Friction of an electron-hole drop moving at near-sonic speed

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(Submitted 30 June 1978)
Zh. Eksp. Teor. Fiz. 75, 1943–1951 (November 1978)

The friction force acting on a moving electron-hole drop is considered. It is shown that at low temperatures the friction due to incoherent emission and absorption of acoustic phonons increases as the speed of the drop approaches that of sound. At supersonic speeds an additional much stronger friction is produced by the Cerenkov emission of sound by the drop as a whole. The coherent interaction of the electrons and holes of the drop with the lattice-deformation field leads also to the appearance of an additional "deformation" mass of the electron-hole drop.

PACS numbers: 71.35. + z

1. INTRODUCTION

The motion of electron-hole drops in a semiconductor under the influence of external forces was investigated in a number of papers.[1-4] The drop velocity is determined by the viscous-friction force applied to the drop by the crystal lattice. The friction force due to emission and absorption of acoustic phonons by the electrons and holes of the drop was calculated by Keldysh[5] for velocities $v \ll s$, where s is the sound velocity. Since the experimentally observed drop velocities reach values $v \sim s$ (Refs. 3, 4, 6) it is of interest to examine the friction force and its velocity dependence near the speed of sound. The deceleration mechanism proposed by Keldysh[5] was investigated at high velocities by Manoliu and Kittel,[7] who calculated numerically the friction force for germanium at drop velocities $(0.5-2.5)\times 10^5$ cm/sec falling short of the sound velocity. According to them, the viscous-friction coefficient changes insignificantly in this interval.

It is shown in the present paper that the friction force increases when the drop approaches the speed of sound. This increase is due to the fact that at $v \geqslant s$ the moving drop, as a macroscopic body, can emit Cerenkov radiation, i.e., the entire aggregate of electrons and holes of the drop can radiate coherently. The wavelength of the radiated sound is in this case of the order of the drop dimension.

Second, incoherent emission of phonons by individual electrons and holes is possible; this phenomenon was considered in a number of studies.[5-7] At low temperatures the intensity of the incoherent emission (which determines the friction force at $v < s$) should increase strongly as $v \to s$. In fact, in the presence of degeneracy and at $v \ll s$, the energy sq lost by an individual particle as it emits a phonon of momentum $\mathbf{q}$ should be of the order of kT. At low temperatures $kT \ll \hbar k_F s$ ($\hbar k_F$ is the Fermi momentum) the momenta of the emitted phonons are therefore small, $q \ll \hbar k_F$, so that the particle loses a small fraction of its momentum. For this reason, at low temperatures and at $v \ll s$, the friction force contains the small factor $(kT/\hbar k_F s)^5$ (see Ref. 5).

With increasing drop velocity the characteristic momentum of the emitted phonons increases, and at $v \geqslant s$ this momentum turns out to be of the order of the Fermi momentum, so that the friction should increase strongly.

This statement can be explained in the following manner: the energy conservation law for phonon emission by an electron can be written in the form

$$\varepsilon(\mathbf{p}+m_e\mathbf{v})-\varepsilon(\mathbf{p}+m_e\mathbf{v}-\mathbf{q})=sq,$$

where $\mathbf{p}$ is the electron momentum in a reference frame co-moving with the drop at velocity v, and m_e is the electron effective mass. The energy lost by the electron in the moving coordinate frame is

$$\varepsilon(\mathbf{p})-\varepsilon(\mathbf{p}-\mathbf{q})=sq-\mathbf{q}\mathbf{v}.$$

At subsonic velocities this is a postive quantity and, as a result of degeneracy, should be of the order of kT. The absolute value of the momentum of a phonon emitted in the direction of the motion (it is these phonons which make the largest contribution to the deceleration force) is of the order of $kT/(s-v)$ and increases with increasing velocity.

At $v \geqslant s$ phonon emission becomes possible even at $T=0$. The internal energy of the drop then increases, in contrast to the case $v<s$, and the restriction on the phonon momentum is lifted. This should lead to an increase of the friction force.

