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$$f_{\mathrm{even}}(l) = -\frac{1}{2k^2}\int [\rho(t) + \sigma(t)]\, Q_l(1 + \frac{t}{2k^2})\, \mathrm{d}t \ ,$$

which satisfies conditions (8) (replacing (6) and (7)) and coincides with the physical scattering amplitudes for even integral l values, and

$$f_{\mathrm{odd}}(l) = -\frac{1}{2k^2}\int [\rho(t) - \sigma(t)]\, Q_l(1 + \frac{t}{2k^2})\, \mathrm{d}t$$

which also satisfies (8) and coincides with the physical scattering amplitudes for odd integral l values.

These two interpolating functions are bounded by $A+B|l|$ for Re $l > m$, they are free of poles, they are the only interpolations

$$|f(l)| < \exp\,[\,|l|\,(\frac{\pi}{2} - \epsilon)]$$

at infinity in Re $l > -\frac{1}{2}$, and they satisfy *automatically* the unitarity condition.

One may now think of making the prolongation for complex values of the energy. $\rho(t)$ and $\sigma(t)$ are actually functions $\rho(s,t)$ and $\sigma(s,t)$ analytic in the s plane, with a cut from $s = 4\mu^2$ to $-\infty$. This holds for equal masses μ. Now $Q_l(z)$ has a cut from $-\infty$ to $+1$ and, as long as $|z + \sqrt{z^2 - 1}| > 1$ (defining $\sqrt{z^2 - 1}$ as positive when z is real > 1) it behaves at infinity like

$$\sqrt{\pi}\,\frac{\Gamma(l+1)}{\Gamma(l+3/2)}\,(z + \sqrt{z^2 - 1})^{-l-1} \ ,$$

so that in Im $s > 0$, Im $l > 0$ the integral converges as well as before and is analytic both in s and in l for Re $l > m$, if m is taken as the upper limit of the number of subtractions in t and in u in the original Mandelstam representation.

We realise that these results are very incomplete because they do not say anything about the interesting region in the l plane, namely the one where you expect to find poles and, possibly, cuts [5]. Another disagreeable feature is that so far we have treated the regions Im $l > 0$ and Im $l < 0$ on the same footing, whilst, in potential scattering, the region Im $l < 0$ does not contain any singularity. One may hope that these difficulties might be overcome by playing on the double analyticity in l and E, using more intensively the unitarity condition, and trying to use the analytic properties of $\rho(t)$ considered as the difference between two analytic functions in two different Riemann sheets.

I wish to thank Professor K. Symanzik for a discussion which initiated this work, and Dr. E.Leader for stimulating discussions.

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PIONIC DECAY OF HYPERONS IN THE POLE APPROXIMATION

AKIHIKO FUJII
Frascati National Laboratory, Frascati (Rome), Italy

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We have correlated the pionic decay rates and the asymmetry parameters of the hyperons Σ^+, Σ^-, and Λ^0 in a simple model with the pole approximation described in fig. 1. The K' particle is the experimentally observed $K\pi$ resonance at 885 MeV, which is assumed to be a stable vector particle [1]. We may also consider the pole diagram shown in fig. 2, where the intermediate boson X may be either the η [2], ρ [3] or ω [4] particle. If we require $|\Delta I| = \frac{1}{2}$ for weak interactions and $\Delta I = 0$ for strong interactions, then X must possess unit isospin, which is then the ρ particle. However, the ρ particle decays into two pions, hence it cannot decay into a single pion strongly. Isospin may not be a

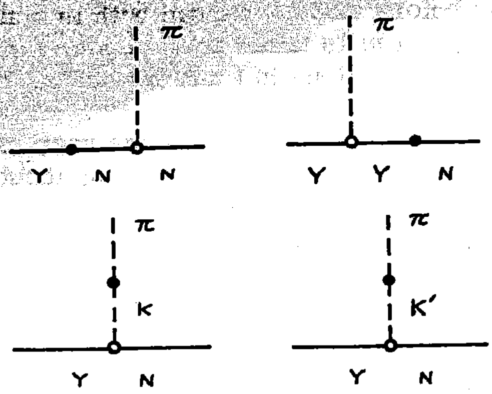

Fig. 1. Decay diagrams in the pole approximation. The open circle represents a strong interaction, and the black circle a weak interaction.

