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[7]Michael Nauenberg (unpublished) has pointed out that, in general, the observed branching ratios for a resonant state strongly coupled to several channels may depend sensitively upon whether the resonant state is produced in association with other particles or is formed directly from the interaction of two specific particles. In the former case, the final state may be considered to be the result of a set of three-body production processes (e.g., $\Lambda\pi\pi$, $\Sigma\pi\pi$, $\bar{K}p\pi$), followed by coherent rescattering of some groups of particles formed. Since only one of the possible intermediate states will be analogous to formation in the two-body situation, the observed final states may have markedly different properties.

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K^--p INTERACTIONS NEAR 760 MeV/c†

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We present the results of a study of the K^--p system at incident K^- laboratory momenta of 620, 760, and 850 MeV/c (center-of-mass energies $E_{\text{c.m.}}$ of 1616, 1681, and 1723 MeV, respectively). Only the most important features of the interactions have been obtained at each of these momenta. At 620 MeV/c the system is dominated by strong $S_{1/2}$ absorption. At 760 MeV/c, effects due to $Y_1^*(1660)$ are observed.[1] Here the presence of large $\cos^2\theta$ terms and the absence of large $\cos^3\theta$ terms in the angular distributions suggest $\frac{3}{2}$ as a plausible spin assignment for the resonance. Finally, at 850 MeV/c, large $\frac{3}{2}$ or $\frac{5}{2}$ amplitudes have set in.

The Lawrence Radiation Laboratory's 15-in. hydrogen bubble chamber was exposed to a separated K^- beam capable of operating at either 760 or 850 MeV/c.[2] A setting at 620 MeV/c was obtained by degrading the 760-MeV/c beam. A total of 8000 interactions, representing all the available data, were analyzed.

In Table I we summarize the observed total cross sections, having determined the path length at each momentum by counting τ decays.[3] The only significant bias occurs when a Σ^+ decays via the protonic mode; at our energies the laboratory angle between the Σ^+ and its decay proton is usually too small to be detected with good efficiency. For this reason only, Σ^+ decaying via the pionic mode were used to establish both the total and differential cross sections. On the other hand, ambiguities in the interpretation of the events arise when one wants to distinguish between the $\Lambda\pi^0$, $\Sigma^0\pi^0$, $\Lambda\pi^0\pi^0$, and $\Sigma^0\pi^0\pi^0$ final states. The method of separation used in our experiment is the same as that described by Ferro-Luzzi, Tripp, and Watson.[4]

All differential cross sections were fitted to a series of the form

$$(4/\lambda\!\!\!^{-2})d\sigma/d\Omega = A_0 + A_1\cos\theta + A_2\cos^2\theta + \cdots, \quad (1)$$

Table I. Total cross sections in millibarns.

