CHIRAL UNITARY APPROACH TO HADRON SPECTROSCOPY

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Abstract

The s-wave meson-baryon interaction in the $S = -1$, $S = 0$ and $S = -2$ sectors is studied by means of coupled channels, using the lowest-order chiral Lagrangian and the N/D method or equivalently the Bethe-Salpeter equation to implement unitarity. This chiral approach leads to the dynamical generation of the $\Lambda(1405)$, $\Lambda(1670)$ and $\Sigma(1620)$ states for $S = -1$, the $N^*(1535)$ for $S = 0$ and the $\Xi(1620)$ for $S = -2$. We look for poles in the complex plane and extract the couplings of the resonances to the different final states. This allows identifying the $\Lambda(1405)$ and the $\Lambda(1670)$ resonances with $\bar{K}N$ and $K\Xi$ quasibound states, respectively. Our results are found to be incompatible with the measured properties of the $\Xi(1690)$ resonance, thus ruling this state out as the remaining member of this octet of dynamically generated resonances. We therefore assign $1/2^-$ for the spin and parity of the $\Xi(1620)$ resonance as the $S = -2$ member of the lowest-lying $1/2^-$ octet.

The low-energy K^-N scattering and transition to coupled channels is one of the cases of successful application of chiral dynamics in the baryon sector. The studies of Refs.[1, 2] showed that one could obtain an excellent description of the low-energy data starting from chiral Lagrangians and using the multichannel Lippman-Schwinger equation to account for multiple scattering and unitarity in coupled channels. By including all open channels above threshold and fitting a few chiral parameters of the second-order Lagrangian one could obtain a good agreement with the data at low energies. This line of work was continued in Ref.[3], where all coupled channels that could be arranged from the octet of pseudoscalar Goldstone bosons and the entire baryon ground state octet were included . In Ref.[3] it was demonstrated that using the Bethe-Salpeter equation (BSE) with coupled channels and using the lowest-order chiral Lagrangians, together with one cut off to regularize the intermediate meson-baryon loops, a good description of all low-energy data was obtained. One of the novel features with respect to other approaches using the BSE is that the lowest-order meson-baryon amplitudes, playing the role of a potential, could be factorized on shell in the BSE, and thus the set of coupled-channels integral equations became a simple set of algebraic equations, technically simplifying the problem. The justification of this procedure was developed in the treatment of meson-meson interactions using chiral Lagrangians and the N/D method[4]. One uses dispersion relations and shows that neglecting the effects of the left-hand singularity (also shown to be small there) one needs only the on-shell scattering matrix from the lowest-order Lagrangian, and the eventual effects of higher-order Lagrangians are accounted for in terms of subtractions in the dispersion integrals. The N/D method has also been recently applied to study pion-nucleon dynamics[5].

The work of Ref.[3] was reanalyzed recently[6] from the point of view of the N/D method and dispersion relations, leading formally to the same algebraic equations found in Ref.[3]. There are also technical novelties in the regularization of the loop function, which is done using dimensional regularization in Ref.[6], while it was regularized with a cut off in Ref.[3].

One of the common findings shared by all the theoretical approaches is the dynamical generation of the $\Lambda(1405)$ resonance which appears with the right width, and at the correct position, with the choice of a cut off of natural size. This natural generation from the interaction of the meson-baryon system with the lowest-order Lagrangian allows us to identify that state as a quasibound meson-baryon state. This would explain why ordinary quark models have had so many problems explaining this resonance[7].

In ordinary quark models the $\Lambda(1405)$ resonance would mostly be a SU(3) singlet of $J^P = 1/2^-$ and there would be an associated octet of s-wave excited $J^P = 1/2^-$ baryons that would include the N*(1535), the $\Lambda(1670)$, the $\Sigma(1620)$ and a Ξ^* state. In the chiral approach one would also expect the appearance of such a nonet of resonances. In fact, it appears naturally in the approach of Ref.[3], with a degenerate octet, when setting all the masses of the octet of stable baryons equal on one side and the masses of the octet of pseudoscalar mesons equal on the other side. Yet, to obtain this result it is essential that the coupled channels do not omit any of the channels that can be constructed from the octet of pseudoscalar mesons and the octet of stable baryons.