The increased role of the phonon emission by a moving drop at low temperatures leads at subsonic velocities to the drop cooling considered by us earlier.[8]

In Sec. 2 of this paper we calculate the friction force due to incoherent emission and absorption of phonons by electrons and holes of the drop, as a function of the drop velocity and of the lattice temperature. The growth, discussed above, of the friction force at subsonic velocities turns out to be large only at very low temperatures $T \ll T_0$, where $T_0 = 2\hbar k_F s$.

In Sec. 3 is calculated the friction force that appears at supersonic velocities as a result of Cerenkov emission of sound by the drop as a whole. The value of this force at the Cerenkov-radiation threshold exceeds substantially the low-temperature friction force at $v<s$.

It appears that the Cerenkov friction force makes it impossible to attain in experiment velocities higher than that of sound.

Coherent interaction of the entire aggregate of the electrons and holes of the drop with the lattice-deformation field produces, besides the Cerenkov friction force, an additional "deformation" mass of the electron-hole drop. This deformation mass and its dependence on the drop velocity are obtained in Sec. 4.

2. FRICTION FORCE DUE TO INCOHERENT INTERACTION WITH PHONONS

We obtain in this section the friction force due to incoherent emission and absorption of acoustic phonons by electrons and holes of a drop moving with velocity $v<s$. The expression for the friction force $\mathbf{F}_1$ (per electron-hole pair) is of the form[5,9]

$$\mathbf{F}_1=-\frac{\pi}{2\hbar\rho s n_0}\sum_{i=e,h}D_i^2\int\frac{d\mathbf{p}d\mathbf{q}}{(2\pi\hbar)^6}\mathbf{q}q\delta(\varepsilon_{i\mathbf{p}}-\varepsilon_{i\mathbf{p}-\mathbf{q}}-sq+\mathbf{q}\mathbf{v})$$
$$\times[f_{i\mathbf{p}}(1-f_{i\mathbf{p}-\mathbf{q}})(N_\mathbf{q}+1)-f_{i\mathbf{p}-\mathbf{q}}(1-f_{i\mathbf{p}})N_\mathbf{q}]. \tag{1}$$

Here ρ is the crystal density, n_0 is the density of the electron-hole pairs in the drop, D_e and D_h are the deformation potentials of the electrons and holes, N_q is the equilibrium distribution function of the phonons for the lattice temperature T, $f_{e\mathbf{p}}$ and $f_{h\mathbf{p}}$ are the Fermi distribution functions of the electrons and holes at the drop temperature T_d, which can differ from the lattice temperature.

The stationary temperature T_d of the moving drop must be determined from the condition $dE/dt=0$, where E is the internal energy of the drop. An explicit expression for dE/dt is obtained by replacing the factor q under the integral sign in (1) by $sq-\mathbf{q}\cdot\mathbf{v}$ (Ref. 5). The dependence of the drop temperature T_d on the drop velocity and on the lattice temperature was investigated by us earlier.[8] It was shown that at subsonic velocities the drop becomes heated at high lattice temperatures and is cooled at low ones ($T\ll T_0$). This dependence of T_d on the velocity must be taken into account in the calculation of the friction force (1).

It is easy to verify that when account is taken of the condition $dE/dt=0$ the expression for the absolute value of the friction force is given by the integral of (1), in which the factor $\mathbf{q}$ must be replaced by $|\mathbf{q}|s/v$. At $\varepsilon_F\gg kT$ and $\varepsilon_F\gg ms^2$ (ε_F is the Fermi energy) this expression can be reduced to[1)]

$$F_1=\frac{1}{(2\pi)^3}\frac{m_e^2D_e^2+m_h^2D_h^2}{\hbar^2\rho n_0}(2k_F)^5I(\beta), \tag{2}$$

where m_e and m_h are the masses of the electron and hole, k_F is the wave vector corresponding to the Fermi momentum, $\beta=v/s$, and

$$I(\beta)=\frac{1}{\beta^2\xi^5}\int_0^\xi z^4dz\int_{1-\beta}^{1+\beta}tdt\left[\frac{1}{e^{xzt}-1}-\frac{1}{e^z-1}\right]; \tag{3}$$

here $\xi=T_0/T$ and $x=T/T_d$.

