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Alternate Modes of Decay of Neutral K Mesons*

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Consideration is given to the possibility that the θ^0 particle can decay via the "anomalous" modes $(e\pi\nu)$ and $(\mu\pi\nu)$. It is shown that, as a function of time, the decay rate for θ^0 disintegration into a final state of this type, having definite signs of the charges, is the sum of three terms. Two of these are the expected exponentially decaying functions, characterized by the rates λ_1 and λ_2 for decay of the particles denoted as θ_1 and θ_2 by Gell-Mann and Pais. The third term is the product of an exponential having the decay rate $(\lambda_1+\lambda_2)/2$ and an oscillating factor whose frequency depends on the difference in mass between the θ_1 and θ_2 particles. It arises from interference between θ_1 and θ_2 decay and has opposite signs for the $(e^+\pi^-\nu)$ and $(e^-\pi^+\bar{\nu})$ modes. Similar effects occur for the muon decay modes. This interference may have the effect that the rate of disintegration into any one such final state *increases* with time during the early stages of decay. The detailed shape of the decay curve may be quite sensitive to the (θ_1,θ_2) mass difference.

On the basis of presently available data on "anomalous θ^0" decay events, it is estimated that $1/\lambda_2>5\times10^{-9}$ sec.

I. INTRODUCTION

It has been suggested by Gell-Mann[1] and Nishijima[2] that, in addition to the now well-established θ^0 meson, there exists another neutral particle, the $\bar{\theta}^0$ meson, which is charge conjugate to the θ^0 and therefore has opposite sign for its "strangeness" quantum number. Gell-Mann and Pais[3] have remarked that the normal two-pion decay mode of these particles is best discussed, not in terms of θ^0 and $\bar{\theta}^0$, but rather in terms of the neutral particles θ_1 and θ_2 whose state vectors are respectively the real and imaginary parts of the θ^0 state vector. Thus, if we denote the state vector by the particle symbol, we have

$$\theta^0=2^{-\frac{1}{2}}(\theta_1+i\theta_2), \tag{1}$$

$$\bar{\theta}^0=2^{-\frac{1}{2}}(\theta_1-i\theta_2), \tag{1'}$$

so that θ_1 is even and θ_2 odd under charge conjugation.

Only one of the particles θ_1 or θ_2 is able to decay into two pions, because two pions in a state of definite relative angular momentum l have definite signature $(-1)^l$ under charge conjugation. It is of course assumed that the interactions involved in the decay are invariant under charge conjugation. If the spin of the θ^0 is taken to be even, as is indicated by the apparent existence of a mode of decay into two neutral pions,[4] then it is the θ_1 rather than the θ_2 which decays into two pions. It was pointed out by Gell-Mann and Pais that the θ_2 may disintegrate into two pions plus a photon; the partial lifetime for this process should be considerably larger than the θ_1 lifetime. They suggested that two distinct lifetimes could be used as a basis for the experimental separation of the two particles. This suggestion was followed by the specific proposal of Pais and Piccioni[5] for an experimental method to test these ideas.

Snow[6] has called attention to the possibility that similar considerations would apply to the τ^0 meson, should it and its charge conjugate $\bar{\tau}^0$ exist. Here too one would describe weak decay processes in terms of τ_1 and τ_2 particles whose state vectors are the real and imaginary parts of the τ^0 state vector. In this case both τ_1 and τ_2 can undergo three-pion decay, but a difference in lifetimes is still expected, since the τ_1 and τ_2 have three-pion states of different symmetries available to them.