The lowest-order Lagrangian involving the octet of pseudoscalar mesons and the $1/2^+$ baryons is given in Refs.[8, 9, 10, 11].

At lowest order in momentum, that we will keep in our study, the interaction Lagrangian reads

$$L_1^{(B)} = < \bar{B} i \gamma^\mu \frac{1}{4f^2} [(\Phi \partial_\mu \Phi - \partial_\mu \Phi \Phi) B - B (\Phi \partial_\mu \Phi - \partial_\mu \Phi \Phi)] > , \qquad (1)$$

where Φ and B are the SU(3) matrices for the mesons and baryons, respectively and the symbol $<>$ stands for the trace of the resulting SU(3) matrix. The Lagrangian of Eq. (1) leads to a common structure of the type $\bar{u} \gamma^u (k_\mu + k'_\mu) u$

for the different channels, where $u, \bar{u}$ are the Dirac spinors and k, k' the momenta of the incoming and outgoing mesons.

The lowest-order amplitudes for these channels are easily evaluated from Eq. (1) and are given by

$$V_{ij} = -C_{ij} \frac{1}{4f^2} (2\sqrt{s} - M_{Bi} - M_{Bj}) \left(\frac{M_{Bi} + E}{2M_{Bi}} \right)^{1/2} \left(\frac{M_{Bj} + E'}{2M_{Bj}} \right)^{1/2}, \quad (2)$$

with E, E' the energy of the initial, final baryon, and the matrix C_{ij}, which is symmetric, is given in Ref.[3].

Note that the use of physical masses in Eq. (2) effectively introduces some contributions of higher orders in the chiral counting. In the standard chiral approach one would be using the average mass of the octets in the chiral limit and higher order Lagrangians involving SU(3) breaking terms would generate the mass differences. By introducing the physical masses one guarantees that the phase space for the reactions, thresholds and unitarity in coupled channels are respected from the beginning.

Ref.[6], using the N/D method[4] for this particular case, proved that the scattering amplitude could be written by means of the algebraic matrix equation

$$T = [1 - V\,G]^{-1}\,V \quad (3)$$

with V the matrix of Eq. (2) evaluated on shell, or equivalently

$$T = V + V\,G\,T \quad (4)$$

with G a diagonal matrix given by the loop function of a meson and a baryon propagators.

One can see that Eq. (4) is just the Bethe-Salpeter equation but with the V matrix factorized on shell, which allows one to extract the scattering matrix T trivially, as seen in Eq. (3).

The analytical expression for G_l can be obtained from Ref.[12] using a cut off and from Ref.[6] using dimensional regularization. In this latter case one obtains

$$\begin{aligned} G_l &= i2M_l \int \frac{d^4q}{(2\pi)^4} \frac{1}{(P-q)^2 - M_l^2 + i\epsilon} \frac{1}{q^2 - m_l^2 + i\epsilon} \\ &= \frac{2M_l}{16\pi^2} \Big\{ a_l(\mu) + \ln \frac{M_l^2}{\mu^2} + \frac{m_l^2 - M_l^2 + s}{2s} \ln \frac{m_l^2}{M_l^2} + \\ &\quad + \frac{\bar{q}_l}{\sqrt{s}} \left[\ln(s - (M_l^2 - m_l^2) + 2\bar{q}_l\sqrt{s}) + \ln(s + (M_l^2 - m_l^2) + 2\bar{q}_l\sqrt{s}) \right. \\ &\quad \left. - \ln(-s + (M_l^2 - m_l^2) + 2\bar{q}_l\sqrt{s}) - \ln(-s - (M_l^2 - m_l^2) + 2\bar{q}_l\sqrt{s}) \right] \Big\} \quad (5) \end{aligned}$$

which has been rewritten in a convenient way to show how the imaginary part of G_l is generated and how one can go to the unphysical Riemann sheets in order to identify the poles. The dimensional regularization scheme is preferable if one goes to higher energies where the on-shell momentum of the intermediate states is not much smaller than the cut off.

