proportional to $m^{-5/2}$ [1], although the fraction of the light holes is determined by the mass ratio $(m_1/m_2)^{3/2}$ and amounts to only 4%. Second, the light holes interact strongly with the optical phonons, the number of which decreases exponentially with temperature, leading to a stronger temperature dependence of the mobility than the predicted $\mu(T) \sim T^{-1.5}$.

From the fact that the light holes makes a large contribution to the conductivity it can be concluded that their quantization should also lead to the appearance of longitudinal magnetoresistance, as was indeed observed by us in the experiment.

Figure 2 shows plots of the relative change of conductivity, $-\Delta\sigma/\sigma_0$, in a magnetic field H for two temperatures, 77 and 155°K. We see that $-\Delta\sigma/\sigma_0$ first changes rapidly with increasing magnetic field, and the change slows down when H > 60 kOe, but the behavior of the curves in the region of rapid variation of $-\Delta\sigma/\sigma_0$ is different for 77 and 155°K. Doubling the temperature causes the rate of the rapid change of $-\Delta\sigma/\sigma_0$ as a function of H to be approximately half as large for T = 155°K as for T = 77°K; this should indeed be the case if the change of the conductivity is due to the quantization of the light-hole band and is a function of $\eta = (\hbar\omega_c/kT)$.

The slow change of $-\Delta\sigma/\sigma_0$ when H > 60 kOe can be attributed to a certain nonsphericity of the constant-energy surface of the heavy holes [4].

The authors are grateful to V.A. Chuenkov for a discussion of the results.

[1] E.M. Conwell, High Field Transport in Semiconductors, Academic Press, 1967.
[2] V.G. Veselago, L.P. Maksimov, and A.M. Prokhorov, Prib. Tekh. Eksper. No. 4, 192 (1968).
[3] S.V. Kironov, E.G. Rudashevskii, and V.I. Chernykh, ibid. No. 5, 192 (1969).
[4] C. Kittel. Quantum Theory of Solids (Russ. trans.) pp. 285-286, Nauka, 1967 [Wiley, 1963].

PARAMAGNETIC STRUCTURE DUE TO EXCHANGE INTERACTION OF ELECTRONS IN A METAL

V.P. Silin
P.N. Lebedev Physics Institute, USSR Academy of Sciences
Submitted 16 March 1970
ZhETF Pis. Red. 11, No. 8, 419 - 421 (20 April 1970)

The theoretical prediction of diamagnetic domains and periodic diamagnetic structures in normal metals, and the subsequent experimental observation of their manifestations [1, 2] has resulted by now in a large number of theoretical papers (cf., e.g., the review [3]) providing a definite understanding of the nature of such a structure. An important feature of the domains discussed in this case is that their nature is governed to a considerable degree by the quantization of the electron levels in the magnetic field. Therefore, the diamagnetic structures are connected with quantum de Haas - van Alphen magnetization oscillations. In this paper we indicate another possibility, of a different physical nature, of the existence of inhomogeneous magnetic structures in normal metals. Such structures are possible under conditions when the quantization of the orbital motion of the electrons is insignificant, and consequently the magnetization oscillations are likewise insignificant. At the same time, the magnetic structures discussed below can occur in sufficiently strong fields.

An indication of the existence of a spatially inhomogeneous static state of a metal with different values of magnetization at different points of space can be obtained with the aid of the dispersion equation for the electromagnetic waves in a metal. Namely, the presence of solutions of such an equation in the

limiting case of zero wave frequency ω, but at a finite wave vector $\vec{k}$, corresponds to the conditions of the existence of periodic structures, as in the case, for example, in the theory of the antiferromagnetism of chromium [4 - 6]. We, on the other hand, are interested in magnetic structures arising under the influence of a sufficiently strong magnetic field. In other words, we must consider waves in a metal situated in a magnetic field. We can therefore use the results of the theory of waves in normal metals, based on the theory of a degenerate electron fluid [7, 8].

We shall use henceforth a spin-wave dispersion equation in the approximation with one constant B_0 characterizing the exchange interaction. We neglect here the electron collisions; this is equivalent to stipulating a sufficiently strong magnetic field in which the gyroscopic rotation frequency Ω of the electron is much higher than the electron collision frequency. We then get for waves whose magnetization is perpendicular to the homogeneous magnetic field [8]

$$\frac{B_0\omega}{1+B_0}\frac{1}{2}\int_0^{\pi}\sin\theta\,d\theta\sum_{n=-\infty}^{+\infty}I_n^2\left(\frac{k_\perp v}{\Omega}\sin\theta\right)\frac{1}{\omega\pm\Omega_0-n\Omega-k_z v\cos\theta}=1.$$

Here v is the Fermi velocity of the electron, Ω the effective frequency corresponding to the spin levels of the electron, and k_z and $k_\perp$ are resp ctively the projections of the wave vector on the direction of the homogeneous magnetic field and on a plane perpendicular. Finally, the ± sign corresponds to the two different circular polarizations of the magnetic field of the spin wave.

