The total deformation of the samples to which Fig. 2 pertains was about 4% at the end of the experiments.

The spectral behavior of the effect can be seen from the plots of Fig. 3. The effect was observed in the wavelength region coinciding with the band of intrinsic absorption of light by CdS.

The effect is maximal when the illumination wavelength is close to the intrinsic-absorption maximum (5300 Å). The spectral dependence of the effect was plotted at 75°C at a constant illumination of 150 lux.

The sole purpose of the present note is to report the experimental facts. We assume that the observed phenomenon is due to the change of the conditions for the motion of dislocations during the time of illumination. This can pertain either to the change in the density of the free electrons interacting with the moving dislocations [1] or with the change in the form of the potential relief along which the dislocation moves (Peierls barriers) during the course of photoionization of the atoms constituting the crystal lattice of the CdS. The presence of a maximum of the effect in the region of the intrinsic absorption gives grounds for assuming that the nature of the observed phenomenon differs from the change in the state of the local centers capable of pinning the moving dislocation, as is observed in colored ionic crystals illuminated with light in the F-band [2].

Further experiments will be aimed at a detailed study of the nature of the observed phenomenon.

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CRITICAL VELOCITIES IN He II

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As is well known, at velocities higher than critical ($v_c \equiv v_{c1}$), i.e., at velocities corresponding to complete dragging of the He II by the vessel wall (in the case of translational or rotational motion of the vessel), the fountain effect and the second sound remain in force, while the velocity of the second sound remains the same as in the precritical mode [1]. This is evidence that when $v > v_{c1}$ the He II does not go over into the normal phase, and moreover, the ratio of the density ρ_n of the normal component to that of the superfluid one (ρ_s) does not change noticeably. However, such a behavior of He II does not contradict Landau's known point of view that when $v \gtrsim v_c$ the number of excitations produced is unlimited, owing to the interaction of the He II with the vessel walls (i.e., all the He II goes over to the normal phase), since the usual identification of the Onsager-Feynman vortex filaments, which determine v_{c1}, with microscopic excitations is not fully correct. It is known that: (i) the vortex filaments represent macroscopic motions of the superfluid component and the length of the vortex filaments and rings has macroscopic dimensions, (ii) the velocity perturbation due to the filament ($v_s(r) = \hbar(m|r - r_v(t)|)$) extends over large distances, (iii) as a result, the minimum of the "free" energy $\tilde{E} = E - M\omega$ corresponds in rotary motion to an average filament density $n = n_0 \equiv 2\omega/\kappa = m\omega/\pi\hbar$ (i.e., to an average distance between filaments $b = b_0 \equiv (\pi\hbar/m\omega)^{1/2}$, where E is the kinetic energy of the superfluid component, M its angular momentum, $\kappa = 2\pi\hbar/m$ the circulation for one filament, and ω the angular velocity of the vessel). Consequently, the vortex filaments are effectively repelled from one another, so that the distance between them cannot exceed b and an unlimited number of filaments is produced. It is known that when $n = m\omega/(\pi\hbar)$ the angular momentum is equal to its solid-body value [1,2], i.e., the superfluid component, without going over into the normal phase, imitates the rotation of the liquid as a whole, corresponding to the normal phase, and in the case of translational motion it imitates the flow of the liquid as a whole (with solid-body total momentum at an average distance between vortex rings $b = [2\pi\hbar R/(3mv)]^{\frac{1}{2}}$). In only one respect is there a similarity between the vortex filaments and the excitations. According to Ginzburg-Pitaevskii [3], we have $\rho_s = 0$ on the vortex axis ($\rho_s(r) \sim r$ as $r \to 0$), and only when $r \gg a$ do we have $\rho_s = \rho_{s0}(T)$ (a - effective radius of the core of the vortex, ρ_{s0} is the value of ρ_s in the absence of vortices). Inasmuch as $\rho_s + \rho_n = \text{const}$, it is natural to propose that the core of the vortex will be filled with the normal component, i.e., when a vortex system is produced with a density

