

Spin, Parity, and Nature of the $\Xi(1620)$ Resonance

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Using a unitary extension of chiral perturbation theory with a lowest-order s -wave SU(3) chiral Lagrangian we study low-energy meson-baryon scattering in the strangeness $S = -2$ sector. A scattering-matrix pole is found around 1605 MeV which corresponds to an s -wave Ξ resonance with $J^P = 1/2^-$. We identify this resonance with the $\Xi(1620)$ state, quoted by the Particle Data Group with $I = 1/2$ but with unknown spin and parity. The addition of the $S = -2$ state to the recently computed $\Lambda(1670)$, $\Sigma(1620)$, and $N(1535)$ states completes the octet of $J^P = 1/2^-$ resonances dynamically generated in this chiral unitary approach.

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Since the discovery of the first baryon resonance, the $\Delta(1232)$, over 50 years ago, there has been a persistent question regarding the nature of such resonances: Are these genuine states that appear as bare fields in the driving term of the meson-baryon scattering matrix? Or can they be generated dynamically, i.e., by iterating an appropriate nonpolar driving term to all orders? Early attempts by Chew and Low to generate the $\Delta(1232)$ by iterating a crossed nucleon pole term were eventually overtaken by the success of the SU(3) quark models which established the $\Delta(1232)$ as part of the SU(3) ground-state decuplet. The subsequent discovery of over 100 additional baryon resonances and their mostly successful incorporation into quark models appeared to have settled the question in favor of treating them as genuine fields. However, one persistent exception for many years was the lowest-lying $S = -1$ resonance, the SU(3) singlet $\Lambda(1405)$, which appeared quite naturally as a dynamical pole in the K^-p scattering matrix using a variety of approaches [see the discussion of the Particle Data Group [1] about the history of the $\Lambda(1405)$]. The advent of chiral Lagrangians combined with unitarization techniques placed these efforts on more solid theoretical grounds [2–5]. Next came the suggestion that the $N^*(1535)$, which in some works is considered the $J^P = 1/2^-$ chiral partner of the nucleon and hence would be degenerate with the nucleon were it not for chiral symmetry breaking [6,7], can also be generated dynamically, extending the same techniques into the $S = 0$ sector [2,3,8–11]. Invoking the same framework but extending it to higher energy, Ref. [12] returned the search to the $S = -1$ sector and identified the $\Lambda(1670)$ and $\Sigma(1620)$ resonances as two additional members of the same $J^P =$

$1/2^-$ octet that appear naturally within the chiral unitary framework.

The purpose of this paper is to extend this chiral approach with unitarization yet again to the $S = -2$ sector and demonstrate that the last remaining member of the lowest-lying $J^P = 1/2^-$ octet, an s -wave Ξ resonance, can indeed be found as a pole in the appropriate meson-baryon scattering matrix. The experimental situation of the $S = -2$ low-lying Ξ resonances is rather unclear. There is a well-known p -wave state, the $P_{13}\Xi(1530)$ resonance, with $I(J^P) = 1/2(3/2^+)$, which is 4-star rated by the Particle Data Group [1]. Next come the $\Xi(1620)$ and $\Xi(1690)$ resonances which are rated with one and three stars, respectively. While an isospin value of $I = 1/2$ is quoted for both resonances, the spin and parity of these states have not been measured. We note that the $\Xi(1620)$ and $\Xi(1690)$ resonances are of too low mass to be accommodated in most quark models [13,14], which produce Ξ states at energies of the order or above 1800 MeV. In the present paper we present arguments that the $\Xi(1620)$ corresponds to the lowest-lying $J^P = 1/2^-$ octet of baryon resonances generated dynamically through multiple scattering of meson-baryon pairs in a coupled-channels approach.

