may be treated as a perturbation. On the other hand, such "forbidden" crossings may provide information on the $n=2$ fine structure; see M. Leventhal, Bull. Am. Phys. Soc. 11, 327 (1966).

[8] W. E. Lamb, Jr., and R. C. Retherford, Phys. Rev. 79, 549 (1950).

[9] This confirms the assumed value of the $2P$ lifetime, $\tau=1.595\times10^{-9}$ sec, to better than 1%.

[10] For example, we previously used an incorrect value of the free-electron–free-proton g-value ratio; see footnote 55 in Ref. 1. Correcting this, and incorporating the latest atomic constants, we find that $\mathcal{S}$ from H(605) should be lowered by 0.01 MHz.

[11] The "beam-tilt" effect is mainly responsible for the "beam notch" discussed in Sec. IIIA of Ref. 1. Considering this effect, we can now account for better than 99% of the β beam at the crossing point. This places an upper limit on stray-field quenching; we calculate that stray fields cannot be larger than $|E_S|=0.1$ V/cm, as compared with the applied quenching field $|E_Q|=0.6$ V/cm.

[12] We have also measured such a difference in deuterium. See R. T. Robiscoe and B. L. Cosens, Bull. Am. Phys. Soc. 11, 62 (1966).

SYMMETRY OF PARAELECTRIC DEFECTS IN ALKALI HALIDES*

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(Received 22 April 1966)

For an interpretation of many measurements on paraelectric impurities, such as OH^-, CN^-, NO_2^-, and Li^+ in alkali halides,[1] it is important to know the symmetry of the lattice imperfection created by the impurity. One of the most direct methods of determining the symmetry is by sound velocity and absorption measurements. The results of such studies on Li^+ in KCl are reported here.

For the inelastic properties, an "elastic dipole" can be defined (in analogy to the electric dipole for the dielectric properties) which is a measure of the elastic distortion of the lattice produced by the lattice imperfection and which is a second-rank tensor and can be represented by an ellipsoid (instead of a vector as for the electric moment). For symmetry reasons, the dipoles of the Li^+ (which forms an electric dipole because it is displaced from the center of the vacancy left by the replaced K^+ ion) are ellipsoids of revolution with the dipolar axis as axis of rotation. The stresses resulting from an acoustic wave in a cubic crystal can be represented by "symmetric components" which reflect the symmetry of the crystal.[2] One of these components is, e.g., the hydrostatic pressure $\frac{1}{3}(\sigma_1+\sigma_2+\sigma_3)$ which is totally symmetric (A_{1g} representation). Each of these stress components produces a strain component of the same symmetry, related by the corresponding elastic modulus (again written in symmetric form). In the cubic crystal there are 3 such moduli: $C_{11}+2C_{12}$, $C_{11}-C_{12}$, and C_{44}. The velocity of a wave of given directions of propagation and polarization can then be expressed in terms of these constants; they are given in Table I, column 3. Column 4 indicates the irreducible representations involved. The different symmetric stress components interact with a defect depending on its symmetry.[2] For instance the A_{1g} component (hydrostatic component) does not induce transitions

Table I. Interaction (indicated by "x") of various ultrasonic stresses with defects of different symmetries in a cubic host lattice. The velocity is expressed in terms of the appropriate moduli C_{ij} and the density ρ. "Stress symmetry" indicates the representations of the cubic group according to which the components of the stress tensor for a given acoustic mode transform.

Propagation direction	Polarization direction	Velocity	Stress symmetry	Equilibrium orientation of the defect $\langle 100\rangle$	$\langle 111\rangle$	$\langle 110\rangle$
[100]	[100]	$\{[(C_{11}+2C_{12})+2(C_{11}-C_{12})]/3\rho\}^{1/2}$	E_g, A_{1g}	x		x
[100]	[010]	$(C_{44}/\rho)^{1/2}$	T_{2g}		x	x
[110]	[001]	$(C_{44}/\rho)^{1/2}$	T_{2g}		x	x
[110]	$[1\bar{1}0]$	$[(C_{11}-C_{12})/2\rho]^{1/2}$	E_g	x		x

of defects from one equilibrium position to another; thus the corresponding modulus is not changed by the presence of the defects. In columns 5-7 in the table, the possibilities of interaction and thus change of velocity for the different defect symmetries are indicated. For example, if the dipoles are oriented in the $\langle 111\rangle$ direction, a wave propagating in the $[110]$ direction will interact strongly if the polarization is in the $[001]$ direction, and will not interact at all if it is in the $[1\bar{1}0]$ direction.

