sensitivity to the assumptions about the asymptotic behavior of the cross sections at high energy is in progress, so that a critical comparison of theory and experiment can be attempted.

The authors wish to thank the Brookhaven National Laboratory Accelerator Department for valuable cooperation in providing desired beam characteristics, magnet measurements, etc.

---

*Work performed under the auspices of the U. S. Atomic Energy Commission.

†Visitor from Rutherford High Energy Laboratory.

[1]K. J. Foley, R. S. Gilmore, R. S. Jones, S. J. Lindenbaum, W. A. Love, S. Ozaki, E. H. Willen, R. Yamada, and L. C. L. Yuan, Proceedings of the International Conference on High-Energy Physics, Dubna, 1964 (to be published); S. J. Lindenbaum, Rapporteur's Report, Proceedings of the International Conference on High-Energy Physics, Dubna, 1964 (to be published).

[2]K. J. Foley, R. S. Gilmore, R. S. Jones, S. J. Lindenbaum, W. A. Love, S. Ozaki, E. H. Willen, R. Yamada, and L. C. L. Yuan, Phys. Rev. Letters 14, 74 (1965).

[3]The total correction for form factor is quite small. The effect on the real part was less than 0.01.

[4]K. J. Foley, S. J. Lindenbaum, W. A. Love, S. Ozaki, J. J. Russell, and L. C. L. Yuan, Phys. Rev. Letters 11, 425 (1963).

[5]W. Galbraith, E. W. Jenkins, T. F. Kycia, B. A. Leontić, R. H. Phillips, A. L. Read, and R. Rubenstein, Phys. Rev. 138, B913 (1965).

[6]V. S. Barashenkov and V. I. Dedyu, Dubna Joint Institute for Nuclear Research Report No. JINR P-1598, 1964 (unpublished).

[7]G. Hohler, G. Ebel, and J. Giesecke, Z. Physik 180, 430 (1964).

[8]In general, this corresponds to several standard deviations from the best fit, and also would imply the physically unexpected result that the range of the interaction responsible is of the order of five times larger than that corresponding to the exchange of a pion.

BREAKING OF $U_{\mathcal{L}}(12)$ SYMMETRY*

Reinhard Oehme

The Enrico Fermi Institute for Nuclear Studies and the Department of Physics, The University of Chicago, Chicago, Illinois

(Received 15 April 1965)

In a recent paper[1] we have proposed a spurion theory of broken $U_{\mathcal{L}}(12)$ symmetry. In this scheme the formal, initial $U_{\mathcal{L}}(12)$ invariance[2] of an amplitude is not only broken by the imposition of Bargmann-Wigner equations,[3] which define the asymptotic "states" of physical particles, but also by spurions.[4] It is the purpose of this note to give a more detailed discussion of these spurions and to report certain consequences of the spurion theory for meson-baryon amplitudes.

We note that the spurion theory of $U_{\mathcal{L}}(12)$ should be considered as an approximation scheme for S-matrix elements and vertex functions, and not as an exact supermultiplet theory which complies a priori with all the axioms of field theory or of dispersion theory. It is well known that a relativistic theory which incorporates the successful features of SU(6) models encounters difficulties because of the noncompact nature of the Lorentz group.[5] In $U_{\mathcal{L}}(12)$ schemes, the impact of these difficulties is softened because the transformations are not directly applied to the physical amplitudes; rather, they act on formal amplitudes which are generalizations of the "spinor space amplitudes" used in dispersion theory.[6] There are, however, difficulties with the unitarity condition, even after imposition of the Bargmann-Wigner equations.[7,8]

