will give $[h'(a)/h(a)]$ (and hence the phase shift) correct to a higher order than the v and h. If v and h depend on free parameters, these may be found in the usual way by using the stationary character of $[h'(a)/h(a)]$.

If the Born approximation conditions are valid for $r>a$, we may take $v\cong h$. In that case Eq. (A14) becomes

$$[h'(a)/h(a)]\cong[h(a)]^{-2}\int_a^\infty h^2 W dr+(A-k^2B). \quad \text{(A15)}$$

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Σ-Λ Relative Parity and the $\Sigma^0\rightarrow\Lambda^0+\gamma+\gamma$ Decay

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It is shown that the π^0-pole term predicts a large difference by nearly two to three orders of magnitude for the branching ratio of the $\Sigma^0\rightarrow\Lambda^0+2\gamma$ decay mode, depending upon the value of the $\Sigma-\Lambda$ relative parity. It is further argued that this difference is not masked, even if we include other diagrams. It is thus suggested that a study of the branching ratio of the $\Sigma^0\rightarrow\Lambda^0+2\gamma$ decay may serve to determine the $\Sigma-\Lambda$ relative parity.

THE determination of $\Sigma-\Lambda$ relative parity is at present an urgent problem for the development of the theory of elementary particles. Suggestions have been made[1] to determine this parity unambiguously from correlation effects in the Dalitz decay, $\Sigma^0\rightarrow\Lambda^0+e^++e^-$. In this letter we wish to point out that a study of the branching ratio of the $\Sigma^0\rightarrow\Lambda^0+\gamma+\gamma$ decay mode could also serve this purpose.

We will denote even and odd $\Sigma-\Lambda$ parities by $P=\pm1$, respectively. The corresponding parameters will often be denoted by the superscripts $\pm$. First of all, we notice that under the assumption of charge-independent strong interactions and minimal electromagnetic interactions the decay $\Sigma^0\rightarrow\Lambda^0+2\gamma$ is strictly forbidden for either parity, as long as we switch on only those interactions which involve the particles with integral isotopic spin (Σ, Λ, and π) and the photon. This is because the above-mentioned class of interactions is invariant under the isotopic rotation $e^{i\pi T_2}$ (under which $\Sigma^0\rightarrow-\Sigma^0$, $\Lambda^0\rightarrow\Lambda^0$) together with $\gamma\rightarrow-\gamma$ (γ from the isovector current). The decay takes place only when[2] the strong and electromagnetic (minimal) interactions involving the particles with half-integral isotopic spin (N, Ξ, and K) are switched on. Thus if we neglect the K-meson cloud, the decay has to occur only through baryon (half-integral isotopic spin) loops, the main contribution of which may be expected to be given by the π^0-pole term, shown in Fig. 1.

Let us, therefore, first study the contribution of Fig. 1, hoping that it dominates. The possible importance of other intermediate states will be discussed at the end. The $\pi^0\rightarrow2\gamma$ vertex in Fig. 1, involving the baryon (half-integral isotopic spin) loops, can be estimated from the observed rate of π^0 decay. We denote the $\pi^0\rightarrow2\gamma$ matrix element by

$$(2\pi)^4\delta^4(p_\pi-k_1-k_2)F_\pi\epsilon_{\alpha\beta\gamma\delta}e_{1\alpha}e_{2\beta}k_{1\gamma}k_{2\delta},$$

where four-vector $e_{1,2}$ denotes the polarizations of the photons with four-momenta $k_{1,2}$, respectively. We may safely assume that the form factor F_π is nearly a constant. Then in the rest frame of Σ^0, the rate of

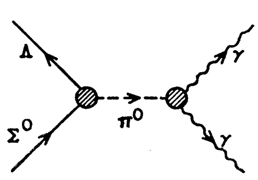

Fig. 1. The π^0 pole diagram for $\Sigma^0\rightarrow\Lambda^0+2\gamma$ decay.

