TABLE IV. Summary of predictions of IVB Model A for $\Delta Y=1$ hadron decays emitting lepton pairs.

$\Delta Y=1$ decay process	Order of semiweak plus electromagnetic matrix element	Predicted ratio of $CP=-1$ to $CP=+1$ matrix elements	Predicted branching ratio due to dominant matrix element
$K^+\to\pi^+l\bar{l}$	g^2e^2	0.3	0.5 to 8.5×10^{-7}
$K_L^0\to\mu\bar{\mu}$	g^2e^4	~2	$\sim5\times10^{-9}$
$K_L^0\to\pi^0l\bar{l}$	g^2e^4	30	0.5 to 1.1×10^{-6}
$\Sigma^+\to pe\bar{e}$	g^2e^2	0.03	3×10^{-9}

the third column of Table IV require some explanation. The only decay in Table IV which is explained in a straightforward fashion is $\Sigma^+\to pe\bar{e}$ decay, where we expect the ratio of the g^3 to g^2e^2 matrix elements to be of order $g/e^2\sim\frac{1}{10}$ and this is confirmed. On the other hand, the estimated ratio of the two matrix elements for $K^+\to\pi^+l\bar{l}$ decay is larger by a factor of 10, an effect which follows from the relative suppression of the $K^+\pi^+\gamma$ vertex with respect to the $\Sigma^+p\gamma$ vertex (the former vertex vanishes for a real photon, whereas the latter does not). The ratio of the matrix elements for $K_L^0\to\pi^0ll$ is also "abnormal," since one expects $g^3/g^2e^4\sim1$. The enhancement of the g^3 with respect to the g^2e^4 matrix element finds its explanation in the fact that the $CP=-1$ and $+1$ matrix elements are associated with the $^1S_0$ and $^3P_0$ states, respectively. The explanation of the number for $K_L^0\to\mu\bar{\mu}$ decay has already been given in terms of a hypothesized amount ($\sim10\%$) of broken SU_3 symmetry.

Quite apart from explaining the numbers listed in Table IV, the most favorable candidates for the detection of gross CP violation appear to be the first two decays. The last two decays seem equally unfavorable, except that to measure electron-pair correlation from polarized Σ^+ decay may be easier than to measure muon-pair spin correlation from $K_L^0\to\pi^0l\bar{l}$ decay. As a practical test of the strong cubic IVB model A, it is suggested that, first, the existence of the induced neutral-lepton-current decay $K^+\to\pi^+\nu\bar{\nu}$ must be established at the 10^{-8}–10^{-7} level, i.e., at the g^3 rather than the g^4 level in the matrix element, in order to justify the laborious search for gross CP violation in $K^+\to\pi^+l\bar{l}$ or $K_L^0\to\mu\bar{\mu}$ decay. Failure to observe any of the $\Delta Y=1$ induced neutral-lepton-current decay of hadrons in order g^3 would eliminate IVB model A but not B. The latter could then be tested by searching for gross CP violation in direct W production by neutrinos on nucleons. The observation of gross CP violation either in accordance with the predictions of IVB model A or B, would serve as a striking confirmation of the idea that CP violation is not a small accident, but must be incorporated into the basic structure of weak-interaction theory at the semiweak level through the mechanism of a set of strong self-interacting intermediate vector bosons (the "fifth force" in nature).

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Field-Theoretic Model for Low-Energy $J=\frac{3}{2}^+$ K^+p Scattering

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Assuming broken SU_3 Yakawa coupling between the octet baryons and pseudoscalar octet mesons, we discuss low-energy K^+p scattering. We obtain a broad ($\Gamma=1000$ MeV) resonance in the $J=\frac{3}{2}^+$ channel with mass 1750 MeV. The phase shift does not reach $\pi/2$ in the resonance region. The calculated cross section is in qualitative agreement with the experimental data near the resonance energy.

I. INTRODUCTION

In two recent papers[1,2] (subsequently referred to as I and II), we have treated low-energy elastic π^+p and $\bar{K}\Xi$ scattering, assuming Yukawa coupling between the octet baryons and pseudoscalar octet mesons and using the SU_3 values for the coupling constants[3] with the phenomenologically determined value[4] $\alpha=0.75$ for the F/D mixing parameter. We used the physical values for the masses. Virtual baryon-antibaryon pair-creation processes were neglected. An approximate expression for the off-energy-shell T matrix was obtained by summing to infinite order those terms in the perturbation expansion which did not contain more than two

[1] L. B. Rédei, Nucl. Phys. **B10**, 419 (1969).