The purpose of this note is to call attention to other possible features of neutral K-meson decay which suggest an additional experimental method for testing some of the ideas discussed above. The method to be discussed is based on the existence of the "anomalous θ^0" processes,[7,8] which are evidently three-body disintegrations of neutral K mesons. Some of the observed events of this kind may be interpreted in terms of the decay scheme $\tau^0\to\pi^++\pi^-+\pi^0$. However, many of the events are kinematically inconsistent with this interpretation[7,8] but are consistent with the following schemes which naturally suggest themselves:

$$K_{\mu3}{}^0\to\mu^\pm+\pi^\mp+\nu, \tag{2}$$

$$K_{e3}{}^0\to e^\pm+\pi^\mp+\nu. \tag{3}$$

The basis for this interpretation of at least some of the "anomalous θ^0" events is that there is good evidence[9] that the charged $K_{\mu3}$ decay proceeds according to

$$K_{\mu3}{}^\pm\to\mu^\pm+\pi^0+\nu. \tag{4}$$

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[1] M. Gell-Mann, Phys. Rev. **92**, 833 (1953).
[2] K. Nishijima, Progr. Theoret. Phys. (Japan) **13**, 285 (1955).
[3] M. Gell-Mann and A. Pais, Phys. Rev. **97**, 1387 (1955).
[4] Osher, Moyer, and Parker, Bull. Am. Phys. Soc. Ser. II, **1**, 185 (1956).
[5] A. Pais and O. Piccioni, Phys. Rev. **100**, 1487 (1955).
[6] G. Snow (to be published).
[7] R. W. Thompson, Progr. Theoret. Phys. **III** (to be published), contains a summary of published anomalous V^0 events.
[8] Ballam, Grisaru, and Treiman, Phys. Rev. **101**, 1438 (1956).
[9] Hoang, Kaplon, and Yekutieli, Phys. Rev. **101**, 1834 (1956).

It may also be noted that there has been direct experimental identification of electron secondaries in several of the anomalous events.[10]

We now make the supposition that, aside from differences in charge and sign of the strangeness number, there are at most two distinct kinds of K mesons, θ and τ. Then the reactions (2) and (3) are interpreted as alternate decay modes of either θ^0 or τ^0, or both. For the sake of definiteness we shall assume that they are alternate decay modes of the θ^0. The phenomena to which attention is directed here will not be qualitatively different if the other possibilities prevail.

Since the states $(e^+\pi^-\nu)$, etc., appearing in Eqs. (2) and (3) do *not* have definite signature under charge conjugation, it will in general be possible for both θ_1 and θ_2 to decay via these modes. This means that in the time dependence of θ^0 decay via the anomalous modes there will be an effect of interference between the θ_1 and θ_2 contributions. It is this effect which will be our principal concern here.

II. DECAY RATES

Taking into account the two possibilities for the signs of the final charges in reactions (2) and (3), and including the possibility that in each case we may replace neutrino by antineutrino, we have a total of eight possible "anomalous" decay modes for θ^0 (and for $\bar{\theta}^0$), in addition to the two-pion and two-pion plus photon decay modes. Let us fix our attention in particular on, say, the electron decay mode and associate with it a definite neutrino character. The generalization following from inclusion of the other modes will be evident. In terms of the state vectors, we may describe the decay schemes as follows:

$$\theta_1 \to a_1(e^+\pi^-\nu + e^-\pi^+\bar{\nu}) + b(\pi^+\pi^-) + \cdots, \tag{5}$$

$$\theta_2 \to -ia_2(e^+\pi^-\nu - e^-\pi^+\bar{\nu}) + c(\pi^+\pi^-\gamma) + \cdots, \tag{5'}$$

where the presence of additional terms corresponding to other modes of decay is to be understood. The quantities a_1, a_2, b, c, etc., are proportional to the transition amplitudes for the indicated modes. The assumption of invariance under charge conjugation has been used in an obvious way, e.g., in equating the amplitudes for $(e^+\pi^-\nu)$ and $(e^-\pi^+\bar{\nu})$ in Eq. (5). It follows from the invariance under time reversal that a_1 and a_2 may be taken to be real numbers (see Appendix I). The transition probabilities for decay of θ_1 and θ_2 are given by

$$\lambda_1 = 2a_1^2 + |b|^2 + \cdots, \tag{6}$$

$$\lambda_2 = 2a_2^2 + |c|^2 + \cdots, \tag{6'}$$

if the amplitudes are suitably normalized.