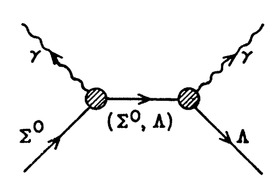

Fig. 2. The (Σ^0,Λ)-pole diagram for $\Sigma^0\rightarrow\Lambda^0+2\gamma$ decay.

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[1] J. Sucher and G. A. Snow, Nuovo cimento 18, 195 (1960); N. Byers and H. Burkhardt, Phys. Rev. 121, 281 (1961); L. Michel and H. Rouhaninejad, Phys. Rev. 122, 242 (1961); and S. Chiba (to be published).

[2] In this case one photon comes from the isoscalar and the other from the isovector part of the current.

$\Sigma^0 \rightarrow \Lambda^0+2\gamma$ is given by

$$W(\Sigma^0\rightarrow\Lambda^0+2\gamma)=W(\pi^0\rightarrow 2\gamma)[(g^{\pm}{}_{\Sigma\Lambda\pi})^2/4\pi]m_\pi^{-3}\int_{m_\Lambda}^{E_\Lambda{}^{\max}} dE_\Lambda \frac{(E_\Lambda{}^2-m_\Lambda{}^2)^{\frac{1}{2}}(E_\Lambda\mp m_\Lambda)(m_\Sigma{}^2+m_\Lambda{}^2-2m_\Sigma E_\Lambda)^2}{[m_\Sigma{}^2+m_\Lambda{}^2-m_\pi{}^2-2m_\Sigma E_\Lambda]^2}, \quad (1)$$

where $W(\pi^0\rightarrow 2\gamma)=|F_\pi|^2 m_\pi{}^3/64\pi$ is the rate of $\pi^0\rightarrow 2\gamma$ decay and $E_\Lambda{}^{\max}=(m_\Sigma{}^2+m_\Lambda{}^2)/2m_\Sigma$. $g^{\pm}{}_{\Sigma\Lambda\pi}$ denotes[3] the renormalized pseudoscalar ($P=+1$) and scalar ($P=-1$) coupling constants for $\Sigma\leftrightarrows\Lambda+\pi$, respectively. The upper and lower signs in Eq. (1) correspond, respectively, to $P=\pm1$. If we use[4]

$$W(\pi^0\rightarrow 2\gamma)\approx 0.5\times 10^{16}\ \text{sec}^{-1},$$

we obtain

$$W(\Sigma^0\rightarrow\Lambda^0+2\gamma) \approx\frac{(g^{\pm}{}_{\Sigma\Lambda\pi})^2}{4\pi}\times\begin{cases}9\times10^{9}\ \text{sec}^{-1} & \text{for } P=+1\\ 2.4\times10^{13}\ \text{sec}^{-1} & \text{for } P=-1.\end{cases} \quad (2)$$

Thus a large difference between the branching ratios of the $\Sigma^0\rightarrow\Lambda^0+2\gamma$ mode for the two parity cases is expected, insofar as Fig. 1 is the dominant mechanism.

Let us now consider the normal decay mode $\Sigma^0\rightarrow\Lambda^0+\gamma$. In the absence of an experimental[5] value for the absolute rate of this mode, we have to resort to theoretical prediction for this rate at present. Denoting the effective interaction for $\Sigma^0\rightarrow\Lambda^0+\gamma$ by

$$\tfrac{1}{2}\mu^{\pm}{}_{\Sigma\Lambda}\bar{\Lambda}\sigma_{\mu\nu}\begin{pmatrix}1\\ i\gamma_5\end{pmatrix}\Sigma^0 F_{\mu\nu}, \quad (3)$$

where

$$\mu^{\pm}{}_{\Sigma\Lambda}\equiv(e/2m_p)X^{\pm}{}_{\Sigma\Lambda},\quad \sigma_{\mu\nu}\equiv(1/2i)(\gamma_\mu\gamma_\nu-\gamma_\nu\gamma_\mu),$$

