is the addition of relatively slowly varying form factors at the $\omega\pi\gamma$ and πNN vertices.

We believe that plane-polarized photons will provide the most decisive measurement of the relative contributions of the two processes.

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Current Algebra and the Suppression of Resonances with Baryon Number 2

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The members of the $\overline{10}$ $SU(3)$ multiplet, into which the deuteron (d) would have been classified in the limit of unitary symmetry, were not observed in experiments (apart from the deuteron itself). We show that current algebra and the hypothesis of partially conserved axial-vector current (PCAC) imply a suppression of the λ^0K^-d strong vertex, where λ^0 is a dibaryon state in the aforementioned $\overline{10}$. We therefore suggest that the other particles of that multiplet were not observed because of their small coupling to the states being used in the experiments.

I. INTRODUCTION

In the last two years, the hypothesis of partially conserved axial-vector current (PCAC)[1,2] and current-algebra[3] (CA) have been extensively used to calculate weak-current renormalization constants,[4] electromagnetic properties of baryons,[5,6] decay widths of vector mesons,[7,8] and scattering lengths of meson-hadron interactions,[9,10] and to relate processes which differ by the emission or absorption of soft pions.[11–13]

Those calculations generally confirm the $SU(3)\times SU(3)$ structure of the vector and axial-vector weak currents, and the PCAC assumption for pions. However, the PCAC hypothesis for kaons shows only a qualitative agreement with experiment, while quantitatively the results obtained by using this assumption do not agree with experiment so well as those obtained from PCAC for pions. Thus it is desirable to find further tests of the generalized PCAC.

In this work, we wish to make use of CA and PCAC for kaons, in order to explain the experimental nonexistence of bound states or resonances with $B=2$, $S=-1$, which would have to exist in the limit of unitary symmetry because of the existence of the deuteron. Our calculations serve at the same time as a test for the PCAC assumption for kaons.

In Sec. II we define our problem, describing the relations between the deuteron and the possible resonances mentioned above. We then derive an Adler-Weisberger–type sum rule which relates $K^\pm$-on-deuteron cross sections with a coupling constant measuring the symmetry-breaking effects which cause the nonappearance of the resonances.

In Sec. III we describe the numerical evaluation of the sum rule, and the various models used in the energy regions where no experimental data are available. In Sec. IV we discuss our results.

II. DERIVATION OF THE SUM RULE

In the limit of unitary symmetry for the strong interactions, the deuteron (d) has to be a member of an $SU(3)$ multiplet of states with baryon number 2. Since the $SU(3)$ content of the d wave function is a two-nucleon state, d has to be classified in a multiplet which is obtained by reduction of the product 8×8. Since the only state with $Y=2$ and $T=0$ contained in 8×8 is the upper vertex of the $\overline{10}$, the $SU(3)$ classification of d is unique.

Applying the $SU(3)$ generator that has the quantum numbers of the K^- on $|d\rangle$, we arrive at a state which we shall denote by $|\lambda^0\rangle$, and which has an $SU(3)$

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content

$$|\lambda^0\rangle=(\sqrt{\tfrac{1}{6}})\{(\sqrt{\tfrac{3}{2}})[\Lambda^0(1)n(2)-n(1)\Lambda^0(2)] +(\sqrt{\tfrac{1}{2}})[\Sigma^0(1)n(2)-n(1)\Sigma^0(2)] -[\Sigma^-(1)p(2)-p(1)\Sigma^-(2)]\}. \quad (2.1)$$

In the limit of unitary symmetry, λ^0 would have been a bound state of a hyperon-nucleon system, with the same mass and binding energy as those of d.

The symmetry-breaking part of the strong interactions may change the binding energy of λ^0, turn it into a resonance (move it off the real axis), change its coupling constants to other states from the values appropriate to unitary symmetry, or even make it nonexisting.