When the particle θ_i ($i=1$ or 2) is produced at time $t=0$, the amplitude of the state decays as $\exp(-\lambda_i t/2)$. However, because of the rules of associated production, it is the θ^0 that is produced in association with a Λ^0 or Σ hyperon. Hence the state vector at $t=0$ has the form indicated in Eq. (1) and its time dependence is given by

$$\psi(t) = 2^{-\frac{1}{2}}\{\theta_1 \exp(-\lambda_1 t/2 - i\omega_1 t) + i\theta_2 \exp(-\lambda_2 t/2 - i\omega_2 t)\}, \tag{7}$$

where ω_1 and ω_2 are the De Broglie frequencies corresponding to the rest energies of the θ_1 and θ_2 particles. Under these circumstances the rate at which decay into two charged pions occurs is, according to Eq. (5),

$$R(\pi^+\pi^-) = \tfrac{1}{2}|b|^2 \exp(-\lambda_1 t),$$

as expected. On the other hand, the rate at which $(e^+\pi^-\nu)$ decay occurs is, as demonstrated in Appendix II,

$$R(e^+\pi^-\nu) = \tfrac{1}{2}|a_1 \exp(-\lambda_1 t/2) + a_2 \exp(-\lambda_2 t/2) \exp(i\Delta\omega t)|^2,$$

where $\hbar\Delta\omega/c^2$ is the θ_1, θ_2 mass difference. Thus

$$R(e^+\pi^-\nu) = \tfrac{1}{2}a_1^2 \exp(-\lambda_1 t) + \tfrac{1}{2}a_2^2 \exp(-\lambda_2 t) + a_1 a_2 \cos(\Delta\omega t) \exp[-(\lambda_1+\lambda_2)t/2]. \tag{8}$$

Similarly, the negative electron rate is

$$R(e^-\pi^+\bar{\nu}) = \tfrac{1}{2}a_1^2 \exp(-\lambda_1 t) + \tfrac{1}{2}a_2^2 \exp(-\lambda_2 t) - a_1 a_2 \cos(\Delta\omega t) \exp[-(\lambda_1+\lambda_2)t/2]. \tag{9}$$

We have then the interesting result that the electron decay mode (and similarly for the muon decay mode), for given signs of the final charges, does not have the form merely of a sum of two exponential terms with rates λ_1 and λ_2. There is, in addition, an interference term containing an exponential factor characterized by the average rate $(\lambda_1+\lambda_2)/2$ and an oscillating factor which depends on the θ_1, θ_2 mass difference. Furthermore, the interference term has opposite sign for the $(e^+\pi^-\nu)$ and $(e^-\pi^+\bar{\nu})$ modes; and it could happen that the rate of decay into one or another of these final states will show an *increase* with time during the early stages of decay. The occurrence of this latter effect depends on the relative magnitudes of a_1 and a_2. It would be particularly marked if $a_1 \approx \pm a_2$. This can be understood by noting that if, for example,

$$a_1 = a_2 \tag{10}$$

then according to Eqs. (1) and (5) the θ^0 does not decay into $(e^-\pi^+\bar{\nu})$, so that initially this process does not occur. However, at later times the state vector contains an increasing admixture of $\bar{\theta}^0$, as can be seen from Eq. (7), and $\bar{\theta}^0$ can decay into $(e^-\pi^+\bar{\nu})$. This striking difference in decay rates into final states with different signs of the charges may of course also occur in the muon decay mode.

These effects are illustrated in Fig. 1 which shows decay curves for the case where $a_1 = a_2$. It has been assumed here that $\lambda_2 \ll \lambda_1$. Two choices for the θ_1, θ_2 mass difference, corresponding to $\Delta\omega = 0$ and $\Delta\omega = \lambda_1$,

[10] Slaughter, Block, and Harth, Bull. Am. Phys. Soc. Ser. II, 1, 186 (1956).

are illustrated for comparison. Note that in the latter case the mass difference is about 4×10^{-6} electron volts. The extreme sensitivity of the shapes of the curves to the mass difference suggests that a measurement of the shapes might be used to determine the mass difference. For this purpose it would be necessary to determine the ratio a_1/a_2 (taken to be unity in Fig. 1). This could be accomplished by finding the initial values of the two rates, since

$$\left[\frac{R(e^-\pi^+\nu)}{R(e^+\pi^-\nu)}\right]_{t=0}=\left(\frac{1-a_2/a_1}{1+a_2/a_1}\right)^2. \tag{11}$$

Of course, if it happens that the mass difference is large, i.e., if $\Delta\omega\gg\lambda_1$, then the interference effect would be entirely washed out by the rapid oscillations and there would be no experimental distinction between the electron and positron decay rates. The interference effect would also cease to be observable if the amplitudes a_1 and a_2 have very different magnitudes.