1 Strangeness $S = -1$ sector

We take the $K^- p$ state and all related channels using SU(3) mesons and baryons within the chiral approach, namely $\bar{K}^0 n$, $\pi^0 \Lambda$, $\pi^0 \Sigma^0$, $\pi^+ \Sigma^-$, $\pi^- \Sigma^+$, $\eta \Lambda$, $\eta \Sigma^0$, $K^0 \Xi^0$ and $K^+ \Xi^-$. Hence we have a problem with ten coupled channels. The coupled set of Eqs. (3) were solved in Ref.[3] using a cut off momentum of 630 MeV in all channels. Changes in the cut off can be accommodated in terms of changes in μ, the regularization scale in the dimensional regularization formula for G_l, or in the subtraction constant a_l. In order to obtain the same results as in Ref.[3] at low energies, we set μ equal to the cut off momentum of 630 MeV (in all channels) and then find the values of the subtraction constants a_l such as to have G_l with the same value with the dimensional regularization formula (Eq. (5)) and the cut off formula of [3] at the $\bar{K}N$ threshold. This determines the values

$$\begin{array}{lll} a_{\bar{K}N} = -1.84 & a_{\pi\Sigma} = -2.00 & a_{\pi\Lambda} = -1.83 \\ a_{\eta\Lambda} = -2.25 & a_{\eta\Sigma} = -2.38 & a_{K\Xi} = -2.52 \ . \end{array} \qquad (6)$$

This guarantees that we obtain the same results at low energies as in Ref.[3] and we find indeed that this is the case when we repeat the calculation with the new G_l of Eq. (5). Then we extend the results at higher energies, looking for the possible appearance of new resonances.

For the purpose of this study let us recall that Ref.[3] obtained the $\Lambda(1405)$ resonance which we show in Fig. 1 obtained from the $\pi\Sigma$ spectrum. Next we go to higher energies and search for new resonances.

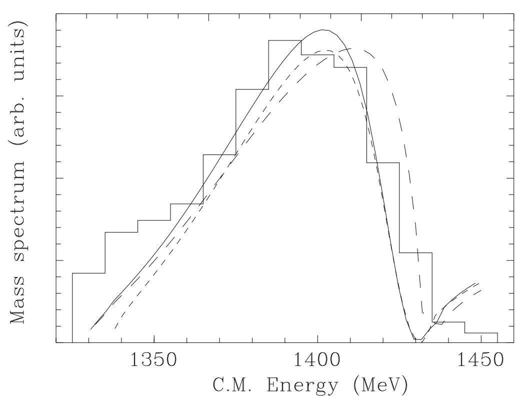

Figure 1: The $\Lambda(1405)$ resonance obtained from the invariant $\pi\Sigma$ mass distribution, with the full basis of physical states (solid line), omitting the η channels (long-dashed line) and with the isospin-basis (short-dashed line).

Fig. 2 shows the real and imaginary parts of the $I = 0$ scattering amplitude, obtained in Ref.[13] normalized as in the Partial Wave Analysis of Ref. [14]. Remarkably, the amplitudes shown by the solid lines, which are obtained using

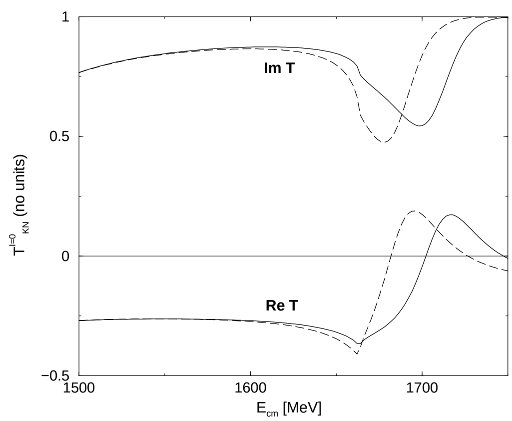

Figure 2: Real and imaginary parts of the $\bar{K}N$ scattering amplitude in the isospin $I = 0$ channel in the region of the $\Lambda(1670)$ resonance.