	P_{K^-} (MeV/c)	620 ± 15	760 ± 7	850 ± 10
	$E_{\text{c.m.}}$ (MeV)	1616	1680	1723
Reaction	Topology[a]	50 events/mb	124 events/mb	145 events/mb
K^-p	2 prong	16.0 ± 1.0	16.7 ± 1.0	22.4 ± 2.0
$\bar{K}^0 n$	0 prong + V	2.3 ± 0.4	3.8 ± 0.3	4.0 ± 0.3
$\Sigma^+\pi^-$	2 prong (+ decays)[b]	4.6 ± 0.6	2.8 ± 0.3	2.0 ± 0.2
$\Sigma^-\pi^+$	2 prong (- decays)	2.0 ± 0.3	3.3 ± 0.2	1.3 ± 0.1
$\Sigma^0\pi^0$	0 prong + V	2.1 ± 0.3	1.4 ± 0.2	0.8 ± 0.1
$\Lambda\pi^0$	0 prong + V	1.9 ± 0.3	2.6 ± 0.2	2.7 ± 0.2
$\Lambda\pi^+\pi^-$	2 prong + V	1.7 ± 0.3	4.3 ± 0.3	3.5 ± 0.3
$\Sigma^0\pi^+\pi^-$	2 prong + V	0.3 ± 0.15	0.8 ± 0.15	0.7 ± 0.15
$\Sigma^-\pi^+\pi^0$	2 prong (- decays)	0.3 ± 0.1	0.8 ± 0.1	0.7 ± 0.1
$\Sigma^+\pi^-\pi^0$	2 prong (+ decays)	0.1 ± 0.05	0.8 ± 0.2	0.5 ± 0.1
$(\Lambda\Sigma^0)\pi^0\pi^0$	0 prong + V	1.1 ± 0.2	1.9 ± 0.2	1.3 ± 0.2
$\Lambda\eta$ ($\eta \to$ neut)	0 prong + V	...	0.5 ± 0.15	0.15 ± 0.1
$\Lambda\pi^+\pi^-\pi^0$	2 prong + V	0.0 ± 0.03	0.25 ± 0.05	0.15 ± 0.05
$\bar{K}^0 p\pi^-$	2 prong + V	0.0 ± 0.03	0.04 ± 0.03	0.10 ± 0.06
$K^-p\pi^0$	2 prong	0.0 ± 0.03	0.15 ± 0.1	0.3 ± 0.1
$K^-\pi^+n$	2 prong	0.06 ± 0.06	0.0 ± 0.05	0.2 ± 0.1
$\bar{K}^0 n\pi^0$	0 prong + V	0.0 ± 0.03	0.0 ± 0.05	0.3 ± 0.1
Total		32.4 ± 1.5	40.1 ± 1.3	40.6 ± 2.1
$(\Sigma\pi)_{I=0}$		6.4 ± 0.8	4.2 ± 0.5	2.3 ± 0.3
$(\Sigma\pi)_{I=1}$		2.3 ± 0.8	3.3 ± 0.4	1.3 ± 0.3
$\pi\lambda^2$		9.50	6.80	5.75

[a]Topological appearance of the event in the bubble chamber.
[b]This means 2 prong, and the positive prong decays.

where $4/\lambda\!\!\!^-{}^2$ is an arbitrary normalization factor. The results of such fits were the dimensionless coefficients A_i and their errors. A fit of order n means a fit up to and including the term $A_n \cos^n\theta$.

In Figs. 1 and 2 we present the energy dependence of the coefficients A_i for all the two-body reactions except $\Sigma^0\pi^0$. In this latter channel statistics are small, and, for angular distributions at least, it is difficult to determine the amount contributed to a given bin by the $\Lambda\pi^0\pi^0$ and $\Sigma^0\pi^0\pi^0$ final states. Significant amounts of $\cos^3\theta$ are required only at 850 MeV/c as is shown in Fig. 3.

At 620 MeV/c the angular distributions are remarkably similar to those obtained at 510 MeV/c by Ferro-Luzzi, Tripp, and Watson[4]: No appreciable changes in the A_i seem to have taken place between these two momenta, and the system is probably well described by a large absorptive $S_{1/2}$ amplitude with small amounts of $P_{1/2}$, $P_{3/2}$, and $D_{3/2}$. At 850 MeV/c the large amounts of $\cos^3\theta$ noticeable in every distribution indicate that $J = \frac{3}{2}$ or $\frac{5}{2}$ amplitudes are now present.

We now discuss the 760-MeV/c data more thoroughly and examine the effect of the newly discovered $Y_1^*(1660)$ on our total cross sections and angular distributions. From now on we shall refer to this resonance as $\Sigma(1660)$; 1660 MeV corresponds to a laboratory momentum of 715 MeV/c, and its half-width, $\frac{1}{2}\Gamma = 20$ MeV, corresponds to ± 43 MeV/c. Significant effects due to $\Sigma(1660)$ therefore should be observed at 760 MeV/c, provided it is coupled reasonably strongly to the K-nucleon system. Assuming that this resonance is adequately described by a pure Breit-Wigner form, we can obtain the strength of the coupling by using the formula for the energy dependence of the total cross section in such a case, namely,

$$\sigma_T = 2\pi\lambda\!\!\!^-{}^2 (J + \tfrac{1}{2}) x_K / (\epsilon^2 + 1), \tag{2}$$

where the normalized channel width $x_K \equiv \Gamma_{\text{elastic}}/\Gamma$, and $\epsilon \equiv 2(E_R - E_{\text{c.m.}})/\Gamma$.