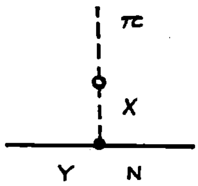

Fig. 2. The decay diagram via pion-pion resonance.

good quantum number for strong decay [5], but we exclude this diagram in the rest of this communication.

Complete symmetry between Σ and Λ is adopted for strong interactions, namely, the coupling of these hyperons to the pion, $i\, g_{Y\pi}\, \overline{Y}\, \gamma_5\, Y \cdot \pi$, or kaon, $i\, g_K\, \overline{N}\, \gamma_5\, Y \cdot K$, is of the same strength as well as charge independent, and their masses are degenerate. The weak "vertices" are assumed to transform as a spinor in charge space so as to yield an overall $|\Delta I| = \frac{1}{2}$ rule. Except for the numerical factors infered from the charge space structure, the weak interactions are described by the following phenomenological hamiltonians:

(YN)-vertex $\quad m_\pi\, \overline{N}(f + f'\,\gamma_5)\, Y + \text{h.c.}$

(K π)-vertex $\quad m_\pi^2\, f_K\, K\, \pi + \text{h.c.}$

(K'π)-vertex $\quad m_\pi\, f_{K'}\, K'_\mu\, \frac{\partial \pi^*}{\partial x_\mu} + \text{h.c.}$

The virtual K' can transform into π through the non-vanishing fourth component. The strong interaction of K' with baryons is simply taken as $i\, g_{K'}\, \overline{N}\, \gamma_\mu\, Y \cdot K'_\mu$, neglecting all other possible "moment" type couplings. The dimensionless coupling constants f's and g's are all real, provided that time reversal invariance is valid.

The effective transition operator for the process $Y \to N + \pi$ may be written in the general form

$$T = A + B\,\gamma_5\ ,$$

which yields the decay rate w and the asymmetry parameter α as follows:

$$w = \frac{1}{4\pi} \cdot \frac{P_N}{m_N} \cdot \left[|A|^2 \frac{(m_Y + m_N)^2 - m_\pi^2}{2m_Y^2} + |B|^2 \frac{(m_Y - m_N)^2 - m_\pi^2}{2m_Y^2} \right] \cdot m_N\ ,$$

where P_N is the magnitude of the three-momentum of the final nucleon,

$$\alpha = \frac{2\,\mathrm{Re}\, x^*}{1 + |x|^2}$$

$$x = \frac{P_N}{m_N + \epsilon_N}\left(\frac{B}{A}\right) = \sqrt{\frac{(m_Y - m_N)^2 - m_\pi^2}{(m_Y - m_N)^2 - m_\pi^2}}\left(\frac{B}{A}\right)\ ,$$

where x is the ratio of the p-wave to s-wave decay amplitude in the non-relativistic approximation, and ϵ_N is the total energy of the final nucleon.

When the corresponding quantities in the channel $\Sigma^+ \to n + \pi^+$, $\Sigma^+ \to p + \pi^0$, $\Sigma^- \to n + \pi^-$ are distinguished by the superscript +, o, - respectively and without a superscript for the channel $\Lambda^0 \to p + \pi^-$, the pole approximation gives

$$\begin{cases} A^{(+)} = -\,2f'(g_{N\pi} + g_{Y\pi})\,\dfrac{m_\pi}{m_Y + m_N}\ , \\ B^{(+)} = 2f\,(g_{N\pi} - g_{Y\pi})\,\dfrac{m_\pi}{m_Y - m_N}\ , \end{cases} \qquad (1)$$

$$\begin{cases} A^{(-)} = 2f_{K'}\, g_{K'}\,\dfrac{m_\pi(m_Y - m_N)}{m_{K'}^2}\ , \\ B^{(-)} = 2f_K\, g_K\,\dfrac{m_\pi^2}{m_K^2 - m_\pi^2}\ , \end{cases} \qquad (2)$$

$$\begin{cases} \sqrt{2}\, A^{(o)} = A^{(+)} - A^{(-)}\ , \\ \sqrt{2}\, B^{(o)} = B^{(+)} - B^{(-)}\ , \end{cases} \qquad (3)$$

$$\begin{cases} \sqrt{2}\, A = A^{(+)} + A^{(-)}\ , \\ \sqrt{2}\, B = B^{(+)} + B^{(-)}\ . \end{cases} \qquad (4)$$

