with respect to the triton beam. The efficiency of the counter was obtained by measuring the yield from a standard RaBe source placed at the target. Errors up to 20 percent in total yield measurements may arise from a difference in sensitivity of the counter to neutrons from Li+T and from the RaBe source. The neutron energy-spectrum from Li+T is not known.

Data were obtained in a series of eight runs using beams of mass 6 (TT^+), mass 4 (HT^+), and mass 3 (principally T^+), which covered a triton energy range from 0.258 to 2.080 Mev. One of each of the following targets was used: lithium on aluminum, lithium on tungsten, aluminum blank, tungsten blank.

Background measurements were obtained from runs on the metal blanks. For triton energies below 1.0 Mev, backgrounds were negligible. Above 1.5 Mev, tritons on aluminum produced backgrounds which were twice as large as those for tritons on tungsten. At angles greater than 90°, backgrounds were at least twice as large as those at 0°.

For the machine energies available, probable neutron-producing reactions are:

$$Li^6+T\rightarrow 2He^4+n+16.02\text{ Mev}, \quad (1)$$
$$Li^6+T\rightarrow Be^8+n+15.97\text{ Mev}, \quad (2)$$
$$Li^7+T\rightarrow Be^9+n+10.52\text{ Mev}, \quad (3)$$
$$Li^7+T\rightarrow 2He^4+2n+8.88\text{ Mev}, \quad (4)$$
$$Li^7+T\rightarrow Be^8+2n+8.83\text{ Mev, and} \quad (5)$$
$$Li^7+T\rightarrow He^5+He^4+n+8.08\text{ Mev}. \quad (6)$$

The compound nucleus from Li^7+T is Be^{10} with an excitation energy of 17.324 Mev. Contributions from the 7.4 percent Li^6 present in the targets are assumed to be small.

The neutron yield at 0°, expressed as a cross section, is shown in Fig. 1. Apparent resonances occur at triton energies of 0.84 and 1.70 Mev. One may assume these peaks to indicate virtual states in Be^{10} of 17.91 Mev and 18.51 Mev.

Figure 2 shows two angular distributions of neutron yields, expressed in terms of differential cross sections. Numerical integration of the data at 0.728 Mev gives a total cross section of

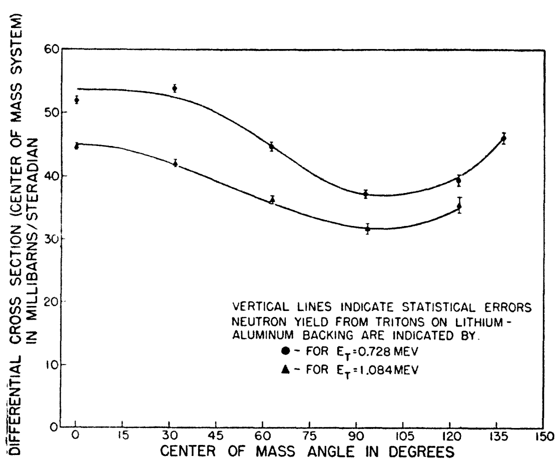

FIG. 2. Angular distribution of neutron yields in the center-of-mass system, expressed as cross section values, for the Li+T reactions.

approximately 0.5 barn. An average total cross section ~0.7 barn is obtained for a triton energy range of 0.25 to 2.0 Mev. Almqvist[3] reports an average total cross section of 1.5 barns for a thick target and a maximum triton energy of 2.6 Mev.

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On the Spin of the Nuclear Ground State in the *jj*-Coupling Scheme

IGAL TALMI*
Physikalisches Institut der Eidgenössischen Technischen Hochschule, Zürich, Switzerland
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RECENTLY, some interest has arisen in the problem concerning the spin of the ground state of the j^n configuration predicted for Majorana forces in the jj-coupling scheme.

Racah[1] showed, on the basis of the "long range" approximation, in which one assumes that the potential between a pair of nucleons can be approximated by a square well in the region where the amplitudes of the wave functions are important, that the ground state has the minimum value of spin J allowed by the Pauli principle. On the other hand, results of Mayer[2] show that in the case of δ-type interaction between the nucleons, the spin of the ground state is j if n is odd and 0 if n is even. This case should be referred to as the "short range" approximation, because here one considers an interaction potential which differs from zero only in a region so small that the change of the wave function in it is small compared to the wave function itself.