Examining in greater detail the behavior of our cross sections in various channels, shown in Table I, we notice first a 2-mb bump in the $\Sigma^-\pi^+$ channel at 760 MeV/c. This final state is a mix-

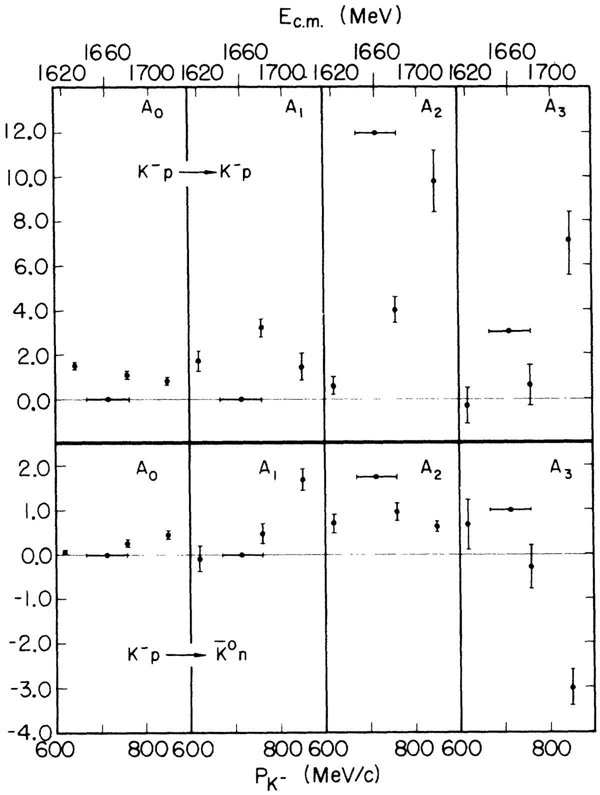

FIG. 1. Coefficients A_i for $K^-p \to K^-p$ and $K^-p \to \bar{K}^0n$ with data fitted up to third order. Cosθ is defined as $\hat{K}_{inc}^- \cdot \hat{K}_{scattered}$. The A_3's are large at 850 MeV/c and consistent with zero at the lower momenta. For reference we have drawn horizontal bars indicating the position and width of $Y_1^*(1660)$ for all coefficients.