[2] L. B. Rédei, Nucl. Phys. **B17**, 38 (1970).

[3] See, e.g., S. G. Gasiorowicz, *Elementary Particle Physics* (Wiley, New York, 1966).

[4] G. Ebel *et al.*, in *Proceedings of the Lund Conference on Elementary Particles*, edited by G. von Dardel (Berlingska Boktryckeriet, Lund, Sweden, 1969).

mesons in their intermediate states. It was found that the attractive $J=\frac{3}{2}^+$ amplitude was only logarithmically divergent (because we have taken into account the recoil of the baryon) and this made it possible to obtain a finite result in the infinite-cutoff limit after coupling-constant renormalization. The other nonresonant amplitudes vanished in this limit owing to the vanishing of the unrenormalized coupling constant (alternatively, one could say that they remained cutoff-dependent even after renormalization). We adjusted the pion-nucleon coupling constant g to the experimental value for the π^+p $J=\frac{3}{2}^+$ scattering length, giving $g^2/4\pi=23.25$. Having fixed the value of g, we obtained 1368 and 1976 MeV for the masses of the $N^*_{\frac{3}{2}\frac{3}{2}}$ and Ω^- which appeared, respectively, as a $Y=1$, $I=\frac{3}{2}$, $J=\frac{3}{2}^+$ resonance and a $Y=-2, I=0, J=\frac{3}{2}^+$ bound state, to be compared with the experimental values $g^2/4\pi=15$, $M_{N^*}=1238$ MeV, and $M_{\Omega^-}=1675$ MeV.

The purpose of the present paper is to treat the elastic $J=\frac{3}{2}^+$ K^+p exotic channel by the same method. We use $g^2/4\pi=23.25$ and $\alpha=0.75$ as in I and II. There will be no other parameters in our calculations.

The question of the possible existence of an exotic K^+p resonance has aroused considerable interest in recent years. As is well known, such a resonance could not be thought of as built up from three quarks and would require a modification of the quark model, at least in its simple form. For a recent discussion of the implications of the possible existence of such exotic resonances for the quark model and for the concept of exchange degeneracy see, e.g., the review talk by Lipkin.[5]

Let us first summarize the experimental features relevant for this work. A bump in the total cross section for K^+p scattering is now well established[6,7] at $q_{\mathrm{lab}}=1.2$ GeV/c meson momentum. The height of the bump is about 18.5 mb. The cross section falls off rather slowly to the right of the maximum. The elastic cross section has also been measured, its maximum value being about 14 mb. The position of the maximum does not seem to be settled yet, owing to the experimental uncertainties in the low-energy cross-section data.[8] The differential cross section in the region $0.7\leq q_{\mathrm{lab}}\leq 0.9$ GeV/c is symmetric around $\cos\theta=0$ and shows a fairly strong $\cos\theta$ dependence indicating the presence of a substantial $(\cos\theta)^2$ term.[9] There seem to be no polarization data available at low energies.[10]

A number of theoretical works have previously predicted the existence of a $J=\frac{3}{2}^+$ K^+p resonance. These are, e.g., the strong-coupling model,[11] the N/D bootstrap method,[12] and a calculation based on a relativistic Schrödinger equation[13] where the potential is taken to be an off-energy-shell extrapolation of the SU_3 Yukawa-Lagrangian-Born term. However, in all these calculations it is implied that the $J=\frac{3}{2}^+$ phase shift passes through $\pi/2$ at resonance energy, and since they work within the elastic approximation they fail to provide an explanation for the fact that the experimental cross section in the resonance region is much smaller than the value given by the elastic unitarity limit $8\pi/q^2$.

The main result of our calculations is that the present model gives a rapidly rising phase shift in the $J=\frac{3}{2}^+$ elastic K^+p channel, but the phase flattens off before reaching $\pi/2$. The derivative of the phase shift has a maximum indicating the existence of a very-short-lived resonance structure giving rise to a broad maximum in the elastic cross section.