and $F_{\mu\nu}=\partial_\mu A_\nu-\partial_\nu A_\mu$, the rate of $\Sigma^0\rightarrow\Lambda^0+\gamma$ for either parity case is given by

$$W(\Sigma^0\rightarrow\Lambda^0+\gamma)=\frac{e^2}{4\pi}\frac{(X^{\pm}{}_{\Sigma\Lambda})^2\omega^3}{m_p{}^2} \approx (X^{\pm}{}_{\Sigma\Lambda})^2\times 5.53\times10^{18}\ \text{sec}^{-1}, \quad (4)$$

where ω denotes the energy of the photon in the rest frame of Σ^0 ($\omega\approx 76$ Mev). By (2) and (4), the branching ratio for the $\Sigma^0\rightarrow\Lambda^0+2\gamma$ mode is given by

$$R\equiv\frac{W(\Sigma^0\rightarrow\Lambda^0+2\gamma)}{W(\Sigma^0\rightarrow\Lambda^0+\gamma)} \simeq\frac{(g^{\pm}{}_{\Sigma\Lambda\pi})^2/4\pi}{(X^{\pm}{}_{\Sigma\Lambda})^2}\times\begin{cases}1.6\times10^{-9} & \text{for } P=+1\\ 4.3\times10^{-6} & \text{for } P=-1.\end{cases} \quad (5)$$

For the even-parity case, hyperfragment analysis[6] indicates that $(g^{+}{}_{\Sigma\Lambda\pi})^2/4\pi\simeq 15$, consistent with global symmetry,[7] while for the odd-parity case, the same analysis,[8] as well as the recently proposed Nambu-Sakurai,[9] method, indicates that $(g^{-}{}_{\Sigma\Lambda\pi})^2/4\pi\simeq0.5$. We therefore shall use[10] these values in Eq. (5). The values of $X_{\Sigma\Lambda}{}^{\pm}$ are also unknown and depend on the structure of strong interactions. From the results of various symmetry hypotheses[11] and other calculations,[12,13] however, it seems plausible to take $X^{+}{}_{\Sigma\Lambda}\simeq 1$–2. For the odd-parity case, there is no symmetry principle to guide us. At any rate, rough perturbation calculations[12] and the recent calculation by Dreitlein and Lee,[13] based on the hypothesis of the dominance of the 2π resonance in the ($T=1$, $J=1$) state, indicate that $|X^{-}{}_{\Sigma\Lambda}|\approx 1$–0.3. Using these values, we obtain from Eq. (5)

$$R\approx\begin{cases}(0.6\text{–}2.4)\times10^{-8} & \text{for } P=+1\\ (0.2\text{–}2.4)\times10^{-5} & \text{for } P=-1.\end{cases} \quad (6)$$

Apart from the absolute decay rates, the energy

[3] We explicitly assume the direct Yukawa-type interaction without derivative.

[4] H. Ruderman, S. Berman, R. Gomez, A. V. Tollestrup, and R. Talman, Bull. Am. Phys. Soc. **5**, 508 (1960); R. F. Blackie, A. Engler, and J. H. Hulvey, Phys. Rev. Letters **5**, 384 (1960); R. G. Glasser, N. Seeman, and B. Stiller, Bull. Am. Phys. Soc. **6**, 39 (1961).

[5] The value quoted by W. H. Barkas and A. H. Rosenfeld [*Proceedings of the 1960 Annual International Conference on High-Energy Physics at Rochester* (Interscience Publishers, Inc., New York, 1960)] for the $\Sigma^0\rightarrow\Lambda^0+\gamma$ rate is $W(\Sigma^0\rightarrow\Lambda^0+\gamma)\gtrsim10^{11}$ sec^{-1}. Recently methods have been suggested by J. Dreitlein and H. Primakoff (to be published) to determine the Σ^0 radiative transition rate.

[6] D. B. Lichtenberg and Marc Ross, Phys. Rev. **107**, 1714 (1957); F. Ferrari and L. Fonda, Nuovo cimento **9**, 842 (1958).