A priori, we introduce the spurions in such a way that they preserve the Lorentz invariance of the amplitudes, transform as SU(3) singlets, and do not split the masses of particles with different spin values which are contained in the same SU(6) supermultiplet. We write the general spurion in the form

$$S=(\Gamma+\Gamma_5\gamma_5+\Gamma_\mu\gamma_\mu+\Gamma_{5\mu}i\gamma_5 +\tfrac{1}{2}\Gamma_{\mu\nu}\sigma_{\mu\nu})\otimes 1, \qquad (1)$$

and if used in first order, the coefficients are Lorentz covariants which must be constructed out of the independent momenta available in the amplitude into which S is being inserted. In the second and higher order, we can proceed in different ways. If we continue to construct the coefficients in Eq. (1) out of external momenta only, then the masses of the SU(6) supermultiplets remain degenerate.[1] On the other hand, in second order we may also construct terms of the form

$$S_2 = (s1\otimes 1 + s_p\gamma_5\otimes\gamma_5 + s_v\gamma_\mu\otimes\gamma_\mu + s_A i\gamma_\mu\gamma_5\otimes i\gamma_\mu\gamma_5 + s_T\sigma_{\mu\nu}\otimes\sigma_{\mu\nu})\otimes 1\otimes 1. \quad (2)$$

These spurions can give rise to mass differences between the SU(3) multiplets which are contained in a given SU(6) supermultiplet.

Liberal use of the spurions (1) and (2) will eliminate all restrictions due to the broken $U_{\mathfrak{L}}(12)$-invariance scheme, and we are left only with SU(3) and Lorentz invariance. On the other hand, if we use only the spurion S with momentum-dependent coefficients, then we obtain a scheme which reproduces many of the successful results of SU(6) models.[1]

As an example, let us consider briefly the amplitude for quark-quark scattering. In the limit of formal $U_{\mathfrak{L}}(12)$ symmetry, there are two terms:

$$\bar{u}_A(p')u^A(p)\bar{u}_B(k')u^B(k), \quad (3)$$

and

$$\bar{u}_A(p')u^A(k)\bar{u}_B(k')u^B(p), \quad (4)$$

where $k+p=k'+p'$. With the spurion S in first order, we have irreducible insertion $[i\gamma\cdot(k+k')\otimes 1 + 1\otimes i\gamma\cdot(p+p')]$ in Eq. (3) and an analogous one in Eq. (4). In second order the terms with $i\gamma\cdot(k+k')\otimes i\gamma\cdot(p+p')$ for Eq. (3) and $i\gamma\cdot(p+k')\otimes i\gamma\cdot(k+p')$ for Eq. (4) are added, and this is all we can get with the momentum-dependent spurion (1). The remaining insertions are $\gamma_5\otimes\gamma_5$ and $i\gamma_5\gamma\cdot(k+k')\otimes i\gamma_5\gamma\cdot(p+p')$ for Eq. (3) and corresponding forms for Eq. (4).[6] They can only be obtained with the second-order spurion S_2 given in Eq. (2), the inclusion of which leaves us with $P\otimes$ SU(3), where P is the Poincaré group.

It is important to note that our amplitude satisfies the elastic unitarity condition only in the SU(3) limit. We find that unitarity mixes the amplitudes of type (3) and of type (4), and, using the Fierz relations, we see that also those invariants are generated which are not present in the original amplitudes obtained from the momentum-dependent spurions (1).

In general, we cannot hope to get complete elastic unitarity with the spurion S alone. However, this should not detract from the possible usefulness of this spurion method as an approximate scheme. We may recall that most dispersion-theoretical approximations violate unitarity in some way, especially the pole-dominance models. In the spurion theory, we have the choice of various paths from formal $U_{\mathfrak{L}}(12)$ invariance to $P\otimes$SU(3) with mass splitting, and the usefulness for a specific amplitude of a given chain of approximations may well depend upon dynamical considerations.

Let us now consider meson-baryon amplitudes within the framework of the spurion scheme. We restrict ourselves to the octet of baryons, and sometimes also to pseudoscalar mesons.

