where $e_n = -Ze/A$, $e_p = Ne/A$, e—proton charge, and Z(N) —number of protons (neutrons).

Let us estimate B(E1) in the quasi-classical approximation for the low-lying 1^- levels ($\omega \ll 2\Delta_{n,p}$). In this case we can obtain, for example (see [4]),

$$\sum_{\lambda,\lambda'} |(q_{30})_{\lambda\lambda'}|^2 \frac{E_\lambda E_{\lambda'} - \varepsilon_\lambda \varepsilon_{\lambda'} + \Delta^2}{2E_\lambda E_{\lambda'}(E_\lambda + E_{\lambda'})^3}$$

$$\approx \frac{1}{4\Delta^2} \sum_{\lambda\lambda'} |(q_{30})_{\lambda\lambda'}|^2 \varphi\left(\frac{\varepsilon_\lambda - \varepsilon_{\lambda'}}{2\Delta}\right) \delta(\varepsilon_\lambda) \sim \rho_0 R^6 \Delta^{-2}, \quad (3)$$

where $\varphi(x) = x^{-2} - \ln(x + \sqrt{1+x^2}) \cdot [x^3\sqrt{1+x^2}]^{-1}$, ρ_0 —energy level density near the Fermi surface. The estimate in (3) is valid if the energy differences $\epsilon_\lambda - \epsilon_{\lambda'} \lesssim 2\Delta$. By examining the Nilsson scheme we see that at the observed values of quadrupole nuclear deformation β_0 there exist levels λ and λ' for which this condition is satisfied. Estimating in similar fashion the numerator of (2), we obtain

$$B(E1, 1^- \to 0^+) \sim \left(\frac{N-Z}{A}\right)^2 (eR)^2 \beta_0^2 \rho_0 \Delta \frac{2\Delta}{\omega}. \quad (4)$$

Formula (2) and the estimate (4) are derived without introducing the static octupole deformation (see [6] in this connection).

The reduced probability for the excitation of the 3^- level (rotational satellite of the 1^- level) is in this model

$$B(E3, 0^+ \to 3^-) = \frac{1}{2\omega} \left\{\sum_{\lambda\lambda'p} |(q_{30})_{\lambda\lambda'}|^2 \frac{E_\lambda E_{\lambda'} - \varepsilon_\lambda \varepsilon_{\lambda'} + \Delta^2}{2E_\lambda E_{\lambda'}} \times \frac{E_\lambda + E_{\lambda'}}{(E_\lambda + E_{\lambda'})^2 - \omega^2}\right\}^2 \left\{\sum_{\lambda\lambda' n,p} |(q_{30})_{\lambda\lambda'}|^2 \frac{E_\lambda E_{\lambda'} - \varepsilon_\lambda \varepsilon_{\lambda'} + \Delta^2}{2E_\lambda E_{\lambda'}} \times \frac{E_\lambda + E_{\lambda'}}{[(E_\lambda + E_{\lambda'})^2 - \omega^2]^2}\right\}^{-1}. \quad (5)$$

We note that (5) is derived without introducing the effective nucleon charge for the E3 transitions. Estimating $B(E3, 0^+ \to 3^-)$ in analogy with (3), we obtain

$$B(E3, 0^+ \to 3^-) \sim B_{s.p.}(E3)\, \rho_0 \Delta \frac{2\Delta}{\omega}, \quad (6)$$

where $B_{s.p.}(EL) \sim e^2 R^{2L}$ is the reduced probability for the single-particle transition. Comparing (6) with the estimated $B_{hydr}(E3)$ in the hydrodynamic model, $B_{hydr} \sim B_{s.p.}(E3) \cdot A^{2/3}$ [2] we obtain the inequality

$$B_{hydr}(E3) > B(E3) > B_{s.p.}(E3). \quad (7)$$

It is seen from (5)—(7) that measurement of the cross section for the Coulomb excitation of the 3^- level is an essential check on the proposed model.