[7] M. Gell-Mann, Phys. Rev. **106**, 1296 (1957).

[8] F. Ferrari and L. Fonda, Nuovo cimento **9**, 842 (1958).

[9] Y. Nambu and J. J. Sakurai, Phys. Rev. Letters **6**, 377 (1961).

[10] Since we are interested in establishing the large value of the branching ratio R for the odd-parity case as compared to that for the even-parity case, we feel that we are only taking maximum odds against our argument by such choice of coupling constants. Any mistake in this choice is, therefore, expected only to strengthen our argument.

[11] For instance, the hypothesis of global symmetry ($P=+1$) yields $X^{+}{}_{\Sigma\Lambda}=-X_{\text{neutron}}\approx1.9$, while that of unitary symmetry [M. Gell-Mann, California Institute of Technology Synchrotron Report No. CTSL-20, 1961 (unpublished)] gives $X^{+}{}_{\Sigma\Lambda}\approx(\sqrt{3}/2)\times1.9$ [S. Coleman and S. L. Glashow, Phys. Rev. Letters **6**, 423 (1961)]. In this connection see also G. Feinberg (to be published).

[12] A perturbation calculation [H. Katsumori, Prog. Theor. Phys. (Kyoto) **24**, 1371 (1960)], taking the contribution of pionic current only into account, gives $X^{+}{}_{\Sigma\Lambda}\approx1.8$ for $g^2{}_{\Sigma\Sigma\pi}/4\pi\simeq(g^{+}{}_{\Sigma\Lambda\pi})^2/4\pi\simeq15$; while for the odd-parity case it gives $X^{-}{}_{\Sigma\Lambda}\approx-1.0$ to -0.45 for $g^2{}_{\Sigma\Sigma\pi}/4\pi\simeq15$–3 and $(g^{-}{}_{\Sigma\Lambda\pi})^2/4\pi\simeq0.5$.

[13] J. Dreitlein and B. W. Lee (to be published). On the hypothesis of the dominance of the two-pion resonance in the ($T=1$, $J=1$) state, these authors obtain $W(\Sigma^0\rightarrow\Lambda^2+\gamma)\approx(g_{\Sigma\Sigma\pi}g_{\Sigma\Lambda\pi}/g^2{}_{NN\pi})^2\times9.1\times10^{18}$ sec^{-1} for either parity. So, if we choose $g_{\Sigma\Sigma\pi}=g^{+}{}_{\Sigma\Lambda\pi}=g_{NN\pi}$ for the even-parity case, we get [cf. Eq. (4)] $|X^{+}{}_{\Sigma\Lambda}|\approx1.3$, while if we choose $(g_{\Sigma\Sigma\pi}g^{-}{}_{\Sigma\Lambda\pi}/g^2{}_{NN\pi})^2\approx1/15$ for the odd-parity case, we get $|X^{-}{}_{\Sigma\Lambda}|\approx0.33$.

spectra of the Λ hyperon and the photons are also expected to be different for the even- and odd-parity cases. These spectra, calculated from Fig. 1, are shown in Figs. 3(a) and (b).

We have shown [Eq. (6)] that the π^0-pole term predicts a large difference by nearly two to three orders of magnitude for the branching ratio R depending upon the value of P. We have, however, to study the effects of other diagrams, especially for the even-parity case, since the π^0-pole term leads to a very low branching ratio $\sim 10^{-8}$ due to the P-wave pion emission in this case. The first set of candidates, in order of increasing mass, after the one-pion-pole term, would appear to be terms involving a higher number of pions. Fortunately the 2π, 4π, $\cdots$ configurations cannot contribute to the process by the G $(=e^{i\pi T_2}\times$charge conjugation) invariance of strong interactions, and we may reasonably neglect the contributions from 3π, $5\pi\cdots$ intermediate states. We also neglect the intermediate states involving nucleon (cascade) and $\bar{K}$ (K) meson, partly because of the higher mass involved and partly because the K-meson coupling constant seems to be smaller than that of the π meson. Rough perturbation-theoretic calculation justifies this omission.