We replace t in the first term of (3) by a new integration variable $y=zxt$ and integrate by parts. We then obtain

$$I(\beta)=(3\beta^2\xi_d^2)^{-1}[\Phi(\xi_d(1+\beta),\xi)-\Phi(\xi_d(1-\beta),\xi)], \tag{4}$$

where

$$\Phi(\gamma,\xi)=\frac{1}{\gamma^3}\int_0^\gamma\frac{z(\gamma^3-z^3)}{e^z-1}dz-\frac{3}{2}\frac{\gamma^2}{\xi^5}\int_0^\xi\frac{z^4dz}{e^z-1};\quad \xi_\kappa=\frac{T_0}{T_d}. \tag{5}$$

The drop temperature T_d is determined by its velocity and by the lattice temperature, so that the parameter ξ_d in (4) is a function of β and ξ.

We now investigate a few limiting cases. At low velocities ($\beta\ll1$) we expand the function Φ in powers of β up to terms of order β^3. We must take into account here the quadratic dependence of the parameter ξ_d on β (Ref. 8). We then get for $I(\beta)$

$$I(\beta)=\frac{2}{3}\frac{\beta}{\xi^5}\int_0^\xi\frac{dz\,z^5e^z}{(e^z-1)^2}. \tag{6}$$

Expression (2) for the friction force then coincides with Keldysh's result.[5]

In the case of high temperatures ($\xi\ll1$) we have[8]

$$\xi_d=\xi(1+\beta^2/3)^{-1}.$$

At low values of γ and ξ the expansion $\Phi(\gamma,\xi)=3\gamma/4-3\gamma^2/8\xi$ is valid. We thus obtain in this case the expression

$$I(\beta)=\beta/6\xi, \tag{7}$$

which is valid at arbitrary velocity and at $\xi\gg1$. In fact, as shown by numerical calculation, formula (7) is valid, accurate to 10%, up to $\xi=1$.

We turn now to the most interesting case of low temperatures: $T\ll T_0$. We obtain the friction force at $v=s$. In this case $I(\beta)=\Phi(2\xi_d,\xi)/3\xi_d^2$. Using the asymptotic value $\Phi(\infty,\infty)=1.6$ and the expression[8] $\xi_d=0.26\xi^{5/3}$, which is valid at $\xi\gg1$ and $\beta=1$, we get

$$I(1)=8\xi^{-10/3}. \tag{8}$$

Numerical calculation shows that this formula is accurate enough at $\xi\gtrsim10$. We note that cooling the drop ($\xi_d>\xi$) decreases the friction force (without allowance for the cooling we would have $I(1)=0.55/\xi^2$ at $\xi\gg1$).

It follows from Ref. 8 that at $T\ll T_0$ the temperature of the drop is given in the entire velocity interval $v<s$, with the exception of a small region near the speed of sound, by the expression

$$T_d=(1+\beta^2/3)^{1/3}(1-\beta^2)^{2/3}T,\quad 1-\beta\gg\xi^{-5/3}. \tag{9}$$

Using (9) we can obtain the following expression for $I(\beta)$, valid at $\xi\gg1$, $1-\beta\gg\xi^{-5/3}$:

$$I(\beta)=\frac{83}{\xi^5}\frac{1+\beta^2/15}{1-\beta^2}\beta. \tag{10}$$

It should be noted that formulas (9) and (19) are accurate enough in the region $1-\beta\ll1$ only if $\xi\gtrsim20$, owing to the slow approach of the integrals in (5) to their asymptotic values.

Formula (10) describes the growth of the friction force when the drop velocity approaches that of sound. Without allowance for the drop cooling, this growth would be

