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Anomalous Propagation of Ultrasound in Metals by Open-Orbit Electrons

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(Received 30 January 1978)

Open-orbit electrons are shown to generate an ultrasonic wave throughout a metal when the conditions for an open-orbit resonance are satisfied.

We have observed anomalous propagation of 30-MHz shear waves through copper and silver crystals by open-orbit electrons. At certain values of the applied magnetic field, acoustic energy from a wave packet of 1.0-μsec duration spreads out over about 4.0 μsec, reaching the receiver transducer long before the main body of the wave packet arrives. We first present the experimental data which suggested the effect, and then propose a simplified theoretical model which agrees with the principal features of the data.

Open-orbit electrons traveling parallel to the propagation vector $\vec{q}$ of an ultrasonic wave absorb energy resonantly[1] when the period of the open orbit, $D = \hbar K/eB$, is some integral multiple n of the sound wavelength λ. Here $\hbar K$ is the crystal momentum corresponding to the period of the open orbit in $\vec{k}$ space, e is the electronic charge, and B the applied magnetic field. Hence, the resonance condition is $B_n = \hbar K/ne\lambda$.

While studying these effects with shear waves in copper, we decided in mount the quartz transducers with their polarization vectors midway between the fast- and slow-mode directions so that we could measure both modes in the experiment. Figure 1(b) shows the attenuation for the fast mode. Open-orbit resonances (OOR) are seen at the expected field values, but superimposed on the curve are small "beat" patterns at field values which correspond to OOR for the slow mode, as shown in Fig. 1(a). In other words, it appears that some of the slow-mode signal arrives at the receiver transducer simultaneously with the fast-mode signal, but only when B corresponds to a slow-mode OOR. Furthermore, the phase of the anomalous slow-mode signal must be linearly proportional to B/B_n in order to produce the observed beats.

Direct observation of the signal from the receiver transducer shows that the slow-mode wave packet does indeed spread out at OOR, as shown in Fig. 2. As B is swept through a resonance, not only does the amplitude of the packet diminish, but a signal at the sound frequency begins to spread out far in advance of the packet. In some cases the spreading effect is much more evident than the attenuation resonances. In our copper and silver specimens the $n = 4$ and 5 OOR are not visible in the attenuation, but we can easily see the wave packet spread out at values of B_n which correspond to those resonances.

We use a highly simplified model calculation to show that open-orbit electrons pick up energy while traversing the packet, and then transfer it coherently back to the lattice outside the packet.

We first consider the situation where the wave

$u_0 e^{iqx}$ exists only on the negative semiaxis. (Since the Fermi velocity $\vec{v}$ is much greater than the sound velocity, we ignore the time dependence.) The effective force on an open-orbit electron is $\vec{F}_e(\vec{k}) u_0 e^{iqx}$ for $x<0$, and zero for $x>0$. The current produced by $\vec{F}_e$ is computed using Chambers's path-integral method.[2] The excess energy of an electron arriving at x_0 when $t=0$ is

$$\Delta\epsilon = \int_{-\infty}^{0} \vec{F}\cdot\vec{v}\exp[(x-x_0)/l_x]dt, \tag{1}$$

where l_x is an effective mean free path. For $\vec{B}$ along the z axis, $dt=-(D/Kv_x)dk_y$. Using $\vec{F}_e$ we obtain

$$\Delta\epsilon = -(D/K)u_0\exp(iqx_0)\exp(iQk_y)\int_{k_y^\infty}^{k_y^0}(\vec{F}_e\cdot\vec{v}/v_x)\exp(-iQk_y')dk_y', \tag{2}$$

where $Q=qD/K-iD/Kl_x$, $k_y^\infty = k_y(t\to\infty)$, and $k_y^0=k_y+Kx_0/D$.

