$$

where $\hat{\tau}$ is the absorption time of a nonequilibrium cold phonon of frequency ω by the Planck distribution with the temperature T. We then calculate the energy flux

$$\mathbf{w} = \int_0^\infty d\omega\, \rho(\omega)\, \omega\, [-D(\omega)\nabla n], \tag{2}$$

where ρ is the density of states. Substituting w in the form of a functional of T in the energy conservation law

$$\partial\epsilon/\partial t + \operatorname{div}\mathbf{w} = 0, \tag{3}$$

we obtain the sought-for equation.

If $T \ll \omega_D$ and the scattering by defects is a Rayleigh scattering, then

$$D(\omega) \sim \omega^{-4}, \quad 1/\hat{\tau} \sim \omega T^4, \quad \rho \sim \omega^2, \quad \epsilon \sim T^4,$$

and Eqs. (1)-(3) have self-similar solutions from which all the physical parameters of the problem can be excluded. We shall determine the temperature T_∞ by the condition $P = \epsilon(T_\infty)\, l(T_\infty)$ and introduce the dimensionless variables $\bar{z} = z/l_\infty$, $\bar{t} = t/\tau_\infty$, and $\bar{\omega} = \omega/T_\infty$, where $l_\infty \equiv l\, T_\infty$ and $\tau_\infty \equiv \tau(T_\infty)$. It is easy to check that $T_\infty = T_0 (l_0/d)^2 \ll T_0$, so that $l_\infty \gg l_0$ and $\tau_\infty \gg \tau_0$. If we substitute in Eqs. (1)-(3)

$$n(\omega, z, t) = \bar{t}^{-\lambda} g(\zeta, \eta) \qquad (\lambda = 1/21), \tag{4}$$

$$T(z, t) = T_\infty \bar{t}^{-\mu} f(\zeta) \qquad (\mu = 5/21), \tag{5}$$

with self-similar variables

$$\eta = \bar{\omega}\bar{t}^{a}, \quad \zeta = \bar{z}\bar{t}^{-\beta} \quad (a = 4/21,\ \beta = 20/21), \tag{6}$$

then the equations obtained for g and f will not have any parameters, and hence the most important results can be obtained without knowing the explicit form of these

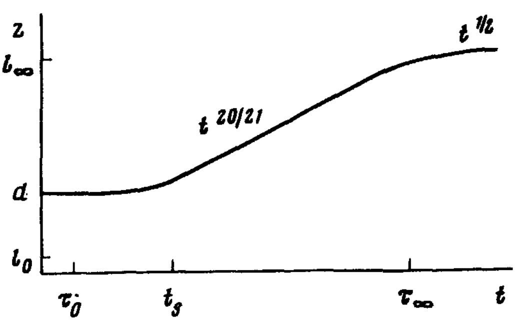

FIG. 1. Time dependence of the depth of the "heated" layer.

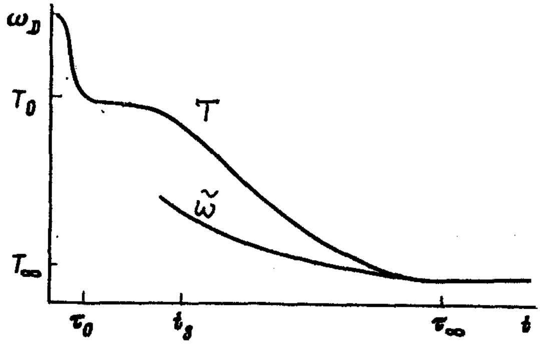

FIG. 2. Time dependence of the temperature (of the average phonon frequency) and of the frequency of phonons that transfer the energy.

functions. Thus the energy of this system of equations is concentrated in the region which is determined by the condition $\zeta \approx 1$, i.e., for $z \approx l_\infty (t/\tau_\infty)^{20/21}$. The characteristic temperature in this region is $T \approx T_\infty (t/\tau_\infty)^{-5/21}$. The frequency of phonons that transfer the energy is determined by the condition $\eta \approx 1$, i.e., $\tilde{\omega} \approx T_\infty (t/\tau_\infty)^{-4/21}$. These relations include the time dependence of the main physical parameters (Figs. 1 and 2) and the pumping power (at l_∞ and τ_∞).

The self-similar solution of Eqs. (4) and (5) can be obtained for these times t_s when the characteristic temperature, which is determined from (5), is of the order of T_0. It is easy to see that $t_s \approx \tau_0 (d/l_0)^{8/5}$ and to verify that $\tau_0 \ll t_s \ll \tau_\infty$. At the time t_s, which is characteristic for the solution of (5), $z_s \approx d$. On the other hand, t_s coincides in order of magnitude with the time during which the heated layer of thickness d doubles in width due to nonlocal thermal conductivity.[4] This means that the self-similar regime at $t \gg t_s$ matches the non-self-similar stage at $t \lessdot t_s$. Thus we have the following physical picture. First the temperature T_0, which remains constant to $t \approx t_s$, is established in the layer d during the time τ_0. During this time the heated layer begins to expand, and the propagation is converted to a self-similar regime which persists to $t \approx \tau_\infty$. At this time, the temperature spreads to a depth l_∞ as it decreases to T_∞. At this moment the frequency of phonons that transfer the energy is equal to this temperature; because of this, the mechanism of nonlocal thermal conductivity is no longer in effect. It can be shown that at $t \gg \tau_\infty$ the anharmonic processes are unessential and that the phonons of different frequencies diffuse independently of each other, thereby breaking down the Planck distribution; however, their average frequency remains at the level $\omega \approx T_\infty$.

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