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PHONON SCATTERING ON IMPURITY IONS IN SODIUM CHLORIDE CRYSTALS

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LOCAL distortions in the sodium chloride crystal lattice in the environment of different impurity ions were recently successfully investigated with the aid of nuclear magnetic resonance. In particular, the radii of the distorted zones around Ag^+, Br^-, and K^+ ions were shown to be related as 1:1.4:1.9. [1] The distorted zone apparently plays an important role in the scattering of phonons. The phonon scattering cross section of an impurity ion can therefore be expected to be proportional to the square of the radius of the distorted zone and, consequently, the cross section for scattering on Ag^+, Br^-, and K^+ ions could be expected to be related as 1:2.0:3.5. The purpose of the present project was to verify this assumption.

At a low concentration of impurity ions, when the elastic constants and the heat capacity of a crystal with an impurity practically do not differ from the corresponding properties of a pure crystal, it can be asserted that $\Delta R/R_0 = f(l_0/l_{imp})$, where R_0 is the thermal resistance of the pure

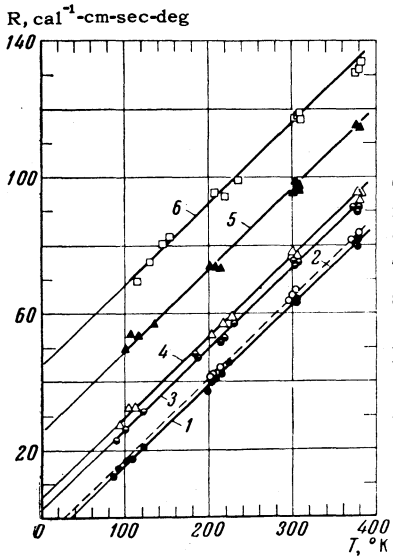

FIG. 1. Temperature dependence of thermal resistance for specimens with admixture of Br^-. The curves refer to specimens with the following quantities of NaBr: 1 – 0; 2 – 0.0012; 3 – 0.0030; 4 – 0.0045; 5 – 0.0150; 6 – 0.0310

crystal, ΔR is the added thermal resistance due to the impurity; l_0 and l_{imp} are the phonon mean free paths for scattering on phonons and on impurity ions, respectively.

In accord with the elementary theory of thermal conductivity we have $l_0 \sim 1/R_0\bar{v}C_V$, where $\bar{v}$ is the mean velocity of sound and C_V is the heat capacity per unit volume. From dimensionality considerations it follows in turn that $l_{imp} \sim 1/SN$, where S is the cross section for phonon scattering an impurity ion, and N is the number of impurity ions per unit volume.

Thus, if S is the cross section of the distorted zone, and if the assumption mentioned earlier is correct, then $\Delta R/R_0$ is a single-valued function of the dimensionless parameter η

$$\Delta R/R_0 = f(\eta), \qquad \eta = SN/R_0\bar{v}C_v \tag{1}$$

independent of the type of impurity ion.

The thermal conductivity was measured on the same NaCl monocrystals with AgCl, NaBr and KCl impurities as before.[1] The measurements were made with the apparatus described in [2] (setup A) in the temperature interval 100—380 deg K. Because of the transparency of the NaCl specimens, heat transfer can take place directly through the specimen from the heater to the

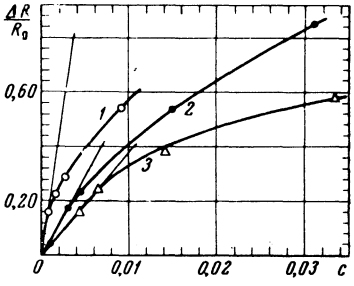

FIG. 2. Dependence of $\Delta R/R_0$ on the molecular concentration of the impurity; o – K^+; • – Br^-; Δ – Ag^+.

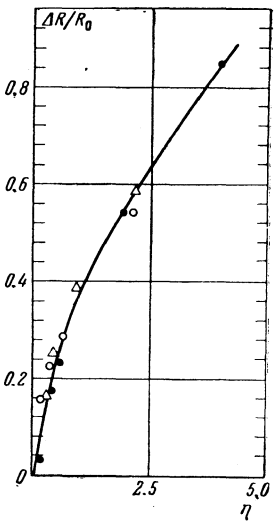

FIG. 3. Relation (1) (designation of points same as in Fig. 2).

cooler.[3] The error due to his, however, is less than 0.4% of the basic heat flow in the given temperature interval.

The measurements have shown in all cases that the temperature dependence of the thermal resistance is linear, R = a + bT, the parameter a being determined by the type and concentration of the impurity and the coefficient b being independent of them. Figure 1 gives, as an example, the corresponding experimental curves for specimens with bromine impurities. Similar familes of curves were obtained for specimens with silver and potassium.