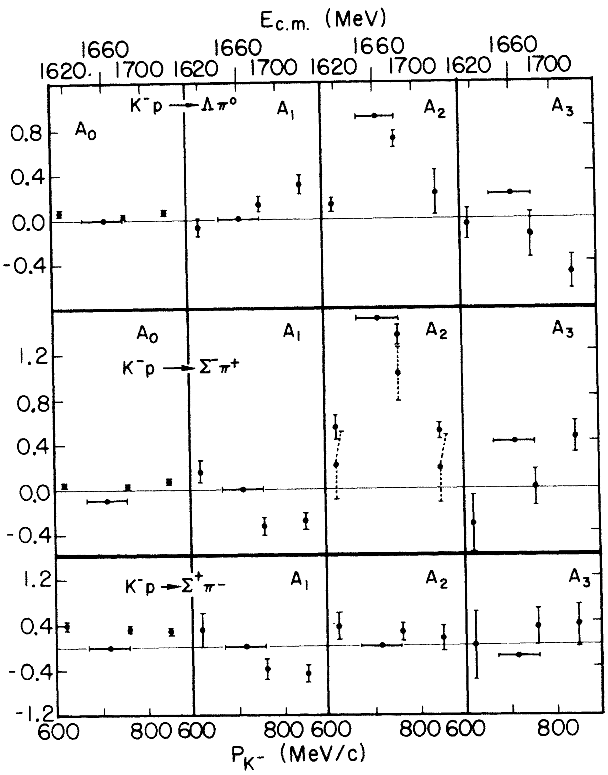

FIG. 2. Coefficients A_i for $K^-p \to \Lambda\pi^0$, $K^-p \to \Sigma^-\pi^+$, and $K^-p \to \Sigma^+\pi^-$ with data fitted up to third order. Here cosθ is defined as $\hat{K}_{inc}^- \cdot \hat{Y}$. In the A_2's, one notices significant bumps in $\Sigma^-\pi^+$ and $\Lambda\pi^0$, whereas in $\Sigma^+\pi^-$ the nonresonant background interferes destructively with the resonant amplitude at 760 MeV/c. The A_3's are again consistent with zero at 760 MeV/c. The dashed curves in the $A_2(\Sigma^-\pi^+)$ are the values of the coefficients of $\cos^2\theta$ when $\Sigma^-\pi^+$ is fitted up to fourth order. The bump in $\cos^2\theta$ is still significant, whereas the coefficients of $\cos^4\theta$ (not shown) are consistent with zero.

ture of isospin 0 and 1, but the isospin of the effect can be established through the following formulas:

$$\sigma(\Sigma\pi, I=0) = 3\sigma(\Sigma^0\pi^0), \tag{3}$$

$$\sigma(\Sigma\pi, I=1) = \sigma(\Sigma^-\pi^+) + \sigma(\Sigma^+\pi^-) - 2\sigma(\Sigma^0\pi^0). \tag{4}$$

It is clear from the results of this separation (Table I) that the rise is indeed due to the isospin-1 part of the amplitude. We therefore attribute the effect to $\Sigma(1660)$. One notices also a small rise at 760 MeV/c in every one of the $\Lambda\pi\pi$ and $\Sigma\pi\pi$ cross sections. These are also interpreted as manifestations of the three-body decay modes of the resonance. Adding these cross sections, we estimate x_K of Eq. (2) to be 0.25 ± 0.10 for spin $\frac{3}{2}$, and appropriate fractions of that value for other spin assignments. The $\Sigma\pi$ and $\Lambda\pi$ normalized channel widths are estimated to be $x_\Sigma = 0.35 \pm 0.15$ and $x_\Lambda = 0.15 \pm 0.10$. Finally, since it is not possible to separate the $\Lambda\pi\pi$ and $\Sigma\pi\pi$ final states into isospin without more information than that given by the K-proton system alone, we cannot calculate partial widths for these last two channels. Because the elastic cross sections are proportional to $(x_K)^2$, the small values of x_K that we have obtained lead, for all possible spin assignments, to negligible rises in the K-nucleon total cross sections at 760 MeV/c. Therefore, no lower limit on the spin of the resonance can be set by this method. Since Table I shows that at 760 MeV/c the cross sections due to the resonance are small compared to the background, we conclude that the nonresonant amplitudes dominate over the resonant amplitude.

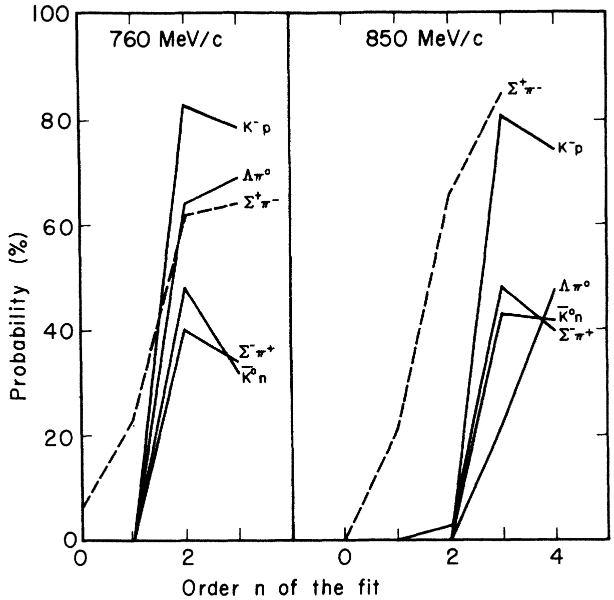

FIG. 3. The χ^2 probability of describing two-body angular distributions with a fit of order n. The probability of a fit is defined as