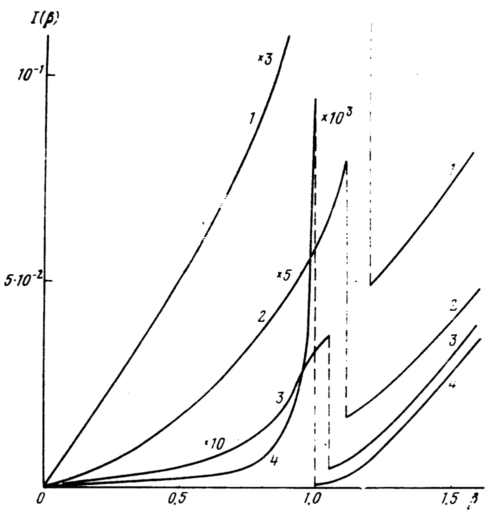

FIG. 1. Results of numerical calculation of the function $I(\beta)$ (formula (4)) for different values of the parameter $\xi = T_0/T$. 1) $\xi=3$, 2) 6, 3) 10, 4) 30. The values of $I(\beta)$ on the initial sections of curves 1 – 4 are magnified 3, 5, 10, and 10^3 times, respectively.

much faster $[I(\beta)\sim(1-\beta)^{-3}]$.

Finally, at velocities somewhat higher than that of sound we have[7] $\xi_d=2.4(\beta-1)^{-1}$ and then we get from (4) and (5), using the condition $\xi\gg1$,

$$I(\beta)=0.26(\beta-1)^2,\quad \xi^{-2/3}\ll\beta-1\ll1. \tag{11}$$

The friction force is independent of the lattice temperature in this case.

Figure 1 shows the results of a numerical calculation of the function $I(\beta)$ at different values of the parameter $\xi=T_0/T$. The plots of the kinematic friction coefficient $\gamma(\beta)=F_1(m_e+m_h)^{-1}v^{-1}$ (referred to the $\gamma(0)$ against the drop velocity for various temperatures are shown in Fig. 2. Up to values $\beta\approx0.5$ the friction coefficient grows slowly. At low temperatures and close to the speed of sound, the friction coefficient grows rapidly. At $\xi\gg1$ we have in accord with (10)

$$\frac{\gamma(\beta)}{\gamma(0)}=\frac{1+\beta^2/15}{1-\beta^2},\quad 1-\beta\ll\xi^{-2/3}.$$

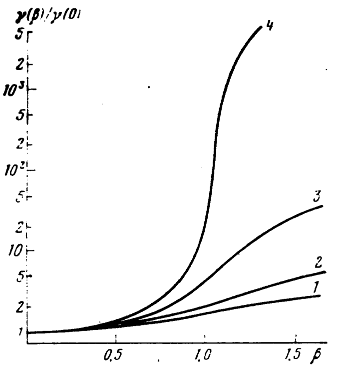

FIG. 2. The kinematic friction coefficient $\gamma(\beta)$ referred to the value $\gamma(0)$ vs. the drop velocity at various temperatures. Values of the parameter $\xi=T_0/T$: 1) $\xi=3$, 2) 6, 3) 10, 4) 30.

Curve 4 of Fig. 2 at $\xi=30$ fits this formula almost exactly up to $\beta=0.97$. At not too low temperatures, however, the growth of the friction coefficient $\gamma(\beta)$ from the value $\beta=0$ to the value $\beta=1$ is small. Thus, at $\xi=6$ (Fig. 2, curve 2), which corresponds to $T\approx2$ K for germanium, we have $\gamma(1)/\gamma(0)\approx2$.

We now estimate the absolute value of the friction force F_1 in germanium at $v=w$. The force F_1 is calculated in this section for the simplest model of a single-valley semiconductor with nondegenerate bands. It appears that the results obtained above do not make it possible to determine with sufficient accuracy the friction force in real semiconductors, but give a qualitatively correct dependence of this force on the drop velocity and on the lattice temperature. We shall therefore use the experimental data for the estimate of the absolute magnitude of the force. In germanium at $v\ll s$ and $T=2$ K the experimental value of the kinematic coefficient was found by Aleksseev *et al.*[10] to be $\gamma\approx1.6\cdot10^8$ sec^{-1}, in agreement with the results of the detailed calculation.[7] Therefore at $m_e+m_h\approx0.5\cdot10^{-27}$ g and $s=5\times10^5$ cm/sec, using the results of the numerical calculation (Fig. 1, curve 2) we get at $v=s$ the values $F_1\approx2.5$ and ≈0.7 meV/mm for $T=2$ K and 1 K, respectively.

We note that the region of supersonic velocities is apparently not attainable in experiment, in view of the strong dragging produced in this region by the coherent Cerenkov radiation of sound.

