We note that the logarithm of the dislocation velocity is half as large as ln w, a most important fact in the estimate of exponentially small effects.

Attempts to elucidate the role of quantum effects in the overcoming of Peierls barriers by dislocations were made in [6] on the basis of non-rigorous assumptions. These estimates are not confirmed by our calculations, however.

The authors are grateful to S.I. Anisimov and S.V. Iordanskii for useful discussions.

[1] I.Lothe and I.P. Hirth, Phys. Rev. 115, 543 (1959).
[2] A.P. Kazantsev and V.L. Pokrovskii, Zh. Eksp. Teor. Fiz. 58, 677 (1970) [Sov. Phys.-JETP 31, 362 (1970)].
[3] I.M. Lifshitz and Yu.M. Kagan, ibid. 62, 385 (1972) [35, No. 1 (1972)].
[4] S.V. Iordanskii and A.M. Finkel'shtein, ibid. 62, 403 (1972) [35, No. 1 (1972)].
[5] P. Guyot and J.E. Dorn, Can. J. of Phys. 45, 983 (1967).
[6] A.C. Arko and J. Weertman, Acta Met. 17, 687 (1969).

NEW METHOD OF AMPLIFYING MONOCHROMATIC ULTRASOUND IN A SEMIMETAL

S.Ya. Rakhmanov
L.D. Landau Institute of Theoretical Physics, USSR Academy of Sciences
Submitted 22 November 1971
ZhETF Pis. Red. 15, No. 1, 66 - 69 (5 January 1972)

In [1] there were discussed the conditions under which an intense sound wave of high frequency can strongly drag the carriers of one of the valleys of a semimetal or of a degenerate semiconductor, i.e., produce a particle drift with an average velocity on the order of the sound velocity s. In this communication we assume that a constant electric field E is applied along the wave propagation direction. Within the framework of the model (see below) assumed for the description of this situation, the following result is obtained: the initial wave is amplified (or is damped if the direction of E is reversed), and the entire power jE acquired by the carriers from the field is transferred to the wave (j is the total current of the particles of one valley, generated both by the sound wave and by the field).

Just as in [1], it is assumed that the following conditions are satisfied for the carriers of one valley: 1) A strong magnetic field, sufficient to obtain a one-dimensional spectrum [2], $\hbar\Omega \gtrsim \varepsilon_F$, is applied along the wave propagation direction (Ω is the Larmor frequency and ε_F the Fermi energy). 2) The sound intensity is high enough, $\Delta >> \hbar\omega$, where Δ is the energy of the interaction of the carriers with the wave and ω is the sound frequency. 3) The sound wavelength λ satisfies the condition $(2\pi\hbar/\lambda) - 2p_F \lesssim p_F(\Delta/\varepsilon_F)$, where p_F is the particle momentum on the Fermi surface. 4) The temperature is low, $kT \lesssim \Delta$. We assume also that $\omega > \nu_n$ (ν_n is the frequency of the collisions with the impurities) and $\Delta << \varepsilon_F$.

We use the following model: a one-dimensional gas of charged particles interacts (for concreteness, via a deformation potential with constant Λ) with a classical sound wave and is elastically scattered by the impurities. The corresponding Hamiltonian is

$$\hat{H} = \int dx \left\{ \frac{d}{2} [U_t'^2 + S^2(U_x')^2] + n [\hat{\psi}^+ \Lambda \frac{\partial U}{\partial x} \hat{\psi} + \hat{\psi}^+ \frac{\hat{p}^2}{2m} \hat{\psi} + \hat{\psi}^+ V_{\text{np}} \hat{\psi} + \hat{\psi}^+ eE\hat{x}\hat{\psi}] \right\}. \quad (1)$$

Here $U(x, t)$ describes the wave, d is the crystal density, m the effective carrier mass, ψ^+ and $\hat{\psi}$ the creation and annihilation operators, n the carrier-number density, and V_i the impurity potential. In this model, at $E = 0$, the total energy of the particles and of the sound wave is conserved, and consequently, when the sound is attenuated the particle energy increases. At $E \neq 0$, the carriers acquire energy from the electric field and transfer part of this energy (determined by the parameters of the problem) to the sound wave. In the cases of interest to us (see conditions 1 - 4) the carrier spectrum has a gap located somewhat above the Fermi surface and almost impenetrable to the particles, so that the carriers cannot absorb energy: at $E = 0$ the sound attenuation is very small, and at $E \neq 0$ the entire power jE is transferred to the wave.

To explain the foregoing statements, let us describe briefly the solution of the problem. We assume that there exists an undamped sound wave. Substituting $U(x, t) = A \sin (kx - \omega t)$, we obtain from (1) an equation for the carrier density matrix[1)]

$$i\hbar\dot{\hat{\rho}} = [\hat{H}_0\ \rho] + [eE\hat{x}, \hat{\rho}] + W\hat{\rho}. \tag{2}$$

The operator W (the collision integral) describes the averaged scattering by the impurities (see [1]). It is convenient to solve Eq. (2) in a coordinate system moving with the wave, where H_0 takes the form

$$\hat{H}_0 = \frac{\hat{p}^2}{2m} + S\hat{p} + \Delta\cos(k\ \ x), \quad \Delta = \Lambda k A . \tag{3}$$

The spectrum of the operator H_0 has a gap Δ with a center at the point $\varepsilon_0 = \hbar^2k^2/8m$. In the coordinate system coupled to the wave, the impurities move with velocity s, and therefore the carrier scattering is accompanied by a change of energy. The most frequent (with frequency ν_n) is scattering with an energy change $\sim\hbar\omega$ (see the figure in [1]). Transitions through the shell (with an energy change $\sim\Delta$) are very rare, and have a frequency $\nu_n(\Delta/\varepsilon_F)(2\Delta/\hbar\omega)$.

