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Theoretical Evidence for $I = 0$ Z^*'s

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We present strong theoretical evidence for the existence of highly inelastic $I = 0$ K-N (Z^*) resonances in the $S_{1/2}$ and $D_{3/2}$ states at c.m. energies near 1830 MeV with widths of about 100 MeV.

We present strong theoretical evidence for the existence of $I=0$, $Y=+2$ exotic baryon resonances (Z^*'s) in the K-N system. For some time now, the question of the existence of such Z^* resonances has been vigorously debated.[1] From a theoretical standpoint this debate seems unwarranted, since simple model-independent dynamical arguments almost certainly predict the existence of $I=0$ K-N resonances. In this Letter we present these arguments, illustrating them with a particular dynamical calculation of the properties of these resonances. We find highly inelastic $S_{1/2}$ and $I_{3/2}$ resonances in the $I=0$ K-N channel near 1830 MeV c.m. energy with widths of about 100 MeV.[2-5] Our calculation fits the total $K\pi N$ inelastic cross section, and is also consistent with recent $KN \rightarrow K^*N$ angular distributions obtained by Hirata et al.[6]

The dynamical mechanism producing the resonances is the reflection—essentially through unitarity—of the large, very rapidly opening K^*-N production near its threshold. As is well known,[7] a rapidly opening inelastic threshold can affect the elastic channel and drive it resonant. Since the K^*-N production can only open rapidly in a relative S wave, only the K-N waves coupled to K^*-N S waves can experience this effect, and they are the $S_{1/2}$ and $D_{3/2}$ states. A resonance produced by such a mechanism can in general occur above or below the inelastic threshold, but if it occurs above, it will necessarily be a highly inelastic resonance. Such an inelastic resonance will not necessarily show up as a clear bump in the elastic K-N scattering, but it certainly will appear in a phase-shift analysis.

The isospin of the resonance depends on the mechanism responsible for the inelastic channel. In the case of $KN \rightarrow K^*N$, the dominant mechanism is thought to be π exchange as shown in Fig. 1(a). This reflects back on K-N scattering through the box diagram of Fig. 1(b), where it provides a kind of potential. The isospin factor for the box is 1 for the $I=0$ K-N system, but $\frac{1}{9}$ for $I=1$. Hence we expect K^*N production via π exchange to be important only for $I=0$.[8] All these arguments are much more familiar in the π-N system, where the rapid onset of ρ production drives the $D_{13}(1520)$ resonance, as shown by Cook and

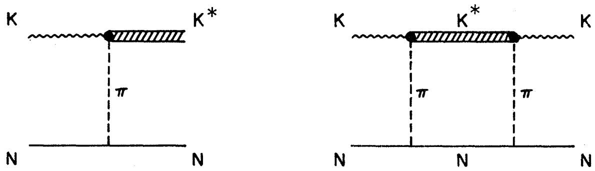

FIG. 1. Diagrammatic representation of basic mechanisms responsible for the existence of Z_0^*'s.

Lee,[9] and also drives the $S_{11}(1700)$ resonance.[10] Taking π exchange as the dominant mechanism leads to isospin factors favoring the $I=\frac{1}{2}$ state over the $I=\frac{3}{2}$ state by 4 to 1 in this case.

Any dynamical calculation which includes the effects of unitarity and allows for the large, rapidly opening K^*-N channel will undoubtedly give resonances in the elastic channel. Standard methods like the multichannel N/D have the disadvantage of treating the K^* as stable. We have therefore chosen to perform a three-body calculation of K-N scattering in the $I=0$ channel using the relativistic three-body equations first proposed by Blankenbecler and Sugar[11] and already applied with considerable success to π-N scattering.[12,13] The equations are essentially covariant generalizations of the Lippmann-Schwinger equations and the solutions satisfy two- and three-body unitarity. For the K-N system we take Fig. 1(b) as the "potential." Other potential terms or crossed cuts are no doubt of consequence in the K-N system, but since they are slowly varying in the K^* threshold region, they will not affect the existence or gross properties of the resonances as long as the inelastic cross section is correctly calculated.

Since the K-N system is not extremely relativistic at these energies and since multipion production does not seem to be important in this region, we expect the Blankenbecler-Sugar method to be particularly appropriate. The input into the calculation is the K^*-box "potential" of Fig. 1(b). It includes coupling constants, widths, and masses, all of which are known. In addition it is necessary to introduce form factors at all vertices and these bring in an additional parameter —a form-factor scale or cutoff. We have found that in this model no reasonable value of the cutoff parameters will make the K-N resonances disappear, although their detailed position and widths does depend on the cutoff, as one might expect.