It should be remarked that a determination of the θ_1, θ_2 mass difference may be of value in connection with questions concerning the nature of the weak interactions responsible for θ decay. Presumably the mass difference is caused by virtual processes associated with weak interactions, since strong interactions acting differently on the θ_1 and θ_2 would make possible the rapid transition $\theta^0\rightleftarrows\bar{\theta}^0$, in violation of the accepted selection rules. It is well known that the mass difference produced by virtual processes involves an integral over energy of the square of the matrix element to all virtual states. Hence a determination of the mass difference would give some insight into the energy dependence of those weak interactions which distinguish between the θ_1 and θ_2, such as the interaction producing the two-pion decay.

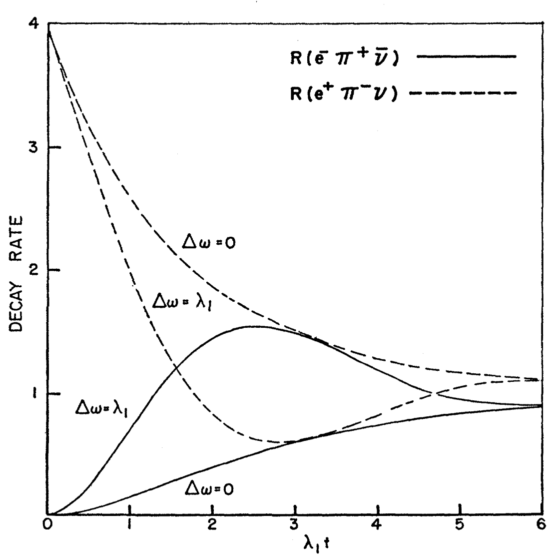

FIG. 1. Decay rates of the θ^0 into the modes $(e^+\pi^-\nu)$ (dashed curves) and $(e^-\pi^+\bar{\nu})$ (solid curves) under the assumption $a_1/a_2=1$. The abcissa is the time in units of the θ_1 mean life. The two curves in each case correspond to different choices for the mass difference.

We have remarked that there are eight possible decay modes of θ_1 and θ_2, of the types represented by Eqs. (2) and (3). One may ask whether there is some general principle leading to relationships among the amplitudes for the various modes. There does not appear to be any well-established principle of this kind. However, if it is assumed that the strangeness concept can be extended to the light fermions,[11] then it turns out that there are very severe restrictions on the amplitudes of the different modes. If we follow the reasoning of reference 11, we must assign half-integral strangeness numbers to the light fermions and characterize the decay into these particles by $\Delta S=0$. Then, of the decay processes of the types represented by Eqs. (2) and (3), only the following are allowed for θ^0:

$$\theta^0\rightarrow e^++\pi^-+\nu, \tag{12}$$

$$\theta^0\rightarrow\mu^++\pi^-+\nu. \tag{12'}$$

The charge conjugate modes are allowed to the $\bar{\theta}^0$. It follows that $a_1=a_2$ for the electron decay mode (and similarly for the muon mode). But this is just the condition, Eq. (10), which leads to the prediction of a striking growth in time, during the early stages of decay, for the e^- (and μ^-) mode. This conclusion would also hold if the τ^0 particle contributes to $K_{\mu3}{}^0$ and $K_{e3}{}^0$ processes, since the τ^0 presumably is to be assigned the same strangeness as the θ^0.

