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Part I

Elementary Particles and Interaction Symmetries

The Nucleon–Nucleon Interaction and the Interaction of Nucleons with Isobars[1]

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Abstract

The nucleon–nucleon interaction and the interaction of nucleons with isobars. G. E. Brown (Nordita, Copenhagen, Denmark, and State University of New York, Stony Brook, New York, USA).

The origin of the nucleon–nucleon interaction from the exchange of the π- and ω-mesons and from two-pion exchange, is discussed. The two-pion exchange term is described in a dispersion theoretical formalism, and it is shown how the π–π scattering phase shift $\delta_{\pi\pi}^{J=0}$ governs the strength of the σ-degrees of freedom. Going over to dynamical models, the interaction of nucleons with the N_{33}^*-isobar is discussed. A model for the two-pion exchange involves the virtual isobar as intermediate state. The presence of virtual isobars is seen, perhaps most clearly, in the process $n+p\rightarrow d+\gamma$. The isobar also enters into various exchange current corrections to magnetic moments and β-decay. Using any one of several simple models for the isobar, a consistent picture of the various phenomena in the two- and three-nucleon systems can be achieved.

1. Introduction

Calculations of the nucleon–nucleon interaction have been proceeding on several levels, two of which I wish to discuss here. The penultimate way in which the nucleon–nucleon interaction will be found, it seems to me, is through dispersion of theoretical calculations, following the early work of Amati, Leader and Vitale [1]. In some sense, one can say that this is not a calculation, in that the dispersion relations are used to relate quantities determining the nucleon–nucleon force to observables in pion–nucleon and pion–pion scattering. I shall return to this in § 2. The other level at which I wish to discuss the forces, is that of simple dynamical calculations. It seems to me that rigorous, detailed dynamical calculations are probably much harder to carry through than the dispersion-theoretical ones, but one can get a long way in a low-brow approach, provided that one builds in low-energy theorems, etc., which keep one from going too far astray.

2. Dispersion theoretical calculations

I shall begin by discussing the more high-brow way of "calculating" the nucleon–nucleon interaction, while my audience is fresh, and ask you to bear with me, in case this seems rather complicated; it is complicated in detail but not in concept.

Everyone agrees that the long-range part of the nucleon–nucleon interaction is given by the one-pion-exchange term, and this is so familiar that I shall not bother to write it down. I shall also assume that the short-range repulsion is given by the exchange of ω-mesons. The coupling constant of the ω-meson to nucleons has recently been determined [2] rather accurately from the nucleon–nucleon forward scattering dispersion relations to be $g^2/4\pi \cong 7$; I shall not quote errors in this somewhat preliminary work, but this gives us the general order-of-magnitude, and gives a repulsion which seems adequate to describe the low-energy nucleon–nucleon scattering, as you will see later from Fig. 6. This ω-coupling strength is about 12 times the coupling strength of the ρ-meson, and somewhat exceeds the maximum limit of this ratio (9) which one could get in SU(3). This larger coupling strength may indicate that our ω-exchange potential is also summarizing some of the three-pion continuum.

[1] Supported under US Atomic Energy Commission contract AT-(11-1)-3001.

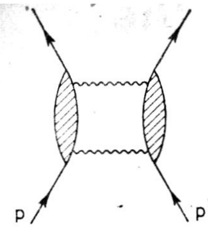


Fig. 1. The two-pion-exchange term. The wavy lines represent pions. The cross-hatched areas represent all possible strong interactions.

From this point onwards, many people have introduced other mesons in the spirit of the one-boson-exchange models, and obtained tolerable fits to the nucleon–nucleon scattering data. All of these models require the exchange of a spin-zero, isoscalar meson of intermediate mass, 400 to 800 MeV, which we shall call the σ-meson. It is now quite clear that no such meson exists in the form required, in nature, and what is being summarized by such a meson is the exchange of two-pion systems, in which these two-pions are coupled to angular momentum zero. This is the situation which is convenient to discuss in the framework of dispersion relations.

The two-pion-exchange term, in all generality, is shown in Fig. 1.

In the dispersion-theoretical treatment, it is convenient to go to the centre-of-mass system of the two pions, considering the process, Fig. 1, to be made up out of the product of two graphs, such as shown in Fig. 2.