3. FRICTION FORCE DUE TO CERENKOV RADIATION OF SOUND FROM A MOVING DROP

Motion of a drop faster than sound is accompanied by Cerenkov radiation of phonons having a wavelength of the order of the drop radius, and this leads, in particular, to an additional friction force. This radiation is due to coherent interaction of the entire aggregate of electrons and holes of the drop with the deformation field. The effects connected with this interaction can be treated classically on the basis of the macroscopic equations of motion of the drop in an elastic medium. These equations can be obtained from the Lagrange function (we regard here the crystal as an isotropic elastic medium and disregard by the same token, just as in Sec. 2, the interaction with the transverse oscillations)

$$L=\frac{M\dot{\mathbf{R}}^2}{2}+\int d\mathbf{r}\left[\frac{1}{2}\rho\dot{u}_i^2-\frac{\lambda}{2}u_{ii}u_{kk}-\mu u_{ik}u_{ik}-Dn(\mathbf{r}-\mathbf{R})u_{ii}\right]. \tag{12}$$

Here M is the mass of the drop, $\mathbf{R}$ is the radius vector of the center of mass of the drop, ρ is the density of the crystal, $\mathbf{u}(\mathbf{r})$ is the lattice displacement vector, u_{ik} is the strain tensor, λ and μ are the Lamé coefficients, $D=D_e+D_h$ is the total deformation potential of the electrons and holes, and $n(\mathbf{r})$ is the density of the electron-hole liquid ($n=n_0$ inside the drop and $n=0$ outside).

The equations of motion are

$$M\ddot{R}_i=-D\frac{\partial}{\partial R_i}\int d\mathbf{r}\,n(\mathbf{r}-\mathbf{R})\frac{\partial u_j}{\partial x_j}, \tag{13}$$

$$\rho\ddot{u}_i-(\lambda+\mu)\frac{\partial}{\partial x_i}\frac{\partial u_j}{\partial x_j}-\mu\frac{\partial^2u_i}{\partial x_j\partial x_j}=D\frac{\partial}{\partial x_i}n(\mathbf{r}-\mathbf{R}). \tag{14}$$

The right-hand side of (13) is the force exerted on the drop by the lattice. This force can contain a term connected with the external action on the lattice (in which case $\partial u_j/\partial x_j$ in (13) is determined by solving Eq. (14) without the right-hand side, with suitable boundary conditions). We consider the force $\mathbf{F}^*$ due to the deformation of the lattice by the drop itself.

It is convenient to express the force $\mathbf{F}^*$ in terms of the Fourier components of the displacement vector $\mathbf{u}(\mathbf{k})$:

$$\mathbf{F}=D\int\frac{d\mathbf{k}}{(2\pi)^3}\mathbf{k}(\mathbf{k}\mathbf{u}(\mathbf{k}))^* n(\mathbf{k})e^{-i\mathbf{k}\mathbf{R}(t)}, \tag{15}$$

where $n(\mathbf{k})$ is the Fourier component of the density $n(\mathbf{r})$. We obtain $\mathbf{u}(\mathbf{k},t)$ from (14) and substitute it in (15). Then

$$\mathbf{F}=-\frac{D^2}{\rho s}\int\frac{d\mathbf{k}}{(2\pi)^3}\mathbf{k}k|n(\mathbf{k})|^2\int\limits_0^\infty d\tau\sin ks\tau\sin(\mathbf{k}[\mathbf{R}(t)-\mathbf{R}(t-\tau)]), \tag{16}$$

where $s=((\lambda+2\mu)/\rho)^{1/2}$ is the longitudinal sound velocity.

The force $\mathbf{F}$ produced by the action of the drop in itself via the deformation field has a complicated dependence on the character of the drop motion (of the function $R(t)$). In uniform subsonic motion we have $\mathbf{F}=0$. If the drop acceleration is small enough, the force $\mathbf{F}$ can be expanded in a series of components proportional to $\ddot{\mathbf{R}}$, $\dddot{\mathbf{R}}$, etc. The term proportional $\ddot{\mathbf{R}}$ is then the increment to the inertia force, i.e., it determines the change of the mass of the drop (see Sec. 4). The term containing $\dddot{\mathbf{R}}$ yields the radiation-friction force.