We explicitly calculate the contributions from Σ^0- and Λ-pole terms (Fig. 2), which give roughly the same contribution for even- and odd-parity cases. This contribution is strongly dependent upon the magnitudes of the transition moments[14] $X_{\Sigma^0\Sigma^0}$ and $X_{\Lambda\Lambda}$, as well as upon $X_{\Sigma\Lambda}$. However, $X_{\Sigma^0\Sigma^0}$ and $X_{\Lambda\Lambda}$ receive contributions only from the isoscalar part.[15] If the isoscalar part is small for (Σ,Λ) hyperons as in the case of nucleons ($\mu_S \approx -0.06\ e/2m_p$ for nucleons), it is plausible[16] that

$$|X_{\Sigma^0\Sigma^0}| \quad \text{and} \quad |X_{\Lambda\Lambda}| \lesssim 0.2.$$

[14] We define $X_{\Sigma^0\Sigma^0}$ and $X_{\Lambda\Lambda}$ in the same way as $X_{\Sigma\Lambda}$ [see Eq. (3)], i.e., $\mu_{ij} \equiv (e/2m_p)X_{ij}$.

[15] This is easy to see by writing $\mu = \mu_s + T_3\mu_v$ in obvious notation.

[16] This is true under the hypotheses of many symmetry schemes. For instance in the limit of strict global symmetry $(P=+1)$, or symmetries somewhat weaker than global symmetry, the magnetic moments of Σ^0 and Λ vanish. In this connection, see G. Feinberg and R. E. Behrends, Phys. Rev. **115**, 745, 1959; K. Tanaka, Phys. Rev. **122**, 705 (1961), and G. Feinberg (to be published). There are also perturbation-theoretic calculations [see W. G. Holladay, Phys. Rev. **115**, 1331 (1959)], which indicate that the Λ and Σ^0 magnetic moments are quite small ($<0.1\ e/2m_p$) as long as the K-meson coupling constants are not as large as the π-meson ones. The hypothesis of unitary symmetry,[11] on the other hand, gives $\mu_{\Sigma^0\Sigma^0} = -\mu_{\Lambda\Lambda} = -\frac{1}{2}\mu_{\text{neutron}}$. This symmetry scheme, as emphasized by Gell-Mann, must however be badly broken, since experiments indicate that $KN\Lambda$ and $KN\Sigma$ coupling constants are much smaller than the πNN coupling constant in contradiction with unitary symmetry. It is hard to tell, therefore, how much one has to rely upon the conclusions of this scheme about the magnetic moments in the presence of other unknown strong symmetry-breaking interactions. At any rate, we hope a choice between the above results can be made by a direct measurement of the magnetic moment of the Λ hyperon. The magnetic moment of Σ^0 can be obtained indirectly from a measurement of the magnetic moments of Σ^+ and Σ^- by using the Marshak, Okubo, and Sudarshan relation $\mu_{\Sigma^+} + \mu_{\Sigma^-} = 2\mu_{\Sigma^0}$, which is based on charge independence alone.

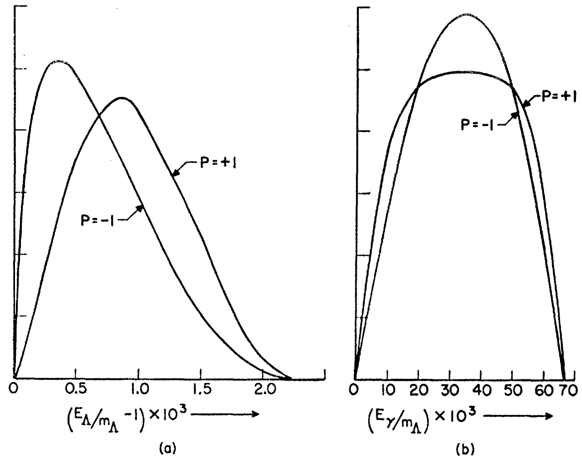

FIG. 3. The energy spectra of the Λ hyperon (a) and the photons (b), for even $(P=+1)$ and odd $(P=-1)$ relative $\Sigma-\Lambda$ parities. The spectra for $P=\pm 1$ are normalized to have the same area.