II. DESCRIPTION OF MODEL

In this section we describe briefly the chain of approximations we used to get an approximate expression for the elastic $J=\frac{3}{2}^+$ K^+p scattering amplitude. For details Refs. 1 and 2 can be consulted.

We start from the SU_3 interaction Hamiltonian[3]

$$H_{\mathrm{int}}=i\sqrt{2}g_0\!:\int d^3x \times\{(2\alpha-1)\bar{B}_j{}^i\gamma_5B_k{}^j\phi_i{}^k+\bar{B}_k{}^j\gamma_5B_j{}^i\phi_i{}^k\}:\,, \quad (1)$$

where g_0 is the unrenormalized pion-nucleon coupling constant; the $B_j{}^i$ and $\phi_i{}^k$ represent the octet 3×3 SU_3 matrices built out of the octet baryon octet pseudoscalar meson fields, respectively. α is the F/D mixing parameter for which we used the experimental value[4] $\alpha=0.75$. The physical masses are used throughout, except that for the Λ and Σ particles we used the mean value of their masses.

For the description of low-energy phenomena it seems reasonable to neglect those terms on the right-hand side of Eq. (1) which give rise to virtual baryon-antibyron pair-creation processes. This way we obtained the model interaction Hamiltonian V:

$$V=V_++V_-\,,\quad V_-=(V_+)^\dagger\,, \quad (2)$$

the operator V_+ being given by

$$V_+=i\sqrt{2}g_0\int d^3x[(2\alpha-1)\bar{B}^{(+)}{}_j{}^i\gamma_5B^{(-)}{}_k{}^j\phi^{(+)}{}_i{}^k +\bar{B}^{(+)}{}_k{}^j\gamma_5B^{(-)}{}_j{}^i\phi^{(+)}{}_i{}^k]\,, \quad (3)$$

where the $B^{(-)}$ ($\bar{B}^{(+)}$) annihilates (creates) one baryon

[5] H. Lipkin, rapporteur's talk, in Ref. 4.
[6] For a recent measurement see R. L. Cool *et al.*, Phys. Rev. D **1**, 1887 (1970).
[7] D. R. O. Morrison, rapporteur's talk, in Ref. 4.
[8] For a recent summary of the available experimental data on low-energy K^+p scattering, see, e.g., A. D. Martin and B. Perrin, Nucl. Phys. **B10**, 125 (1969).
[9] See, e.g., Martin and Perrin, Ref. 8, and also Morrison, Ref. 4.
[10] For a review of the experimental situation and references see D. R. O. Morrison, Ref. 7. A survey of phenomenological phase-shift analysis can be found in R. Levi Setti, rapporteurs' talk, in Ref. 4. For low energies see Martin and Perrin, Ref. 8.
[11] C. J. Goebel, Phys. Rev. Letters **16**, 1130 (1966).
[12] J. C. Pati and K. V. Vasavada, Phys. Rev. **144**, 1270 (1966).
[13] J. Katz and S. Wagner, Phys. Rev. **188**, 2196 (1969).

and $\phi^{(+)}$ creates a meson. The V_+ (V_-) raises (lowers) the meson number by one while leaving the number of baryons unchanged. We make use of the identity

$$\langle b|S|a\rangle=\langle b|a\rangle -2\pi i S(W_a-W_b)\lim_{\epsilon\to 0}\langle b|T(W_a+i\epsilon)|a\rangle \quad (4)$$

for the S-matrix element $\langle b|S|a\rangle$ for elastic K^+p scattering, where $|a\rangle$ and $|b\rangle$ are one-meson–one-proton states and

$$\langle b|T(z)|a\rangle=\langle b|V[1+R_0(z)V]^{-1}|a\rangle,\quad R_0(z)=(H_0-z)^{-1} \quad (5)$$