In an external electric field, Zener tunneling through the gap takes place [3], but when the condition $eE << \Delta^2k/\varepsilon_F$ is satisfied, its probability is exponentially small[2)]. By virtue of the foregoing circumstances, the carriers initially located below the gap (assume $T = 0$, $\varepsilon_0 - \varepsilon_F \gtrsim \Delta/2$) will continue to stay there. The change of the particle distribution over the levels is described by Eq.(2), which reduces, when the condition $\Delta >> \hbar\omega$ is satisfied, to the diffusion equation if the transitions through the gap are neglected and the integral operator W is expanded in terms of the small parameter. Within a time on the order of $\nu_n^{-1}(\Delta/\hbar\omega)^2$ a stationary distribution is established, together with a large stationary current

$$j \sim neS , \tag{4}$$

which depends on E ($j_{st}(E) \geq j_{st|E=0}$, and the current j_{st} in (4) is calculated in the laboratory frame). Subsequently the total energy of the carriers remains unchanged, and the absorbed power $j_{st}E$ is transferred entirely to the sound and

---

1)Owing to the interaction with the carriers, the wave amplitude changes within a time longer than the characteristic times of Eq. (2), and consequently can be assumed constant in (2) if $\alpha = n\Lambda^2/\varepsilon_F dS^2 << 1$. In bismuth, $\alpha \sim 10^{-3}$.

2)This condition is satisfied with a large margin, for example, by the fields customarily employed for sound amplification in semiconductors, $E \sim mS/e\tau$, where τ is the free-path time, if $\tau \sim \nu_n^{-1}$.

goes mostly to increase the amplitude of the initial wave (the harmonics are generated at a lower rate, $\sim(\Delta/\varepsilon_F)^{\ell}$, where ℓ is the number of the harmonic). The growth of the wave amplitude is proportional to the nondiagonal element of the density matrix $\rho_{p,p-k}$, which should be determined from (2). We can, however, use the energy conservation law directly.

In the particular case when all the levels below the gap are filled (i.e., $\varepsilon_F = \varepsilon_0 - \Delta/2$), all the calculations become exceedingly simple. In this case, j_{st} = neS and is independent of E, the power transferred to the sound is neSE and the wave amplitude increases like

$$A(t) = \sqrt{A^2(0) + \frac{2enSE}{dS^2k^2}t}\,.$$

We disregard the processes that limit the growth of A(t).

The model used by us does not reflect all the processes in real semimetals or semiconductors. We have neglected the interaction of the sound with carriers of other valleys (bearing in mind that it is possible to produce conditions when requirements 1 - 4 are satisfied for only one valley: for example, requirement 1 in bismuth is satisfied for holes in stronger fields than for electrons). No account was taken of carrier scattering by thermal phonons and zero-point lattice oscillations, or of intervalley scattering (these processes are less effective than intravalley scattering by impurities, and do not change the results). An important role can apparently be played by interparticle interactions. In spite of the relative weakness, they influence both the rate of the particle transition through the gap, and the distribution under the gap. These questions will be considered in another paper.

The author thanks M.Ya. Azbel' and E.I. Rashba for a discussion of the work.

[1] S.Ya. Rakhmanov, ZhETF Pis. Red. 12, 266 (1970) [JETP Lett. 12, 181 (1970)].
[2] N.B. Brandt and E.A. Svistova, Usp. Fiz. Nauk 101, 249 (1970) [Sov. Phys.-Usp. 13, 370 (1970)].
[3] J.M. Ziman, Theory of Solids, Cambridge Univ. Press, Chap. 6, Sec. 8.

POSSIBILITY OF GENERATING ULTRASHORT LIGHT PULSES IN LASERS WITH SMALL LUMINESCENCE LINE WIDTH OF THE MEDIUM

V.N. Lugovoi and A.M. Prokhorov
P.N. Lebedev Physics Institute, USSR Academy of Sciences
Submitted 25 November 1971
ZhETF Pis. Red. 15, No. 1 70 - 72 (5 January 1972)

A theory of stimulated Raman emission (SRE) in an optical resonator excited by an external longitudinal monochromatic light beam was developed in [1 - 3]. It was shown, in particular, that if the resonator is filled with a liquid or a solid the phases of the oscillations of the different SRE components are arbitrary, i.e., they are not interconnected. A theoretical investigation [4] of the case of a gas-filled resonator has shown that the phases of the SRE components are interconnected, and the output radiation is a sequence of ultrashort pulses with a spectrum width determined by the number of generated components. An investigation [5] of stationary stimulated Mandel'shtam-Brillouin emission (SMBE) under similar conditions has shown that the phases of the different SMBE components are also interconnected if the resonator contains