estimate the Fermi energy of the observed d-band holes to be about 0.11 ± 0.03 eV and, if spin-orbit coupling is neglected, this is also the Fermi energy for the heavy d-band holes. The temperature variation of the amplitude of the fast period (that from the s band) at $\theta = 14°$ gave $m^*/m_0 = 2.1 \pm 0.3$.

Using the observed effective masses to calculate the contribution of the s band and the light d holes to the electronic specific heat, we find that the observed surfaces account for less than 10 % of the observed value[1] of 9.3 mJ deg^{-2} mole^{-1}. If we neglect many-body effects, the remainder is contributed by the heavy d holes. These must therefore have a very high effective mass, which makes them unobservable in the de Haas-van Alphen effect at 1°K.

There remains the question of whether any other d-band levels at symmetry points are above the Fermi level. Most d-band calculations[8,13] indicate that the energy separation of X_5 and W_1' is very small (~0.01 eV), whereas the other d levels are appreciably lower. Our experiments suggest that W_1' is above the Fermi level, because otherwise it is very difficult to accommodate the 0.36 hole per atom on the square faces of the Brillouin zone. However, it is the galvanomagnetic data[6] which provide conclusive evidence that W_1' is above the Fermi level and that the d-band levels at K and L lie below it. The heavy d-band holes then form a surface open in [100] directions which accounts for the observed stereogram. In addition, this model of the Pd band structure predicts correctly the sign and the approximate magnitude of the Hall constant[6] for H[100], the only direction in which geometric discompensation is observed.

*Work supported by the National Science Foundation and by the Advanced Research Projects Agency.

†National Science Foundation Predoctoral Fellow.

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SUPERCONDUCTING SURFACE SHEATH OF A TYPE-II SUPERCONDUCTOR BELOW THE UPPER CRITICAL FIELD H_{c2}†

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We have calculated the order parameter of a bulk type-II superconductor near its surface (bounded by a vacuum) for magnetic fields parallel to the surface and close to the upper critical field H_{c2} when the Ginzburg-Landau parameter $\kappa \gg 1$. We find that the energy gap near the surface is larger than about 70 % of its zero-field value for fields just below H_{c2}. The maximum value of the order parameter is in general not at the surface but is field dependent and occurs within approximately one coherence length from the surface. The superconducting properties (order parameter, energy gap, superconducting electron concentration) of the metal near its surface do not show any abrupt changes near H_{c2}.

Consider a semi-infinite superconducting half-space with the boundary surface at $x = 0$ and vacuum for $x < 0$. The applied magnetic field H_0 is parallel to the z direction. Assume that the Ginzburg-Landau parameter $\kappa = \lambda/\xi \gg 1$ (λ is the penetration depth and ξ the coherence length) and also that $(H_{c2} - H_0)/H_{c2} \ll 1$. We know that to this approximation Abrikosov's normalized solution[1] for the magnetic field in the bulk of the material near the transition temperature is

$$H_B = H_0 - (1/2\kappa)|\Psi_B|^2, \tag{1}$$

where the bulk order parameter Ψ_B is normalized with respect to the order parameter in zero field, the magnetic field with respect to $\sqrt{2}H_C$, the distances with respect to λ, and H_C is the bulk thermodynamic critical field.

Between the magnetic field at the boundary surface H_0 and the bulk field H_B, there must be a region in which the magnetic field goes over continuously from H_0 to H_B. For a square vortex lattice the maximum value of the order parameter is $|\Psi_B|_{\max}^2 = 2\sqrt{2}\kappa(\kappa-H_0)/(2\kappa^2-1) = \sqrt{2}\beta\langle|\Psi_B|^2\rangle$, where $H_{c2}\equiv\kappa$ in normalized units and $\beta = 1.18$. Therefore $H_B = H_0$ for $\kappa \gg 1$ and $(\kappa-H_0)/\kappa \ll 1$. We may therefore assume to the first approximation that for large-κ materials the magnetic field near the surface of the metal is also H_0, provided (as it turns out) the "surface layer" is smaller than the penetration depth λ.