The states are generated in the following way. One constructs the set of coupled channels from the octets of ground-state baryons and pseudoscalar mesons and, using the SU(3) version of the chiral Lagrangians for mesons and baryons [15–18], one implements any of the unitary extensions of chiral perturbation theory to generate the scattering amplitudes connecting the various channels. Recent unitary extensions of chiral perturbation theory (U χ PT) rely upon the inverse amplitude

method of [19] extended to coupled channels [20] or the N/D method including explicit fields of genuine resonances (which would survive in the large N_c limit) [21], also used in [22] to study πN scattering. This latter method was also used in [5] to study the low-energy interaction of antikaons with baryons and it was found to be equivalent to the formalism used in [4] to study the same problem using the Bethe-Salpeter equation with coupled channels, with the difference that the loops were regularized in [5] with dimensional regularization, while a cutoff was applied in [4].

In order to search for a Ξ resonance we follow closely the work of Ref. [12], where the $\Lambda(1670)$ and $\Sigma(1620)$ resonances were discovered in the $S = -1$ channel. Here we focus on the $S = -2$ sector for which, as an example, the zero-charge states of the coupled-channels framework

$$V_{ij} = -C_{ij} \frac{1}{4f^2} (2\sqrt{s} - M_i - M_j) \left(\frac{M_i + E_i}{2M_i} \right)^{1/2} \left(\frac{M_j + E_j}{2M_j} \right)^{1/2} \quad (2)$$

is obtained from the chiral Lagrangian for the meson-baryon interaction at lowest order and C_{ij} are SU(3) coefficients with i and j denoting channel indices. The (diagonal) matrix G in Eq. (1) accounts for the loop integral of a meson and a baryon propagator and depends on the regularization scale, μ , and a subtraction constant for each channel, a_l , that comes from a subtracted dispersion relation. The explicit expression of G can be found in Refs. [5,12]. The regularization scale μ is of course arbitrary but the subtraction constants depend on it.

Our study relies upon coupled channels of one meson and one baryon. Although there is phase space for the production of one extra pion, these channels proved to be rather unimportant in the study of πN and coupled channels in $I = 1/2$ around the $N^*(1535)$ resonance region [11], and for the states studied here the only experimental evidence of their contribution is in a negligible branching ratio of the $\Xi(1690)$ into $\Xi\pi\pi$ [1].

In the present work we use the isospin basis, which contains the states $\pi\Xi$, $\bar{K}\Lambda$, $\bar{K}\Sigma$, and $\eta\Xi$ for isospin $I = 1/2$ and the states $\pi\Xi$ and $\bar{K}\Sigma$ for isospin $I = 3/2$. For the particular isospin $I = 1/2$ case needed in the present study, the coefficients C_{ij} have the values shown in Table I. We have four different subtraction constants, namely $a_{\pi\Xi}$, $a_{\bar{K}\Sigma}$, $a_{\bar{K}\Lambda}$, and $a_{\eta\Xi}$. Assuming that the regularization scale μ gives the size of the maximum momentum in the cutoff regularization method, then the values of a_l can be deduced approximately from [5] and they are found to be of the order of -2 when a cutoff of 630 MeV is chosen as done in [4]. This value for the subtraction constant is what is called a magnitude of natural size in [5], since it corresponds to using cutoffs of the order of 1 GeV which are considered natural in the chiral approaches. We point out that the choice of the coefficients a_l accounts for contributions of higher-order Lagrangians to the process, as shown in [20].

In the present $S = -2$ study we set, as a trial run, the four values of the subtraction constants to a value of -2

TABLE I. Coefficients C_{ij} of the meson baryon amplitudes for isospin $I = 1/2$ ($C_{ji} = C_{ij}$).

