Nonlinear restriction imposed on the coherent acoustic fields which counteract the spreading of electron-hole plasma at transonic velocities

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(Submitted 5 December 1986)
Pis'ma Zh. Eksp. Teor. Fiz. **45**, No. 6, 288–291 (25 March 1987)

The initial concentration of electron-hole pairs is estimated theoretically. This concentration must be known in order for the expanding-plasma shock front to break the sound barrier.

Experimental results of many studies indicate that the electron-hole plasma photoproduced near the surface of a semiconductor spreads at supersonic[1,2] velocities and even Fermi[3,4] velocities. However, the theoretical models,[5,6] which describe the spreading of plasma from the surface of a crystal due to the internal pressure forces, ignore the coherent acoustic fields that resist this motion. We know,[7] on the other hand, that the inverse effect of the strain-induced waves, which are effectively excited when electron-hole drops move at transonic velocities, so far could not be used to accelerate the drops to supersonic velocities. In the present letter we show that if the electron-hole plasma layer is expanded, the sound barrier can be broken by imposing a nonlinear restriction on the amplitude of the quasisynchronously excited acoustic waves.

We will use a step-by-step approach to analyze the interaction, in terms of the strain energy, of the compressed electron-hole plasma with the coherent acoustic fields. In the first step we will consider the generation of acoustic waves by an incompressible plasma shock front. The distribution of the concentration n of nonequilibrium carriers is assumed to be

$$n = n_{\rm fr} \, [1 - \Theta(x - V_{\rm fr} \, t)] , \qquad (1)$$

where x is the coordinate directed along the axis perpendicular to the semiconductor surface, t is the time, $V_{\rm fr} \simeq c_0$ is the velocity of the shock front, c_0 is the velocity of the longitudinal sound, $n_{\rm fr}$ is the plasma density at the shock front, and $\Theta(x)$ is a unit function $[\Theta(x \leqslant 0) = 0,\ \Theta(x > 0) = 1]$. When the sources are moving at transonic velocities, they excite most effectively (quasisynchronously) the acoustic waves that are propagating in the same direction. We can write a simplified equation which describes the generation of strain-induced waves that accompany the electron-hole plasma front in the form (compare with the equation in Ref. 8)

$$D_t - \Delta D_\xi - \epsilon c_0 D D_\xi = -d\,(2\rho_0 c_0)^{-1} n_\xi \, . \qquad (2)$$

Here D is the deformation of the crystal, $\xi = x - V_{\rm fr} t$ is the accompanying coordinate, $\Delta = V_{\rm fr} - c_0$ is the deviation from an exact linear matching condition, ϵ is a

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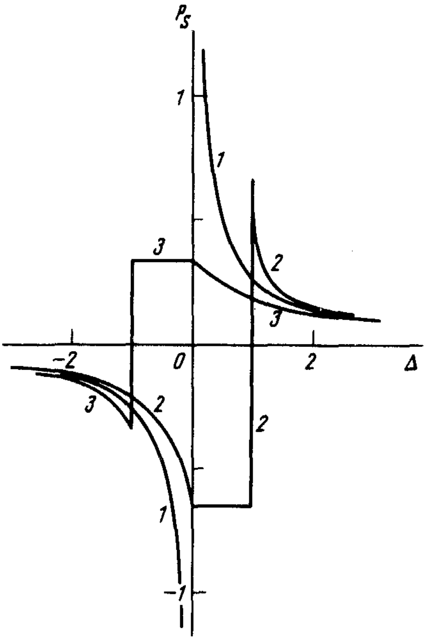

FIG. 1. The pressure exerted on the electron-hole plasma front by the acoustic field versus the velocity of the front. 1—P_S ($\epsilon = 0$); 2—P_S ($\epsilon > 0, d < 0$); 3—P_S ($\epsilon > 0, d > 0$).

nonlinear acoustic parameter,[9] ρ_0 is the equilibrium crystal density, and d is the strain-energy constant of an electron-hole pair.[7]