Therefore the vector potential in the surface layer is approximately $\vec{A} = (0, H_0x, 0)$ plus the gradient of an arbitrary function of the coordinates which we have equated to zero. Because the Ginzburg-Landau[2] equations (on which Abrikosov's solution is based) are gauge invariant, the phase of the order parameter is then determined only to within an arbitrary function of the coordinates. We therefore assume that near the surface the order parameter is of the form

$$\Psi = e^{iky}D(x), \tag{2}$$

where k is as yet undetermined, and D is a function of x. With the above vector potential and Eq. (2), the first normalized Ginzburg-Landau equation,

$$[(i/\kappa)\vec{\nabla} + \vec{A}]^2\Psi - \Psi + |\Psi|^2\Psi = 0, \tag{3}$$

reduces to

$$-\frac{1}{\kappa^2}\frac{d^2D}{dx^2} + \left(H_0x - \frac{k}{\kappa}\right)^2 D - D + |D|^2 D = 0. \tag{4}$$

We introduce the parameter $\mu^2 = \kappa/H_0$, which is equal to the ratio of the upper critical field $(H_{c2}\equiv\kappa)$ to the applied field, and the new variables

$$\eta = \zeta - \zeta_0 = (\kappa/\mu)x - k(\mu/\kappa).$$

Then Eq. (4) becomes, with $D \equiv D(\zeta)$,

$$\frac{d^2D}{d\zeta^2} + [\mu^2(1-D^2) - (\zeta-\zeta_0)^2]D = 0. \tag{5}$$

At the boundary surface the Ginzburg-Landau boundary condition[2]

$$\left(\frac{i}{\kappa}\frac{d}{d\zeta} + A_\zeta\right)\Psi(y,\zeta) = 0 \tag{6}$$

must be satisfied. For our particular choice of the vector potential, Eq. (6) reduces to

$$dD/d\zeta = 0 \text{ at } \zeta = 0. \tag{7}$$

The as yet undetermined constant $k[=(\kappa/\mu)\zeta_0]$ can be obtained in the following way. The free-energy difference between the superconducting state in an applied magnetic field and that of the normal state in zero field integrated over the total volume is[2]

$$\int dV(F_{SH} - F_{N0}) = \int dV\left(-|\Psi|^2 + \frac{1}{2}|\Psi|^4 + \frac{H^2}{8\pi} + \left|\frac{i}{\kappa}\vec{\nabla}\Psi + \vec{A}\Psi\right|^2\right). \tag{8}$$

For our particular choice of Ψ and $\vec{A}$ the free-energy difference per unit area in the yz plane (at constant magnetic field) is

$$\Delta F = B + C\int_{\eta_1(\zeta_0)}^{\infty} d\eta\left[\left(\frac{dD}{d\eta}\right)^2 + \{\eta^2 - \mu^2\}D^2 + \frac{\mu^2}{2}D^4\right], \tag{9}$$

where B and C are constants and D is now $D(\eta, \zeta_0)$. By minimizing ΔF with respect to ζ_0, one obtains with the boundary condition Eq. (7), with Eq. (5), and the condition that the slope and the amplitude of $D(\zeta)$ is zero for large values of ζ,

$$\zeta_0^2 = \mu^2[1 - \tfrac{1}{2}D^2(0)], \tag{10}$$

where $D(0) \equiv D(\zeta = 0) \neq 0$ is the amplitude of the order parameter at the surface of the superconductor. Equation (10) shows that the phase of the order parameter is related to the applied magnetic field and also to the amplitude of the order parameter on the boundary surface between superconductor and vacuum.

Equation (5) can now be solved numerically with the boundary condition Eq. (7) and with the help of Eq. (10) for a given magnetic field (μ^2 = constant) by selecting a value of $D(0)$ which makes the function $D(\zeta)$ for large values of ζ (in the bulk of the metal) become very small compared to $D(0)$. This is obviously the correct solution for $H_0 \geq H_{c2}$ because no superconductivity exists in the bulk of the metal for magnetic fields at and above H_{c2}. Below H_{c2} this is justified only if

$$D^2(0) \gg \langle|\Psi_B|^2\rangle = \frac{\kappa - H_0}{\kappa\beta(1-1/2\kappa^2)}. \tag{11}$$

In turns out that $D(0)$ is of the order of unity, and therefore the computed solutions of Eq. (5) are correct within our approximations which are $\kappa \gg 1$ and $(\kappa - H_0)/\kappa = (\mu^2-1)/\mu^2 \ll 1$.