The experimental evidence concerning the existence of λ^0 as a bound state or resonance is not definite.[14,15] In K^--nucleus capture reactions several peaks in $\Lambda^0 N$ effective mass have been observed, but most may be explained by effects other than the existence of λ^0 or λ^+ (the particle obtained by acting with T_+ on λ^0; this λ^+ should have been a member of the same $\overline{\mathbf{10}}$). In hyperon-nucleon final-state interactions no evidence for λ^0 is found. The low-energy Λ^0-nucleon scattering parameters show no existence of resonance or bound state in this system (though an apparent bump in the low-energy Λ^0-p cross section, which might as well be a statistical fluctuation, has recently been observed.[16])

It is clear, therefore, that even if λ^0 exists, it is weakly coupled to other states, since otherwise it would manifest itself in the low-energy scattering of systems where it can contribute (as the d does in p-n scattering). We wish to show that this situation is necessitated by CA and PCAC.

For definiteness we look at the coupling constant of the strong-interaction vertex $\lambda^0 K^- d$, $g_{\lambda^0 K^- d}$. This coupling constant is directly responsible for the formation of λ^0 in K^-d interactions, which seems to be the natural place to look for λ^0 in experiment, λ^0 and d being in the same multiplet, and connected by the K^- field. The sum rule has, however, a more basic importance, since via PCAC the K^- field is related, roughly speaking, to the axial charge Q^{A,K^-}; in the limit of $SU(3)\times SU(3)$ chiral symmetry, $Q^{A,K^-}|d\rangle$ is just one state in the $SU(3)\times SU(3)$ multiplet in which d is a member. Thus the sum rule measures the leakage of $Q^{A,K^-}|d\rangle$ from the above multiplet, i.e., the breaking of the above symmetry.

To obtain a sum rule for $g_{\lambda^0 K^- d}$ we start from the $SU(3)\times SU(3)$ commutator

$$[Q^{A,K^+},Q^{A,K^-}]=\tfrac{3}{2}Y+T_0\,, \quad (2.2)$$

and from the PCAC assumption

$$\partial^\mu J_\mu{}^{A,K^-}=c\phi^{K^-}. \quad (2.3)$$

Using the method of Fubini and Furlan,[17] i.e., taking the matrix element of (2.2) between d states of momenta p, p' and polarization vectors ϵ, ϵ', and then separating explicitly the λ^0 intermediate-state contribution in the commutator, we arrive at the sum rule

$$3=\frac{2}{3}\frac{1}{M_d M_{\lambda^0}}(g_{\lambda^0 K^- d})^2\frac{C^2}{m_K{}^4}G^2(2(M_d{}^2+M_{\lambda^0}{}^2),0) +\frac{C^2}{m_K{}^4}\frac{2}{\pi}\int' dw\frac{w}{w^2-M_d{}^2}[\sigma_0{}^-(w)-\sigma_0{}^+(w)]. \quad (2.4)$$

In this sum rule, M_d and M_{λ^0} are the masses of d and λ^0; $g_{\lambda^0 K^- d}$ and G are the coupling constant and form factor defined for the $\lambda^0 K^- d$ vertex as follows:

$$\langle d(p_d,\epsilon_d)|\phi^K|\lambda^0(p_\lambda,\epsilon_\lambda)\rangle =\left(\frac{1}{2p_d{}^0}\right)^{1/2}\left(\frac{1}{2p_{\lambda^0}{}^0}\right)^{1/2}\frac{1}{q^2-m_K{}^2}g_{\lambda^0 K^- d} \times G(Q^2,q^2)\frac{1}{(M_d M_{\lambda^0})^{1/2}}\epsilon_{\mu\nu\rho\sigma}q^\mu Q^\nu \epsilon_d{}^{*\rho}\epsilon_{\lambda^0}{}^\sigma, \quad (2.5)$$