In this case Fig. 3 leads to a branching ratio for the $\Sigma^0 \to \Lambda^0 + 2\gamma$ mode less than 3.8×10^{-8} for either parity case with any choice of signs for $X_{\Sigma^0\Sigma^0}$ and $X_{\Lambda\Lambda}$. This is still much smaller than the contribution from the π^0-pole term for the odd-parity case [cf. Eq. (6)]. The situation does not change by considering the interference term between Fig. 1 and Fig. 2 either.

Thus we have shown that the large difference between the branching ratios for even and odd relative parity, obtained by the consideration of the π^0-pole term, remains true even if we consider other diagrams. We therefore feel safe to conclude, allowing a rather large margin for the uncertainties in the coupling constants and the hyperon-electromagnetic form factors, that if we can see one $\Sigma^0 \to \Lambda^0 + 2\gamma$ event out of less than 10^6 Σ^0 decays, the $\Sigma-\Lambda$ relative parity must be odd; while if we cannot see any such event in more than 10^7 Σ^0 decays, the relative parity must be even. Of course, if one is optimistic enough, one may hope that the value of $X_{\Sigma\Lambda}$, or equivalently the rate of $\Sigma^0 \to \Lambda^0 + \gamma$ decay, is somewhat smaller than that chosen in this article. For instance, if $X_{\Sigma\Lambda}$ is less than 0.2, one may hope to see (if $P=-1$) one event of $\Sigma^0 \to \Lambda^0 + 2\gamma$ out of even 2×10^4 or less Σ^0 events. This, if seen, would unambiguously determine that the $(\Sigma\Lambda)$ relative parity is odd. An accurate determination[5] of Σ^0 lifetime would therefore be particularly interesting in this connection.[17]

As to the experimental possibility, we hope that, with enough intensity, a study of the $\Sigma^0 \to \Lambda^0 + 2\gamma$ decay may be feasible, for instance, in a heavy liquid chamber.

[17] Alternatively, we may note that, if the branching ratio of $\Sigma^0 \to \Lambda^0 + 2\gamma$ is found to indicate odd relative parity, then the same data can be utilized, either to infer the value of $X_{\Sigma\Lambda}$, assuming that of $g^-{}_{\Sigma\Lambda\pi}$, obtained by other [8,9] methods, or to deduce the value of $g^-{}_{\Sigma\Lambda\pi}$, knowing $X_{\Sigma\Lambda}$ by an independent method. The same cannot, however, be done for the even-parity case, since in addition to Fig. 1 there exist other competing diagrams (for example, Fig. 2) with somewhat uncertain contributions.

One process[18] that may be looked for is: $K^-+\text{nucleus}\rightarrow \Sigma^0+\pi^-+(\text{nucleus})^*$; $\Sigma^0\rightarrow\Lambda^0+2\gamma$, with the subsequent pair production by the two photons and the $(p\pi^-)$ production by the Λ hyperon. This may be feasible in a heavy liquid chamber, specially since the photons are expected to be energetic.

In conclusion, except for the large[19] intensity requirement, which will be a drawback, the method described above seems to have certain advantages over the other[1] methods, since it does not need any polarization of Σ^0, and does not involve the somewhat difficult task of studying correlation effects such as between the spin of Λ^0 and the plane of the pair in the Σ^0 Dalitz decay.

[18] This was suggested to one of us (J.C.P.) by G. A. Snow.

[19] The methods suggested in reference 1 involving Dalitz decay of Σ^0 require nearly 10^4 to 10^5 polarized Σ^0 events for an unambiguous determination of P. The present method, on the other hand, needs nearly 10^5 to 10^6 Σ^0 events without any restriction on their polarization.