H_0 being the free Hamiltonian. In order to obtain an approximate expression for the S-matrix element, we have expanded the right-hand side of Eq. (5) in powers at V and neglected those terms in the expansion which contain more than two mesons in their intermediate states. It is then possible to resum the series, and one obtains[1,2] the approximate expression

$$\langle b|T(z)|a\rangle=\langle b|K(z)[1+R_0(z)K(z)]^{-1}|a\rangle, \quad (6)$$

where

$$-K(z)=V_+R_0(z)V_-+V_-R_0(z)V_+. \quad (7)$$

The operator $K(z)$ can be looked upon as a generalized energy-dependent meson-nucleon potential. It becomes non-Hermitian above the production threshold, corresponding to the possibility of meson production. The $K(z)$ in general contains a diagonal part $K_{\mathrm{dn}}(z)$:

$$\langle b|K(z)|a\rangle=\langle b|\tilde{K}(z)|a\rangle+K_{\mathrm{dn}}(z)\langle b|a\rangle. \quad (8)$$

The presence of $K_{\mathrm{dn}}(z)$ is a self-energy effect due to meson-cloud formation and, in the spirit of Ref. 14, it is to be discarded (mass renormalization). This amounts to replacing $K(z)$ by $\tilde{K}(z)$ in Eq. (6).

Using Eqs. (3), (7), and (8), standard commutation rules, and partial-wave decomposition in the c.m. system, we obtained the following low-energy expression (for the details of the low-energy approximations see Ref. 1) for $\langle k'|\tilde{K}(z)|k\rangle$ in the p-wave $J=\frac{3}{2}^+$ elastic K^+p channel:

$$\langle k'|\tilde{K}(z)|k\rangle=-\frac{g_0^2}{4\pi}\frac{1}{6\pi}[\tfrac{1}{3}(3-2\alpha)^2+(1-2\alpha)^2] \times\frac{k'k}{(E_p'E_p\omega'\omega)^{1/2}}\frac{1}{E_\Sigma(k')}\frac{1}{E_\Sigma(k')+\omega'+\omega-z}, \quad (9)$$

where $E_p=(M_p^2+k^2)^{1/2}$, $\omega=(M_K^2+k^2)^{1/2}$, $E_\Sigma=(M_\Sigma^2+k^2)^{1/2}$, k is the c.m. meson momentum; M_p, M_Σ, and M_K are the proton, the average $\Lambda\Sigma$ mass, and the K-meson mass, respectively. The constant α is the F/D ratio, for which we use the experimental value[15] $\alpha=0.75$. The evaluation of the energy-shell T matrix requires the inversion of the integral kernel $[1+R_0(z)\tilde{K}(z)]$ on the K^+p subspace. With expression (9) for $\langle k'|\tilde{K}(z)|k\rangle$ this cannot be done analytically. In order to proceed, we have therefore neglected the influence of the inelastic production channel. This was done by using the identity

$$\langle k'|R_0(z)\tilde{K}(z)|k\rangle=\langle k'|R_0(z)V_{\mathrm{eff}}|k\rangle+\langle k'|F(z)|k\rangle, \quad (10)$$

where

$$\langle k'|V_{\mathrm{eff}}|k\rangle =-\frac{g_0^2}{4\pi}\frac{1}{6\pi}\frac{k'k}{E_p'E_p\omega'\omega}\frac{1}{E_\Sigma(k')}\frac{1}{\omega+M_\Sigma-M_p} \quad (11)$$

[we have put $\alpha=0.75$ and $\omega+E_\Sigma(k)-E_p(k')\sim\omega+M_\Sigma-M_p$] and where $\langle k'|F(z)|k\rangle$ is regular for $\mathrm{Re}z<M_\Sigma+2M_K$, and by neglecting $\langle k'|F(z)|k\rangle$. This can be thought of as being justified by the principle of nearest singularity. In this approximation we obtain

$$\langle k'|T(z)|k\rangle=\langle k'|V_{\mathrm{eff}}[1+R_0(z)V_{\mathrm{eff}}]^{-1}|k\rangle, \quad (12)$$

where the kernel $\langle k'|V_{\mathrm{eff}}|k\rangle$ is given by Eq. (11). The effective potential V_{eff} is separable and the inversion can be done analytically. One gets the following result:

$$S(k)=e^{2i\delta(k)}=\frac{1-(g_0^2/4\pi)(1/6\pi)\Lambda^{(0)}(W-i\epsilon)}{1-(g_0^2/4\pi)(1/6\pi)\Lambda^{(0)}(W+i\epsilon)}, \quad (13)$$

where $\delta(k)$ is the $J=\frac{3}{2}^+$ phase shift, W is the c.m. total energy, and the function $\Lambda^{(0)}(z)$ is given by

$$\Lambda^0(z)=\int_0^\infty dq\frac{q^4}{E_p(q)E_\Sigma(q)\omega(q)} \times\frac{1}{\omega(q)+M_\Sigma-M_p}\frac{1}{E(q)+\omega(q)-z}. \quad (14)$$

