TABLE I.

Ratio: $\frac{\text{RdTh neutrons}}{\text{Ga neutrons}}$	Ratio: $\frac{\text{RdTh protons}}{\text{Ga protons}}$	Ratio: $\frac{\text{Photo-proton ratio}}{\text{Photo-neutron ratio}}$
0.96 ± 0.05	0.98 ± 0.04	1.02 ± 0.08
1.41 ± 0.04	1.43 ± 0.04	$1.01_4 \pm 0.06$

For the photo-proton ratio, the number of counts from the ionization chamber was recorded at the same discriminator bias voltage for each source. A correction for "wall effect" in the chamber has to be considered. If the range of a proton were proportional to its energy, a constant fraction of the photo-protons would be unrecorded. As it is not, the wall effect corrections for the two sets of photo-protons (estimated to be 150 kv and 210 kv) are different. It is possible to adjust the deuterium pressure ($2\frac{1}{2}$ atmos.) so that the total correction is only 10 percent, in which case the uncertainty in the ratio of counts is unlikely to exceed 1 percent.

To integrate the photo-neutron flux, the sources were placed at the center of a sphere of heavy water of 8-cm radius, contained in a thin copper vessel. This was suspended at the center of a cylindrical tank, 60 cm in diameter and 60 cm deep, containing an aqueous solution of $MnSO_4$. After irradiation for a known time, the whole solution was well mixed by a circulating pump and the mean activity determined in an aliquot fraction with a pair of thin glass-walled immersion type Geiger-Müller counters.

Results of two series of measurements are shown in Table I. Clearly the low energy γ-quanta from gallium are not producing many photo-neutrons.

By combining the experimental value[6] of the ratio

$$\frac{\text{2.51-Mev quanta}}{\text{2.21-Mev quanta}}$$

with the theoretical variation in cross section for the photo-disintegration of deuterium near the threshold,[8, 9] one can calculate the expected value of the ratio we have measured for any value of Q_γ. For $-Q_\gamma = 2.185$ Mev the value of the ratio would be 1:65; if, as an extreme case, one supposes our results are consistent with a value of 1:1 for the ratio, then Q_γ and the quantum energy of the Ga^{72} γ-ray cannot differ by more than 3 kev, and therefore the binding energy of the deuteron cannot be less than 2.207 Mev.

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The Magnetic Moments of Sn^{117}, Sn^{119}, and Pb^{207}*

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July 18, 1949

USING the nuclear induction spectrometer described in a previous letter,[1] the magnetic moments of the nuclei Sn^{117}, Sn^{119}, and Pb^{207} have been remeasured. The results are in good agreement with those from hyperfine structure measurements through which they were previously known, and which indicated for both tin nuclei a spin of $\frac{1}{2}$[2] and a magnetic moment of -0.89 nuclear magneton,[3] and for the nucleus Pb^{207} a spin $\frac{1}{2}$[4] and a magnetic moment of $+0.6$ nuclear magneton.[5]

The tin resonances were observed in the neighborhood of 7 MC, with a magnetic field near 4500 gauss, using a sample consisting of an aqueous solution of 5.3 molar $SnCl_2$ with 1.0 molar $MnCl_2$. These frequencies were compared to that of the resonance of Na^{23} when the above sample was replaced with one consisting of an aqueous solution of 0.69 molar NaCl with 1.0 molar $MnCl_2$. The resonances of Sn^{117}, Sn^{119}, and Na^{23} were observed in that order until six of each had been recorded. After correcting for magnetic field drifts, the ratios of the resonant frequencies in the same field were found to be

$$\nu(Sn^{117})/\nu(Na^{23}) = 1.3464 \pm 0.0004,$$
$$\nu(Sn^{119})/\nu(Na^{23}) = 1.4090 \pm 0.0005,$$
$$\nu(Sn^{119})/\nu(Sn^{117}) = 1.0465 \pm 0.0003.$$

Making a correction for electronic diamagnetism,[6] which here has been assigned an error of 15 percent, and taking $\frac{3}{2}$ for the spin and $+2.217 \pm 0.002$ nuclear magnetons for the magnetic moment of Na^{23},[7] one obtains

$$\mu(Sn^{117}) = -(1.000 \pm 0.001)\mu_N,$$
$$\mu(Sn^{119}) = -(1.047 \pm 0.001)\mu_N.$$

The ratio of this magnetic moments is, of course, given by the ratio of their resonant frequencies. Sn^{117} and Sn^{119} have respectively a natural abundance of 9.1 percent and 9.8 percent,[8] or a ratio of abundance of 0.93; the relative signal amplitudes, when a correction for their different gyromagnetic ratios is made, was 0.86 ± 0.12, which served to identify the isotopes. By comparing the recorded tin signals with the Na^{23} signals, the sign of the magnetic moments of both tin nuclei was seen to be negative, and the signal amplitudes were consistent with spin $\frac{1}{2}$.