In order to see how these results depend on the approximations used, it is useful to develop the interaction potential in a series of Legendre polynomials.[3]

$$V(|r_1-r_2|)=\sum_{k=0}^{\infty} f_k(r_1, r_2)P_k(\cos\omega).$$

The energy levels are then given as finite sums of the radial integrals

$$F^k=\int_0^\infty\int_0^\infty R_l^2(r_1)R_l^2(r_2)f_k(r_1, r_2)dr_1dr_2,$$

the coefficients of which are integrals of products of spherical harmonics. The term values of the Majorana interaction for the levels in the jj-coupling are easily found by the method of sums in a way similar to that used for ordinary forces,[4] and the relative order can be determined as a function of the parameters F^k.

It is easy to see that in the "long range" limit

$$f_k(r_1, r_2)=\frac{2k+1}{2}\int_{-1}^{+1} V(|r_1-r_2|)P_k(\cos\omega)d\cos\omega$$
$$=\frac{2k+1}{2}V\int_{-1}^{+1} P_k(\cos\omega)d\cos\omega,$$

where V is the (fixed) depth of the potential well; therefore, only F^0 does not vanish. This approximation is accomplished by setting

$$F^k \ll F^0 \quad k\geqslant 1.$$

In the approximation of Mayer we simply have the expansion of the δ-function in a series of orthogonal functions

$$\delta(\cos\omega-1)=\sum_{k=0}^{\infty}\Theta_{k0}(\cos\omega)\Theta_{k0}(1)=\tfrac{1}{2}\sum_{k=0}^{\infty}(2k+1)P_k(\cos\omega),$$

and the limit is reached by putting

$$F^k=(2k+1)F^0.$$

The simplest configuration in which the two approximations give different order of levels is $(d_{5/2})^3$, where we obtain

$$E_{5/2}=4/5F^0-(56/245)F^2-(217/735)F^4,$$
$$E_{3/2}=3/5F^0+(12/245)F^2-(209/735)F^4.$$

Assuming a special form of potential and of wave functions, one is able to calculate the F^k. Kurath[5] took a gaussian potential and harmonic oscillator wave functions without nodes and obtained the result that a "cross-over" of these levels occurred at a value of r_l, which lies in a region of physical interest. The gaussian potential yields the two extreme limits: if r_0 tends to infinity, the potential approximates a square well; and if r_0 tends to zero and V is appropriately increased (for the well to have a fixed volume), the δ-potential is obtained.

The case of the Yukawa potential

$$V(|r_1-r_2|)=V(r_{12})=\exp(-r_{12}/r_0)/(r_{12}/r_0),$$

is different. For a vanishing r_0 (and appropriately increased V) it gives the Mayer limit; but if r_0 tends to infinity, not a square well but the coulomb potential is obtained, in which case the values of the F^k (for $k\geqslant 1$), while being smaller than that of F^0, are not negligible compared to it.

For the Yukawa potential, it is possible to calculate the F^k by elementary integrations, taking the coulomb-field wave functions without nodes

$$R_l(r)=N_l\, r^{l+1}\exp(-r/r_l),$$

with the aid of the addition theorem

$$\exp\left(-\frac{(r_1^2+r_2^2-2r_1r_2\cos\omega)^{\frac{1}{2}}}{r_0}\right)\Big/\frac{(r_1^2+r_2^2-2r_1r_2\cos\omega)^{\frac{1}{2}}}{r_0}=\sum_{k=0}^{\infty}(2k+1)\frac{K_{k+\frac{1}{2}}(r_1/r_0)I_{k+\frac{1}{2}}(r_2/r_0)}{(r_1/r_0)^{\frac{1}{2}}(r_2/r_0)^{\frac{1}{2}}}\quad r_1\geqslant r_2.$$

The result, however elementary, is complicated: all the F^k are expressed as the quotient of two polynomials in the ratio $\lambda=r_0/r_l$. The numerator is a polynomial of the degree $4l+2$ (with coefficients that are complicated functions of k and l) multiplied by λ^3, and the denominator is $(1+2\lambda)^{4l+4}$.

Returning to the configuration $(d_{5/2})^3$, we find that the cross-over occurs at a value $\lambda=10$. We compute r_l in terms of the nuclear radius R by the relationship $R=\frac{1}{2}r_l[(4l+4)(4l+3)]^{\frac{1}{2}}$, which gives for d-nucleons $r_l=R/3.74$ and determines that the cross-over occurs at $R=r_0\times 3.74/10=\frac{3}{8}r_0$.