the low-energy parameters in Eq. (6), show the resonant structure of the $\Lambda(1670)$ appearing at about the right energy and with a similar size compared to the experimental analysis[14]. The position of the resonance is quite sensitive to the parameter $a_{K\Xi}$ and moderately sensitive to $a_{\eta\Lambda}$. Hence, without spoiling the nice agreement at low energies, which is not sensitive to $a_{K\Xi}$, we exploit the freedom in the parameters by choosing $a_{K\Xi} = -2.70$, moving the resonance closer to its experimental position (dashed lines).

SU(3) symmetry, partly broken here due to the use of physical masses, demands a singlet and an octet of resonances. Within $S = -1$, we have already identified the singlet $\Lambda(1405)$ and the $I = 0$ member of the octet, the $\Lambda(1670)$. Since we found the partial decay widths and couplings of the $\Lambda(1405)$ to $\bar{K}N$ states and the $\Lambda(1670)$ to $K\Xi$ states to be very large, one is naturally led to identify these two resonances as a "quasibound" $\bar{K}N$ and $K\Xi$ state, respectively.

Searching for the $I = 1$ member of the octet, we find that the $I = 1$ amplitudes in our model are smooth and show no trace of resonant behavior, in line with experimental observation. To explore this issue further we conducted a search for the poles of the $\bar{K}N \to \bar{K}N$ amplitudes in the second Riemann sheet and find two poles in the $I = 0$ amplitude $(1426 + \mathrm{i}16,\ 1708 + \mathrm{i}21)$, corresponding to the $\Lambda(1405)$ and the $\Lambda(1670)$, and one in the $I = 1$ amplitude $(1579 + \mathrm{i}296)$, corresponding - most likely - to the resonance $\Sigma(1620)$. The large width found for this resonance may explain why we saw no trace of it in the scattering amplitudes.

2 Strangeness $S = 0$ sector

The strangeness $S = 0$ channel was also investigated using the Lippmann-Schwinger equation and coupled channels in Ref.[1, 15]. The $N^*(1535)$ resonance was found to be generated dynamically within this approach. Subsequently, work was done in this sector using the procedure of Ref.[5] with subtraction constants in Ref.[16], and the $N^*(1535)$ resonance, as well as the low-energy scattering observables, were well reproduced. The exception was the isospin 3/2 channel which was not reproduced in this approach nor in [1, 15], and neither in [18] where the $N^*(1535)$ resonance was generated together with the $N^*(1650)$ at the expense of using more free parameters.

Ref.[17] continued and further improved work along these lines by introducing the $\pi N \to \pi NN$ channels, which proved essential in reproducing the isospin 3/2 part of the πN amplitude.

For total zero charge one has six channels in this case, $\pi^- p$, $\pi^0 n$, ηn, $K^+\Sigma^-$, $K^0\Sigma^0$, and $K^0\Lambda$. The subtraction parameters $a_i(\mu)$ for the meson-baryon propagators that we obtain in the new fit to the data, are

$$\mu = 1200 \text{ MeV}, \quad a_{\pi N}(\mu) = 2.0, \quad a_{\eta N}(\mu) = 0.1, \quad a_{K\Lambda}(\mu) = 1.5, \quad a_{K\Sigma}(\mu) = -2.8 \tag{7}$$

The results obtained for the phase shifts and inelasticities of πN scattering are shown in Fig. 3, where the continuous line corresponds to the calculation while the dotted line is the experimental analysis.