III. THE θ_2 LIFETIME

It is of some interest to consider here what can be said, on the basis of present experimental information, about θ^0 decay via anomalous modes. We denote as "anomalous" all neutral K-meson decay processes other than two-pion decay, so that all θ_2 disintegration modes are anomalous. In view of the present scarcity of quantitative experimental data on such decay processes, our remarks can obviously have only a rough qualitative significance.

From Eqs. (8) and (9) we note that if one observes $(e\pi\nu)$ decay events without regard to the signs of the charges, the interference terms cancel and one has

$$R(e\pi\nu)\equiv R(e^+\pi^-\nu)+R(e^-\pi^+\bar{\nu})$$
$$=a_1^2\exp(-\lambda_1 t)+a_2^2\exp(-\lambda_2 t).$$

Including similar contributions from the other decay modes, we have for the total rate of anomalous events

$$R(\text{anom})=\Lambda_1\exp(-\lambda_1 t)+\lambda_2\exp(-\lambda_2 t), \tag{13}$$

where Λ_1 is the partial rate for θ_1 decay via anomalous modes ($\Lambda_1<\lambda_1$). The total rate of θ_1 decay is fairly well known experimentally[12]: $1/\lambda_1\approx1.8\times10^{-10}$ sec.

In typical cloud chamber experiments one can ob-

[11] R. G. Sachs, Phys. Rev. **99**, 1573 (1955).
[12] Snyder, Chang, and Gupta, Phys. Rev. **100**, 1264(A) (1955).

serve decay processes over a time interval, T, (measured in the K-meson rest frame) which on the average may be of the order 5×10^{-10} sec, or perhaps somewhat longer. Assuming that the θ^0 particles are produced very close to the visible region of the cloud chamber so that they come under observation at $t\approx0$ (this is not always the case experimentally and obvious corrections must be made), we have for the fraction of all decay events which are anomalous

$$f=\frac{(\Lambda_1/\lambda_1)[1-\exp(-\lambda_1T)]+[1-\exp(-\lambda_2T)]}{[1-\exp(-\lambda_1T)]+[1-\exp(-\lambda_2T)]}. \quad (14)$$

Experimentally, the frequency of anomalous events having charged secondaries is about 0.1 relative to the total number of neutral K-meson decay events with charged secondaries. But if the θ^0 has even spin, θ_1 may sometimes decay into two neutral pions and thus be missed. On this account, and also because some of the anomalous events may come from τ^0 decay, we have that $f\lesssim0.1$. (The possible decay mode $\theta_2\rightarrow\pi^0+\pi^0+\gamma$ works in the opposite direction as regards the above inequality, but we shall assume this effect can be neglected.)

With our estimate $T\approx5\times10^{-10}$ sec, we have $\lambda_1T\approx3$; and since $f\lesssim0.1$ we see that $\lambda_2T\ll1$. Equation (14) can therefore be written in the approximate form

$$\Lambda_1/\lambda_1+\lambda_2T\lesssim0.1. \quad (15)$$

This leads to the inequality

$$\lambda_2T\lesssim0.1,\quad 1/\lambda_2\gtrsim5\times10^{-9}\text{ sec}, \quad (16)$$

a result which is not appreciably changed even if Λ_1 is as large as λ_2. There is of course no way to set an upper limit on the θ_2 lifetime on the basis of the present analysis, since we have no idea what fraction of anomalous events are contributed by τ^0 rather than θ^0 decay. Another result may be obtained from Eq. (15); namely, the branching ratio for θ_1 decay via anomalous modes satisfies the inequality

$$\Lambda_1/\lambda_1\lesssim0.1.$$

The work reported here was undertaken after a stimulating discussion with R. Dalitz, and it has also been influenced by a number of interesting comments by C. N. Yang.