Using as a sample an aqueous solution of 1.0 molar $Pb(C_2H_3O_2)_2$ and 0.8 molar $Mn(C_2H_3O_2)_2$, the frequency of the lead resonance which occurred near 5.8 MC at 6600 gauss was repeatedly compared with that of the Na^{23} resonance from the sodium sample above. A resonance frequency ratio of

$$\nu(Pb^{207})/\nu(Na^{23}) = 0.7901 \pm 0.0002$$

was found which leads, in analogy to the above results from the magnetic moments of the tin nuclei, to

$$\mu(Pb^{207}) = +(0.588 \pm 0.001)\mu_N.$$

The sign of the magnetic moment was found to be positive, and the signal amplitudes indicate a spin of $\frac{1}{2}$.

The author wishes to express his appreciation for the frequent and helpful advice of Professor Felix Bloch during the period of these measurements.

* Assisted by the Joint Program of the ONR and the AEC.

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The Application of Dyson's Methods to Meson Interactions

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June 13, 1949

BY developing a combination of the electrodynamic theories of Schwinger[1] and Feynman,[2] Dyson[3] has given comparatively simple rules whereby the matrix element of the Hamiltonian for an effect such as the Lamb-Retherford shift[4] can be written down directly in terms of invariant functions $D_F(x)$ and $S_F(x)$.

It will be shown here that this technique can be applied even to the more general meson interactions containing derivatives of the field variables which lead to Hamiltonians in the interaction representation which depend explicitly on the direction of the general space-like surfaces.

When there is no external field, the S-matrix can be expressed in the form (reference 3, Section III, Eq. (32))

$$U(\infty)=\sum_{n=0}^{\infty}U_n(\infty)=\sum_{n=0}^{\infty}(-i/\hbar c)^n[1/n!]\int_{-\infty}^{\infty}dx_1\cdots \times\int_{-\infty}^{\infty}dx_n P(H(x_1),\cdots,H(x_n)). \quad (1)$$

As an example the vector interaction of neutral scalar mesons with nucleons will be considered. The Hamiltonian is[5]

$$H(x)=(1/c\kappa)j_\mu(x)\partial\phi(x)/\partial x_\mu+(1/2c^2\kappa^2)(j_\mu(x)n_\mu(x))^2, \quad (2)$$

where $j_\mu(x)=igc\bar{\psi}(x)\gamma_\mu\psi(x)$ and $n_\mu(x)$ is the normal to a one-parameter family of space-like surfaces $\sigma(x)$.

Now

$$[\phi(x),\phi(x')]=i\hbar c\Delta(x-x') \quad (3)$$

and

$$\langle\{\phi(x),\phi(x')\}\rangle_0=\hbar c\Delta^{(1)}(x-x'). \quad (4)$$

Also, in the "natural" coordinate system (with time axis in the direction $n_\mu(x)$), $\partial\Delta(x)/\partial x_4=i\delta(x_1)\delta(x_2)\delta(x_3)$. And for the function $\epsilon(x)$ defined by Schwinger,[1] $\partial\epsilon(x)/\partial x_4=-2i\delta(x_0)$. Hence

$$j_\mu(x)j_\nu(x')\langle P(\partial\phi(x)/\partial x_\mu,\partial\phi(x')/\partial x_\nu')\rangle_0 =\tfrac{1}{2}\hbar c j_\mu(x)j_\nu(x')\left(\frac{\partial^2}{\partial x_\mu\partial x_\nu'}\Delta^{(1)}(x-x')+i\epsilon(x-x')\frac{\partial^2}{\partial x_\mu\partial x_\nu'}\Delta(x-x')\right).$$

The factor $\epsilon(x-x')$ can be taken under the differential sign and the additional terms subtracted explicitly. These can be evaluated in the "natural" coordinate system and then expressed in general coordinates, giving

$$j_\mu(x)j_\nu(x')\langle P(\partial\phi(x)/\partial x_\mu,\partial\varphi(x')/\partial x_\nu')\rangle_0 =\tfrac{1}{2}\hbar c j_\mu(x)j_\nu(x')(\partial^2/\partial x_\mu\partial x_\nu')\Delta_F(x-x') +(\hbar c/i)(j_\mu(x)n_\mu)(j_\nu(x')n_\nu')\delta_4(x-x').$$

The terms of order g^n in $U_n(\infty)$ are

$$\left(\frac{-i}{\hbar c^2\kappa}\right)^n\frac{1}{n!}\int_{-\infty}^{\infty}dx_1\cdots\int_{-\infty}^{\infty}dx_n P\left(j_\mu(x_1)\frac{\partial\varphi}{\partial x_{1,\mu}},\cdots,j_\nu(x_n)\frac{\partial\phi}{\partial x_{n,\nu}}\right). \quad (5)$$