$$\int_{\chi_0^2}^{\infty} f(\chi^2, \nu) d\chi^2,$$

where χ_0^2 is the least-squares-fit chi-square and ν the number of degrees of freedom for that distribution. The numbers of events in each distribution are comparable at both momenta. However, we had much less statistics in the $\Sigma^+\pi^-$ channel (dashed lines), since only the pionic decays of the Σ^+ were used. All channels at 850 MeV/c clearly require $\cos^3\theta$, whereas no significant amounts of that power are required anywhere at 760 MeV/c. A small amount of $\cos^4\theta$ is needed for $\Lambda\pi^0$ at 850 MeV/c.

To obtain information on the spin of $\Sigma(1660)$, we examine the two-body angular distributions at 760 MeV/c. Usually the spin of a resonance is determined by looking at the energy dependence of the coefficients A_i at closely spaced momentum intervals throughout the resonance region. Here we have only a single point one half-width above the resonance, and so a quantitative analysis is not possible. However, by considering the A_1, A_2, and A_3 in Fig. 2, we wish to show that our data seem to be more consistent with $J=\frac{3}{2}$ for $\Sigma(1660)$ than with any other spin assignment.

Examining first A_2 in Fig. 2, we shall show that spin $\frac{1}{2}$ is unlikely. We observe a large enhancement at 760 MeV/c in both $A_2(\Sigma^-\pi^+)$ and $A_2(\Lambda\pi^0)$. We attribute these bumps to interferences between the small resonant amplitude and the nonresonant background[5]; as we have seen, the nonresonant background is dominant and is composed only of $S_{1/2}$, $P_{1/2}$, and $P_{3/2}$ amplitudes, since large amounts of other amplitudes would lead to $\cos^3\theta$ interferences, which, as is clearly shown in Fig. 3, are not present. To examine the consistency of our data with the spin-$\frac{1}{2}$ hypothesis, we look at the form of A_2 in a partial-wave expansion up to $J=\frac{3}{2}$, namely,

$$A_2 = 3|P_{3/2}|^2 + 3|D_{3/2}|^2 + 6\mathrm{Re}(P_{1/2}{}^*P_{3/2} + S_{1/2}{}^*D_{3/2}). \quad (5)$$

For $J=\frac{1}{2}$, the large interference effects observed in the A_2 must be due to the $S_{1/2}D_{3/2}$ or $P_{1/2}P_{3/2}$ terms. An $S_{1/2}$ resonance seems unlikely, since $D_{3/2}$ is also small. Again, if the resonance is $P_{1/2}$, then $P_{3/2}$ must be large, and values of the A_2 considerably bigger than what we observe are expected because of the term $3|P_{3/2}|^2$. Spin $\frac{1}{2}$ therefore seems improbable.

By considering the A_1 and the A_3, we next show that $J=\frac{5}{2}$ is unlikely for $\Sigma(1660)$. The values of the A_1 in Fig. 2 again show interferences in all three distributions. If the spin is $\frac{5}{2}$, we would have expected variations in the A_3 of the order of 15/9 those observed in the A_1, if the nonresonant background is assumed to vary slowly throughout the region. To see this, note that in a partial-wave analysis up to angular momentum $\frac{5}{2}$, the expressions for A_1 and A_3 are

$$\begin{aligned} A_1 = \mathrm{Re}[&2S_{1/2}{}^*P_{1/2} + 4S_{1/2}{}^*P_{3/2} + 4P_{1/2}{}^*D_{3/2} \\ &-9S_{1/2}{}^*F_{5/2} - 9P_{1/2}{}^*D_{5/2} - 10P_{3/2}{}^*D_{3/2} \\ &+(45/2)D_{5/2}{}^*F_{5/2}], \end{aligned} \quad (6)$$

$$\begin{aligned} A_3 = \mathrm{Re}(&12P_{3/2}{}^*D_{5/2} + 12D_{3/2}{}^*F_{5/2} \\ &+15S_{1/2}{}^*F_{5/2} + 15P_{1/2}{}^*D_{5/2} + 18P_{3/2}{}^*D_{3/2} \\ &-117D_{5/2}{}^*F_{5/2}). \end{aligned} \quad (7)$$

No such variations are seen in the A_3, which instead are all consistent with zero. It also seems unlikely that, when made to interfere with reasonable values of the nonresonant background, a $\frac{5}{2}$ resonant amplitude with our normalized channel widths would not have led to sizable A_3's in at least one of the angular distributions.