APPENDIX I

In order to establish that the coefficients a_1 and a_2 appearing in Eqs. (5) may be chosen to be real, we make use of the behavior of the θ^0 field under Wigner-type time reversal.[13] For the sake of simplicity we consider θ mesons with zero linear momentum. Denote by θ^m the state of the θ meson corresponding to spin magnetic quantum number m. Then the phases may be chosen in such a manner that under time reversal K

$$K\theta^m=i^{2m}\theta^{-m}. \quad \text{(A–1)}$$

Since the operation of time reversal includes complex conjugation we have, from Eqs. (1) and (A–1),

$$K\theta_1{}^m=i^{2m}\theta_1{}^{-m}, \quad \text{(A–2)}$$

$$K\theta_2{}^m=-i^{2m}\theta_2{}^{-m}. \quad \text{(A–3)}$$

We now apply K to both sides of Eqs. (5) and (5′). Since final state interactions are expected to be very small, we are free to choose the phases such that

$$K(e^+\pi^-\nu)^m=i^{2m}(e^+\pi^-\nu)^{-m}$$

(and similarly for the other states). Then we find

$$i^{2m}\theta_1{}^{-m}\rightarrow i^{2m}a_1{}^*(e^+\pi^-\nu+e^-\pi^+\bar{\nu})^{-m}+\cdots, \quad \text{(A–4)}$$

$$-i^{2m}\theta_2{}^{-m}\rightarrow i^{2m+1}a_2{}^*(e^+\pi^-\nu-e^-\pi^+\bar{\nu})^{-m}+\cdots. \quad \text{(A–5)}$$

But according to Eqs. (5)

$$\theta_1{}^{-m}\rightarrow a_1(e^+\pi^-\nu+e^-\pi^+\bar{\nu})^{-m}+\cdots, \quad \text{(A–6)}$$

$$\theta_2{}^{-m}\rightarrow -ia_2(e^+\pi^-\nu-e^-\pi^+\bar{\nu})^{-m}+\cdots, \quad \text{(A–7)}$$

since the amplitudes are independent of m. Comparison of Eq. (A–6) with Eq. (A–4) and of Eq. (A–7) with Eq. (A–5) leads to the conclusion

$$a_1{}^*=a_1,\quad a_2{}^*=a_2, \quad \text{(A–8)}$$

i.e., that a_1 and a_2 are real.

APPENDIX II

To obtain Eq. (8), we consider the approximate solution of the time-dependent Schrödinger equation for the state vector which at $t=0$ is just θ^0. In the usual way, we expand the state vector in terms of the nearly stationary states θ_1, θ_2, $(e^+\pi^-\nu)$, $(e^-\pi^+\bar{\nu})$, $(\pi^+\pi^-)$, etc. If the amplitudes of these states are denoted by $A(s)$, their time dependence is determined by

$$i\hbar\dot{A}(s)=\sum_{s'}A(s')\exp(i\omega_{ss'}t)V_{ss'}, \quad \text{(A–9)}$$

where $V_{ss'}$ is the matrix element of the interaction between states s and s' and $\hbar\omega_{ss'}=E_s-E_{s'}$. The initial conditions under which these equations are to be solved are : at $t=0$, $A(\theta_1)=A(\theta_2)=1/\sqrt{2}$, all other $A(s)=0$. Following Weisskopf and Wigner,[14] we anticipate the time dependence

$$A(\theta_i)=2^{-\frac{1}{2}}\exp(-\gamma_it/2), \quad \text{(A–10)}$$

[13] E. P. Wigner, Nachr. Akad. Wiss. Göttingen, Math. physik Kl. (1932), p. 546.

[14] V. Weisskopf and E. P. Wigner, Z. Physik **63**, 54 (1930); **65**, 18 (1930).

and solve Eq. (A–9) under the assumption that all other amplitudes $A(s)$ may be dropped from the right-hand side. We then find for states other than θ_1 and θ_2

$$A(s)=\frac{2^{-\frac{1}{2}}}{i\hbar}\left\{V_{s1}\frac{\exp[(i\omega_{s1}-\frac{1}{2}\gamma_1)t]-1}{i\omega_{s1}-\frac{1}{2}\gamma_1}+V_{s2}\frac{\exp[(i\omega_{s2}-\frac{1}{2}\gamma_2)t]-1}{i\omega_{s2}-\frac{1}{2}\gamma_2}\right\}. \quad \text{(A–11)}$$