where $Q=p_d+p_{\lambda^0}$, $q=p_d-p_{\lambda^0}$, and $G=1$ when $\mathbf{p}_d=\mathbf{p}_{\lambda^0}=\mathbf{0}$. We let G depend on Q^2 and q^2 only, since in the derivation we encounter only p_{λ^0} and p_d's that satisfy $\mathbf{p}_d=\mathbf{p}_{\lambda^0}$. $\sigma_0{}^\pm(w)$ are the total cross sections for scattering of zero-mass $K^\pm$ on d at c.m. energy w. The integration over w for $\sigma_0{}^+(w)$ ranges from $w=m_K+M_d$ to $+\infty$, for $\sigma_0{}^-(w)$ from $w=M_N+M_{\Lambda^0}$ to $+\infty$, and not including the δ-function contribution at $w=M_{\lambda^0}$, a fact denoted by the prime on the integral sign. The mass factors in the definition (2.5) were chosen to make $g_{\lambda^0 K^- d}$ have the same dimensions as $g_{\Lambda^0 KN}$ (which is the coupling constant of the strong-interaction Yukawa Hamiltonian $g_{KN\Lambda}\bar{N}\gamma_5 K\Lambda+$H.c.) so that they can be compared; they are coupling constants of the same K^- field within different $SU(3)$ multiplets.

We now put[18]

$$\frac{C}{m_K{}^2}=\frac{g_A{}^\Lambda(M_N+M_{\Lambda^0})}{g_{KN\Lambda}F_{KN\Lambda}(0)}, \quad (2.6)$$

where $g_A{}^\Lambda$ is the weak-coupling constant of p to Λ^0, $g_{KN\Lambda}$ and $F_{KN\Lambda}$ are the strong-coupling constant and form factor for the vertex $KN\Lambda^0$. Note that this is independent of any assumption on the existence of a definite D/F ratio.

Inserting (2.6) into (2.4) and using numerical values for the masses and for $g_A{}^\Lambda$ (we assume $M_{\lambda^0}\sim 2.1$ BeV),

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we obtain

$$3=0.45\left(\frac{g_{\Lambda^0 K^- d}}{g_{\Lambda KN}}\right)^2 \frac{G^2(2(M_d^2+M_{\Lambda^0}^2),0)}{F_{\Lambda KN}^2(0)} +\left(\frac{g_A{}^{\Lambda}(M_N+M_{\Lambda^0})}{g_{\Lambda KN}F_{\Lambda KN}(0)}\right)^2 \frac{2}{\pi}\int' \frac{dww}{w^2-M_d^2} \times[\sigma_0{}^-(w)-\sigma_0{}^+(w)]. \quad (2.7)$$

In order to evaluate numerically the second term of the right-hand side of (2.7), we need the value of $g_{KN\Lambda^0}$, which is not well known from experiment. Assuming at this stage a definite D/F ratio f, we have[19]

$$g_{KN\Lambda}=-\frac{1+2f}{\sqrt{3}}g_{\pi NN}. \quad (2.8)$$

We take the value $f=\frac{1}{4}$,[19] but bear in mind that from calculations of f using CA and PCAC[18,20] some higher values were obtained, so that using these values for f, we get for the second term of (2.7) values smaller by about 15%.

With $f=\frac{1}{4}$ we have

$$3=0.45\left(\frac{g_{\Lambda^0 K^- d}}{g_{\Lambda KN}}\right)^2 \frac{G^2(2(M_d^2+M_{\Lambda^0}^2),0)}{F_{\Lambda KN}^2(0)} +\frac{3.2\times10^{-2}}{F_{\Lambda KN}^2(0)}\frac{2}{\pi}\int' dw\frac{w}{w^2-M_d^2} \times[\sigma_0{}^-(w)-\sigma_0{}^+(w)], \quad (2.9)$$

where masses are measured in BeV and $\sigma_0{}^{\pm}$ in mb.

III. NUMERICAL EVALUATION

Experimental total cross sections for $K^{\pm}$-d scattering are available in the range 0.6–20-BeV/c lab momentum of the kaon[21–25] and two points for K^+d at 0.37 and 0.53 BeV/c.[26] We have then to use models for $\sigma^{\pm}$ in three regions: high energy, near threshold, and in the unphysical region for the $\sigma_0{}^-$ integration.

We make no dynamical extrapolation off the mass shell for the physical cross sections, but we do make a kinematical correction. Using the optical theorem, one can write

$$\sigma_0{}^{\pm}(w)=(4\pi/K_0)\,\mathrm{Im}T_0{}^{(\pm)}(w,\theta=0), \quad (3.1)$$

where K is the c.m. momentum at c.m. energy w of the Kd system, and the subscript 0 denotes quantities that refer to 0-mass kaons. The factor $1/K$ is of a kinematical nature, resulting from expressing the incident flux and final phase space in the elastic channel in terms of K, while T is the dynamical amplitude obtained, say, by the use of the Lehmann-Symanzik-Zimmermann (LSZ) contraction scheme.