	$\pi\Xi$	$\bar{K}\Lambda$	$\bar{K}\Sigma$	$\eta\Xi$
$\pi\Xi$	2	$-\frac{3}{2}$	$-\frac{1}{2}$	0
$\bar{K}\Lambda$		0	0	$-\frac{3}{2}$
$\bar{K}\Sigma$			2	$\frac{3}{2}$
$\eta\Xi$				0

are $\pi^+\Xi^-$, $\pi^0\Xi^0$, $\bar{K}^0\Lambda$, $K^-\Sigma^+$, $\bar{K}^0\Sigma^0$, and $\eta\Xi^0$. We solve the coupled-channels Bethe-Salpeter equation for the scattering amplitude,

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

where the driving (kernel) V matrix

and we discover a pole in the second Riemann sheet of the $I = 1/2$ $\pi\Xi \rightarrow \pi\Xi$ amplitude 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.

Demanding the subtraction constants to be of natural size is important. Indeed, our claims that the $\Xi(1620)$ resonance is dynamically generated rely upon the ability of the present approach to generate it from the lowest-order Lagrangian using natural size cutoffs or subtraction constants. An alternative method would use the lowest-order Lagrangian and bare resonances, whose couplings to the physical channels would be given by fits to the data. This is done in [22] for πN scattering and in [21] in meson-meson scattering. In such a case, the bare resonances, such as the Δ in [22] and the ρ in [21], require finite couplings to the physical channels; such couplings would be compatible with zero for the dynamically generated resonances. If no bare resonant fields are included explicitly, using a lowest-order Lagrangian may either not generate such resonances at all, or they would require unphysical cutoff or subtraction constant (a cutoff of the order of 1 TeV for the ρ according to [21,23]).

Allowing the subtraction constants a_l to change within a reasonable natural range, we obtain the results shown in Table II. Only $a_{\pi\Xi}$ and $a_{\bar{K}\Lambda}$ are varied, since we find the couplings of the resonance to the $\bar{K}\Sigma$ and $\eta\Xi$ states to be very weak. The values of the couplings, calculated from the residue of the diagonal scattering amplitudes [12], are also shown in Table II. The second and third columns 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 fifth 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. 1, shows a smaller apparent width compared to the one obtained at the pole position. For the values of the subtraction constants in the fifth column of Table II we see in Fig. 1 (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 the Flatté effect [24], 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 II, 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 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-

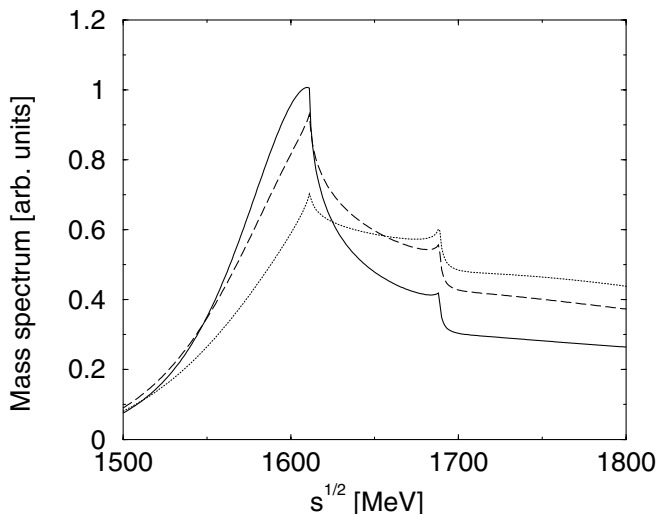


FIG. 1. The $\pi\Xi$ invariant mass distribution as a function of the center-of-mass energy, for several sets of subtraction constants. Solid line: $a_{\pi\Xi} = -3.1$ and $a_{\bar{K}\Lambda} = -1.0$; Dashed line: $a_{\pi\Xi} = -2.5$ and $a_{\bar{K}\Lambda} = -1.6$; Dotted line: $a_{\pi\Xi} = -2.0$ and $a_{\bar{K}\Lambda} = -2.0$. 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.

off for the amplitude, leading to an apparent width much smaller than the actual width at the pole.

We have also analyzed effects of flavor symmetry breaking by replacing the f^2 factor appearing in Eq. (2), where f is taken as an average value $f = 1.123f_\pi$, by the product of physical decay constants $f_i f_j$, where $f_\pi = 93$ MeV, $f_K = 1.22f_\pi$, and $f_\eta = 1.3f_\pi$. The results are displayed in the last column of Table II and should be compared to those in the first column. The changes observed are moderate and do not alter the analysis and conclusions of the present work.