$$n_1 = b_1^{-2} = \left(\frac{\hbar\pi}{m\omega} + \pi a^2\right)^{-1}$$

ρ_n increases by an amount $\Delta\rho = \rho_s \pi a^2/b_1^2$ (and ρ_s decreases by $\Delta\rho$) (the change of n is connected with the fact that the solid-body angular momentum is now equal to the sum of the momentum of the superfluid component and of the normal component dragged in the vortex core by the vessel, i.e., rotating with angular velocity ω). Thus the vortices also make a certain contribution to ρ_n, and in the presence of vortices there is produced in He II a mixed phase analogous to the mixed phase for superconductors of the second kind. Its formation, however, is connected not with the negative surface tension σ_{sn}, as in the latter case, but with repulsion between the vortex filaments. The upper critical velocity v_{c2} (ω_{c2}) at which the He II goes over entirely into the normal phase can be determined either (i) by the fact that the distance between the filaments becomes of the order of a - v'_{c2} (in this case $\rho_s = 0$; the upper critical field in superconductors of the second kind is determined in this manner) [4], or (ii) by the microscopic excitations v''_{c2} after Landau. $v'_{c2} = \omega'_{c2}R \sim R\hbar/ma^2$; when $\tau \ll 1$ we have $a = a_0\tau^{-2/3}$; $\tau = T_\lambda - T/T_\lambda$ [5]. $v''_{c2} = \Delta(t)/p_0$ (Δ and p_0 are the energy and momentum corresponding to the roton minimum in the energy spectrum of He II). Of course, v_{c2} is equal to the smaller of the quantities v'_{c2} or v''_{c2}. Usually $v''_{c2} < v'_{c1}$; only when $\tau \ll 1$ and r is small can $v'_{c1} < v''_{c2}$ occur. It would be of interest to set up experiments aimed at observing the fountaining and second sound at $v > v_{c2}$.

2. In the calculation of v_{c1} (ω_{c1}) one usually neglects the energy of the normal component in the vortex cores. We shall show later that near T_λ this can lead to appreciable

errors. Let us consider ω'_{c1} corresponding to the presence of a system of vortices with density $n = n_1$ ($\omega'_{c1} \simeq \omega_{c1}$; calculation with $n = n_0$ leads to the same results). The free energy E_1 per unit volume is equal to $\tilde{E}_1 = n_1[(\epsilon + \alpha\omega^2) - \omega(2\beta + 2\alpha\omega)]$;

$$\epsilon = \frac{\pi\rho_s \hbar^2}{m^2} \ln\frac{b}{a} + 2\pi a \sigma_{sn} + \pi a^2 (F_n - F_s); \; \alpha = \frac{\pi\rho_s a^2 R^2}{4}; \; \beta = \frac{\pi\hbar\rho_s R^2}{4m};$$

the term $\alpha\omega^2$ is equal to the kinetic energy of the normal component in the cores of vortices rotating with angular velocity ω, the first term in ϵ is connected with the ordinary kinetic energy of the vortices, the second with the surface tension between the normal and superfluid components on the surfaces of the vortex cores, and the third with the additional energy lost when the superfluid component is replaced by the normal ones in the vortex cores ($F_n - F_s$ is the difference between the specific free energies of the normal and superfluid components at the given temperature). The expression for M takes into account the rotation of the normal component in the cores of the vortices. Putting $\tilde{E} = 0$, we obtain a quadratic equation for ω_{c1}; ω_{c1} corresponds to the positive root of the equation

$$\omega_{c1} = [(\alpha\epsilon + \beta^2)^{1/2} - \beta]/\alpha.$$

When $\alpha \ll \beta^2$ we have $\omega_{c1} \simeq \epsilon/(2\beta)$. The second and third terms in ϵ are small both when $T = 0$ and $T \simeq T_\lambda$; when $T \simeq T_\lambda$ they are close to $\tau^{2/3}$. If we neglect these terms, we obtain the usual equaltion $\omega_{c1} \approx 2\hbar(mR^2)^{-1} \ln b/a$. On the other hand, if $\alpha\epsilon \gg \beta^2$, i.e., $R \ll R_0$, where $R_0 = 2a(\ln b/a)^{1/2}$ ($R_0 = 2a_0\tau^{-2/3}(\ln b/a)^{1/2}$ near T_λ), then $v_{c1} = \omega_{c1}R \simeq 2\hbar\tau^{2/3}(\ln b/a)^{1/2}/ma_0$. This is precisely the dependence on τ observed in the experiments for pores with dimensions $R = 0.2\ \mu$ at $T_\lambda - T = 10^{-2} - 10^{-4}$ deg. (In the comparison with the experimental data it is necessary to consider v_c and not ω_c, since in [6] they rotated the entire cylinder, which was covered with porous material and had an outside radius 2.5 cm.) It is easy to see that the condition $R \ll R_0$ is satisfied if $a_0 = (0.7 - 2) \times 10^{-7}$ cm. For pores with $R = 10\ \mu$, the dependence of ω_{c1} on τ is weaker, but approaches $\sim\tau^{2/3}$ at $T_\lambda - T \sim 10^{-4}$ deg, and $\omega_{c1} = \text{const}$ for pores with $R = 150\ \mu$. Indeed, in this case $R \gg R_0$. We note that when the condition $R \ll R_0$ is satisfied we get $\omega_{c1} \sim R^{-1}$, and not $\omega_{c1} \sim R^{-2}$ as in the case when $R \gg R_0$.