In summary, the most natural way to interpret the data is to assume that a small $P_{3/2}$ or $D_{3/2}$ resonant amplitude is interfering with large $S_{1/2}$ and $P_{1/2}$ nonresonant background. However, only an analysis of more angular distributions throughout the whole region will provide a final answer to the spin question and information on the parity. In particular, the energy dependence of the $\Sigma^-\pi^+$ differential cross sections must be established carefully, because we cannot rule out the possibility that the large value of $A_2(\Sigma^-\pi^+)$ at 760 MeV/c is due mainly to nonresonant terms. If this is the case, the value of x_K would be much smaller, and this, in turn, would decrease our confidence in the spin assignment.

Finally, let us say a few words about the three-body final states. If $\Lambda\pi^+\pi^-$ is formed predominantly through the chain $K^-p \to \Sigma(1660) \to \Sigma(1385) + \pi \to \Lambda\pi^+\pi^-$, one could hope to say something about the spin and parity of $\Sigma(1660)$, taking $\Sigma(1385)$ to be $\frac{3}{2}^+$.[6-8] However, the observed cross section for $K^-p \to \Lambda\pi^+\pi^-$ at 760 MeV/c is 4.3 mb, and $\Sigma(1660)$ can contribute only ~0.4 mb to that value since the normalized channel width for $\Sigma(1660) \to \Lambda\pi^+\pi^-$ is approximately 0.10. Therefore, $Y_1^*(1660)$ cannot affect the $\Lambda\pi^+\pi^-$ distributions significantly. Similar negative conclusions are drawn about the $\Sigma\pi\pi$ channels.

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[4]M. Ferro-Luzzi, R. D. Tripp, and M. B. Watson, Phys. Rev. Letters 8, 28 (1962).

[5]Both the square of the resonant amplitude and interference terms between the resonant amplitude and the nonresonant background appear in the coefficients of the even powers of $\cos\theta$. With the observed normalized partial widths for $\Sigma(1660)$, the former are of the order of 0.02 for all channels in our dimensionless units, whereas the latter can be one or two orders of magnitude larger (see reference 4).

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EIGHTFOLD-WAY ASSIGNMENTS FOR $Y_1^*(1660)$ AND OTHER BARYONS†

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In the preceding papers establishing the existence of a new $T=1$, $Y=0$ hyperon resonance at 1660 MeV, $Y_1^*(1660)$,[1] it is noted that the baryons and the many low-lying baryon-meson resonances may be included within just four unitary-symmetry multiplets[2]: the $j=\frac{1}{2}^+$ α octet, the $j=\frac{1}{2}^-$ β singlet, the $j=\frac{3}{2}^-$ γ octet, and the $j=\frac{3}{2}^+$ δ decuplet.[3] All members of these unitary multiplets are seen except Ω_δ^- (sometimes called Z^-) and Ξ_γ, whose discoveries would give crucial tests of the eightfold way. Each of the four unitary multiplets may be viewed as the "ground" (i.e., lowest angular momentum) state of a different Regge trajectory, and each such trajectory may have several manifestations with different physical values of angular momentum. These "Regge recurrences" (should they occur at all) must characterize all the members of a given unitary multiplet, and must be spaced by two units of angular momentum. Identifying the $N_{3/2}^*(1920)$ resonance with Δ_δ ($j=\frac{7}{2}^+$), the $Y_0^*(1815)$ with Λ_α ($j=\frac{5}{2}^+$), and the $N_{1/2}^*(1688)$ resonance with N_α ($j=\frac{5}{2}^+$), we find that all presently known baryon states are described with only four unitary-multiplet Regge trajectories. Because the first Regge recurrences of some