We use the periodic properties of $\vec{F}_e$ and $\vec{v}$ on an open orbit,

$$\vec{F}_e\cdot\vec{v}/v_x = \sum_n F_n(k_z)\exp(i2\pi n k_y/K), \tag{3}$$

to obtain

$$\Delta\epsilon = u_0\exp(-x_0/l_x)l_x\sum_n\{F_n\exp(i2\pi n k_y/K)\exp(i2\pi n x_0/D)[1+iql_x(1-n\lambda/D)]^{-1}\}. \tag{4}$$

The current is given by $\vec{J}=(e/4\pi^3\hbar)\int(\vec{v}/v)\Delta\epsilon\, dS$ or, with use of (4) and $dS/v=dk_y\,dk_z/v_x$,

$$\vec{J}=u_0\exp(-x_0/l_x)l_x\sum_n\{\vec{J}_n\exp(i2\pi n x_0/D)[1+iql_x(1-n\lambda/D)]^{-1}\}, \tag{5}$$

where $\vec{J}_n=(e/4\pi^3\hbar)\int F_n\vec{V}_n^*\,dk_z$ and $\vec{V}_n$ is a Fourier coefficient of $\vec{v}/v_x$.

A resonance, $D_n=n\lambda$, leads to a current maximum. For a small deviation from resonance, $D_n=n\lambda/(1+\delta)$, where $\delta=\Delta B/B_n$ and $\Delta B=B-B_n$, the current is approximately

$$\vec{J}=u_0\exp(-x_0/l_x)\exp[iq(1+\delta)x_0]l_xJ_n/(1-iql_x\delta). \tag{6}$$

This current will exert a collision-drag force on the ions causing a displacement $\vec{u}=C\vec{J}$, where C is the coupling coefficient.[3] Thus we expect an exponentially decaying sound wave having a phase proportional to $(\Delta B/B_n)x_0/\lambda$ to be generated at a distance x_0 in front of the wave packet near resonance.

In order to study the field dependence more closely, we used a coherent detector gated at a time Δt before the arrival of the wave packet. The detector employs a double-balanced mixer with the received signal as one input, and a delayed, pulsed signal from the transmitter oscillator as the second. The output is thus proportional to the amplitude and cosine of the phase of the receiver signal during the gate pulse. Figure 3 shows the results in silver for $\Delta t=1.2$ μsec, i.e., $x_0=0.14$ cm in front of the 0.5-μsec wave packet, and for $\Delta t=1.0$ μsec. The observed beat period for the first case is $\Delta B/B_n=0.03$, in good agreement with the value calculated, $\lambda/x_0=0.025$.

Although the simple model gives the proper field-dependent phase for the generated wave, the observed decay length appears to be several times the value of l_x (about 0.03 cm in the copper specimen) computed from the OOR width. This is not surprising, since we ignored the fact that the sound field is not zero in front of the packet, but is $\vec{u}=C\vec{J}$. We now repeat the calculation for nonzero field, letting the amplitude u_1 and decay length L be variable parameters, and seek a self-consistent solution of the form $u(x)=u_1e^{-x/L}e^{iqx(1+\delta)}$ for $x>0$.

We find that the decay length of the generated wave is $L=l_x/(1-Cl_xJ_n)$, while the relative amplitude is $u_1/u_0=Cl_xJ_n/(1-iql_x\delta)$. Combining these results we obtain

$$L=l_x/[1-|u_1/u_0|(1+q^2l_x^2\delta^2)], \tag{7}$$

which qualitatively agrees with the observed behavior of the spreading of the wave packet near a resonance. Quantitative comparison would require that the model take proper account of the attenuation.

We have found that wave-packet spreading can be used to distinguish between various magnetoacoustic resonances caused by closed and open or extended orbits. The latter always cause spreading, whereas the former produce no observable change in the packet. We have not looked at the Kjeldaas geometry[4] ($\vec{B}\|\vec{q}$), but the effect should be large since virtually all of the electrons would be traveling on extended orbits along $\vec{q}$.