We define a kinematically corrected cross section $\tilde{\sigma}_0$ to be used in the sum rule by

$$\tilde{\sigma}_0{}^{\pm}(w)=(K/K_0)\sigma^{\pm}(w), \quad (3.2)$$

assuming that $T(w)\sim T_0(w)$. Such a correction is essential near the K^-d threshold, since the existence of open exothermic channels causes a finite T matrix even at threshold, so that the physical $\sigma^-(w)$ is divergent at $w\rightarrow M_d+M_{K^-}$. Our correction makes $\tilde{\sigma}_0{}^-$ finite there, as it should.

To describe the low-energy K^-d scattering we use the model of Dalitz and Chand.[27,28] The model assumes independent interactions of K^- with the two nucleons in d, but takes into account interference effects and multiple scattering, and also the internal structure of the deuteron. For the $\bar{K}N$ interactions at low energies, a boundary-condition model is used[29] with a zero-range approximation. As proved by Dalitz and Tuan,[30] this description of the scattering is equivalent to a K-matrix formalism with the standard energy dependence of the K elements.

We have used a modified version of the model which takes into account the K^--$\bar{K}^0$ mass difference, and which is described in the Appendix of Ref. 28. For that we solved the 6×6 system of equations given in Ref. 28 by Kramer's method, so that angular integrations could be performed analytically, and the remaining radial integration was done by a computer. We have also not used the zero-range approximation, but rather the energy-dependent scattering lengths of Kittel.[31] We checked that if we use Kim's[32] parameters I instead, the results do not change significantly. Keeping the kinematical correction in mind, we used the physical m_K when calculating the amplitude $f(\theta)$ by Chand's model, but for the $1/K$ in the optical theorem we took $1/K_0$.

One expects this model to be valid only at the low-energy region, and in fact, Chand obtained agreement both for capture of K^- on d at rest and for scattering at lab momentum 0.2–0.3 GeV/c, while we obtain disagreement with experiment at 0.6 GeV/c. In the region 0.3–0.6 GeV/c the K^-p cross sections are smooth, apart from the narrow $Y^*(1520)$ which occurs at 0.4 GeV/c.

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Thus in the region 0.3–0.6 GeV/c we made a linear extrapolation between the predicted σ_{total} at 0.3 GeV/c and the experimental σ_{total} at 0.6 GeV/c to serve as background, and the $Y^*(1520)$ contribution was calculated in the zero-width approximation described below.

To calculate the low-energy K^+d scattering, we used once more Chand's model with the following changes: (1) Since the charge exchange in K^+d occurs in the K^+n interaction while in K^-d it occurs in the K^-p interaction, the charge-exchange threshold is at different lab momentum of the K. (2) The isospin content of the initial and possible final wave functions is different in the two cases, and therefore the numerical values of coefficients in the 6×6 system of equations in the Appendix of Ref. 28 had to be recalculated. We used for the calculations the KN scattering lengths quoted by Stenger *et al.*[26]

In K^+d interactions, one finds a p-wave contribution even at low energies,[26] and thus we used Chand's model only up to 0.2 GeV/c. We then made interpolations using these predictions and the experimental points at 0.37, 0.53, and 0.6 GeV/c. The results are essentially equal to those obtained by making an extrapolation using the experimental points only.