Solution of this dispersion equation for the frequency, in the limit $k_z v \to \pm\Omega_0$, corresponds to vanishing of the spin-wave frequency in accordance with the equation

$$\omega=\pm\left\{\Omega_0\mp k_z v\left(1+2\exp\left[\frac{1+B_0}{B_0}\frac{2\Omega_0}{\Omega_0\mp k_z v}+f\left(\frac{k_\perp v}{\Omega}\right)\right]\right)\right\},$$

where

где

$$f(x)=2x\int_{-1}^{+1}z\,dz\,\frac{\ln(1-z)}{\sqrt{1-z^2}}\,I_0(x\sqrt{1-z^2})\,I_1(x\sqrt{1-z^2})\pm$$

$$\pm\Omega_0\sum_{n\neq 0}\int_{-1}^{+1}\frac{dz}{\pm\Omega_0(1-z)-n\Omega+i0}\,I_n^2(x\sqrt{1-z^2}).$$

It is necessary here to satisfy the inequality $1+B_0^{-1}<0$, which indeed holds for sodium and potassium.

Consequently, we can speak of a spatially inhomogeneous magnetic state with a period along the direction of the external magnetic field equal to $2\pi v/\Omega_0$. We emphasize that the periodicity is determined by the spin frequency of the electrons. In addition, unlike the diamagnetic structures discussed in [3], there is periodicity along the direction of the external magnetic field. It should be noted that the spin-wave dispersion equation employed by us corresponds to singularities of the spin paramagnetic susceptibility. In this sense, one can speak of a paramagnetic static structure that can occur at negative values of the exchange constant B_0.

[1] J.H. Condon, Phys. Rev. 145, 526 (1966).
[2] J.H. Condon and R.E. Walstedt, Phys. Rev. Lett. 21, 612 (1968).

[3] M.Ya. Azbel', Usp. Fiz. Nauk 98, 601 (1969) [Sov. Phys.-Usp. 12, No. 4 (1970)].
[4] W.M. Lomer, Proc. Phys. Soc. 80, 489 (1962).
[5] A.W. Overhauser, Phys. Rev. 128, 1437 (1962).
[6] T.M. Rice, A.S. Barker, Jr., D.I. Halperin, and D.B. McWhan, J. Appl. Phys. 40, 1337 (1969).
[7] V.P. Silin, Zh. Eksp. Teor. Fiz. 35, 1243 (1958) [Sov. Phys.-JETP 8, 870 (1959)].
[8] V.P. Silin, Fiz. Metal. Metallov. 29, No. 4 (1970).

"SUPERHETERODYNE" AMPLIFICATION OF ULTRASOUND IN SEMICONDUCTORS

Yu.V. Gulyaev and P.E. Zil'berman
Radio Electronics Institute, USSR Academy of Sciences
Submitted 20 February 1970; resubmitted 18 March 1970
ZhETF Pis. Red. 11, No. 8, 421 - 424 (20 April 1970)

We consider in this paper the effect of distributed "superheterodyne" amplification of sound by a drift current of electrons in a semiconductor in the presence of an intense sound wave playing the role of a heterodyne. It will be shown that when the signal (carrier frequency ω_s) interacts nonlinearly with the heterodyne (frequency ω_h), the information in it is converted along the length of the crystal into another frequency (intermediate, ω_i), and vice versa. As a result, if the growth increment possessed by the intermediate frequency is sufficiently large, it is transferred to the signal frequency.

As is well known, [1 - 3], the synchronism condition causes separation of a series of waves that interact most strongly via the pump (the ω_h) wave. These, for example, may be the waves with frequencies ω_s, $\omega_i = \omega_h - \omega_s$, $\omega_p' = \omega_h + \omega_s$, and other combination frequencies of the type $n\omega_h \pm \omega_s$ (n is the integer). It is important that this results in a system of coupled monochromatic waves. Each such monochromatic wave is not a natural wave of the crystal. A natural wave represents, generally speaking, a mixture of monochromatic waves of all the interacting frequencies and is characterized by a single growth increment and by a distinct matching of the phase velocities of the individual monochromatic components. In the general case this increment describes two effects, namely parametric amplification [1 - 3] and the usual linear amplification [4]. The natural wave, which goes over in the absnece of the pump ω_h into a wave with frequency ω_i, will be called for brevity "natural wave of type ω_i" (and similarly for waves of the type ω_s, ω_i', etc.).

Assume that elastic oscillations of frequency ω_s are excited at the entrance into the crystal. Then, owing to the parametric interaction with the pump ω_h, a superposition of different natural waves will be excited in the crystal. Let us assume further that the growth increment of one such wave (say ω_i) is much larger than of the others. This condition is easiest to realize far from parametric resonance ($\omega_s >> _i$ or $\omega_s << \omega_i$). It may then turn out that the admixture of the frequency ω_s in the natural wave of type ω_i will have a much larger amplitude at the exit from the crystal than the oscillations of frequency ω_s in the other natural modes. In this case amplification of sound of frequency ω_s will be determined by the maximum growth increment α_i of the wave of type ω_i.

The described amplification mechanism may have nothing in common with ordinary parametric amplification. Indeed, let us consider a case in which the