Since the radii of the distorted zones around the Ag^+, Br^-, and K^+ impurity ions were determined by measurement of the nuclear magnetic resonance only for 300 deg K, we must limit ourselves in verifying the applicability of Eq. (1) to data on the thermal conductivity at that temperature. Figure 2 shows the dependence of the corresponding variation in thermal resistance on the impurity concentration in molar percentages for 300 deg K. These data were employed to plot Eq. (1) in Fig. 3.* From the figure it is evident that the data pertaining to different impurities fit well the common curve. From this it is to be concluded that phonon scattering on impurity ions is really determined by the cross sections of the corresponding distorted zones.

We note at this point that this assertion is valid not for room temperature alone. All other data, obtained at other temperatures in the 100—380 deg K interval also fit the curve in Fig. 3 (if, of course, the values of R_0, $\bar{v}$, and C_V are taken for the corresponding temperature). This fact indicates that the radius of the distorted zone around the impurity ion is practically independent of the temperature.

The authors wish to thank V. V. Lemanov for his discussion of the paper.

*In the calculation of parameter η, the following values were used. $R_0 = 63$ cm-sec-deg-cal^{-1}; $C_v = 0.42$ cal-cm^{-3}; $\bar{v} = 3.2 \times 10^5$ cm-sec^{-1}; $N = cN_0$, where c is the molar concentration of the impurity, $N_0 = 2.23 \times 10^{22}$ cm^{-3}; for Ag^+, Br^-, and K^+ ions, in accordance with [1], S is equal respectively to 2.48, 4.85 and 8.75×10^{-14} cm^2.

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REMARKS ON THE PAPER BY DEMIVOV, SKACHKOV, AND FANCHENKO ENTITLED CHANNEL EXPANSION IN INTENSE SMALL SPARKS

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1. Using a high-sensitivity electron-optical converter Demidov, Skachkov, and Fanchenko [1] were able to photograph the space-time development of a spark channel excited by discharging low-capacity condensers; on the basis of these photographs these authors conclude that the rate of expansion of the channel is 60—80 km/sec at the initial stage of the discharge.

It is our contention that these photographs are subject to a somewhat different interpretation than that given by the above authors. First, it should be noted that the extremely high rate of expansion of the channel is determined in a time interval of 0.25—0.5 nsec, in the course of which the channel radius is no greater than 15—30 μ. There is reason to believe that in this time period one is not observing the stage corresponding to hydrodynamic expansion of the channel, but rather the preceding stage, i.e., the streamer discharge. For example, according to the data of Saxe and Chippendale, [2] who obtained frame photographs of the development of a streamer with an electron-optical converter, propagation of the streamer in a discharge gap 1 cm in length continues for approximately 4 nsec, in which time the diameter of the luminous streamer remains smaller than 100 μ.

It is reasonable to assume that in the discharge studied in [1] the streamer diameter is finite as it moves from one electrode to the other and that the "channel expansion" observed in the photographs in Figs. 3—4 of [1] may actually characterize the geometric shape of the head of the streamer and its expansion.

To resolve definitely the question of the rate of expansion of a channel in a spark discharge in its initial stages of development it will be necessary to carry out additional investigations similar to those in [2], in order to determine the radius of the initial channel at the instant it bridges the interelectrode gap.

2. Demidov, Skachkov, and Fanchenko give data for the curvature of the current growth at the initial time, which is calculated on the basis of the known capacitance, the period of natural oscillations, and the gap voltage; the ohmic resistance, however, was neglected.

Investigations carried out by us have shown that for a hard spark discharge in air with energies of 0.01—1.0 joules at capacity values ranging from 500 to 5500 pF the maximum value of the curvature of the current growth is actually determined to a considerable degree by the resistance of the arc channel. In this case the maximum value of the curvature of the current growth is not reached when the current appears, but starts several nonoseconds later, when the current is approximately 500 A. [3,4] At small capacitance values the curvature of the current growth is still appreciably smaller than the quantity U/L (U is the gap voltage and L is the inductance of the discharge circuit).

On the basis of these data it may be assumed that the values of the curvature of the current growth given in [1] are too high.

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