position, height, and width of the secondary peak, in addition to the expected agreement with the forward peak. In the face of such a simple explanation, it seems unwarranted to search for a more complex mechanism. In particular, the forward peak is not expected to be nearly so informative about possible resonant contributions as the backward scattering. The diffraction contribution to $f(\theta)$ in the forward direction has a weight $(L+1)^2$, where L is the highest partial wave which is appreciably absorbed. Therefore, it will be comparatively difficult for a dynamical contribution, whose weight will be $\lesssim(2l+1)$, to significantly alter the shape of the differential cross section at very forward angles. Outside the forward peak such alteration is to be expected, but it can only be interpreted reliably if the diffraction scattering is understood in detail.

It is to be expected that the secondary maxima will change in character (that is, strength and position) at other energies, since both η_l and L are functions of energy. It is obviously of interest to know whether such maxima exist at other energies, but the data appear to be inconclusive. It is clear from the results of Perl, Jones, and Ting[3] that the rapid decrease of $d\sigma/d\Omega$ with increasing θ does not persist at angles outside the extreme forward peak. Whether there is a second maximum or merely a leveling off is difficult to say. The data of Cook et al.[6] seem to show some secondary bumps, but the statistical errors are rather large. The results of Helland[7] for π^+p scattering at 1555 MeV show a definite bump near $\cos\theta = 0.3$, which would be consistent with the present model. The statement that a secondary maximum is known not to exist at other energies would appear to be unjustified, at present. In any case, the attempt to associate the secondary maximum observed by Damouth, Jones, and Perl with the 2.08-GeV/c π^-p resonance entails considerable difficulty, since this second maximum appears to occur in both π^-p and π^+p scattering.

It is a pleasure to acknowledge the valuable contributions of Professor Peter Carruthers. The author is also indebted to Professor M. L. Perl for the communication of some of his results prior to publication.

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*Work supported in part by the U. S. Office of Naval Research.

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NUCLEON ELECTROMAGNETIC FORM FACTORS AT HIGH MOMENTUM TRANSFER*

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(Received 15 January 1964)

The recently reported experimental results of Chen et al.[1] on the electromagnetic form factors of the proton at large values of the invariant momentum transfer, q^2, are consistent with the relationship

$$\lim_{q^2\to\infty} G_E^{\,p}(q^2)/G_M^{\,p}(q^2) = 1 \tag{1}$$

suggested[2] earlier on the basis of physical arguments concerning the structure of the nucleon. If further experimental work confirms this result, we may draw the conclusion that the nucleon is indeed a fundamental particle rather than a composite generated by some sort of bootstrap mechanism as suggested, for example, by Chew and Frautschi.[3]

On the other hand, the claim by Chen et al.[1] that their data are consistent with

$$\lim_{q^2\to\infty} G_E^{\,p}(q^2) = \lim_{q^2\to\infty} G_M^{\,p}(q^2) = 0, \tag{2}$$

namely, that there is no "core contribution" to the proton form factors, might be interpreted in just the opposite way. It is the purpose of this note to show that there need be no contradiction between the two results and that, in fact, the condition expressed in Eq. (1) probably provides the more reliable test of the fundamental character of the nucleon.

Equation (1) is obtained from reference 2 by combining the two conditions given there:

$$\lim_{q^2\to\infty} G_E(q^2) = Z_2^{(s)}Q \tag{3a}$$

and

$$\lim_{q^2\to\infty} G_M(q^2) = Z_2^{(s)}Q, \tag{3b}$$

where $Z_2^{(s)} = Z_1^{(s)}$ are, respectively, the wave-function and vertex renormalization constants due to strong interactions and Q is the physical charge of the nucleon, $Q=1$ for the proton, $Q=0$ for the neutron. The physical content of Eqs. (3) is simply that at high momentum transfer one sees only the bare nucleon which has a total charge proportional to the probability $Z_2^{(s)}$ [Eq. (3a)] and the simple Dirac moment associated with that charge [Eq. (3b)]. If, in fact, there is no core, as indicated by Eq. (2), then $Z_2^{(s)}=0$ and the ratio of Eq. (3a) to Eq. (3b) is not well defined. Hence the argument for Eq. (1) would collapse.

One way out of this paradox is to assume that $Z_2^{(s)}$ is very small but not zero. In their present stage the experiments certainly cannot rule this out. Furthermore, there is, from the theoretical side, an important related question of the order in which limits are to be taken in writing equations like Eq. (3). As pointed out in reference 2, for the qualitative argument to have meaning, it must be applied to a model in which the strong interactions are subject to a cutoff of some sort. If the cutoff is Λ^2, then the condition $q^2\to\infty$ means simply $q^2 \gg \Lambda^2$. In fact, q^2 must not be so large that electromagnetic radiative corrections become important, since then it would be necessary to include higher order electromagnetic effects, which cannot be described in terms of the two form factors.

In such a model $Z_2^{(s)}$ is a function of Λ^2. Then Eqs. (3a) and (3b) might be replaced by

$$G_E(q^2) = G_M(q^2) = Z_2^{(s)}(\Lambda^2)Q \text{ for } q^2 \gg \Lambda^2. \tag{4}$$

At the same time we may have

$$\lim_{\Lambda^2\to\infty} Z_2(\Lambda^2) = 0, \tag{5}$$

in agreement with the notion that for finite but large Λ^2, $Z_2^{(s)}$ becomes very small.

