comparable with the maser noise, but the present necessity of using a room-temperature isolator, which gives a minimum noise input of about 300°K, prevented this. Work is now in progress to develop a good room-temperature circulator or a low-temperature isolator. Such circuit elements, and the subsequent use of low-temperature noise sources, would also permit the generation of noise temperatures T_1 and T_2 by thermostated matched loads at accurately known temperatures, thus eliminating the calibration of a secondary standard. In the masers now being designed, the replacement of the coaxial input line with a wave guide will eliminate most of the input attenuation and hence the uncertainty arising from noise generation in the input circuit. These improvements should reduce the uncertainty in the measurement of T_M to about 5°K.

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Resonant Absorption of Gamma Rays*

S. S. Hanna and Luise Meyer-Schützmeister

Argonne National Laboratory, Lemont, Illinois

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In recent years several ingenious methods[1–5] have been used to observe the resonant absorption or fluorescence when nuclei are irradiated by gamma rays. In this investigation we have used in essence a source of gamma rays with a discrete but variable energy to observe the resonant absorption by a direct measurement of the transmission through an absorbing sample.

In the radiative capture of particles of fixed energy and direction of motion, the gamma rays emitted at a given angle are monoenergetic. Moreover, the energy of of the radiation varies smoothly with the angle of emission. The energy of the incident particles is fixed either by using a resonant capture process, or, if the capture is nonresonant, by employing a monoenergetic beam of particles and a thin target. It is not at all necessary[4] to use a reaction which is the inverse of the subsequent resonant absorption. It is usually convenient to do so, however, since the available variation in energy is limited but, in the case of reciprocal reactions, sufficient to cover the increase in energy needed to compensate for the energies lost to the recoiling nuclei in emission and in absorption.

The source of radiation was the reaction $C^{13}(p,\gamma)N^{14}$ at the resonant energy $E_p = 1.76$ Mev corresponding to the level in N^{14} at $E_{ex} = 9.18$ Mev. The 9.18-Mev radiation was detected at an angle θ through a slit of angular width $\Delta\theta$, with a NaI counter and a 10-channel analyzer. The inverse absorption $N^{14}+\gamma\rightarrow N^{14*}$ was obtained by interposing between the source and the slit a quantity of liquid nitrogen, 27.0 cm thick, contained in two Dewar flasks. The energy of the emitted radiation was then varied by changing the angle of observation with the results given in Fig. 1. Data are presented for three different widths of the slit. The measurements at the top were obtained with water instead of liquid nitrogen in the Dewar flasks.

The number of gamma rays emitted at the angle θ with energy in the interval between E and $E+dE$ will be represented by the single-level formula

$$N(E,\theta)=\frac{N_R}{[2(E-E_R)/\Gamma]^2+1},$$

where Γ is the intrinsic width of the level and the resonant energy $E_R = E_0 - (E_0^2/2Mc^2) + (E_0 v/c)\cos\theta$, to a good approximation. The energy E_0 is the unmodified energy of the transition, and v is the speed and M the mass of the radiating compound nucleus. The second term in the expression represents the energy of

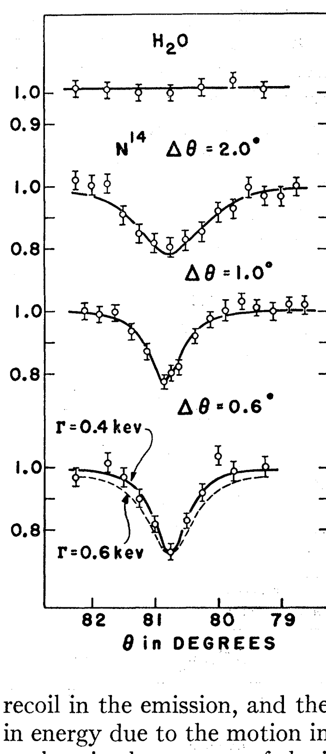

FIG. 1. Resonant absorption in the process $N^{14}+\gamma\rightarrow N^{14*}$ ($E_{ex}=9.18$ Mev), using gamma rays from the inverse reaction $C^{13}(p,\gamma)N^{14}$. The abscissa represents the angle of emission of the radiation. One degree corresponds to a change of 0.7 kev in the energy of the radiation.

recoil in the emission, and the last term gives the shift in energy due to the motion imparted to the compound nucleus in the capture of the incident particle.