Figure 1 represents velocity measurements at 30 Mc/sec on Li^+ in KCl as a function of temperature for the two polarizations just discussed. For the $[1\bar{1}0]$ polarization there is no change of the velocity v within $10^{-3}\%$, while for the $[001]$ polarization a linear decrease with $1/T$ is seen, reaching a change of 0.45% at 1.6°K. Thus these experiments indicate clearly that Li^+ is displaced in the $\langle 111\rangle$ direction, and not in the $\langle 100\rangle$ as had been postulated earlier.[3] For the latter symmetry, the $[001]$ polarization will give zero effects, while the $[1\bar{1}0]$ polarization will show a change of velocity. Experiments with longitudinal and transverse waves in the $[100]$ and the $[111]$ directions are in quantitative agreement with the results of Fig. 1.[4] From the slope in Fig. 1, the dipole moment $(\lambda_1-\lambda_2)$ per defect can be calculated (λ_1 and λ_2 are strains per defect along the two major axes of the dipole ellipsoid) from the formula that applies to this specific defect orientation:

$$2\frac{\delta v}{v}=C_0V_0\frac{4(\lambda_1-\lambda_2)^2}{9kTS_{44}},$$

where C_0 is the concentration of defects in mole fraction, V_0 the atomic volume of the host crystal, and S_{44} the appropriate compliance.

A value of 0.05 for $\lambda_1-\lambda_2$ results from the application of this equation. This compares with a value of ~1 for the stress ordering of C in Fe. This smaller value is reasonable since the small substitutional Li^+ ion produces a much smaller lattice distortion than the interstitial C; also the barriers to orientation are of different orders of magnitude. There is evidence that the orientation in the stress field occurs by tunneling; thus a quantum mechanical treatment has to replace the classical "Langevin-type" calculation. The formula above is valid for $kT>\delta$, where δ is the tunneling-splitting of the ground state. Deviations from the $1/T$ law can be expected at low temperatures; dielectric measurements[5] indicate that δ is of the order of 1°K, and thus deviations should be observable only at temperatures outside the range covered in the present experiments. The absorption is very low and the results are independent of frequency, confirming the idea that the stress ordering occurs by tunneling.

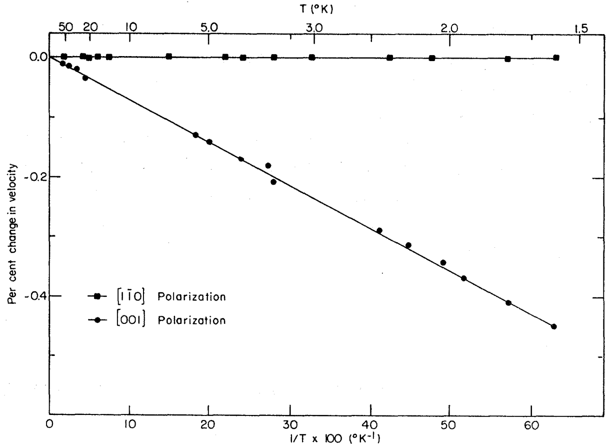

FIG. 1. Temperature dependence of the velocity of sound in KCl with 100 ppm Li^+. The sound propagated in $[110]$ direction at 30 Mc/sec.

For higher concentrations (up to 1500 ppm) there is still no effect on the $[1\bar{1}0]$ polarization, but the $1/T$ dependence for the $[001]$ polarization ceases to be linear. There is first a linear decrease with a slope that scales reasonably well with concentration. It is followed by a curving over into another linear dependence at low temperatures (below 5°K) with a much smaller slope. This behavior is interpreted by interaction between the Li^+ ions which is less strong at higher temperatures and thus does not affect the slope appreciably, but becomes more pronounced at low temperatures. However, there is no increase down to 1.6°K, such as has been observed by Känzig, Hart, and Roberts[1] for the dielectric properties of OH^- in KCl.

Helpful discussions with Professor Kramhansl, Mr. Bowen, and Mr. Gomez are gratefully acknowledged.