If the acoustic-wave profiles which satisfy Eq. (2) are given, the pressure P_S of this sound field on the electron-hole plasma can be calculated:

$$P_S = -d \int_{-\infty}^{\infty} d\xi\, n D_\xi \,. \tag{3}$$

To simplify the notation used below, we will switch, assuming $\epsilon > 0$, to dimensionless variables: $D = D/D_0$ $[D_0 = (|d|n_{fr}/\epsilon\rho_0)^{1/2}/c_0]$, $\Delta = \Delta/\Delta_0$ $(\Delta_0 = \epsilon c_0 D_0)$, $P_S = P_S/P_0$ $(P_0 = |d|n_{fr} D_0)$. We will use, in addition, only the equations for the case $d < 0$, $\Delta \leqslant 0$ and represent the other results graphically (Fig. 1).

If the acoustic nonlinearity ($\epsilon = 0$) is ignored, the solution of problem (1), (2), which satisfies the initial condition $D(\xi, t = 0) = 0$, is

$$D(\Delta \leqslant 0,\ \epsilon = 0) = (2\Delta)^{-1} \{\Theta[\xi] - \Theta[\xi + 2\Delta t]\} \,. \tag{4}$$

Using this solution, we find the pressure which retards the motion of the plasma,

$$P_S(\Delta \leqslant 0,\ \epsilon = 0) = (4\Delta)^{-1} \,. \tag{5}$$

According to (4) and (5), the acoustic fields and their pressure on the electron-hole plasma increase without bound as the sound barrier is approached, rendering it insurmountable. Analysis has shown that allowance for the finite nature of the width of the front in this case does not eliminate (in contrast with the theory of motion of electron-hole drops[10]) the divergence of (4) and (5) in the limit $\Delta \to 0$. It is for this reason that the nonlinear interaction of the strain-induced waves plays the key role.

In the case of ($\epsilon \neq 0$) Eq. (2) describes the excitation of finite-amplitude acoustic waves:

$$D(\Delta \leqslant 0,\ \epsilon > 0) = -(\sqrt{\Delta^2+1}+\Delta)\ \{\ \Theta[\xi] - \Theta[\xi + (\sqrt{\Delta^2+1}+\Delta)^{-1}\ t]\}. \qquad (6)$$

Solution (6), which can be obtained on the phase plane,[8] describes the saturation of quasisynchronously excited acoustic waves: Even when the sources are moving at the velocity of sound, the buildup of acoustic waves occurs only until the change in the propagation velocity of the wave, which is associated with the increase in the amplitude, drives the wave outside the generation region. A nonlinear restriction of the strain fields imposes restrictions on the sound barrier:

$$P_S(\Delta \leqslant 0,\ \epsilon > 0) = -\ \{\Delta + \frac{2}{3}\ [\Delta^3 + (\Delta^2+1)^{3/2}]\ \}. \qquad (7)$$

Relations (5) and (7) and the results of similar calculations for $\Delta > 0$ and $d > 0$ are shown in Fig. 1.

An important feature of solution (6) is that at subsonic velocities of the wave front all acoustic waves that are excited overtake the plasma. This circumstance makes it possible to take into account in the analysis of the expansion of the compressible plasma the interaction of this plasma with the strain waves, by modifying the boundary conditions at the shock front ($x = x_{\rm fr}$).

In the second step we will analyze the gasdynamic problem involving the decay of the initial rupture point.[5,11] We assume that as a result of interband absorption of an extremely short optical pulse near the surface ($x \geqslant 0$), at the initial time we will have a plasma distribution $n = n_0[1 - \Theta(x-L)]$ with a concentration higher than the equilibrium concentration n^* of an electron-hole liquid (L is the thickness of the plasma slab). At $t > 0$ the plasma then begins to expand due to the internal-pressure forces. Before the self-similar expansion wave that travels along the plasma begins to reflect from the crystal surface $[0 \leqslant t \leqslant L/c(n_0)]$, the hydrodynamic velocity v of the electron-hole plasma is described by the relation[5,11]