At supersonic speed the force $\mathbf{F}$ differs from zero even if the drop motion is uniform. This is caused by the Cerenkov radiation of the sound. It is obvious that the force $\mathbf{F}$ is directed in this case opposite to the velocity $\mathbf{v}=\dot{\mathbf{R}}$, and constitutes the friction force. Putting $\mathbf{R}=\mathbf{v}t$ in (16), we get

$$\mathbf{F}=-\frac{\pi D^2}{\rho s}\int\frac{d\mathbf{k}}{(2\pi)^3}\mathbf{k}k|n(\mathbf{k})|^2\delta(\mathbf{k}\mathbf{v}-ks). \tag{17}$$

It is seen that $\mathbf{F}=0$ at $v<s$. At $v>s$ we calculate the integral (17) under the assumption that the drop is a sphere of radius r_0 with constant density n_0. Dividing the result by the number of electron-hole pairs in the drop $N=4\pi r_0^3n_0/3$, we obtain the friction force F_2 per pair[2)]:

$$F_2=\frac{3}{2}\frac{(D_e+D_h)^2n_0}{\rho r_0v^2}\ln(k_Fr_0),\quad v\geqslant s. \tag{18}$$

Thus, when the speed of sound is reached, a Cerenkov friction force F_2 is produced jumpwise, and subsequently decreases like $\sim 1/v^2$. The fact that the force F_2 has a finite value at the threshold of the Cerenkov radiation is peculiar to interactions with longitudinal waves.

We note that in virtue of the macroscopic nature of the Cerenkov friction force F_2, its value is insensitive, unlike the force F_1, to the details of the band structure of the semiconductor. An estimate of the absolute value of the force F_2 for germanium ($D_e+D_h\approx 5$ eV, $n_0\approx 2\times 10^{17}$ cm^{-3}, $\rho=5.3$ g/cm^3, $r_0=10^{-4}$ cm, $v=s=5\times 10^5$ cm/sec) yields $F_2=47$ meV/mm. This is much higher than the low-temperature values of the force F_1 at $v=s$. In experiments on the motion of electron-hole drops, the external forces per pair usually do not exceed several meV/mm. It appears that the Cerenkov force makes the sound barrier unsurmountable.

4. DEFORMATION MASS OF ELECTRON-HOLE DROP

The motion of an electron hole drop in a crystal is accompanied by a displacement of the deformation field produced by the drop. Therefore acceleration of the drop creates a lattice reaction force proportional to the acceleration; this is equivalent to the appearance of a certain additional mass due to the interaction of the drop with the lattice. It is natural to call this additional mass the deformation mass. The increased mass of the drop is a classical analog of the polaron effect. It can be described by using the macroscopic approach developed in the preceding section.

Let us calculate the force $\mathbf{F}$ exerted on the drop by the lattice at subsonic velocities with small accelerations. To this end, we expand the difference $\mathbf{R}(t)-\mathbf{R}(t-\tau)$ in (16) in powers of up to terms of second order inclusive:

$$\mathbf{R}(t)-\mathbf{R}(t-\tau)=\mathbf{v}\tau-\dot{\mathbf{v}}\tau^2/2.$$

The first term of the expansion of the force $\mathbf{F}$ in the small acceleration $\dot{\mathbf{v}}$ is

$$\mathbf{F}=\frac{D^2}{2\rho s}\int\frac{d\mathbf{k}}{(2\pi)^3}\mathbf{k}k(\mathbf{k}\dot{\mathbf{v}})|n(\mathbf{k})|^2\int\limits_0^\infty\tau^2d\tau\sin ks\tau\cos\mathbf{k}\mathbf{v}\tau. \tag{19}$$

This expansion is valid if

$$|\dot{\mathbf{v}}|\ll(1-\beta)^2s^2/r_0. \tag{20}$$

Formula (19) can be represented in the form

$$\mathbf{F}=-\frac{d}{dt}\frac{\partial\tilde{L}}{\partial\mathbf{v}}, \tag{21}$$

$$\tilde{L}=\frac{D^2}{2\rho s}\int\frac{d\mathbf{k}}{(2\pi)^3}k|n(\mathbf{k})|^2\int\limits_0^\infty d\tau\sin ks\tau\cos\mathbf{k}\mathbf{v}\tau. \tag{22}$$

The quantity $\partial\tilde{L}/\partial\mathbf{v}$ is the additional drop momentum due to the field of the deformation that accompanies the drop. It is easy to verify that the function $\tilde{L}(\mathbf{v})$ is the integral term of the Lagrange function (12), calculated by substituting in it the solution of Eq. (14) for the case of uniform motion of the drop with velocity $\mathbf{v}$.