The right-hand side of this equation is logarithmically divergent. We can obtain a cutoff-independent answer by defining the renormalized coupling in the customary manner as

$$g^2=g_0^2\left(1-\frac{g_0^2}{4\pi}\frac{1}{6\pi}\Lambda^0(M_p)\right)^{-1}$$

and expressing the right-hand side of Eq. (13) in terms of g^2 (g is the renormalized pion-nucleon coupling constant; experimentally, $g^2/4\pi=15$), instead of g_0^2. One gets that

$$S(k)=e^{2i\delta(k)} =\frac{1-(g^2/4\pi)(1/6\pi)(W-M_p)\Lambda(W-i\epsilon)}{1-(g^2/4\pi)(1/6\pi)(W-M_p)\Lambda(W+i\epsilon)}, \quad (15)$$

[14] L. Van Hove, N. M. Hugenholtz, and L. P. Howland, *Quantum Theory of Many-Particle Systems* (Benjamin, New York, 1961), p. 136.

[15] See, e.g., A. Messiah, *Quantum Mechanics* (Wiley, New York, 1962); W. Brenig and R. Haag, Fortschr. Physik **7**, 4/5 (1959).

where now

$$\Lambda(z)=\int dq \frac{q^4}{E_p(q)E_\Sigma(q)\omega(q)}\frac{1}{\omega(q)+M_\Sigma-M_p}\times\frac{1}{E_p(q)+\omega(q)-M_p}\frac{1}{E_p(q)+\omega(q)-z} \quad (16)$$

is a finite (cutoff-independent) function. From the two equations above, one gets the final formula for the $J=\frac{3}{2}^+$ K^+p phase shift:

$$\tan\delta=\frac{g^2}{4\pi}\frac{1}{6}\frac{k^3}{E_\Sigma}\frac{1}{W}\frac{1}{\omega+M_\Sigma-M_p}\times\left(1-\frac{g^2}{4\pi}\frac{1}{6\pi}(W-M_p)\,\mathrm{Re}\Lambda(W)\right)^{-1}. \quad (17)$$

III. NUMERICAL RESULTS

We used Eq. (17) to calculate the $J=\frac{3}{2}^+$ K^+p phase shift. The function $\mathrm{Re}\Lambda(W)$ has been evaluated numerically. For the renormalized pion-nucleon coupling constant, we used the value $g^2/4\pi=23.25$, which is the same as that obtained in Ref. 1 by fitting the model to reproduce the experimental value of the $J=\frac{3}{2}^+$ π^+p scattering length. The calculated phase shift is plotted in Fig. 1.

The phase shift rises rapidly in the region $1.6<W<1.9$ GeV, but it does not reach the value $\pi/2$. The derivative $d\delta/dW$ has maximum at $W=W_0=1750$ MeV, the maximum value being $\tau=d\delta/dW|_{W=W_0}=1.8\times10^{-3}$ MeV^{-1}. As is well known from potential scattering,[14] the derivative $d\delta/dW$ gives the delay time and a maximum value is therefore indicative of the existence of a resonance state with maximal delay time. Our calculations thus give a $J=\frac{3}{2}^+$, $Y=2$, $I=1$, K^+p exotic resonance with mass $M_{Z_1}=1750$ MeV and mean lifetime $\tau=1.8\times10^{-3}$ MeV^{-1}. This is a very-short-lived resonance, the total width Γ in the corresponding Breit-Wigner formula being given by $\Gamma=2/\tau\sim1000$ MeV.