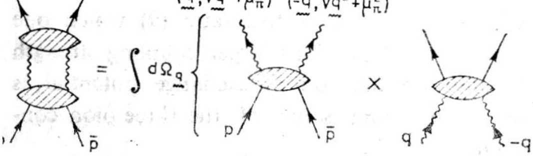
We now go further and decompose each of the constituents into a sum over J, as shown in Fig. 3.

Here

$$t = 4(q^2 + \mu_\pi^2) \tag{1}$$

so that $\sqrt{t}$ is the energy of the two-pion system. It is seen that $f^J(t)$ is the amplitude for emitting two pions of energy $\sqrt{t} = 2\sqrt{q^2+\mu_\pi^2}$, so we can

Fig. 2. Decomposition of the two-pion-exchange term into the product of $p\bar{p} \to \pi+\pi$ processes.



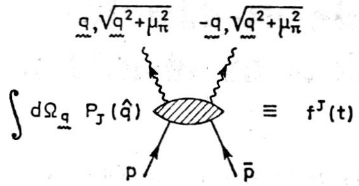


Fig. 3. Decomposition of the $p\bar{p} \to \pi+\pi$ process into terms of definite angular momentum.

expect, roughly, to be able to represent the two-pion-exchange term in the S-matrix as a sum of Yukawas of varying ranges corresponding to the varying masses of the exchanged two-pion systems. For the contribution from $J=0$, which we discuss below,

$$S_{\text{two-pion-exchange}} \simeq \int_{4\mu_\pi^2}^{\infty} q \left|f^{J=0}(t)\right|^2 \frac{e^{-\sqrt{t}r}}{r} \, dt \tag{2}$$

where q is the phase-space factor for the two-pions in an S-state,

$$q = \sqrt{\frac{t}{4} - \mu_\pi^2}. \tag{2.1}$$

Aside from various factors irrelevant for the present discussion, eq. (2) is what emerges from the Amati, Leader and Vitale [1] formalism.

Let us in particular consider this case of $J=0$. Here, the two-pion system behaves as a scalar and we are talking about the σ-degrees of freedom. Now the amplitude for $p+\bar{p} \to \pi+\pi$ shown in Figs. 2 and 3 is the same amplitude as one encounters in $\pi + p \to \pi + p$, but in another regime of the kinematics; i.e., the amplitudes we need in order to calculate $f^J(t)$ are analytic continuations of the pion–nucleon scattering amplitude. In the region of small t, this latter amplitude is determined by the soft-pion theorems and these can be used [3] to calculate the real part of $f^{J=0}(t)$. The functional form calculated in this way seems to agree tolerably well [4, 5] even for relatively large t with that obtained by analytical continuation of the π-nucleon scattering amplitude. We show the various $f^{J=0}(t)$'s in Fig. 4.

We notice that $\mathrm{Re} f^{J=0}(t)$ is relatively flat; certainly the contribution to the integrand on the right-hand side of (2),

$$q(\mathrm{Re} f^{J=0}(t))^2 = \sqrt{\frac{t}{4} - \mu_\pi^2}\,(\mathrm{Re} f^{J=0}(t))^2$$

is relatively flat.