Direct calculation of the function $\tilde{L}$ yields

$$\tilde{L}=N\frac{D^2n_0}{4\rho s^2}\frac{1}{\beta}\ln\frac{1+\beta}{1-\beta}, \tag{23}$$

where N is the number of electron-hole pairs in the drop.

With the aid of (21) and (23) we can express the force $\mathbf{F}$ in the form

$$\mathbf{F}=-\frac{d}{dt}Nm_d(\beta)\mathbf{v}. \tag{24}$$

Here $m_d(\beta)$ is the velocity-dependent deformation mass per electron-hole pair:

$$m_d(\beta)=m_d(0)\frac{3}{2\beta^2}\left(\frac{1}{1-\beta^2}-\frac{1}{2\beta}\ln\frac{1+\beta}{1-\beta}\right), \tag{25}$$

$$m_d(0)=\frac{(D_e+D_h)^2 n_0}{3\rho s^4}. \qquad (26)$$

The quantity $m_d(0)$ is the deformation rest mass. For germanium, $m_d(0)=1.3\cdot 10^{-29}$ g, so that at low velocities the drop mass increases only by 3%. In crystals, where the density of the electron-hole liquid in the drop is much higher, the deformation mass of the drop can be comparable with its ordinary mass.

The increase of the drop mass considered here is connected with coherent interaction of all the particles of the drop with the lattice. Consequently m_d greatly exceeds (by a factor $(\hbar k_F/ms)^3$) the additional mass acquired by an individual pair as a result of the ordinary polaron effect.

According to (25), the deformation mass increases with increasing velocity, and becomes formally infinite at $\beta=1$. It must be borne in mind, however, that the mass m_d can come into play only at sufficiently large drop accelerations, when the inertia force is comparable with the friction force, i.e., $|\dot{\mathbf{v}}| \gtrsim \gamma v$ (γ is the kinematic friction coefficient). On the other hand, introduction of the deformation mass is possible only at sufficiently small accelerations that satisfy the inequality (29). Therefore formula (25) is meaningful for velocities such that $(1-\beta)^2 \gg \beta\gamma r_0/s$.

[1)]Direct calculation of the friction force by formula (1) (without explicit allowance for the condition $dE/dt=0$) leads to a factor $(1-t)$ under the integral sign in formula (3). If we put $x=1$, then formulas (2) and (3) coincide in fact with the corresponding formulas of Manoliu and Kittel,[4] who disregarded the change of the drop temperature in the course of its motion.

[2)]For a drop with sharp boundaries, the integral (17) diverges at large k. Actually the density $n(r)$ drops off to zero at the drop surface in a layer with a thickness of the order of $1/k_F$. The limit integration with respect to k should therefore be k_F, and it is this which leads to the logarithmic factor in (18).

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Translated by J. G. Adashko

Erratum: Inverted hot-electron states and negative conductivity in semiconductors [Sov. Phys. JETP 45, 539 (1977)]

Ya. I. Al'ber, A. A. Andronov, V. A. Valov, V. A. Kozlov, A. M. Lerner, and I. P. Ryazantseva

Zh. Eksp. Teor. Fiz. 75, 1952 (November 1978)

PACS numbers: 72.30. + q, 72.20.Dp, 99.10. + g

Formula (3.7) has the wrong sign and should be corrected to

$$\mathrm{Re}\,\sigma_{\parallel}=\ldots=-\bar{\sigma}F(\theta).$$

Thus, $\mathrm{Re}\,\sigma_{\parallel}\geqslant 0$.

0038-5646/78/110984-01$02.40