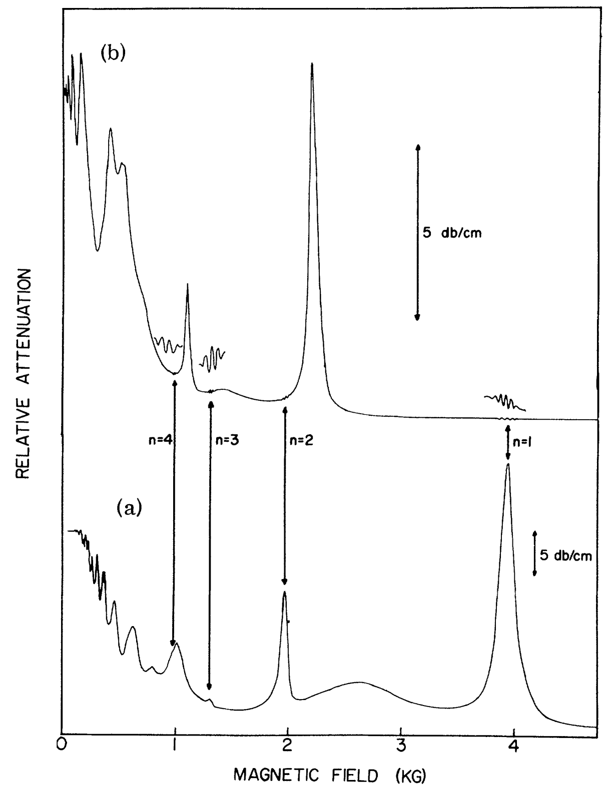

FIG. 1. Attenuation of 34-MHz shear waves in copper at 4.2 K with $\vec{q}\parallel[101]$ and $\vec{B}\parallel[12\bar{1}]$. (a) Slow mode, showing open-orbit resonances. (b) Fast mode, showing open-orbit resonances, but with beat patterns at field values corresponding to slow-mode resonances. Some of the beat patterns are shown expanded.

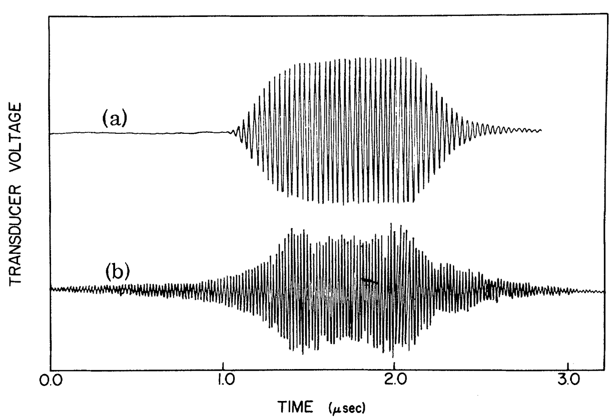

FIG. 2. Receiving transducer voltage for a 1.0-μsec pulse of 34-MHz slow shear waves with $\vec{q}\parallel[101]$ and $\vec{B}\parallel[12\bar{1}]$. (a) $B=4800$ G (off resonance); (b) $B=2000$ G ($n=2$ resonance), with amplification 6 times that in (a). Note that a 34-MHz signal is arriving more than 1.0 μsec before the main packet.

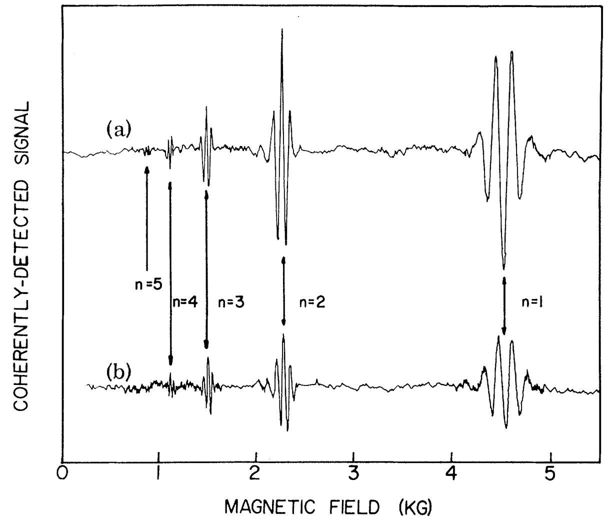

FIG. 3. Output of the gated coherent detector for 34-MHz shear waves in silver with $\vec{q}\parallel[101]$ and $\vec{B}\parallel[12\bar{1}]$. Signals appear only near the open-orbit resonances, indicated by arrows. (a) Detector gated 1.0 μsec before expected arrival of the slow-mode wave packet; (b) detector gated 1.2 μsec before the arrival of the packet.