The matrix element for an effect in which there are $n-2$ free mesons will have one pair of meson operators combining to give a vacuum effect. This pair can be chosen in $n(n-1)/2$ ways and the term is thus

$$\begin{aligned}&\frac{1}{2}\left(\frac{-i}{\hbar c^2\kappa}\right)^n\frac{1}{(n-2)!}\int dx_1\cdots\int dx_n P\left(j_\mu(x_1)\frac{\partial\phi}{\partial x_{1,\mu}}\right),\cdots,\\ &\qquad j_\nu(x_{n-2})\frac{\partial\phi}{\partial x_{n-2,\nu}},\ j_\rho(x_{n-1}),\ j_\pi(x_n)\left\langle P\left(\frac{\partial\phi}{\partial x_{n-1,\rho}},\frac{\partial\phi}{\partial x_{n,\pi}}\right)\right\rangle_0\\ &=\frac{1}{2}\left(\frac{-i}{\hbar c^2\kappa}\right)^n\frac{1}{(n-2)!}\left\{\int dx_1\cdots\int dx_n P\left(j_\mu(x_1)\frac{d\varphi}{\partial x_{1,\mu}}\right),\cdots,\right.\\ &\qquad j_\nu(x_{n-2})\frac{\partial\varphi}{\partial x_{n-2,\nu}},\ j_\rho(x_{n-1}),\ j_\pi(x_n)\frac{1}{2}\hbar c\frac{\partial^2}{\partial x_{n-1,\rho}\partial x_{n,\pi}}\Delta_F(x_{n-1}-x_n)\\ &\quad+\frac{\hbar c}{i}\int dx_1\cdots\int dx_{n-1}P\left(j_\mu(x_1)\frac{\partial\varphi}{\partial x_{1,\mu}}\right),\cdots,\\ &\qquad\left. j_\nu(x_{n-2})\frac{\partial\phi}{\partial x_{n-2,\nu}},\ (j_\rho(x_{n-1})n_\rho(x_{n-1}))^2\right\}. \quad (6)\end{aligned}$$

The only other contribution to this element of this order is from $U_{n-1}(\infty)$. By a similar argument this can be shown to be just the term required to cancel the n_μ-dependent part of (6), leaving an invariant expression independent of the surface σ. For an effect involving $n-4$ free mesons there will be two terms dependent on n_μ from $U_n(\infty)$. These can be shown to cancel with similar terms from $U_{n-1}(\infty)$ and $U_{n-2}(\infty)$ leaving an invariant expression as before. A similar cancelation of the n_μ-dependent terms occurs for any effect and up to any order. The calculation of matrix elements is now no more complicated than in the electrodynamic case. The Hamiltonian can be taken to be

$$H(x)=(1/c\kappa)j_\mu(x)\partial\varphi(x)/\partial x_\mu \quad (7)$$

and effectively

$$\langle P(\partial\phi(x)/\partial x_\mu,\partial\phi(x')/\partial x_\nu')\rangle_0=\tfrac{1}{2}\hbar c(\partial^2/\partial x_\mu\partial_\nu')\Delta_F(x-x'). \quad (8)$$

Using these formulas, the graphical methods of Dyson[3] can be employed with only trivial changes. The Hamiltonian should of course contain two additional terms to allow for the renormalization of both the nucleon and the meson masses.

An exactly similar analysis is possible for all meson interactions with either nucleons or photons. Dyson[6] has shown that the particular case dealt with above leads to no physical effects. It serves here as a simple example of the formalism.

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Electromagnetic Induction in a Superconductor

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July 18, 1949

According to London's formulation of the macroscopic equations for a superconductor, a negligible electromotive force should be induced in a figure of rotation turned about its axis in a magnetic field. The field should penetrate only a very small distance into the rotating body and the electromotive force induced by this penetrating field would be very small.

To test this conclusion a lead ellipsoid was turned about a vertical axis in a cryostat at speeds between 3000 and 9000 r.p.m. A vertical field of 30 to 70 gauss was produced by a solenoid, and sliding contacts enabled the induced electromotive force to be measured with a potentiometer. Above the transition temperature the observed electromotive force had the expected value with an uncertainty of less than 5 percent. This uncertainty was apparently due to thermal and other extraneous electromotive forces.

When the ellipsoid was superconducting, the induced electromotive force was definitely less than 5 percent of its value above the transition temperature. It may therefore be concluded that the description by London's equations, which suggests about 10^{-3} percent of the normal effect, is consistent with this experimental result.

Preliminary results already reported as showing the full induced electromotive force were in error.[1] The rotation of the ellipsoid apparently produced enough heat in the bearings and at the sliding contacts to destroy the superconductivity while a potentiometer setting was being made. This heat was developed in the rotating system and was neither evidenced in the thermometers, nor effectively absorbed by a neighboring bath of liquid helium. It was possible to reduce the rate of development of this heat by reducing the pressure on the sliding contacts and loosening the bearings. In this way the rotor could be kept superconducting for two to five minutes.

In the final measurements, the superconductivity was checked by means of a pick-up coil concentric with and close to the circumference of the rotor. The flux through this coil due to the applied field was measured by a ballistic galvanometer. It dropped to about one-third of its value when the rotor became superconducting. This value was quite consistent with complete exclusion of the field from the rotor, since there was some space between the pick-up coil and the rotor.