The $S_{1/2}$ and $D_{3/2}$ amplitudes for some typical cutoff parameters are shown as Argand plots in Fig. 2. The clear resonance looping is apparent

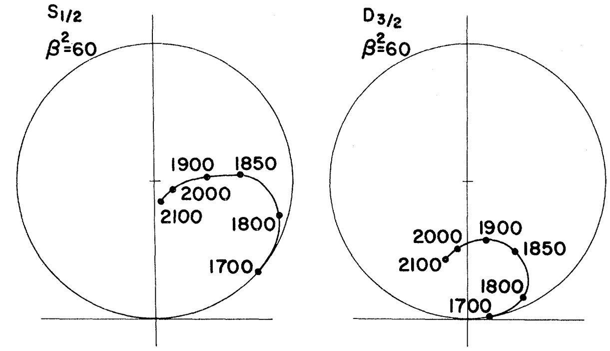

FIG. 2. Argand diagrams for $I=0$ $S_{1/2}$ and $D_{3/2}$ K-N resonances corresponding to a cutoff parameter $\beta^2 = 60m_\pi^2$ in the $K^*K\pi$ form factor.

as well as the highly inelastic nature of the resonances.[14] Both resonances are at 1830 MeV with widths of about 100 MeV. By varying the cutoff parameter it is possible to push the resonance positions below the K^*-production threshold. In that case the inelasticity parameter η stays nearly at 1 and the resonances would then appear clearly in the elastic scattering. This is ruled out by experiment. Further variation of the resonance positions can also come from left-hand or "potential" effects that we have neglected. It is not our purpose to explore these effects here but simply to point out that the K^*-box mechanism will produce these resonances and that it is difficult to imagine some other mechanism that will make them go away.

The inelastic K-N cross section is shown in Fig. 3, where it is compared with the data. We show separately the inelastic cross section coming from the $S_{1/2}$ and $D_{3/2}$ waves and the total inelastic cross section coming from solving all waves with the K^*-box driving mechanism of Fig. 1(b). In the waves other than the $S_{1/2}$ and $D_{3/2}$ there are certainly other important mechanisms that we are neglecting, and hence the total inelastic cross-section calculation is not to be taken literally. Nevertheless the energy dependence and magnitudes are roughly correct. We do not compare with the elastic K-N data since there are many other nonresonant waves that are important and our highly inelastic resonances will not dominate the elastic cross section or angular distributions. There is a well-known bump in the $I=0$ K-N total cross section near the K^*-N threshold,[15] but its unambiguous identification with the resonances we discuss here must await a full phase-shift analysis.

Recently Hirata et al.[6] have obtained angular

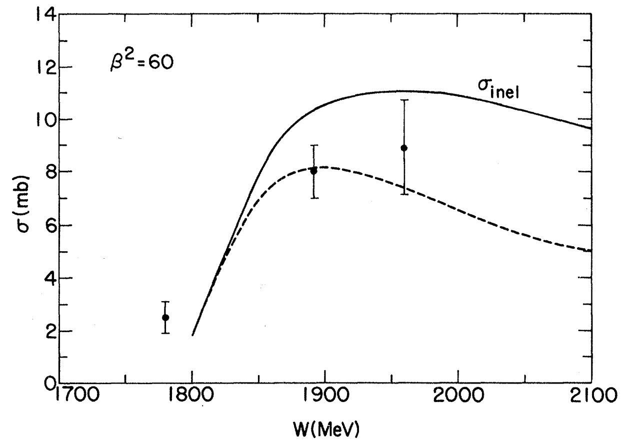

FIG. 3. The solid curve represents our fit to the total $K\pi N$ inelastic cross section. The experimental points are those of Hirata *et al.* (Ref. 6). The dashed curve represents the sum of the $S_{1/2}$ and $D_{3/2}$ contributions.

distribution data for $I=0$, $KN\rightarrow K^*N$. These data are strongly forward peaked even very near threshold. They argue that this is evidence for the π-exchange production mechanism of Fig. 1(a), and against elastic resonances, particularly in the S wave of the K^*-N system. For example at 1900 MeV c.m. energy they find a forward-backward asymmetry (FBA) of about 10/1. In fact, significant forward-peaking is not necessarily incompatible with a large S wave in the total cross section. For example, at 1900 MeV, a single one-pion–exchange denominator with the kinematics appropriate to the case in question leads to a FBA of 36/1 while the cross section is roughly 80% S wave. There are, however, numerator factors in the Born term of Fig. 1(a) that in fact reduce the FBA to 2.5/1 with no form factors and to about 4/1 with the form factors we use. Thus in order to obtain the observed FBA one needs some strong diffraction mechanism to absorb out the low partial waves. Our inelastic resonances produce just such a mechanism and we believe that the angular distribution for K^* production is evidence for—rather than against—our inelastic resonances.

As this Letter was being written, we became aware of a contribution to the Kiev Conference which analyzes K^+d interactions in the momentum range 0.6 to 1.5 GeV/c.[16] The first qualitative features of an $I=0$ K-N phase-shift analysis are in accord with the theoretical picture presented here of highly inelastic $S_{1/2}$ and $D_{3/2}$ resonances. In particular, almost all of the 200 presently acceptable phase-shift solutions of Ref. 16 require strong inelasticity in a D wave.