$$v = \int_n^{n_0} c\ (n')\,(n')^{-1}\ dn'\ . \qquad (8)$$

Here $c(n)$ is the acoustic velocity in the plasma: $c^2(n) = M^{-1}\partial p/\partial n$, the internal pressure $p(n)$ is related to the energy E of the electron-hole pair by the relation $p = n^2 \partial E/\partial n$, and M is the mass of the electron-hole pair.[5,7] The shock front propagates deep into the crystal. The plasma velocity at the shock front ($x = x_{\rm fr}$) is given by Eq. (8):

$$v\,(x_{\rm fr}) = \int_{n_{\rm fr}}^{n_0} c\ (n)\,n^{-1}\ dn\ . \qquad (9)$$

Since in the case at hand[5] the plasma expands into the vacuum $[n(x > x_{\rm fr}) = 0]$, by

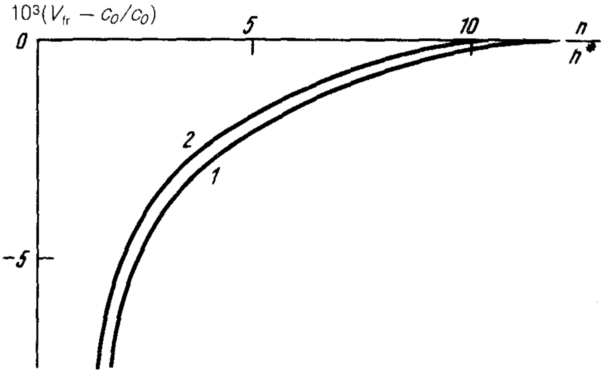

FIG. 2. The velocity of the shock front of the electron-hole plasma versus its concentration in germanium. 1—$V_{fr}(n_0)$; 2—$V_{fr}(n_{fr})$.

virtue of the boundary conditions at the shock front we have

$$V_{fr} = v(x_{fr}), \tag{10}$$

$$p(n_{fr}) = P_0 |P_S|. \tag{11}$$

The second relation, which takes into account the pressure exerted on the plasma by the acoustic field, $P_0 P_S$, is a corollary of the continuity of the momentum flux.[11] According to (11), an excitation of acoustic waves at subsonic plasma-expansion velocities thus accounts for the dependence of the shock-front velocity on the plasma concentration at the front. Figure 2 shows the velocity of the electron-hole plasma shock front plotted as a function of the initial concentration of the plasma and as a function of the plasma concentration at the shock front in germanium ($d < 0$). The two curves were calculated from Eqs. (9)–(11). The analytic functions were found by using the $E(n)$ relation, in which the correlation energy was ignored [$E(n) \cong an^{2/3} - bn^{1/3}$] and the following values of the physical quantities: $|d| \cong 7$ eV, $M \cong 0.4 \times 10^{-27}$ g (Ref. 1), and $\epsilon \cong 14.6$. The kinetic and exchange energies were estimated in the Hartree-Fock approximation, with allowance for the actual band structure of Ge (Ref. 12): $a \cong 1.6 \times 10^{-26}$ erg·cm^2 and $b \cong 1.4 \times 10^{-20}$ erg·cm. According to our calculations (Fig. 2), the initial plasma concentration (n_0) must therefore be only an order of magnitude higher than the equilibrium concentration n^* in order to break the sound barrier. Estimates show that this assertion also applies to silicon ($d > 0$). Equating the flux of electron-hole pairs [$\sim n_{fr}(V_{fr}) V_{fr}$] to the flux of light quanta that are absorbed [$\sim (1 - R) \times (h\nu)^{-1} I_0$, where I_0 is the light intensity, R is the coefficient of the reflection from the surface, and $h\nu$ is the energy of the light quantum], we can crudely calculate the effect of a pulsed laser on the supersonic plasma expansion. Using the results obtained above, we find that when Ge absorbs light with a wavelength of 1.06 μm ($R \cong 0.4$), the plasma front breaks the sound barrier at $I_0 \gtrsim 10^5$ W/cm^2, consistent with the results of the experiment of Tamor *et al.*[1]

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Translated by S. J. Amoretty