The meson-baryon vertex has been discussed in reference 1 in connection with the electromagnetic form factors. Using only the spurion S with momentum-dependent coefficients, we found that almost all the favorable results of the $U_{\mathfrak{L}}(12)$ limit are retained. Furthermore, on the mass shell of the meson-baryon vertex, we have the same unique coupling as with formal $U_{\mathfrak{L}}(12)$ symmetry, except for a new term involving only the singlet vector meson. This term had been left out in reference 1. Evidently, the uniqueness of the couplings and the favorable d/f ratios will be destroyed if we introduce the second-order spurion S_2 given in Eq. (2).

With the spurion S, our results for the electromagnetic form factors are similar to those obtained by Bég and Pais[7] on the basis of a specific modification of the $U_{\mathfrak{L}}(12)$ scheme, at least as long as we do not insist upon a strict pole model. However, the method of these authors makes predictions for pseudoscalar-meson–baryon scattering which are quite different from those of our spurion scheme. We have pointed out in reference 1 that a first-order insertion of the spurion S makes possible polarization effects for a set of reactions for which formal $U_{\mathfrak{L}}(12)$ invariance predicts zero polarization.[8,9] In many cases this is also true for the ξ scheme mentioned by Bég and Pais. However, we have an important exception in the reaction

$$K^- + p \to \Xi^- + K^+. \quad (5)$$

If we write the amplitude for this process in the form

$$\bar{u}_{\Xi^-}(p')[A(s,t)-i\gamma\cdot\tfrac{1}{2}(k+k')B(s,t)]u_p(p), \qquad (6)$$

we find that we have $B(s,t)\equiv 0$ in the $U_{\mathcal{L}}(12)$ limit, whereas for most other reactions only $\mathrm{Im}AB^*\equiv 0$ is obtained.[9] In the ξ scheme mentioned above, the amplitude which contributes to Reaction (5) is given by

$$\bar{\Psi}_{ABC}(p')\Psi^{AB'C'}(p)\bar{\Phi}_{B'}{}^{B}(k')\Phi_{C'}{}^{C}(k), \qquad (7)$$

where

$$\Phi(k)=(\xi-i\gamma\cdot k/\mu)\gamma_5 P(k),$$
$$\bar{\Phi}(k')=(\xi'+i\gamma\cdot k'/\mu)\gamma_5 P(k'). \qquad (8)$$

Even if we assume that $\xi(s,t)$ and $\xi'(s,t)$ are arbitrary complex functions, we still find $B(s,t)\equiv 0$. In our theory with S spurions we have, in first order, the term

$$\bar{\Psi}_{ABC}(p')[i\gamma\cdot(k+k')]_{A'}{}^{A}\Psi^{A'B'C'}(p)$$
$$\times\bar{\Phi}_{B'}{}^{B}(k')\Phi_{C'}{}^{C}(k), \qquad (9)$$

where here $\bar{\Phi}$ and Φ are given by Eq. (8) with $\xi\equiv\xi'\equiv 1$. The amplitude (9) and the corresponding $U_{\mathcal{L}}(12)$-invariant expression (7) together give rise to independent complex functions $A(s,t)$ and $B(s,t)$ for Reaction (5), and this makes it possible to account for the observed polarization effects.[10]

Finally, we consider the annihilation of nucleon-antinucleon pairs at rest into mesons. The consequences of formal $U_{\mathcal{L}}(12)$ invariance have been discussed by several authors.[11] Due to the orthogonality of the spinor wave functions for particle and antiparticle with equal momenta, one finds that

$$\lim_{p\to 0}\bar{\Psi}_{ABC}(-p)\Psi^{AB'C'}(p)=0. \qquad (10)$$

This implies that in the $U_{\mathcal{L}}(12)$ limit there are no annihilations into two mesons. In our spurion scheme, with S used in first order, we have the unique amplitude

