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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Note added in proof. The essentital 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).

scattering amplitude has two cuts in the $\cos\theta$ plane, one from

$$\cos\theta = 1 + (\alpha_{D\min})^2/2k^2$$

to

$$\cos\theta = +\infty,$$

the other one from

$$\cos\theta = -[1 + (\alpha_{E\min})^2/2k^2]$$

to

$$\cos\theta = -\infty,$$

where k is the c.m. momentum. Then, provided $2\alpha_E > \alpha_D$, one can try to fit it by a potential

$$r^{-1}\int_{\alpha_{D\min}}^{\infty} C_D(\alpha)\exp[-\alpha r]d\alpha + P_M r^{-1}\int_{\alpha_{E\min}}^{\infty} C_E(\alpha)\exp[-\alpha r]d\alpha,$$

where P_M is the Majorana exchange operator. Step by by step, one can construct $C_E(\alpha)$ and $C_D(\alpha)$ from the knowledge of the discontinuities across the cuts, and if these are known till

$$\cos\theta = \pm[1 + (\alpha_{\max})^2/2k^2],$$

then $C_E(\alpha)$ and $C_D(\alpha)$ are known for $\alpha < \alpha_{\max}$. Physically this means that the first thing we can get is the external part of the potential; the farther we know the discontinuities, the better we know the inner part of the potential.

Let us apply now these considerations to the nucleon-antinucleon case. The scattering amplitude T for this process can be split into two parts T_a and T_b (Fig. 1).

T_a is a purely elastic amplitude analogous to the nucleon-nucleon amplitude; it satisfies the unitarity condition, which expresses the absence of any possible reaction other than scattering (we are below the meson production threshold). In the absence of anomalous thresholds it is easily seen that the lowest singularities of T_a for fixed energy are given by

$$t = -(N_f - N_i)^2 = \mu^2 \text{ (one-meson pole)},$$

$$t = (2\mu)^2 \text{ (first branch point)},$$

$$\bar{t} = -(\bar{N}_f - N_i)^2 = (2M)^2,$$

where μ is the π-meson mass and M the nucleon mass.

The lowest singularities in T_b lie at

$$t = \bar{t} = (2M)^2.$$

To be more accurate we should everywhere replace $2M$ by the deuteron mass but this is unimportant.

If we return to the variable $\cos\theta$, we see that in the region

$$|\cos\theta| < 1 + 4M^2/2k^2,$$

the only contributions to the discontinuities come from T_a and that the left-hand cut is missing. Hence $C_E(\alpha) = 0$ for $\alpha < 2M$; and $C_D(\alpha)$, for $\alpha < 2M$, is the same as if T_b were absent. We know, however, that T_a satisfies an elastic unitarity condition. Therefore $C_D(\alpha)$ is necessarily real for $\alpha < 2M$. From this we conclude that *the range of the imaginary part of the potential cannot exceed half the nucleon Compton wavelength.*

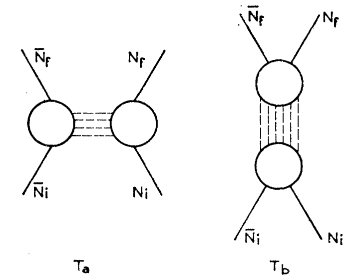

FIG. 1. The two contributions to the scattering amplitude T. The dashed lines represent mesons (π mesons, K, $\bar{K}$ mesons).

Strictly speaking, our analysis is restricted to energies below the meson production threshold. However, as long as the meson production cross section is small, T_a satisfies fairly well an elastic unitarity condition and therefore $C_D(\alpha)$, in the range $\mu \leqslant \alpha \leqslant 2M$, is almost real. Concerning the problem of the energy dependence of the potential, we can guess that the Charap-Fubini-Tausner argument[6] can be applied to T_a and therefore the external part of the real potential is energy-independent in the low-energy region. This is probably not the case for the imaginary potential. We are probably in a situation analogous to that of the optical potential in nuclear physics; the real potential has a weak dependence on energy while the imaginary potential has a strong energy dependence.

As far as we can see, the weak points of our reasoning are the following:

(1) The assumption that the Mandelstam representation holds for T_a and T_b.

(2) The assumption, underlying the work of reference 5, that the potential one obtains eventually is not too singular at the origin.

(3) The use of a potential model. This is, however, the way the experiments have been analyzed up to now and on the other hand, we cannot see any clearer way of defining a range in configuration space.

As a concluding remark we wish to point out that if one takes our argument seriously, it imposes a very severe restriction on the phenomenological analysis of the data on nucleon-antinucleon scattering. At low energies this restriction seems to be compatible with the known experimental facts.[1] At higher energies, it is not yet clear, in our opinion, that it is incompatible with experimental facts.[7]

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[6] J. Charap and S. Fubini, Nuovo cimento **14**, 540 (1959); **15**, 73 (1960); J. Charap and J. Tausner, *ibid.* **18**, 316 (1960).

[7] We think that the tail of the $r^{-1}\exp[-2Mr]$ imaginary potential and the inner part of the real potential might play an important role. This will be discussed in a forthcoming publication.