The mechanism described here may be responsible for the anomalous transmission of microwave electromagnetic (EM) signals reported by Lubzens, Grunzweig-Genossar, and Schultz.[5] The acoustic wave generated at the specimen surface by an EM signal could, at an OOR, be spread throughout the specimen by the open-orbit electrons and generate an EM signal on the opposite side. This is like the situation where a surface EM field sets up "current sheets" within the specimen when a magnetic field is applied.[6] However, the surface current and current sheets have a thickness of the order of the skin depth, whereas the ultrasonic wave sets up an oscillatory current within the bulk; so it should be much more effective in coupling energy to the electrons.

We wish to acknowledge helpful discussions with P. R. Antoniewicz. This research was supported by Grant No. DMR76-11331 from the National Science Foundation.

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ERRATUM

DIFFUSION OF POSITIVE MUONS IN VANADIUM. A. T. Fiory, K. G. Lynn, D. M. Parkin, W. J. Kossler, W. F. Lankford, and C. E. Stronach [Phys. Rev. Lett. 40, 968 (1978)].

In the receipt date line, "received manuscript" should read "revised manuscript."

The value for σ given in the sixth line, second column, page 969, should read $\sigma = 0.370\ \mu s^{-1}$.

An inadvertent transposition of lines in the first column of page 970 rendered a portion of this paper unreadable—lines 9–25 should precede the first eight lines on this page. The text beginning with the last paragraph on page 969 through the end of the first column of page 970 should read as follows:

We have plotted the temperature dependence of a depolarization rate Λ defined through $A(\Lambda^{-1}) = e^{-1}$ in Fig. 1. The depolarization rate of 0.209 μs^{-1} is observed at $T = 10$ K in the 1786-Oe field and is equal, within statistical error, to the 0.208-μs^{-1} depolarization rate measured in a 400 Oe field by Harmannn *et al.*[14] It appears that the Zeeman limit applies to our high-field data.[11] There is a local maximum in Λ in the vicinity of 80 K, where the character of the diffusion process apparently changes.

The temperature dependence for the correlation rate τ^{-1} is shown in Fig. 2. For $T \lesssim 50$ K it is linear, where $\tau^{-1} = (2.4 \pm 0.8) \times 10^4\ s^{-1}\ K^{-1}$ T and the normalized $\chi^2 = 2.1$. The theoretical expression for multiphonon processes, asymptotically approaching $\tau^{-1} \propto T^7$ at low temperatures,[1] does not fit these data, where we find $\chi^2 \sim 10^{11}$. At these temperatures the jump rates for hydrogen (extrapolated[15]) and τ^{-1} for the muon data are of the same order of magnitude. We believe it unlikely that the low-temperature behavior is an effect of trapping and detrapping from impurities, since we expect the binding enthalpy to impurities to be on the order of 0.1 eV, as it is for hydrogen.

In the region $T > 50$ K, on the other hand, the τ^{-1} values are much smaller than the jump rates observed for hydrogen. This is explained by impurity trapping at high temperatures, where τ is close to being the mean time of stay at impurities.[5] We note that for a jump rate on the order of $3 \times 10^7\ s^{-1}$, the muon may diffuse a distance equal to the mean distance between the impurities in 2.2 μs. We therefore associate the structure in Λ and τ^{-1} observed for $T > 50$ K with impurity trapping. In the absence of trapping, positive curvature in τ^{-1} vs T is expected at higher temperatures,[1] where multiphonon processes begin to contribute.