$$G_S\bar{\Psi}_{ABC}(-p)(i\gamma\cdot k)_{C'}{}^{C}\Psi^{A'B'C'}(p)$$
$$\times\bar{\Phi}_{A'}{}^{A}(k')\Phi_{B'}{}^{B}(k), \qquad (11)$$

where $\vec{p}=0$ and $\vec{k}=-\vec{k}'$ is the momentum of a final meson in the c.m. system. Because of $C(i\gamma\cdot k)^{T}C^{-1}=-i\gamma\cdot k$, we find that the amplitude (11) violates charge-conjugation invariance, and hence we must have $G_S=0$. We see that annihilation into two mesons is forbidden in our theory with momentum-dependent spurions S in first order. But a more detailed discussion shows that also in higher order many terms vanish because of charge-conjugation invariance. A suppression of two-meson annihilations seems to be quite compatible with experiments.

In the $U_{\mathcal{L}}(12)$ limit, and as a consequence of Eq. (10), we have a unique amplitude for the annihilation into three mesons; it is given by

$$\bar{\Psi}_{ABC}(-p)\Psi^{A'B'C'}(p)\bar{\Phi}_{A'}{}^{A}(k_1)$$
$$\times\bar{\Phi}_{B'}{}^{B}(k_2)\bar{\Phi}_{C'}{}^{C}(k_3), \qquad (12)$$

with $\vec{p}=0$. For nucleon-antinucleon systems, Eq. (12) gives rise only to the production of nonstrange particles and excludes φ mesons. With a spurion S, we obtain a new category of amplitudes which are of the form

$$\bar{\Psi}_{ABC}(-p)S_{C'}{}^{C}\Psi^{A'B'C'}(p)\bar{\Phi}_{A'}{}^{A}(k_1)$$
$$\times\bar{\Phi}_{D}{}^{B}(k_2)\bar{\Phi}_{B'}{}^{D}(k_3). \qquad (13)$$

These amplitudes could survive for $\vec{p}=0$ and may give rise to the production of strange mesons. However, our theory does not determine the coefficients of the amplitudes (13), and there may well be a suppression for dynamical reasons. One may also speculate that the annihilation process is mainly due to forces of very short range, and that for these forces the formal $U_{\mathcal{L}}(12)$ symmetry is a much better approximation than for those with longer range.

After this note had been written, I was informed about a related note by P. G. O. Freund which deals with the breaking of $U_{\mathcal{L}}(12)$ symmetry due to "kineton emission." I would like to thank Peter Freund for a discussion of his work. Thanks are also due to Chai S. Lai for checking some of my calculations.

---

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

[1]R. Oehme, Phys. Rev. Letters 14, 664 (1965).

[2]K. Bardakci, J. M. Cornwall, P. G. O. Freund, and B. W. Lee, Phys. Rev. Letters 14, 48 (1965); R. Delbourgo, A. Salam, and J. Strathdee, to be published.

[3]A. Salam, R. Delbourgo, and T. Strathdee, Proc. Roy. Soc. (London) A284, 146 (1965); M. A. B. Bég and A. Pais, Phys. Rev. Letters 14, 267 (1965); B. Sakita and K. C. Wali, Phys. Rev. Letters 14, 404 (1965); W. Rühl, Phys. Letters 14, 334 (1965) and to be published. These papers contain further references.

[4]Concerning the use of symmetry-breaking spurions, see also M. Gell-Mann, Phys. Rev. Letters 14, 77 (1965); K. Bardakci, J. M. Cornwall, P. G. O. Freund, and B. W. Lee, Phys. Rev. Letters 14, 264 (1965); R. Oehme, Phys. Letters 15, 284 (1965).

[5]See, for example, L. Michel and B. Sakita, Ann. Inst. Henri Poincaré (to be published). S. Coleman, to be published; S. Weinberg, to be published; L. O'Raifeartaigh, to be published.

[6]M. L. Goldberger, Y. Nambu, and R. Oehme, Ann. Phys. (N.Y.) 2, 226 (1957).