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 $\Xi(1620)$. The experimental situation regarding this 1-star resonance is the following: Three experiments see the resonance in the $\Xi\pi$ spectrum in reactions where K^-p goes to three or four particles in the final state, including $\Xi\pi$ [25–27]. Taking into account the experimental errors of all experiments, the resonance mass is found in the range 1600–1640 MeV and the width in the range 15–50 MeV. The $\Xi(1690)$ state is better known and is rated as a 3-star resonance. Reference [28] even gives quite accurate ratios of partial decay widths. 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 II 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 these a_l parameters beyond this range leads to a disappearance of the pole—and with it the resonance—entirely. The above argument clearly favors identifying the resonance found here with the $\Xi(1620)$ state and leads us to conclude that the $\Xi(1690)$ is most likely an intrinsic or genuine resonance using the terminology of Ref. [29].

The other argument in favor of the $\Xi(1620)$ assignment is the following: The results of Table II 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. [28] 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.

We also point out here that using QCD sum rules one obtains a $J^P = 1/2^-$ octet of excited baryons where the Σ and the Ξ states appear degenerate and with a mass around 1620 MeV [30]. Furthermore, the lowest Ξ resonances have been seen on the lattice, however, these results do not yet allow any quantitative conclusions [31].

In summary, we find a Ξ resonance with an energy around 1606 MeV and a width at the pole position around 100 MeV. Because of a significant threshold effect the apparent (Breit-Wigner) width is much smaller and

TABLE II. Resonance properties for various sets of subtraction constants.

	Set 1	Set 2	Set 3	Set 4	Set 5	Set 1 (physical f_i)
$a_{\pi\Xi}$	-2.0	-2.2	-2.0	-2.5	-3.1	-2.0
$a_{\bar{K}\Lambda}$	-2.0	-2.0	-2.2	-1.6	-1.0	-2.0
$a_{\bar{K}\Sigma}$	-2.0	-2.0	-2.0	-2.0	-2.0	-2.0
$a_{\eta\Xi}$	-2.0	-2.0	-2.0	-2.0	-2.0	-2.0
$ g_{\pi\Xi} ^2$	8.7	7.2	7.4	7.2	5.9	6.7
$ g_{\bar{K}\Lambda} ^2$	5.5	4.6	4.2	5.8	7.0	3.4
$ g_{\bar{K}\Sigma} ^2$	0.68	0.59	0.54	0.74	0.93	0.36
$ g_{\eta\Xi} ^2$	0.36	0.27	0.38	0.14	0.23	0.09
M	1607	1597	1596	1604	1605	1568
$\Gamma/2$	140	117	134	98	66	122

compatible with experimental findings quoted in [1] for the Ξ resonances in question. Of the two relevant Ξ states with isospin 1/2, the 1-star rated $\Xi(1620)$ and the 3-star rated $\Xi(1690)$, we argued that the $\Xi(1690)$ must be ruled out because, on the one hand, it was impossible within our approach to find a pole with an energy close to 1690 MeV and, on the other hand, there are large disagreements between the resonance couplings to meson-baryon states found here and the measured partial decay widths to those states [28]. The findings presented here indicate that overwhelming evidence supports the assignment of the quantum numbers $J^P = 1/2^-$ to the $\Xi(1620)$ resonance. Thus, the computation within the chiral unitary approach of the $\Xi(1620)$, along with its partners, the $N^*(1535)$, the $\Lambda(1670)$, and the $\Sigma(1620)$, plus the $\Lambda(1405)$ which largely accounts for the SU(3) singlet, completes the $J^P = 1/2^-$ nonet of dynamically generated s -wave resonances. Clearly, the special nature of the resonance discussed here calls for renewed experimental efforts, especially for the 1-star rated $\Xi(1620)$, $\Sigma(1620)$ states.