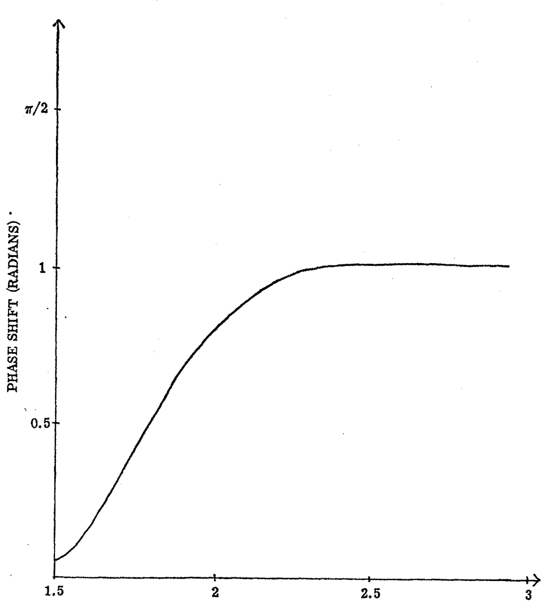

FIG. 1. Calculated $K^+pJ=\frac{3}{2}^+$ phase shift.

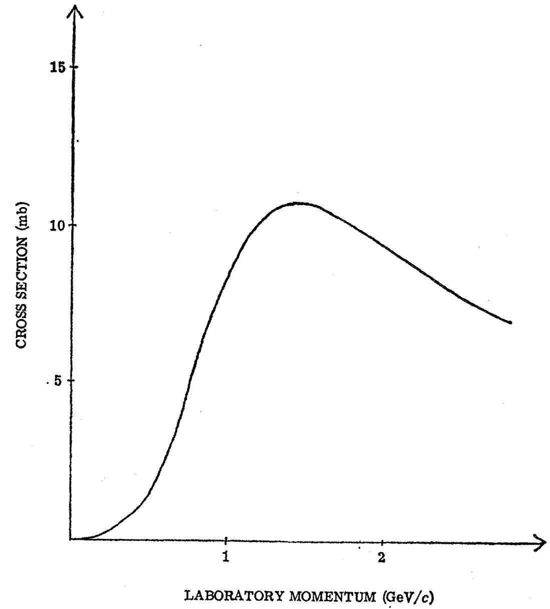

FIG. 2. Calculated elastic cross section $K^+p\rightarrow K^+p$.

In Fig. 2 we have plotted the calculated elastic cross section as a function of the K^+ lab momentum. (In the approximation we have been using, all the phase shifts are zero except the phase shift in the resonant $J=\frac{3}{2}^+$ channel.)

The calculated cross section has its maximum at $W=2$ GeV ($q_{\mathrm{lab}}=1450$ MeV/c), the height at maximum being about 11 mb. This agrees rather well with the experimental value[7] for the height of the elastic cross section, 14 mb, and the position of the bump in the total cross section[6] at 1900 MeV.

It might be worth noticing that according to our calculation the position of the maximum in the cross section lies about 250 MeV to the right of the mass of the resonance.

In Fig. 3 we have plotted the calculated differential cross section at resonant incident meson momentum $q_{\mathrm{lab}}=908$ MeV/c. (Discrete points indicate experimental values taken from Ref. 8 at $q_{\mathrm{lab}}=860$ MeV/c.)

From Eq. (17) there follows the $J=\frac{3}{2}^+$ effective-range formula

$$\frac{k^3}{M_K{}^2\omega}\cot\delta(\omega)=\left(\frac{g^2}{4\pi}\frac{1}{6}\frac{M_K{}^2\omega}{E_\Sigma W}\frac{1}{\omega+M_\Sigma-M_p}\right)^{-1}\times\left(1-\frac{g^2}{4\pi}\frac{1}{6\pi}(W-M_p)\,\mathrm{Re}\Lambda(M_p+M_K)\right)$$

for small ω.

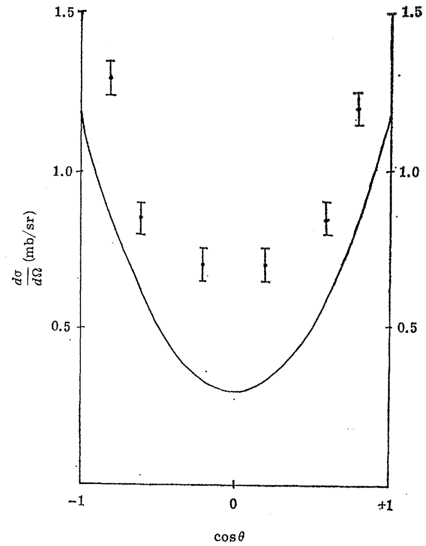

FIG. 3. Calculated differential cross section at $q_{\text{lab}}=908$ MeV/c.