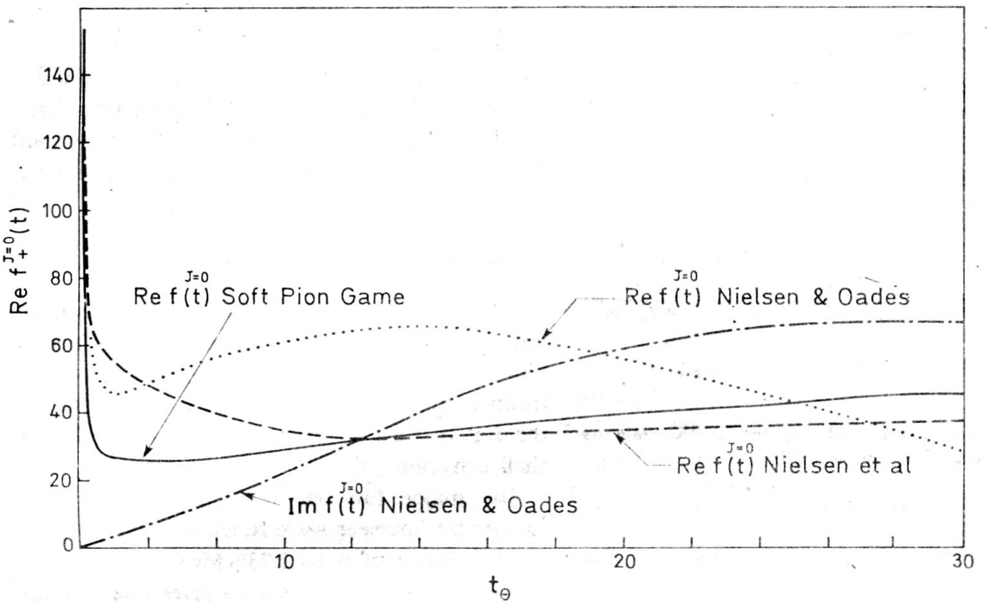


Fig. 4. Forms of $f^{J=0}(t)$ obtained by soft-pion theorems ——, analytic continuation from the pion–nucleon amplitude: ----, Nielsen et al. [4]; ....., Nielsen and Oades [5].

We have still to discuss the imaginary part of $f^{J=0}$, in $f^{J=0}(t)$. In the region $4\mu_\pi^2 < t < 16\mu_\pi^2$ one can rigorously show from unitarity that the phase of $f^J(t)$ is $\delta_{\pi\pi}^J(t)$, the latter being the phase shift for elastic π–π scattering in channel J.

We thus have

$$\operatorname{Im} f^{J=0}(t) = \tan \delta_{\pi\pi}^{J=0}(t) \operatorname{Re} f^{J=0}(t) \tag{3}$$

which is usually assumed to hold for all t up to $\sim 50\ \mu_\pi^2$ because of the smallness of inelastic processes in $\pi\pi$ scattering. Now if $\delta_{\pi\pi}^{J=1}$ were to go through 90° at $t = t_0$, then $f^{J=0}(t)$ would become very large at that point and contributions to the integrand on the right-hand side of eq. (2) would come mainly from $t = t_0$, so that

$$S_{\text{two-pion-exchange}} \simeq C \frac{e^{\sqrt{t_0}\, r}}{r}. \quad (\pi\pi \text{ resonance at energy } t_0)$$

In this case, the description of all $J=0$ exchanges in terms of one meson of mass $m_\sigma = \sqrt{t_0}$ would be a good approximation.

In fact, the $\pi\pi$ phase shift $\delta_{\pi\pi}^{J=0}$ seems to go up to 60–70° in the region 500–700 MeV, and does not reach 90° until higher energies, as shown in Fig. 5.

The behaviour at energies higher than $\sim$700 MeV is not very important, because the strong repulsion from the ω-meson exchange is expected to obscure any other contributions of shorter range $\hbar/\sqrt{t} < \hbar/m_\omega c$.

Thus, it seems necessary to include the exchange of a wide spectrum of masses in the σ-degrees of freedom.

In the case of $J=1$, the two pions have the quantum numbers of the ρ-meson, and whereas the main weighting for these exchanges comes from the region of the ρ-mass, effects from the continuum of lower mass exchanges can be directly seen [6] in the nucleon–nucleon scattering data.

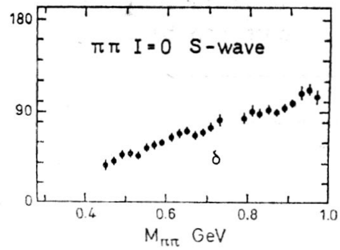


Fig. 5. Behaviour of the $J=0$ ππ-phase shift. This figure is taken from a report by A. D. Martin and P. Estabrooks to the Daresbury Meeting on Pion Exchange, March 23, 1973.

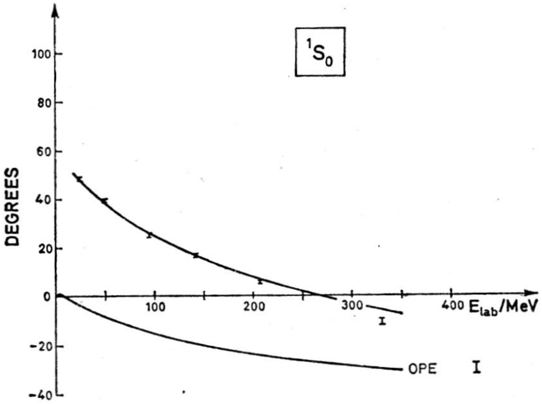


Fig. 6. Nucleon–nucleon scattering in the $^1S_0$ state as calculated by Riska and Verwest [7].