The solutions of Eq. (5) are shown in Fig. 1 as a function of the parameter $1/\mu^2 = H_0/H_{c2}$. At H_{c3} the amplitude $D(0)$ approaches zero. Near H_{c3} the D^2 term can then be neglected with respect to unity in Eq. (5), and the minimum value of μ^2 ($\equiv$ maximum magnetic field) may then be obtained from Eq. (5) with an arbitrary amplitude of $D(\zeta)$ and the condition $\mu^2 = \zeta_0^2$. Saint-James and de Gennes[3] obtained $\mu^2 = 0.590$, which determines H_{c3}.

Recently Abrikosov[4] has calculated the magnetic-field dependence of the superconducting surface sheath between H_{c2} and H_{c3} by assuming a Gaussian trial function for the order parameter whose maximum value is at the boundary surface. There is good agreement between his and our results except that the maximum order parameter is not strictly at the surface except for magnetic fields close to H_{c3}. Abrikosov's order parameters at the surface are about 8 % larger than those shown in Fig. 1.

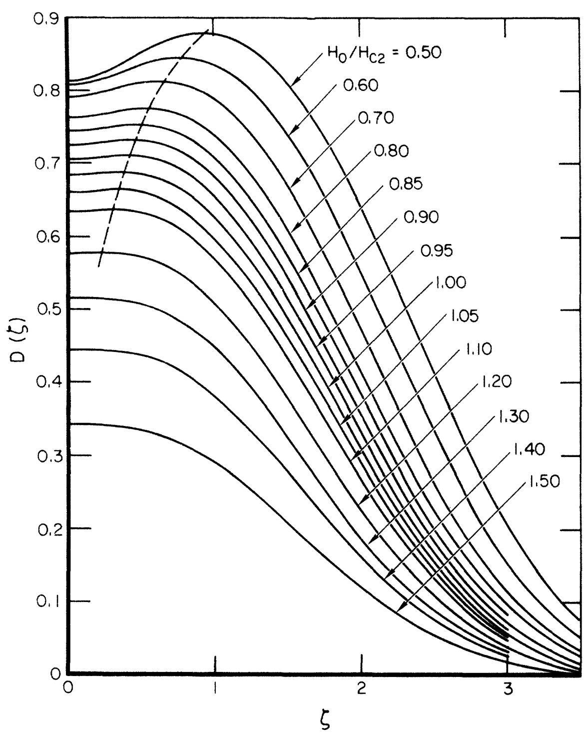

FIG. 1. The normalized amplitude of the order parameter $D(\zeta)$ [Eq. (2)] near the surface of a type-II superconductor for $\kappa \gg 1$ is shown as a function of the normalized distance $\zeta = (x/\xi)(H_0/H_{c2})^{1/2}$ and the parameter H_0/H_{c2}. For $H_0/H_{c2} < 1$ the plotted values are strictly correct only within the approximation $(H_{c2}-H_0)/H_{c2} \ll 1$. The position of the maximum of the order parameter is indicated by the dashed line. For $H_0/H_{c2} > 1.2$ the maximum of the order parameter is located at the surface within the accuracy of the plot.

The order parameters of Fig. 1 stay approximately "constant" over a larger distance for small values of ζ compared to a Gaussian. This makes the superconducting sheath of Fig. 1 thicker compared to a sheath one obtains by assuming a Gaussian order parameter. For magnetic fields below H_{c2} the curves shown in Fig. 1 are a very good approximation for the condition $(\mu^2-1)/\mu^2 \ll 1$ because the maximum order parameter in the bulk of the metal is very small compared to that on the surface. When $(\mu^2-1)/\mu^2$ is only small compared to unity, one would expect that a two-dimensional periodic solution takes over in the bulk of the metal whose maximum value of the order parameter will still be small compared to the maximum value near the surface, and whose effect it will be to modulate spatially the order parameter near, but not at, the surface. From Fig. 1 one would expect that below H_{c2} the bulk periodic solution will establish itself at a distance of about three to four coherence lengths ξ from the surface. An exact calculation which describes simultaneously the "one-dimensional surface sheath" and the two-dimensional solution in the bulk for magnetic fields considerably smaller than H_{c2} seems to be very complex.

From the above calculations one would expect that experiments which test only the surface of a superconductor will show no drastic effects near H_{c2} as long as the magnetic field is applied parallel to the surface of the superconductor. This appears indeed to be the case, if Tomasch's profile parameter in his electron tunneling experiments is interpreted as reflecting the behavior of the energy gap near the surface in a fairly direct way. With this interpretation of the profile parameter it appears also from the experimental results[5] that the "surface sheath" does also exist in magnetic fields small compared to H_{c2}. Some experiments, for example, which are likely to be affected by the superconducting sheath below and above H_{c2} are high-frequency experiments in the upper Gc/sec range and the near infrared, dc measurements if an appreciable amount of current is carried on the surface, electron tunneling experiments, flux entry through the surface, and the Kapitza resistance.