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One of us (J.C.P.) would like to thank Professor M. Gell-Mann, Professor R. L. Walker, and Professor G. A. Snow for many helpful comments and discussions.

Note added in proof. The essentIial content of this work was presented by one of us (S.O.) at the 1961 Spring meeting of the Japanese Physical Society held in Tokyo. After the completion of this work Dr. K. Fujii kindly called our attention to a recent similar work by Okun and Rudik.[20] These authors do not emphasize, however, the importance of the π^0-pole term, and their main interest in the $\Sigma^0\rightarrow\Lambda^0+2\gamma$ decay is not in connection with the determination of the $\Sigma-\Lambda$ relative parity. Instead, they discuss mainly Fig. 2 in connection with the determination of the Σ^0 magnetic moment. See also J. Bernstein and R. Oehme, Phys. Rev. Letters **6**, 639 (1961).

[20] L. B. Okun and A. P. Rudik, Zhur. Eksp. i Teoret. Fiz., **39**, 378 (1960).

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Range of the Nucleon-Antinucleon Annihilation Potential

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It is shown that the assumption of the validity of the Mandelstam representation for nucleon-antinucleon scattering leads to a potential, fitting the data at a given energy, with an imaginary part, the range of which cannot exceed half the nucleon Compton wavelength.

MANY theoreticians state that on the basis of field theoretical arguments, the range of the nucleon-antinucleon annihilation potential must be of the order of the nucleon Compton wavelength.[1] However, this is not obvious because one has to define in a correct way a complex potential describing scattering and disappearance of the nucleon-antinucleon system. To our knowledge this has not been done up to now. Consequently other theoreticians, mainly under the pressure of early experimental results in the low-energy region (these results turned out later to be wrong) and of more recent results in the 1–2 Gev region,[2] tried to construct field-theoretical[3] or phenomenological[4] models in which the annihilation force has a long range.

We do not wish to discuss here the experimental situation. We would like to show that it looks very difficult, in the framework of Mandelstam representation, to have a nucleon-antinucleon annihilation potential with a range larger than half the nucleon Compton wavelength.

In a paper published elsewhere[5] Targonski and the author have indicated a method of construction of an energy-dependent nucleon-nucleon potential, fitting a scattering amplitude at a given energy, when this scattering amplitude has the analytic properties implied by the Mandelstam representation, with respect to the scattering angle. This potential is a superposition of Yukawa potentials. There is no objection to applying this method to the case of nucleon-antinucleon scattering at a given energy. The only change will be that the potential obtained in this way will be complex, since absorption takes place. For simplicity we shall neglect spin complications and assume that the energy at which we try to construct the potential is below the one-meson production threshold. We shall assume that there is only one kind of nucleon, to avoid the troubles due to $p\bar{p}\rightarrow n\bar{n}$ scattering, but this can be easily corrected because one can start from initial states with given isospin.

Let us summarize briefly the general method. The

[1] J. S. Ball and G. F. Chew, Phys. Rev. **109**, 1385 (1958); J. S. Ball and J. R. Fulco, *ibid.* **113**, 647 (1959).

[2] *Proceedings of the 1960 Annual International Conference on High-Energy Physics at Rochester* (Interscience Publishers, Inc., New York, 1960), p. 658.

[3] M. Lévy, Nuovo cimento **8**, 92 (1958); J. Mandelbrojt, Nuovo cimento (to be published).

[4] Z. Koba and G. Takeda, Progr. Theoret. Phys. (Kyoto) **19**, 269 (1958); B. Jancovici, M. Gourdin, and L. Verlet, Nuovo cimento **8**, 485 (1958); M. Lévy, Phys. Rev. Letters **5**, 380 (1960); J. Mandelbrojt (to be published); O. Hara, Phys. Rev. **122**, 669 (1961).

[5] A. Martin and Gy. Targonski, Nuovo cimento **20**, 1182 (1961).