The $N^*(1535)$ peak is visible in the phase shifts and inelasticities of the S_{11} amplitude in the panel. The peak of the $\pi^- p \to \eta n$ cross section is also well reproduced but strength is missing just after the peak indicating the contribution of higher order partial waves. Figure 3 shows that the S_{31} data are also fairly well reproduced once the $\pi N \to \pi\pi N$ channels are introduced in the approach. These amplitudes are fitted simultaneously to the scattering data and the $\pi N \to \pi\pi N$ cross sections and are somewhat different than those determined previously in Ref.[19, 20].

3 Strangeness $S = -2$ sector

Here we focus on the $S = -2$ sector for which,i.e., the zero-charge states of the coupled-channels are $\pi^+\Xi^-$, $\pi^0\Xi^0$, $\bar{K^0}\Lambda$, $K^-\Sigma^+$, $\bar{K^0}\Sigma^0$ and $\eta\Xi^0$.

In the study of $S = -1$ resonances performed in Ref.[13] the a_l parameters were extracted by matching the results to those of Ref.[3] and the range of values obtained, from -1.84 to -2.67, serves as an indication for what we might assume as reasonable natural size parameters in the present $S = -2$ study. We search for poles in the second Riemann sheet of the scattering amplitude, focussing on the elastic $\pi\Xi \to \pi\Xi$ amplitude in the $I = 1/2$ channel. As a trial run, we set the four values of the subtraction constants to a value of -2 and we discover a pole at $1607 + i140$ MeV. This would lead to a width around 280 MeV, unacceptably large compared to those of the two I=1/2 resonances of interest, the $\Xi(1620)$ and the $\Xi(1690)$, which are reported to be of the order of 50 MeV or less. The mass of the particle, around 1607 MeV, would be closer to the $\Xi(1620)$ resonance.

Allowing the subtraction constants a_l to change within a reasonable natural range, we obtain the results shown in Table 1. Only $a_{\pi\Xi}$ and $a_{\bar{K}\Lambda}$ are varied,

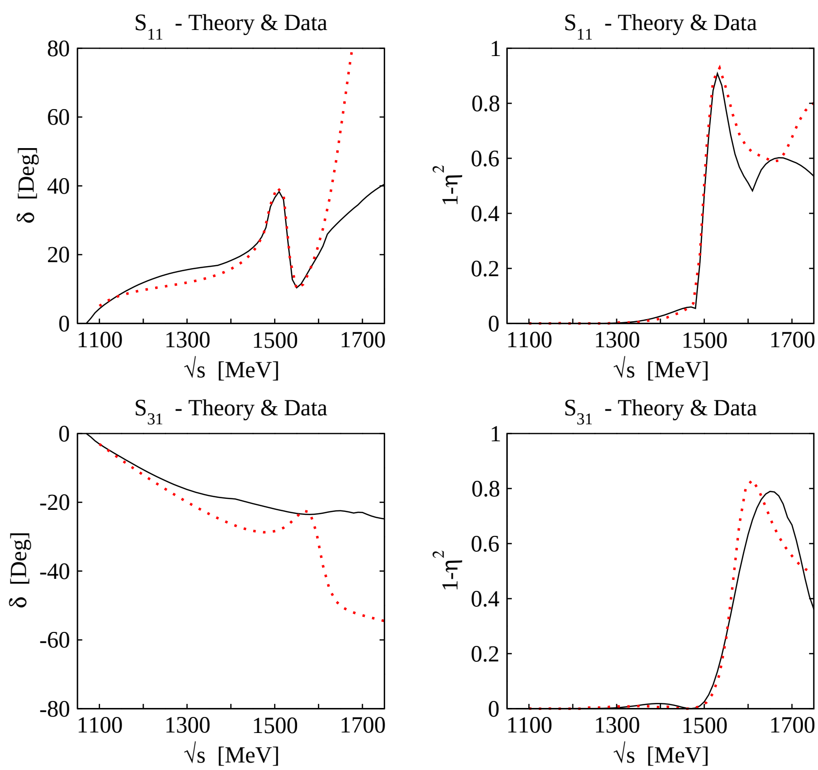

Figure 3: Phase-shifts and inelasticities of S_{11} and S_{31} πN scattering with $\pi\pi N$ channels.

since we find the couplings of the resonance to the $\bar{K}\Sigma$ and $\eta\Xi$ states to be very weak and therefore the results are insensitive to the subtraction constants corresponding to these two channels. The values of the couplings, calculated from the residue of the diagonal scattering amplitudes [13], are also shown in Table 1.