Presumably, the isoscalar Z^*'s discussed here belong to 10* multiplets of SU(3) with spin and parity $\frac{3}{2}^-$ and $\frac{1}{2}^-$. Although they are usually assigned to octets, the πN states $D_{13}(1520)$ and $S_{11}(1700)$ might also be in these 10* multiplets, on the grounds that, as discussed above, they have a similar dynamical origin. If one then assumes the Gell-Mann–Okubo mass formula, the equal-mass splitting for the antidecuplet leads in the $D_{3/2}$ case to bound Y_1^* and $\Xi_{3/2}^*$ states. For the $S_{1/2}$ case, the Y_1^* is of very low mass (1570 MeV) and the $\Xi_{3/2}^*$ is bound [though here the situation may be affected by mixing with the close-by $S_{11}(1535)$ πN state]. The existence of these bound states contradicts experiment. Hence, either the $D_{13}(1520)$ and $S_{11}(1700)$ πN states are not to be grouped with the Z_0^*'s in 10* multiplets or the mass splitting in such multiplets is more complex. The latter is easily possible since the Z^*'s are complex structures from the quark point of view and most of the simple successes of SU(3) have been in cases of correspondingly simple quark arguments. Indeed, "it is intriguing that SU(3) could be so badly broken for exotics that their multiplets are not even complete."[17]

In summary, we believe that the strong opening of K^*-N production in the $I=0$ K-N system drives exotic inelastic K-N resonances. The main rationale for naming these states exotic is given by the quark model. There is no reason, however, to expect such a nonunitary model to give highly inelastic resonances such as ours. If the complete phase-shift analysis should not find these resonances, one can possibly invent competing cross-channel mechanisms that could cancel the resonances, but it would be more of a puzzle to understand why and how they disappear than to understand their existence in the first place.

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that we predict is completely consistent with the experimental elastic and inelastic cross sections.

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Observation of the $\overline{\Omega}$ †

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We have observed an $\overline{\Omega}$ event. The $\overline{\Omega}$ is produced in the reaction $K^+d \to \overline{\Omega}\Lambda\Lambda p\pi^+\pi^-$ at 12 GeV/c, and decays via the mode $\overline{\Omega} \to \overline{\Lambda}K^+$. The fitted mass for this particle is $M_{\overline{\Omega}} = 1673.1 \pm 1.0$ MeV. The $\overline{\Omega}$-production cross section is of the order of 0.1 μb.

We have observed an example of the $\overline{\Omega}$.[1] The decay mode is

$$\overline{\Omega}^+ \to \overline{\Lambda} + K^+.$$

The event is observed in a systematic search for interactions with a charged vee and associated neutral vee. This signature with a positively charged vee is characteristic of the possible decays $\overline{\Omega} \to \overline{\Lambda}K^+$ and $\overline{\Xi} \to \overline{\Lambda}\pi^+$. The experiment is a study of the K^+d interaction at 12 GeV/c carried out in the 82-in. Stanford Linear Accelerator Center (SLAC) bubble chamber. A total of 500 000 pictures were taken. So far we have examined 60% of the film in this systematic search, and have observed the following:

$$\overline{\Xi}^+ \to \overline{\Lambda} + \pi^+ \text{ (45 events)},$$

$$\Xi^- \to \Lambda + \pi^- \text{ (15 events)},$$

and

$$\overline{\Omega}^+ \to \overline{\Lambda} + K^+ \text{ (1 event)}.$$

In Fig. 1 we show the $\overline{\Omega}$ event. The reaction has been fitted uniquely by the hypothesis

$$K^+d \to \overline{\Omega}\Lambda\Lambda p\pi^+\pi^-,$$

where only one of the Λ's decays by the charged mode in the bubble chamber. The $\overline{\Lambda}$ from the $\overline{\Omega}$ decay can be seen on the left, and decays via the mode

$$\overline{\Lambda} \to \overline{p} + \pi^+,$$

in which the antiproton annihilates with a proton of one of the deuterons in the chamber.

Figure 2 shows a plot of $P_\perp$ (transverse momentum) versus α for the $\overline{\Xi}$ and $\overline{\Omega}$ events[2]:

$$\alpha = (P_{\parallel}^{\overline{\Lambda}} - P_{\parallel}^{m})/(P_{\parallel}^{\overline{\Lambda}} + P_{\parallel}^{m}).$$

Here the symbol $P_\parallel$ means longitudinal momentum, and the superscript $\overline{\Lambda}$ refers to the antilambda, while m refers to either the π^+ or K^+ meson. The kinematic ellipses for $\overline{\Xi} \to \overline{\Lambda}\pi$ and $\overline{\Omega} \to \overline{\Lambda}K$ decays, shown in Fig. 2, are calculated for a $\overline{\Xi}$ or $\overline{\Omega}$ momentum of 2708 MeV/c, the momentum of the $\overline{\Omega}$ event. The horizontal axis of each ellipse will shrink slightly with increasing incident momentum, but the effect is small. There is a small region of kinematic ambiguity at the intersection of the two ellipses, but most of the regions are well separated. In Fig. 2 we have plotted the $\overline{\Xi}$ events and the $\overline{\Omega}$ event. In this calculation the $\overline{\Lambda}$ momentum is taken from a three-constraint fit for each $\overline{\Lambda}$ to the charged-vee decay vertex. The meson momentum and the