Further verification of Eq. (1) would seem to argue that this is the correct interpretation, that a primitive Dirac particle with a Dirac moment exists and does underlie the nucleon structure, although the weight that must be given to this "core" might be vanishingly small. Then it would appear to be proper to consider the nucleon to be the manifestation of a fundamental particle.

It is of some interest to note that the equivalent of Eq. (1) cannot be obtained by the same argument for the neutron form factors since then $Q=0$ in Eq. (3). The consequences of Eqs. (3) may, in this case, be brought out more clearly in terms of the Dirac form factor $F_1(q^2)$ and the Pauli form factor $F_2(q^2)$ which are related to G_E and G_M by[4]

$$G_E(q^2) = F_1(q^2) - (q^2/2M)F_2(q^2) \tag{6a}$$

and

$$G_M(q^2) = F_1(q^2) + 2MF_2(q^2), \tag{6b}$$

where M is the physical mass of the nucleon. From Eqs. (3) it follows that

$$\lim_{q^2\to\infty} F_1(q^2) = Z_2^{(s)}Q \tag{7a}$$

and

$$\lim_{q^2\to\infty} q^2F_2(q^2) = 0 \tag{7b}$$

as already noted in reference 2. Therefore $q^2F_2\to 0$ for both the neutron and proton,[5] while it is only for the neutron that $F_1\to 0$ irrespective of the value of $Z_2^{(s)}$. The severe test of these ideas for the neutron lies in the verification of Eq. (7b), which is already established for the proton to the extent that Eq. (1) is established and that G_E does not become very large.

In conclusion, it should be remarked that Eq. (3) may be derived from a generalization of Källén's asymptotic condition[6] for the high-energy limit of the matrix element. He suggests that the matrix element should approach a value equal to the value of the matrix element in lowest order perturbation theory (Born approximation) multiplied by the vertex renormalization constant $Z_1 = Z_2$. In our case we must apply the same condition for $q^2 \gg \Lambda^2$ but with q^2 still satisfying the condition mentioned earlier that it be small enough for electromagnetic radiative corrections to be negligible. Then Z_i is replaced by $Z_i^{(s)}$ and the condition for the matrix element of the current density between nucleon states of four-momentum p and p' reads

$$\langle p'|j_\mu(0)|p\rangle \to Z_1^{(s)} Q\bar{u}(p')\gamma_\mu u(p), \qquad (8)$$

where $u(p)$ is the usual spinor. Since Eq. (8) is a four-vector condition it clearly contains two statements about the form factors, one for the spacelike part of the vector and another for the timelike part. This can be seen most easily in the Breit frame, $\vec{p}'=\vec{p}$. There, Eq. (8) reads

$$\langle p'|j_4(0)|p\rangle \to iZ_1^{(s)}Q, \qquad (9a)$$

$$\langle p'|\vec{j}(0)|p\rangle \to i(\vec{\sigma}\times\vec{q})Z_1^{(s)}Q/2M. \qquad (9b)$$

On the other hand, the general expression for the matrix element in terms of the form factors in the Breit frame is[4]

$$\langle p'|j_4(0)|p\rangle = iG_E(q^2), \qquad (10a)$$

$$\langle p'|\vec{j}(0)|p\rangle = i(\vec{\sigma}\times\vec{q})G_M(q^2)/2M. \qquad (10b)$$

Comparison of Eqs. (9) and (10) leads directly to Eqs. (3).

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

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[6]G. Källén explicitly stated this result for quantum electrodynamics, in Handbuch der Physik (Springer-Verlag, Berlin, 1958), Vol. V, Pt. 1, pp. 362-363. He also generalized the suggestion to other field theories in unpublished lectures. I am indebted to Dr. W. Theis for calling to my attention the connection with Källén's results that is discussed here.

MUON POLARIZATION IN $K_{\mu3}^+$ DECAY

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(Received 5 February 1964)

Investigation of $K_{\mu3}$ decay permits verification of the applicability of the universal V-A interaction to lepton decays of strange particles and also determination of the characteristics of the interaction form factors. For a two-component neutrino, the μ mesons generated in a $\pi\to\mu+\nu$ or $K_{\mu2}\to\mu+\nu$ decay are known to be completely polarized longitudinally, whereas in a three-particle decay the μ-meson polarization depends on space correlation of the decay particles, i.e., on the mode of interaction realized in the decay.

As shown by Pais and Treiman,[1] the matrix element for a K_{l3} decay depends on two scalar form factors and can be written as follows (see, e.g., Brene, Egardt, and Qvist[2]):

$$M=[\tfrac{1}{2}f_+(P_K+P_\pi) + \tfrac{1}{2}f_-(P_K-P_\pi)]\bar{\mu}(P_\nu)\gamma_\lambda(1+\gamma_5)\nu(P_l), \qquad (1)$$

where f_+ and f_- are form factors due to strong interactions, and P is the particle momentum.

According to the theory of the universal V-A interaction, the matrix element describes both the K_{e3} and $K_{\mu3}$ decays. Assuming the form factors to be constant, Gatto,[3] Zachariasen,[4] and Fujii and Kawaguchi[5] estimated the ratio $\xi=f_-/f_+$. The decay-branching ratio, known from experiment,[6] can likewise be determined from the expression (1). It is found to be a quadratic function of the parameter ξ:

$$W(K_{\mu3})/W(K_{e3}) = 0.651+0.126\xi+0.0189\xi^2.$$

A solution of this equation gives two values of ξ, to which correspond different shapes of the μ-meson spectrum in $K_{\mu3}$ decay. This difference, however, is large only in the low-energy region, where measurements are difficult to carry out. It would therefore be more convenient to measure