Eq. (3) is known to be the direct consequence of the $|\Delta I| = \frac{1}{2}$ rule.

Experimentally it is known that

$$\alpha^{(+)} \sim 0\ , \quad \alpha^{(-)} \sim 0\ , \quad w^{(+)} \sim w^{(-)}\ .$$

These facts are satisfied if either one of the following conditions holds,

$$A^{(+)} = B^{(-)} = 0\ ,$$

$$|A^{(-)}|^2 \frac{(m_\Sigma + m_N)^2 - m_\pi^2}{2m_\Sigma^2} = |B^{(+)}|^2 \frac{(m_\Sigma - m_N)^2 - m_\pi^2}{2m_\Sigma^2}$$

or

$$A^{(-)} = B^{(+)} = 0 \ ,$$

$$|A^{(+)}|^2 \frac{(m_\Sigma + m_N)^2 - m_\pi^2}{2m_\Sigma^2} = |B^{(-)}|^2 \frac{(m_\Sigma - m_N)^2 - m_\pi^2}{2m_\Sigma^2}$$

In both of the above cases we obtain

$$\frac{w}{w^{(+)}} \sim \frac{P_N}{P_N^{(+)}} \sim \tfrac{1}{2}$$

$$x = \pm \sqrt{\frac{(m_\Lambda - m_N)^2 - m_\pi^2}{(m_\Lambda + m_N)^2 - m_\pi^2}} \cdot \sqrt{\frac{(m_\Sigma + m_N)^2 - m_\pi^2}{(m_\Sigma - m_N)^2 - m_\pi^2}}$$

$$= \pm 0.54$$

or

$$\alpha = \pm 0.84 \ ,$$

which are in good agreement with experimental findings.

If K' turns out to be a particle of spin zero it must be scalar because it decays strongly into K and π. If this is the case eq. (4) still holds with a trivial modification in eq. (2).

Actually eq. (4) has already been obtained by Treiman [6] and Pais [7] from weak global symmetry, where all kaon effects are omitted. Feldman, Matthews and Salam [8] showed that the pole approximation without K' could give a reasonable result if the coupling constant $g_{\Lambda NK}$ is about $\sqrt{3}$ times the coupling constant $g_{\Sigma NK}$. The present note bridges these two arguments and may suggest the existence of a weak symmetry covering both the pion and kaon effects.

The author expresses his sincere thanks to Professor R. Gatto for his discussions and comments.

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DEUTERONS FROM HIGH-ENERGY PROTON BOMBARDMENT OF MATTER

S. T. BUTLER and C. A. PEARSON
The Daily Telegraph Theoretical Department,
School of Physics, University of Sydney

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In a recent publication we described a mechanism which can account for the large number of relativistic deuterons detected among the secondaries from high energy collisions [1]. These deuterons had been observed among the secondaries from targets exposed to the 25 GeV proton beam of the CERN proton synchroton [2]. They were also known to be responsible for a large proportion of the "grey" tracks in nuclear emulsions [3-5].

It had previously been suggested that these deuterons are formed in elementary nucleon-nucleon collisions [6].

However, the extreme energy dependence of the results from such a process and the complete disagreement with experiment at low energies indicate that this is not the dominant means by which the deuterons are produced.

At the time when our previous letter was published there already seemed good reason to believe that the alternative mechanism proposed by us, a secondary process requiring the presence of nuclear matter, was responsible for the majority of the observed deuterons. Since then, a number of experiments have been performed at CERN and at Brookhaven which confirm this belief.

At Brookhaven the momentum distribution of