In Fig. 6 is shown the nucleon–nucleon scattering in the $^1S_0$ state calculated [7] with the $\mathrm{Re}\, f^{J=0}(t)$ obtained from the soft-pion theorems, $\mathrm{Im}\, f^{J=0}(t)$ obtained from (3) with

$$\tan \delta^J_{\pi\pi}(t) = \frac{q}{(1/a_{\pi\pi})\,(1 - q^2/q_x^2)\, q \cot \delta_x}$$

the parameters being chosen so as to give a scattering length $a_{\pi\pi} = 0.24$ and a maximum in δ of $(\delta_{\pi\pi})_{\max} = 70°$ at 600 MeV. The $f^{J=1}(t)$ were taken from Ref. 3.

As noted earlier, $(g^2/4\pi)$ is taken to be 12 $(g_\rho^2/4\pi)$ and $g_\rho^2/4\pi = 0.6$.

Of course, one has enough parameters at one's disposal, if one assumes the $f^J(t)$ to be variable parameters, to fit almost anything, especially the nucleon–nucleon scattering in only one channel. The point I want to make is that picking behaviours for the $f^J(t)$ obtained from π-nucleon and π-π scattering, one obtains a good fit.

The work of Riska and Verwest has not been extended to all channels because it was shown [8] that the higher helicity amplitudes $f^J(t)$ for $J = 2$ have non-negligible effects, and these had not been properly handled in the earlier work. The program of Riska et al. is now moving towards proper inclusion of these higher helicity amplitudes.

Note, however, that the phenomenological soft-core potential of Reid [9] has a repulsion that would correspond to a $g_\omega^2/4\pi = \sim 70$! However, such a tremendous repulsion does not seem to be necessary.

The dispersion theoretical calculations seem to be moving slowly towards a satisfactory conclusion, showing that the various relations between different processes following from analyticity seem to hold reasonably well. I now want to move on to a discussion of dynamical models.

3. Dynamical models: games people play

The dispersion formalism gives one the nucleon–nucleon scattering only on-shell; if one wishes to obtain it off-shell, one needs models. These introduce arbitrariness, of course, but much of the fun in nuclear physics comes from the off-shell behaviour of the interactions.

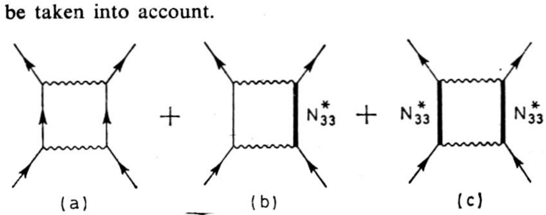
The major fun has come from introducing isobars into nuclear wave functions, especially the $\Delta = N^*_{33}$ isobar of mass 1 236 MeV.

The introduction of isobars gives one a model —certainly a vastly oversimplified one—for the two-pion-exchange interaction, Fig. 1. We show this model in Fig. 7.

It was shown by Riska and Brown [10] following earlier work by Sugawara and von Hippel [11] that these processes gave roughly the correct magnitude for the σ-meson exchange. This work has been greatly improved in a contributed paper to this meeting by Green and Haapakoski [12] who treat the nucleon and isobar in a coupled-channels formalism.

The coupling of NN to NN* is taken to be the one-pion-exchange potential. It does not matter much whether one uses the quark model [13], manufactures a Rarita-Schwinger (3,3)-isobar [10] or describes the isobar in Chew-Low theory [14], the coupling potential $V_{\mathrm{NN,N^*N}}$ has the same functional form, and the strength varies at most ~25% between these models. Just as the OPEP must be cut off in interactions between nucleons, the $V_{\mathrm{NN,N^*N}}$ must also be cut off at

Fig. 7. Model for the two-pion-exchange. In addition to these processes, the crossed pion exchange should also be taken into account.



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