One finds numerically that

$$(g^2/4\pi)(1/6\pi)M_K \operatorname{Re}\Lambda(M_K+M_p)\sim 0.3$$

with $g^2/4\pi=23.25$.

IV. SU_3 PARTNERS

The K^+p resonance belongs to a broken SU_3 27-plet. There are two more single-channel exotic resonances belonging to the same 27-plet which we were able to treat by the same method. We only quote the results for these.

$Y=0, I=2, J=\frac{3}{2}^+$ resonance. This is a $\pi^+\Sigma^+$ state. We obtained 1600 MeV for the mass and 1.6×10^{-3} MeV^{-1} for the mean lifetime. The phase shift stays under $\pi/2$. The maximum of the cross section is at 1800 MeV, and its height is about 15 mb.

$Y=-2, I=1, J=\frac{3}{2}^+$ resonance. This is a $K^-\Xi^-$ state. We obtained 2050 MeV for the mass and 3.1×10^{-3} MeV^{-1} for the lifetime. The phase shift does not go through $\pi/2$. The maximum of the cross section is at 2150 MeV, and its height is 18 mb.

Having calculated the masses of three of the states in the 27-plet, we can calculate the others if we assume the validity of the Gell-Mann–Okubo mass formula which gave in our case

$$M=\{1960-75Y-60[I(I+1)-\tfrac{1}{4}Y^2]\}\ \text{MeV}$$

for the masses in the 27-plet.

V. CONCLUSIONS

For the $J=\frac{3}{2}^+$ K^+p channel, the results of our calculations are in reasonable agreement with experiment in

TABLE I. Summary of numerical results. a is the $J=\frac{3}{2}^+$ π^+p scattering length. M_{Z^*} is here defined as the position of the maximum in the K^+p cross section.

	Calculated	Experimental
a	fitted	0.22
$g^2/4\pi$	23	15
M_{N^*}	1368 MeV	1238 MeV
M_{Ω^-}	1796 MeV	1675 MeV
M_{Z^*}	2000 MeV	1900 MeV

the resonance energy region. This we take as a circumstantial support for the main conclusion of our work that the observed maximum in the total K^+p cross section is due to a $J=\frac{3}{2}^+$ elastic K^+p resonance with mass about 1750 MeV and lifetime of the order of 1.8×10^{-3} MeV^{-1}. It is a member of a broken SU_3 27-plet. It is characteristic for the resonance that it has a large width $\Gamma\sim1000$ MeV and that the phase shift does not go through $\pi/2$ in the resonance energy region.

To summarize the results of Refs. 1 and 2 and the present work, our model Hamiltonian can account for the mass of the $\Delta(1238)$ π^+p $J=\frac{3}{2}^+$ resonance, the mass of the Ω^- as a $\bar{K}\Xi$ bound state, and the position of the maximum in the K^+p cross section, as being due to a long-lived $J=\frac{3}{2}^+$ resonance provided that one fits the pion-nucleon coupling constant to the experimental π^+p scattering length giving $g^2/4\pi=23.25$. This is 50% higher than the experimental value of the squared pion-nucleon coupling constant and this seems fair enough in view of the approximations, such as the neglect of the production channel and the contribution of intermediate states of more than two mesons that had to be made. It may very well be that the main effect of the neglected contributions is to reduce the value of the coupling constant necessary to reproduce the experimental π^+p scattering length. Thus it may well be that our model is quite reliable for the description of low-energy resonant $J=\frac{3}{2}^+$ pseudoscalar meson-octet baryon scattering. For S-wave scattering, the model cannot be expected to be good since direct meson-meson interaction contributions are known to be important.

Apart from the pion-nucleon coupling constant, there are no other parameters in our model.

The numerical results of Refs. 1 and 2 and the present article are summarized in Table I.

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