The probability $P(s)$ for finding the system in state s is obtained by integrating $|A(s)|^2$ over all values of the energy E_s. Under the usual assumption that V_{si} is a well-behaved analytic function of ω_{si} in the complex plane we obtain

$$P(s)=\frac{a_{s1}^2}{2\lambda_1}[1-\exp(-\lambda_1 t)]+\frac{a_{s2}^2}{2\lambda_2}[1-\exp(-\lambda_2 t)]+a_{s1}a_{s2}\,\mathrm{Re}\frac{1-\exp[-(\gamma_1+\gamma_2^*)t/2]}{\gamma_1+\gamma_2^*}, \quad \text{(A–12)}$$

where λ_i is the real part of γ_i and a_{si} is proportional to V_{si}.

Taking the time derivative of $P(s)$, we find for the rate of decay into modes s the expression given by Eq. (8), where $\Delta\omega$ is the imaginary part of $(\gamma_1+\gamma_2^*)/2$ and may be interpreted as a mass difference between the θ_1 and θ_2 particles.

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Neutral Pion Production in Deuterium and Hydrogen: Ratios*

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Neutral photopion production in deuterium and hydrogen has been studied by counting the decay photons singly. Pure gas targets were used. The photon counter consisted of three counters, the first in anticoincidence and the second and third in coincidence. The photon counting rate was obtained by taking the difference between the counting rates with a Pb converter alternately behind and in front of the anticoincidence counter. Data were obtained for 300-, 400-, and 500-Mev bremsstrahlung at laboratory angles of 30, 73, and 140°. An attempt to apply the photon-difference method was only moderately successful owing to poor control of the bremsstrahlung end point and the unfavorable shape of the neutral pion excitation curve. At 73° and 140° the data are consistent with a constant value of 0.90 per nucleon for the deuterium-to-hydrogen neutral photopion production ratio. At 30° and 300 Mev there is some evidence for a larger ratio, which points to constructive interference between production on the proton and neutron in the deuteron. The data are consistent with the predictions of the "symmetrical theory" concerning the coupling of neutral pions to nucleons and with Watson's phenomenological theory of photopion production.

I. INTRODUCTION

THE four basic reactions involving photopion production on nucleons are

$$\gamma+p\rightarrow n+\pi^+, \quad \text{(a)}$$

$$\gamma+p\rightarrow p+\pi^0\rightarrow p+2\gamma, \quad \text{(b)}$$

$$\gamma+n\rightarrow n+\pi^0\rightarrow n+2\gamma, \quad \text{(c)}$$

$$\gamma+n\rightarrow p+\pi^-. \quad \text{(d)}$$

Of these, (a) and (b) are directly accessible for experimental study and considerable information has been accumulated concerning their cross sections.[1,2] The cross sections for (c) and (d), however, can only be inferred from studies made on complex nuclei, the most suitable of which is the deuteron. Because of its loose binding, the neutron in the deuteron can be considered "almost free" for high-energy interactions. In addition, one hopes that the simple structure of the deuteron may make it possible to correct for the effects of binding and thus deduce fairly precise information about the elementary interactions with the neutron. Quite precise studies of the π^-/π^+ ratio in deuterium[3] have also been made which give information about reaction (d). This information has been reasonably well interpreted in conjunction with that on the first two processes by the phenomenological theory of Watson.[4] The theory then leads to the unique prediction that the cross sections for photopion production on neutron and proton should be

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[1] Experiments on positive photopion production in hydrogen are reviewed in papers on this subject by Walker, Teasdale, Peterson, and Vette, Phys. Rev. **99**, 210 (1955) and Tollestrup, Keck, and Worlock, Phys. Rev. **99**, 220 (1950).

[2] Experiments on neutral photopion production are reviewed in papers on the subject by Goldschmidt-Clermont, Osborne, and M. Scott, Phys. Rev. **97**, 188 (1955) and by D. C. Oakley and R. L. Walker, Phys. Rev. **97**, 1283 (1955).

[3] Sands, Teasdale, and Walker, Phys. Rev. **95**, 592 (1954).

[4] Watson, Keck, Tollestrup, and Walker, Phys. Rev. **101**, 1159 (1956).