[7]M. A. B. Bég and A. Pais, Phys. Rev. Letters 14, 509 (1965).

[8]R. Blankenbecler, M. L. Goldberger, K. Johnson, and S. B. Treiman, Phys. Rev. Letters 14, 518 (1965).

[9]J. M. Cornwall, P. G. O. Freund, and K. T. Mahanthappa, Phys. Rev. Letters 14, 515 (1965).

[10]D. Carmony, G. Pjerrou, P. Schlein, W. Slater, D. Stork, and H. K. Ticho, Phys. Rev. Letters 12, 482 (1964).

[11]Y. Hara, Phys. Rev. Letters 14, 603 (1965); R. Delbourgo, Y. C. Leung, M. A. Rashid, and J. Strathdee, Phys. Rev. Letters 14, 609 (1965); Ngee-Pong Chang and J. M. Shpiz, Phys. Rev. Letters 14, 617 (1965). See also H. Harari and H. J. Lipkin, Phys. Letters 15, 286 (1965).

EVIDENCE FOR A NEW 2π RESONANCE AT 700 MeV*

M. Feldman, W. Frati, J. Halpern, A. Kanofsky, M. Nussbaum, S. Richert, and P. Yamin
University of Pennsylvania, Philadelphia, Pennsylvania

and

A. Choudry, S. Devons, and J. Grunhaus

Columbia University, New York, New York
(Received 6 April 1965)

In the course of a study of the production and decay modes of netural bosons produced in the reaction $\pi^- + p \rightarrow n +$ neutral particles, we have obtained evidence for a neutral boson of mass about 700 MeV which decays into two pions. The experiments were made at the AGS at Brookhaven National Laboratory using π beams of momenta up to several BeV/c, and the main evidence for the resonance here reported was observed with π momentum of 1.520 BeV/c. Our experimental procedure was to detect the π^- interacting in liquid hydrogen with no charged-particle production, but with the production of a neutron whose energy is measured by the time-of-flight technique. The two-body reaction is indicated by a characteristic neutron energy which is a function of the π^- energy and the angle of detection of the neutron as well as of the mass of the neutral particle produced. The decay of the latter is observed by a large solid-angle array of γ-sensitive spark chambers.

The apparatus is shown schematically in Fig. 1. The interacting π^- is signaled by the response $\bar{1}23\bar{4}$ in the plastic scintillator, and the neutron by the delayed $\bar{6}5$. The hydrogen target is 7.5

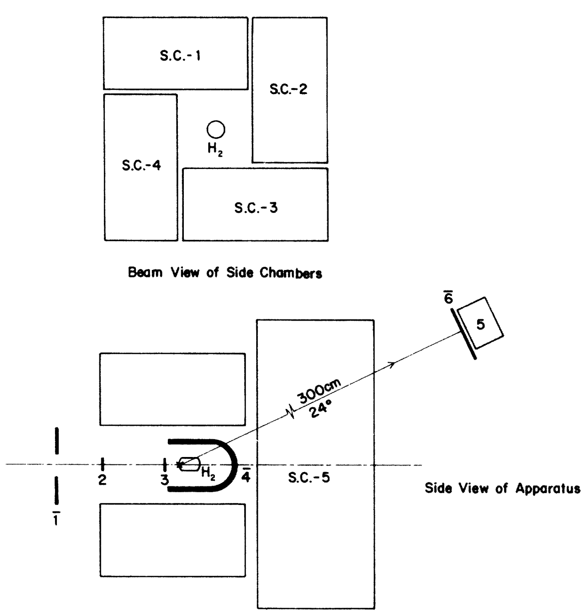

FIG. 1. Schematic cross section of apparatus. The spark chambers have at least 4 radiation lengths of lead plates in all directions, while the neutron counters consist of 20 plastic scintillators about 7 in. in diameter and 4 in. thick coupled to 58AVP photomultipliers.