Table 1: Resonance properties for various sets of subtraction constants

	Set 1	Set 2	Set 3	Set 4	Set 5
$a_{\pi\Xi}$	-2.0	-2.2	-2.0	-2.5	-3.1
$a_{\bar{K}\Lambda}$	-2.0	-2.0	-2.2	-1.6	-1.0
$a_{\bar{K}\Sigma}$	-2.0	-2.0	-2.0	-2.0	-2.0
$a_{\eta\Xi}$	-2.0	-2.0	-2.0	-2.0	-2.0
$\mid g_{\pi\Xi} \mid^2$	8.7	7.2	7.4	7.2	5.9
$\mid g_{\bar{K}\Lambda} \mid^2$	5.5	4.6	4.2	5.8	7.0
$\mid g_{\bar{K}\Sigma} \mid^2$	0.68	0.59	0.54	0.74	0.93
$\mid g_{\eta\Xi} \mid^2$	0.36	0.27	0.38	0.14	0.23
M	1607	1597	1596	1604	1605
$\Gamma/2$	140	117	134	98	66

The second and third columns in Table 1 show that a change of 10% in the subtraction constants $a_{\pi\Xi}$ and $a_{\bar{K}\Lambda}$ modifies the mass of the resonance only slightly but has a larger influence on the width. Investigating the dependence of the results on the values of these two subtraction constants we observe that the mass of the resonance is confined to a range around 1600 MeV. The width, on the other hand, can be reduced considerably by a simultaneous increase of the strength of $a_{\pi\Xi}$ and a decrease of $a_{\bar{K}\Lambda}$, while keeping both of them negative and still reasonably close to the reference value of -2. In the last column we see that the width can be reduced to 130 MeV with acceptable values for the coefficients. While this width might still appear as grossly overestimating the experimental ones, we show below that this is not the case.

Since the $\Xi(1620)$ resonance decays only into $\pi\Xi$ final states, it is experimentally visible through the $\pi\Xi$ invariant mass distribution in reactions leading, among others, to π and Ξ particles. Our calculated distribution, displayed in Fig. 4, shows a very interesting feature, namely a smaller apparent width compared to the one obtained at the pole position. For instance, for the values of the subtraction constants in the last column of Table 1 we see in Fig. 4 (solid line) an apparent Breit-Wigner width of around 50 MeV and a shape for the distribution which resembles the experimental peaks observed. This well-known phenomenon, usually referred to as Flatté effect [21], is due to the presence of a resonance just below the threshold of a channel to which the resonance couples very strongly. In our case the $\bar{K}\Lambda$ channel opens at 1611 MeV and, as shown in Table 1, the resonance couples very strongly to that state. What actually happens is that at an invariant energy close to the resonance mass the amplitude is given essentially by the inverse of the resonance width. As soon as the threshold

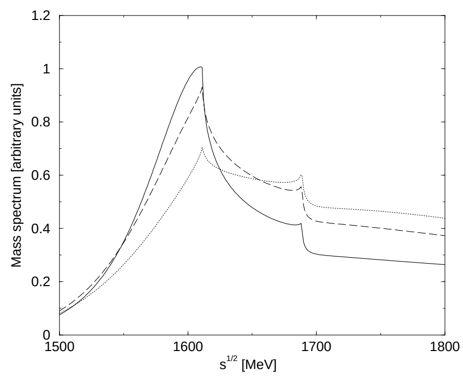

Figure 4: The $\pi\Xi$ invariant mass distribution as a function of the center-of-mass energy, for several sets of subtraction constants. Solid line: $a_{\pi\Xi} = -3.1$ and $a_{\bar{K}\Lambda} = -1.0$; Dashed line: $a_{\pi\Xi} = -2.5$ and $a_{\bar{K}\Lambda} = -1.6$; Dotted line: $a_{\pi\Xi} = -2.0$ and $a_{\bar{K}\Lambda} = -2.0$. The value of the remaining two other subtraction constants, $a_{\bar{K}\Sigma}$ and $a_{\eta\Xi}$, is fixed to -2.0 in all curves.