The author is indebted to R. D. Kessinger for the numerical computation of Fig. 1, and to L. J. Barnes, T. G. Berlincourt, R. R. Hake, A. G.

Pression, and W. J. Tomasch for discussions related to this work.

*Work supported by the U. S. Atomic Energy Commission.

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EXPERIMENTAL EVIDENCE FOR PARITY IMPURITY IN A NUCLEAR GAMMA TRANSITION*

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In this communication we present a preliminary account of an experimental investigation of parity admixture in nuclear states. Our measurements give evidence of a small parity admixture in a nuclear gamma transition of 482 keV in Ta^{181}. The size of the observed effect is in order-of-magnitude agreement with estimates based on the current-current hypothesis of weak interactions,[1] which predicts a strangeness-preserving nonleptonic weak coupling. Some of our results have been presented earlier.[2]

A number of experiments[3] have been done in the last few years with the aim of finding such a parity admixture. With the exception of recent work by Abov, Krupchitsky, and Oratovsky,[4] who have reported a parity-nonconserving term in the angular distribution of neutron-capture gamma rays, no evidence has been found up to now. Recently, Michel[3] has analyzed these experiments and calculated the magnitude of the effects of a parity-nonconserving force as postulated in the weak interaction scheme.

We report here on a measurement of the circular polarization of a gamma transition of 482 keV in Ta^{181}. This transition has been selected since it is found to be a particularly favorable candidate for finding a parity admixture. The transition taking place between a $\frac{5}{2}^+$ and a $\frac{7}{2}^+$ state with asymptotic numbers (402) and (404), respectively, according to the Nilsson scheme[5] has been observed to be strongly hindered. The $M1$ part is hindered by a factor of 3×10^6, and the $E2$ part which is the dominant (97 %) decay mode is hindered by a factor of 35 compared to the Weisskopf estimate. Following Michel's arguments and assuming a term of the type $\vec{\sigma}\cdot\vec{p}$ to cause an admixture of negative-parity states, one can estimate the size of the $E1$ matrix-element and thus the size of the $E1$-$M1$ interference. The Nilsson states $\frac{5}{2}^-$(503 and 303) and $\frac{7}{2}^-$(503 and 303) are presumably mainly responsible for this interference since transitions from and to these states are classified as unhindered. The $E1$-$M1$ interference gives rise to a circular polarization, P, of the gamma ray. This polarization can be expressed in the following way:

$$P=-[2/(1+q^2)]FR.$$

The quantity q is the mixing ratio between competing regular multipoles of lowest order ($M1$ and $E2$ in our case), F is the ratio of the parity-nonconserving potential of the form $H_{\text{int}}=G''\vec{\sigma}\cdot\vec{p}$ to the total (parity-conserving) nuclear potential, and R is a quantity that depends on the nuclear structure only.

For the case of Ta^{181} the quantity F has been estimated by Michel[3] to be $F=8\times10^{-7}$. The quantity R is due to a contribution R^+ from the (503) particle states and R^- from the (303) hole states. For the former we have

$$R^+=\frac{\alpha(\frac{5}{2}^-503,\frac{5}{2}^+402)\mathfrak{M}(E1,\frac{5}{2}^-503\to\frac{7}{2}^+404)+\alpha(\frac{7}{2}^-503,\frac{7}{2}^+404)\mathfrak{M}(E1,\frac{5}{2}^+402\to\frac{7}{2}^-503)}{\mathfrak{M}(M1,\frac{5}{2}^+402\to\frac{7}{2}^+404)},$$

where α characterizes the amplitude of the admixture of the states from the $\vec{\sigma}\cdot\vec{p}$ force. We have, for example,

$$\alpha(\tfrac{5}{2}^-503,\tfrac{5}{2}^+402)=\frac{\hbar}{m_0r_0}\frac{\langle\frac{5}{2}^-503|\vec{\sigma}\cdot\vec{p}|\frac{5}{2}^+402\rangle}{E_{402}-E_{503}}.$$

A detailed calculation was recently performed by Wahlborn[6] for this case with the result

$$P=-1.3\pm0.4(e_1/e)\times10^{-4}.$$

e_1 is the effective charge associated with the