We may neglect the purely electronic absorption, which, although large, is nevertheless essentially constant over the range of energies used in the experiment. Hence, the cross section for absorption is

$$\sigma(E)=\frac{\sigma_R}{[2(E-E_R')/\Gamma_a]^2+1},$$

where E_R' is equal to $E_0+(E_0^2/2Mc^2)$. The second term in this expression represents the energy of recoil in the absorption. The width Γ_a is here distinguished from the intrinsic width because of the Doppler broadening due to the thermal motion of the nuclei in the absorber.[6]

The intensity of the radiation transmitted through the absorber is given by

$$I(\theta)=\int_0^\infty N(E,\theta)e^{-\rho\sigma(E)}dE,$$

where ρ is the number of nuclei per unit area of the absorber. The transmitted intensity has a minimum when E_R is equal to E_R', i.e., when $\cos\theta_R=p_\gamma/p$, where p_γ and p are the momenta of the gamma ray and the radiating nucleus, respectively. With $E_p=1.76$ and $E_\gamma=9.18$ Mev, it follows that $\theta_R=80.8°$.

The function $I(\theta)$ is plotted in Fig. 1 for various values of the parameter Γ. Each curve is normalized to pass through the point of maximum absorption. For the data taken with $\Delta\theta=0.6°$, good agreement is obtained with $\Gamma=0.4$ kev. The Doppler broadening in the absorption process may now be neglected since it is less than 0.1 of this width. After correcting for the width of the slit, we obtain $\Gamma=(0.35\pm0.1)$ kev. This value confirms the recent result of Marion and Hagedorn[7] that $\Gamma<0.4$ kev. The depth of the valley in the plotted curves for $\Delta\theta=0.6°$ corresponds to a cross section $\sigma_R=(0.79\pm0.10)$ barn at resonance for the absorption $N^{14}+\gamma\rightarrow N^{14*}$. This value of the cross section leads to $(2I+1)\Gamma_\gamma=29$ ev, where I is the spin of the excited state and Γ_γ is the partial width for the emission of the 9.18-Mev radiation. With this value of $(2I+1)\Gamma_\gamma$, we obtain a cross section $\sigma_R(p,\gamma)=33$ mb at resonance for the capture process. If this result is increased by 10% to include transitions to the excited state at 6.44 Mev, it may be compared with the result $\sigma_R(p,\gamma)>64$ mb given by Marion and Hagedorn as computed from the measurements of Seagrave.[8]

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General Partial Wave Analysis of the Decay of a Hyperon of Spin $\frac{1}{2}$

T. D. LEE* AND C. N. YANG
Institute for Advanced Study, Princeton, New Jersey
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THIS note is to consider the general problem of the decay of a hyperon of spin $\frac{1}{2}$ into a pion and a nucleon under the general assumption of possible violations of parity conservation, charge-conjugation invariance, and time-reversal invariance. The discussion is in essence a partial wave analysis of the decay phenomena and is independent of the dynamics of the decay. *Nonrelativistic approximations are not made on either of the decay products.*

In the reference system in which the hyperon is at rest there are two possible final states of the pion-nucleon system: $s_{\frac{1}{2}}$ and $p_{\frac{1}{2}}$. Denoting the amplitudes of these two states by A and B, one observes that the decay is physically characterized by *three* real constants specifying the magnitudes and the relative phase between these amplitudes. One of these constants can be taken to be $|A|^2+|B|^2$, and is evidently proportional to the decay probability per unit time. The other two constants are best defined in terms of experimentally measurable quantities. We discuss three types of experiments:

(a) The angular distribution of the decay pion from a completely polarized hyperon at rest.

It has been pointed out before[1] that the distribution is proportional to

$$[1+\alpha\cos\chi]d\Omega, \tag{1}$$

where $d\Omega$ is the solid angle of the pion momentum vector $\mathbf{p}_\pi$ and χ is the angle between $\mathbf{p}_\pi$ and the spin of the hyperon. The constant α is given by

$$\alpha=2\,\mathrm{Re}(A^*B)/(|A|^2+|B|^2), \tag{2}$$

and characterizes the degree of mixing of parities in the decay.

That the distribution is of the form (1) follows immediately from the assumption that the spin of the hyperon is $\frac{1}{2}$. One easily proves (2) by considering the decay probabilities for the cases $\chi=0$ and $\chi=\pi$. In the former case the amplitude of decay is $(A+B)$ and in the latter $(A-B)$. One therefore obtains $(1+\alpha)/(1-\alpha)=|A+B|^2/|A-B|^2$, which results in (2).

Recent experiments[2] have indicated that the absolute value of α is quite large for Λ^0 decay. With improved statistics these experiments can establish beyond doubt that parity nonconservation is not peculiar to neutrino processes. It is, however, not possible to determine the sign of the parameter α through the experiments quoted above, as the sign of the polarization of the Λ^0 in the production process is unknown. Further, it appears that these experiments cannot give an accurate measurement of the magnitude of α because of the difficulty in deter-