Were we to have $\omega_{c1} > \omega'_{c2}$, then a transition from the superfluid phase directly into the normal phase would occur at $\omega = \omega'_{c2}$, bypassing the mixed phase. In this case $\omega'_{c2} \sim \tau^{4/3}$, and not $\omega_{c1} \sim \tau^{2/3}$, if $a = a_0\tau^{-2/3}$, as proposed above. (If the coefficient of ρ_s^2 in the expansion of the thermodynamic potential in [3] were to tend to a constant value as $\tau \to 0$, as in [3], and if $\rho_s \approx \tau^{2/3}$, then $a = a_0\tau^{-1/3}$, i.e., $\omega_{c1} \sim \tau^{1/3}$ when $R \ll R_0$ and $\omega'_{c2} \sim \tau^{2/3}$.) It would be of interest in this connection to set up experiments analogous to [6], but at still lower values of R or τ; a much stronger dependence of v_c on τ was observed in [7] for pores with $R = 3 \times 10^{-7}$ cm, but at $\tau \sim 1/2 - 1/4$. The numerical coefficient in the expression for ω'_{c2} is not known exactly, but apparently in the experiments of [6], $\omega_{c1} < \omega'_{c2}$ even at $R = 0.2\ \mu$, when $\omega_{c1} \sim \omega'_{c2}$, i.e., the mixed phase was realized.

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IS ONE OF THE LEPTON QUANTUM NUMBERS MULTIPLICATIVE ?

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As is well known, a single additive lepton quantum number that assumes only three values (-1, 0, +1) is insufficient for a description of the aggregate of experimental data on weak interaction processes in which leptons take part. We can consider different possibilities, among which we mention the following:

1. There are two different additive lepton charges - muonic and electronic. From the point of view of the experiment proposed in the present note, this alternative does not differ from the case when there is only one additive lepton charge, whose signs are opposite for μ^- and e^- [1].

2. There is one additive lepton charge, the values of which for e^-, ν_e and μ^-, ν_μ are different (say +1 for e^-, ν_e and +2 for μ^-, ν_μ) [2].

3. There is only one additive lepton charge ℓ (say +1 for ν_e, ν_μ, e^-, μ^- and -1 for $\bar{\nu}_e$, $\bar{\nu}_\mu$, e^+, μ^+) and one multiplicative [3] lepton number M equal to +1 for ν_e, e^-, $\bar{\nu}_e$, e^+ and -1 for ν_μ, μ^-, $\bar{\nu}_\mu$, μ^+.

Alternative 3 is the least rigorous of the foregoing possibilities, since it permits, in principle, muonium ⇌ antimuonium transitions [4]. It calls for a value +1 for M in all particles that are not leptons (see, for example, the decay $\pi^+ \to \mu^+ + \nu_\mu$ etc.) and it therefore seems to us quite artificial. In addition, it is not compatible with notions of Feynman and Gell-Mann [5] concerning the interaction of two currents. However, only experiment can answer the question whether alternative 3 is realized in nature.

To clarify the question of the existence of a multiplicative lepton number, it was proposed to study experimentally the oscillations of the muonium ⇌ antimuonium transition [3]. In this case, however, the required experiments are quite difficult. Besides, such oscillations are caused by a second-order interaction or by a rather exotic interaction. We propose below a concrete-experiment formulation free of these shortcomings.

Let us consider the decay of a muon, say a positive one (for reasons of experimental nature, which will be made clear in what follows). According to possibilities 1 and 2, its scheme is $\mu^+ \to e^+ + \nu_e + \bar{\nu}_\mu$. On the other hand, if there exists a multiplicative quantum number M besides the additive charge (alternative 3), then the muon decay proceeds [3] in