Since the SU(3) octet of low-lying $J^P = 1/2^-$ resonances represents the chiral partner of the $J^P = 1/2^+$ ground-state octet, the description of these states demonstrates the extraordinary power of the chiral unitary approach. That these states should appear especially with as simple a driving term as provided by the first-order chiral Lagrangian with very few open parameters is truly remarkable and shows that the nature of these states is well represented by a cloud of meson-baryon components. This would be in principle different to other excited states in higher partial waves or higher excitation energy which might be better represented in terms of their quark constituents.

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- [1] D. E. Groom *et al.*, Eur. Phys. J. C **15**, 1 (2000).
- [2] N. Kaiser, P. B. Siegel, and W. Weise, Nucl. Phys. A **594**, 325 (1995).
- [3] N. Kaiser, T. Waas, and W. Weise, Nucl. Phys. A **612**, 297 (1997).
- [4] E. Oset and A. Ramos, Nucl. Phys. A **635**, 99 (1998).
- [5] J. A. Oller and U. G. Meissner, Phys. Lett. B **500**, 263 (2001).
- [6] C. De Tar and T. Kunihiro, Phys. Rev. D **39**, 2805 (1989).
- [7] D. Jido, M. Oka, and A. Hosaka, Phys. Rev. Lett. **80**, 448 (1998).
- [8] N. Kaiser, P. B. Siegel, and W. Weise, Phys. Lett. B **362**, 23 (1995).
- [9] J. C. Nacher *et al.*, Nucl. Phys. A **678**, 187 (2000).
- [10] J. Nieves and E. R. Arriola, Phys. Rev. D **64**, 116008 (2001).
- [11] T. Inoue, E. Oset, and M. J. Vicente Vacas, Phys. Rev. C **65**, 035204 (2002).
- [12] E. Oset, A. Ramos, and C. Bennhold, Phys. Lett. B **537**, 99 (2002).
- [13] S. Capstick and N. Isgur, Phys. Rev. D **34**, 2809 (1986).
- [14] W. H. Blask *et al.*, Z. Phys. A **337**, 327 (1990), and references therein.
- [15] A. Pich, Rep. Prog. Phys. **58**, 563 (1995).
- [16] G. Ecker, Prog. Part. Nucl. Phys. **35**, 1 (1995).
- [17] V. Bernard, N. Kaiser, and U. G. Meissner, Int. J. Mod. Phys. E **4**, 193 (1995).
- [18] U. G. Meissner, Rep. Prog. Phys. **56**, 903 (1993).
- [19] A. Dobado and J. R. Pelaez, Phys. Rev. D **56**, 3057 (1997).
- [20] J. A. Oller, E. Oset, and J. R. Pelaez, Phys. Rev. D **59**, 074001 (1999); **60**, 099906 (E) (1999).
- [21] J. A. Oller and E. Oset, Phys. Rev. D **60**, 074023 (1999).
- [22] U. G. Meissner and J. A. Oller, Nucl. Phys. A **673**, 311 (2000).
- [23] H. Leutwyler, Nucl. Phys. (Proc. Suppl.) **7A**, 42 (1989).
- [24] S. M. Flatté, Phys. Lett. B **63**, 224 (1976).
- [25] R. T. Ross *et al.*, Phys. Lett. B **38**, 177 (1972).
- [26] E. Briefel *et al.*, Phys. Rev. D **16**, 2706 (1997).
- [27] A. de Bellefon *et al.*, Nuovo Cimento Soc. Ital. Fis. **28A**, 289 (1975).
- [28] C. Dionisi *et al.*, Phys. Lett. B **80**, 145 (1978).
- [29] C. M. Shakin and H. Wang, Phys. Rev. D **63**, 014019 (2000).
- [30] D. Jido and M. Oka, hep-ph/9611322.
- [31] W. Melnitchouk *et al.*, hep-lat/0202022.