To account for the high-energy contribution to the sum rule, we used a Regge-pole parametrization. In the crossed channel $K^-+K^+\to \bar{d}+d$, an ω trajectory can be exchanged, and since the ω has a negative signature, it will contribute to the difference $(\sigma_{K^-d}-\sigma_{K^+d})$ as well. The intercept $\alpha(0)$ for the ω is 0.5,[33] and thus

$$\Delta\sigma(w)=Aw^{2[\alpha(0)-1]}=A/w\,, \tag{3.3}$$

$$\int_{w_0}^{\infty}\frac{\Delta\sigma(w)wdw}{w^2-M_d^2}=\Delta\sigma(w_0)\frac{w_0}{2M_d}\ln\frac{w_0+M_d}{w_0-M_d}. \tag{3.4}$$

The $Y^*(1520)$ and the $Y^*(1405)$, $Y^*(1385)$ (which we assume to dominate in the unphysical region) were calculated using the impulse approximation for K^-d interactions, and the zero-width approximation for the K^-N scattering amplitude. We write for a resonance contribution

$${}_{\text{res}}\sigma_{\text{tot}}{}^{K^-d}=(4\pi/K_0)\,\text{Im}f^{K^-d}(\theta=0)\,, \tag{3.5}$$

$$\text{Im}f^{K^-d}(\theta=0)=\text{Im}f^{K^-n}(\theta=0)+\text{Im}f^{K^-p}(\theta=0)\,. \tag{3.6}$$

The kaon-nucleon amplitudes f^{K^-N} are expanded in partial waves appropriate to spin-$\frac{1}{2}$–spin-0 scattering, and for the partial-wave amplitudes we make the pole approximation

$$\text{Im}f_{l,j}=\tfrac{1}{2}\pi\alpha\delta(w_1-M_{l,j})\,, \tag{3.7}$$

where l, j, are the quantum numbers of the resonance, $M_{l,j}$ its mass, w_1 is the K^-N c.m. energy corresponding to K^-d c.m. energy w, and α is a constant.

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For the resonances with $j=\frac{3}{2}$, α is related to the coupling constant H of the vertex KNY^* via[34]

$$\frac{H^2}{4\pi}=\frac{3(M_{l,j}M_N)^2}{(M_{l,j}+M_N)^2-m_K{}^2}\frac{\alpha}{q^2}, \tag{3.8}$$

where q is the K^-N c.m. momentum when $w_1=M_{l,j}$. For the $Y^*(1405)$ we have

$$\alpha=2M_{l,j}|q|/E\omega\,, \tag{3.9}$$

where q, W, ω are the unphysical c.m. momentum, nucleon energy, and kaon energy, at the resonance energy. We used the coupling constants given by Warnock and Frye.[34]

To calculate the $Y^*(1520)$ contribution, we used directly the above method with the physical mass of K to obtain the amplitude.

For the $Y^*(1405)$ that falls in the unphysical region, we used also the physical mass to calculate the amplitude. Here, however, we have to subtract the "tail" of the resonance that lies in the physical region, since there we took σ from experiment (or the model of Chand) including the $Y^*(1405)$ formation. To do that we assumed a Breit-Wigner form for the resonance, and calculated the part of the integrated cross section that lies in the physical region. One has to remember that the width of the resonance in the K^-d channel is larger than the width in K^-N interactions, since the transformation $w\to w_1$ has a Jacobian different from 1. It can be shown that the width is enlarged by the factor M_dw_1/M_Nw.

For the $Y^*(1385)$ contribution, use of (3.8) with the physical K mass leads to a negative α and therefore negative σ_{tot} in (3.5). We made here, therefore, a dynamical correction to account for the mass extrapolation, by calculating α from (3.8) taking $m_K=0$. Because the energy of the $Y^*(1385)$ was not changed while m_K was set to zero, the resonance energy in the K^-d scattering was shifted. However, since this was the only shift of an intermediate-state energy in our calculations, when we calculated the part of the $Y^*(1385)$ that lies in the physical region in order to subtract it, we treated the resonance as if its energy were the one obtained when the physical mass is used. In the $Y^*(1385)$ case, there are contributions both from K^-n and K^-p reactions (unlike the former two resonances), and we applied a Glauber correction[35] to the cross section. (For a zero-mass K, this resonance is in the high-energy domain).

Collecting all contributions, we obtain for the second term in the right-hand side of (2.9) [assuming $F_{KN\Lambda}(0)\sim 1$] the value 3.02, out of which 1.82 comes from the resonances, 0.90 from the energy region 2.65 to $+\infty$

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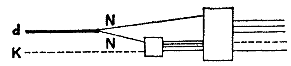

FIG. 1. Processes which give most of the contribution to the sum rule.