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[*] This research was mainly supported by the U. S. Atomic Energy Commission. Additional support was received from the Advanced Research Projects Agency through the use of space and technical facilities of the Materials Science Center.

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[1] W. Känzig, H. R. Hart, and S. Roberts, Phys. Rev. Letters 13, 543 (1964); U. Kuhn and F. Lüty, Solid State Commun. 2, 281 (1964); H. S. Sack and M. C. Moriarty, Solid State Commun. 3, 93 (1965); M. C. Moriarty, Cornell University Materials Science Center Report No. 290, 1964 (unpublished); G. Lombardo and R. O. Pohl, Phys. Rev. Letters 15, 291 (1965); V. Narayanamurti, Phys. Rev. Letters 13, 693 (1964).

[2] For example, A. S. Nowick and W. R. Heller, Advan. Phys. 14, 101 (1965).

[3] J. A. D. Matthew, Cornell University Materials Science Center Report No. 373, 1965 (unpublished); G. J. Dienes, R. D. Hatcher, R. Smoluchowski, and W. Wilson, Phys. Rev. Letters 16, 25 (1966).

[4] This is also confirmed by measurements, by F. Welsh in this laboratory, of Young's modulus and internal friction in the kc/sec range (private communication).

[5] H. Bogardus and H. S. Sack, to be published.

PROPAGATION OF ZERO SOUND IN LIQUID He^3 AT LOW TEMPERATURES*

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(Received 6 June 1966)

We have observed the propagation of sound in liquid He^3 at 0.32 atm and at frequencies of 15.4 and 45.5 MHz down to a temperature T^* of 2 mdeg on the magnetic temperature scale valid for powdered cerium magnesium nitrate (CMN) in the form of a right circular cylinder with diameter equal to height. As the temperature rises the sound attenuation increases, reaches a maximum, and then decreases. At low temperatures the attenuation is proportional to T^{*2} and is independent of frequency. At high temperatures the attenuation is proportional to ω^2/T^{*2}, where ω is the angular frequency. The sound propagation velocity is relatively temperature independent at both high and low temperatures but near the attenuation maximum the velocity changes.

In 1957 Landau[1] predicted that at sufficiently low temperatures a new type of sound, which he called zero sound, could be propagated in liquid He^3. Based on Landau's idea, a more detailed theory of the velocity and attenuation of sound in both the hydrodynamic (first sound) and zero-sound regions was worked out by Khalatnikov and Abrikosov.[2] At temperatures sufficiently high that quantum effects are unimportant, it is predicted that the attenuation of zero sound be proportional to T^2 and independent of frequency. In the first-sound region it is predicted that the attenuation is proportional to ω^2/T^2, corresponding to classical viscous attenuation with viscosity proportional to T^{-2}. Both of these temperature and frequency dependences are observed in the present experiments. In what follows we shall show that there is quantitative agreement with theory on velocity changes and semiquantitative agreement on attenuation, so that the present experiments clearly demonstrate the phenomenon of the propagation of zero sound.

The first experiments to support strongly the existence of zero sound in He^3 are those of Keen, Matthews, and Wilks,[3] who measured the effective acoustic impedance of liquid He^3 and found a transition from a high-temperature to a low-temperature regime at about 0.092°K for a frequency of 1000 MHz. It was not possible for them to measure the propagation characteristics of the sound nor, as it will be shown here, to deduce correctly by means of theory the velocity change between first and zero sound. However, on the basis of the present experiments there is no doubt that the phenomenon observed by them was indeed the transition from first to zero sound.

Only a few brief remarks on the present experiments can be made here. Most details are thoroughly discussed elsewhere.[4] The cell used for the measurements is illustrated schematically in Fig. 1. The CMN acts both as refrigerant and as thermometer. The sound propagation in the He^3 occurs in a small gap between a fused quartz ultrasonic delay line and a receiving crystal of X-cut quartz of fundamental frequency 15 MHz. The delay line serves to allow the transient in the receiving crystal, caused by electrostatic pickup from the initial pulse at the transmitting crystal, to damp sufficiently before the sound pulse reaches the gap. The gap length used for the experiments reported here was 1.38×10^{-2} cm and was provided by three small quartz flats. The princi-