Taking for r_0 the value 1.4×10^{-13} cm, we see that the $J=5/2$ level is lower than the $J=3/2$ level for all nuclei having $R>0.53\times 10^{-13}$ cm which, in fact, applies to all cases where this configuration appears.

In order to compare this result with that cited above we must choose the wave functions so that they will give equal nuclear radii. For the oscillator wave functions, $R=\frac{1}{2}r_l(2l+3)^{\frac{1}{2}}$, and in this case $r_l=R/1.3$. The cross-over occurs at $\lambda=1.32$, or at $R\sim r_0$. The difference in the results can be attributed to the singularity of the Yukawa potential at the origin.

The above calculations support the assumption that the Mayer approximation is more reasonable than the "long range" one for the physically interesting cases.

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* Hebrew University, Jerusalem, Israel.

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Adiabatic Magnetization of a Superconductor

R. L. Dolecek
Naval Research Laboratory, Washington, D. C.
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It has been demonstrated[1,2] that the suppression of superconductivity by a magnetic field is accompanied by the absorption of energy. However, knowledge of possible changes in entropy upon application of subthreshold magnetic fields is necessary for thermodynamic treatment of this phenomenon.[3]

In initiating a study of the adiabatic magnetization of a superconductor, an upper limit of possible entropy change upon application of subthreshold magnetic fields has been established. This was done by measuring the temperature change accompanying the adiabatic magnetization of a superconducting polycrystalline tin sphere by subthreshold fields. For this geometry the transition from the superconducting to the intermediate state should be quite well defined at a threshold field equal to two-thirds the critical field for bulk material, and the application of a magnetic field smaller than this threshold value should not convert any portion of the sample to the intermediate state. A carbon thermometer was employed for the measurement of temperatures. Lack of complete thermal insulation of the sample, and small induction effects upon the application of the magnetic fields, limited the accuracy of measurement of temperature changes to approximately 10^{-4} degree K. To this accuracy, no measurable changes were observed upon the application of subthreshold magnetic fields. The result is in agreement with the observations of Keesom and Kok.[1] Sample data are given in Table I. For analysis, the temperature (T) at which the field was applied is expressed in terms of the transition temperature (T_c) of the sample, the upper limit of possible entropy change (ΔS) is expressed in percent of the theoretical entropy change ($\Delta S'$) that would be expected to accompany the complete suppression of superconductivity,[4] and the magnetic field applied (H) is expressed in terms of the threshold field (H_T) at the temperature (T). The data of de Haas and Engelkes[5] were used in calculating the critical fields. As evidenced by measurable entropy changes, the transition temperature and threshold fields for the sample used were in good agreement with these data.

Table I. Upper limit of entropy change produced by the application of subthreshold magnetic fields to a superconducting tin sphere.

Reduced temperature of sample T/T_c	Magnetic field applied H/H_T	Upper limit percent entropy change $100\times\Delta S/\Delta S'$
0.65	0.76	0.038
0.73	0.92	0.053
0.76	0.72	0.031
0.80	0.85	0.040
0.86	0.96	0.14
0.91	0.43	0.29

Data obtained in this investigation indicate little if any change in the entropy of a superconductor upon application of subthreshold magnetic fields. It is felt that, in thermodynamic treatments of the magnetization of superconductors, the assumption of zero entropy change upon application of subthreshold magnetic fields is valid. Study of the entropy change of a superconductor upon application of magnetic fields which exceed the threshold value is in process.

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Meson to Proton Mass Ratios*

Walter H. Barkas, F. M. Smith, and Eugene Gardner†
Radiation Laboratory, University of California, Berkeley, California
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An extended series of measurements has been made on the ratios of meson masses to the mass of the proton. As in our previous work,[1] momenta and ranges of the mesons were measured, but a number of improvements over our earlier procedures have been devised to reduce systematic errors. Using prior knowledge of the approximate mass ratios, protons and mesons from separate targets in the 184-in. cyclotron were magnetically selected so as to lie in the same 5 percent velocity interval. The particles were stopped in the same nuclear emulsion, which they entered at a small angle to the surface. Relative momenta were calculated with errors of less than one part per thousand for the orbits,[2] which are approximately semi-circular. The rectified ranges in emulsion of