GeV, and 0.29 from the low-energy part from threshold up to 2.65 GeV.

One can argue that another kinematical correction is needed to account for the threshold behavior in the K^+d channel. Near the physical threshold, the physical amplitude T is proportional to K, and therefore vanishes at threshold (no open exothermic channels exist there). However, the zero-mass T_0 is not proportional to K at the physical threshold and does not necessarily vanish there. One can show that eliminating the proportionality to K from T_0 leads to the assumption that at low energies $\sigma^0 \sim \sigma$ rather than $T_0 \sim T$. Thus we have to introduce into the sum rule σ and not $\bar{\sigma}_0$ at low energies, and $\bar{\sigma}_0$ at higher energies. Making the transition between σ and $\bar{\sigma}_0$ at $p_{\text{lab}} = 0.150$ GeV/c, and applying an extrapolation to smooth the integrated σ, we obtain 2.99 for the above term; if we make the transition at $p_{\text{lab}} = 0.450$ GeV/c, we obtain 2.95.

The error estimate is difficult because of the use of models in energy regions that contribute much to the sum rule. In the region $w = 2.65$–8.5 GeV we used experimental data with 1% error. In the high-energy region the K^-d cross section may not yet have reached its asymptotic region (the K^+d clearly has), and we estimate the error in this region by the experimental error in $\Delta\sigma(w_0)$. In the low-energy region we take the differences obtained in using various sets of scattering lengths to estimate the error. For the $Y^*(1520)$ and $Y^*(1405)$ contributions we take 5%, and for the $Y^*(1385)$ we take 10%. Althogether we have an error of 0.20, of which 0.12 comes from the $Y^*(1385)$.

IV. DISCUSSION

One does not have to take very seriously the exact numerical value of the integral which was obtained, since (1) we did not make very good estimates of the cross sections at low energy and for regions below threshold where most of the contribution comes from ; (2) the value of $|g_{\Lambda^0 KN}|$ seems to be somewhat larger than the value appropriate to $f = \frac{1}{4}$ (which we used), making the integral smaller; and (3) extrapolation to zero mass tends to enlarge the value of the integral, as found by Adler[4] in his renormalization calculation.

However, qualitatively it is quite clear that the sum rule is exhausted by the integral on the right-hand side, leaving practically no contribution to the λ^0 term.

To give a quantitative analysis of our results, we list the values of $(g_{\lambda^0 K^- d}/g_{\Lambda KN})^2$ that are obtained for different values of the integral that appears in the sum rule:

Integral	3.0	2.9	2.8	2.7	2.6	⋯	0
$(g_{\lambda^0 K^- d}/g_{\Lambda KN})^2$	0	0.22	0.44	0.66	0.88	⋯	6.6

We then see that if we use the coupling constants as a measure for the strength of the coupling, our results indicate that the $\lambda^0 K^- d$ vertex is suppressed relative to the analogous ΛKN vertex of the K^- field within the baryon octet.

However, since the two coupling constants are related to different types of interaction Hamiltonians, it seems that a better measure of the symmetry breaking is the relative contributions of the λ^0 pole term and of the integral to the value 3 of the sum rule. For this criterion our results clearly indicate suppression of the λ^0.

The essential point about the sum rule is that it limits the total strength of transitions that are induced by the K^- field on the deuteron. In the physical region and just below it those transitions are dominated by impulse-type processes in which the K^- interacts strongly with one of the nucleons (with subsequent final-state interaction) and which are described by the diagrams of the form shown in Fig. 1. The large contributions from such processes arise from the loose binding of the deuteron, leading to a large coupling constant for the vertex dpn, and from the strong K^--nucleon interaction. Those two properties of the K^-d system lead—via the sum rule— to the smallness of the coupling constant of the $\lambda^0 K^- d$ vertex, which is confirmed by experiment.

Alternatively, if the symmetry-breaking effects are large enough to cause nonexistence of λ^0, the right-hand side of the sum rule contains only the integral term, and thus we obtain an Adler-Weisberger sum rule which agrees very well with experiment and serves as a positive test for the PCAC assumption for kaons.

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