is crossed, the new channel leads to an additional energy-dependent contribution for the width which grows very rapidly with increasing energy. This produces a fast fall-off for the amplitude, leading to an apparent width much smaller than the actual width at the pole. This phenomenon has been observed, e.g., in the case of the $a_0(980)$ meson resonance as discussed in Refs. [12, 4].

The question now arises which of the two $I = 1/2$ candidates should be identified with the resonance obtained here. The value found for the mass of the state would suggest identification with the $\Xi(1620)$ which is rated as a one-star resonance in the PDG and has unknown spin and parity. The $\Xi(1690)$ state is better known and is rated as a 3-star resonance. Even if the spin and parity are unknown, there is far more information available for this resonance than for the $\Xi(1620)$ [22]. Ref. [23] gives ratios of partial decay widths having sufficient accuracy for us to draw conclusions from the properties of the Ξ resonance found in this work. We therefore investigate whether the parameters of the theory provide enough flexibility to produce a pole with a real part closer to 1690 MeV, since the results of Table 1 show that, by decreasing the size of $a_{\pi\Xi}$ or $a_{\bar{K}\Lambda}$, one increases the mass of the resonance. However, the presence of the $\bar{K}\Lambda$ threshold leads to mass values that stabilize around the cusp of this threshold for a certain range of the parameters. Continuing to change the a_l parameters beyond this range does not increase the resonance mass but leads to a disappearance of the pole – and with it the resonance. The above argument clearly favors identifying the resonance found here with the $\Xi(1620)$ state.

The other argument in favor of the $\Xi(1620)$ assignment is the following: The results of Table 1 show that the resonance couples strongly to the $\pi\Xi$ and the $\bar{K}\Lambda$ channels but very weakly to $\bar{K}\Sigma$ and $\eta\Xi$. This is opposite to the observed properties of the $\Xi(1690)$ resonance, for which Ref. [23] gives a ratio of branching ratios for $\bar{K}\Sigma$ to $\bar{K}\Lambda$ around 3 and for $\pi\Xi$ to $\bar{K}\Sigma$ of less than 0.09. In our opinion, this argument rules out identifying the resonance found here with the $\Xi(1690)$ state.

In summary, we have demonstrated that the chiral approach to the $\bar{K}N$ and the other coupled channels, which proved so successful at low energies, extrapolates smoothly to higher energies and provides the basic features of the scattering amplitudes, generating the resonances which would complete the states of the nonet of the $J^P = 1/2^-$ excited states. The qualitative description of the data without adjusting any parameters is telling us that the basic information on the dynamics of these processes is contained in the chiral Lagrangians. There is still some freedom left with the chiral symmetry breaking terms. In our formulation they would go into the a_l subtraction constants, and the use of different decay constants for each meson, by means of which one could obtain a better description of the data. The analysis of the poles and the couplings of the resonances to the different channels lead us to identify the strong coupling of the $\Lambda(1405)$ resonance to the $\bar{K}N$ state and the large coupling of the $\Lambda(1670)$ resonance to the $K\Xi$ state, allowing us to classify these resonances as quasibound states of $\bar{K}N$ and $K\Xi$, respectively. In the $S = -2$ sector, the study performed here has allowed us to identify the resonance generated dynamically with the $\Xi(1620)$, hence providing a prediction for the spin and parity of this resonance which is currently not assigned by PDG.

Acknowledgments

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