

International Journal of Modern Physics A
© World Scientific Publishing Company

DYNAMICALLY GENERATED RESONANCES IN THE CHIRAL UNITARY APPROACH TO MESON BARYON INTERACTION

E. OSET, S. SARKAR, M.J. VICENTE VACAS

*Departamento de Física Teórica and IFIC, Centro Mixto Universidad de Valencia-CSIC,
Institutos de Investigación de Paterna, Aptd. 22085, 46071 Valencia, Spain*

A. RAMOS

*Departament d'Estructura i Constituents de la Matèria, Universitat de Barcelona,
Diagonal 647, 08028 Barcelona, Spain*

D. JIDO

Physik-Department, Technische Universität München D-85747 Garching, Germany

J.A. OLLER

Departamento de Física, Universidad de Murcia, 30071 Murcia, Spain

U.G. MEISSNER

*Universität Bonn, Helmholtz-Institut für Strahlen- und Kernphysik (Theorie)
Nußallee 14-16, D-53115 Bonn, Germany*

Received (Day Month Year)

Revised (Day Month Year)

In this talk we report on the use of a chiral unitary approach for the interaction of the octets of meson and baryon and the octet of mesons with the decuplet of baryons. Two octets of $J^\pi = 1/2^-$ baryon states and a singlet are generated dynamically in the first case, resulting in the case of strangeness $S = -1$ in two poles of the scattering matrix close to the nominal $\Lambda(1405)$ resonance. In the second case many resonances are also generated, among them an exotic baryon with $S = 1$ corresponding to a ΔK resonance. We make suggestions of experiments which could show evidence for the existence of these states.

Keywords: Keyword1; keyword2; keyword3.

1. Description of the meson baryon interactions

The introduction of unitarity constraints in coupled channels in chiral perturbation theory has led to unitary extensions of the theory that starting from the same effective Lagrangians allow one to make predictions at much higher energies. One of the interesting consequences of these extensions is that they generate dynamically low lying resonances, both in the mesonic and baryonic sectors. By this we mean that

2 Authors' Names

they are generated by the multiple scattering of the meson or baryon components, much the same as the deuteron is generated by the interaction of the nucleons through the action of a potential, and they are not preexistent states that remain in the large N_c limit where the multiple scattering is suppressed.

Starting from the chiral Lagrangians for the interaction of the octets of meson and baryon ¹ and using the N/D method to obtain a scattering matrix fulfilling exactly unitarity in coupled channels ², the full set of transition matrix elements with the coupled channels is given in matrix form by

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

Here, the matrix V , obtained from the lowest order meson–baryon chiral Lagrangian, contains the Weinberg-Tomozawa or seagull contribution, as employed e.g. in Ref. ³,

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

For strangeness $S = -1$ the C_{ij} coefficients are given in Ref. ⁴, and an averaged meson decay constant $f = 1.123f_\pi$ is used ³, with $f_\pi = 92.4$ MeV the weak pion decay constant.

The diagonal matrix G stands for the loop function of a meson and a baryon and is defined by a subtracted dispersion relation in terms of phase space with a cut starting at the corresponding threshold ². It corresponds to the loop function of a meson and a baryon once the logarithmic divergent constant is removed. The analytical properties of G are properly kept when evaluating the previous loop function in dimensional regularization and an explicit subtraction constant, a_i , appears in the expression.

This meson baryon loop function was calculated in Ref. ⁴ with a cut-off regularization. The values of the a_i constants are found to be around -2 to agree with the results of the cut-off method for cut-off values of the order of the mass of the $\rho(770)$ ², which we call of natural size.

2. Poles of the T-matrix

The study of Ref. ³ showed the presence of poles in Eq. (1) around the $\Lambda(1405)$ and the $\Lambda(1670)$ for isospin $I = 0$ and around the $\Sigma(1620)$ in $I = 1$. The same approach in $S = -2$ leads to the resonance $\Xi(1620)$ ⁵ and in $S = 0$ to the $N^*(1535)$ ^{6,7}. One is thus tempted to consider the appearance of a singlet and an octet of meson–baryon resonances. Nevertheless, the situation is more complicated because indeed in the SU(3) limit there are *two* octets and not just one, as we discuss below. The presence of these multiplets was already discussed in Ref. ² after obtaining a pole with $S = -1$ in the $I = 1$ channel, with mass around 1430 MeV, and two poles with $I = 0$, of masses around that of the $\Lambda(1405)$.

The appearance of a multiplet of dynamically generated mesons and baryons seems most natural once a state of the multiplet appears. Indeed, one must recall that the chiral Lagrangians are obtained from the combination of the octet of

pseudoscalar mesons (the pions and partners) and the octet of stable baryons (the nucleons and partners). The SU(3) decomposition of the combination of two octets tells us that

$$8 \otimes 8 = 1 \oplus 8_s \oplus 8_a \oplus 10 \oplus \overline{10} \oplus 27 . \quad (3)$$

Thus, on pure SU(3) grounds, should we have a SU(3) symmetric Lagrangian, one can expect e.g. one singlet and two octets of resonances, the symmetric and anti-symmetric ones.

The lowest order of the meson–baryon chiral Lagrangian is exactly SU(3) invariant if all the masses of the mesons, or equivalently the quark masses, are set equal. As stated above [see Eq. (2)], in Ref. ³ the baryon masses take their physical values, although strictly speaking at the leading order in the chiral expansion they should be equal to M_0 . For Eq. (2) being SU(3) symmetric, all the baryons masses M_i must be set equal as well. When all the meson and baryon masses are equal, and these common masses are employed in evaluating the G_l functions, together with equal subtraction constants a_l , the T -matrix obtained from Eq. (1) is also SU(3) symmetric.

If we do such an SU(3) symmetry approximation and look for poles of the scattering matrix, we find poles corresponding to the octets and singlet. The surprising result is that the two octet poles are degenerate as a consequence of the dynamics contained in the chiral Lagrangians. Indeed, if we evaluate the matrix elements of the transition potential V in a basis of SU(3) states, we obtain something proportional to $V_{\alpha\beta} = \text{diag}(6, 3, 3, 0, 0, -2)$ taking the following order for the irreducible representations: 1, 8_s , 8_a , 10, $\overline{10}$ and 27, with positive sign meaning attraction.

Hence we observe that the states belonging to different irreducible representations do not mix and the two octets appear degenerate. The coefficients in $V_{\alpha\beta}$ clearly illustrate why there are no bound states in the 10, $\overline{10}$ and 27 representations.

In practice, the same chiral Lagrangians allow for SU(3) breaking. In the case of Refs. ^{4,3} the breaking appears because both in the V_{ij} transition potentials as in the G_l loop functions one uses the physical masses of the particles as well as different subtraction constants in G_l , corresponding to the use of a unique cut-off in all channels. In Ref. ² the physical masses are also used in the G_l functions, although these functions are evaluated with a unique subtraction constant as corresponds to the SU(3) limit. In both approaches, physical masses are used to evaluate the G_l loop functions so that unitarity is fulfilled exactly and the physical thresholds for all channels are respected.

By following the approach of Ref. ³ and using the physical masses of the baryons and the mesons, the position of the poles change and the two octets split apart in four branches, two for $I = 0$ and two for $I = 1$, as one can see in Fig. 1. In the figure we show the trajectories of the poles as a function of a parameter x that breaks gradually the SU(3) symmetry up to the physical values. The dependence of masses and subtraction constants on the parameter x is given by

$$M_i(x) = M_0 + x(M_i - M_0),$$

4 Authors' Names

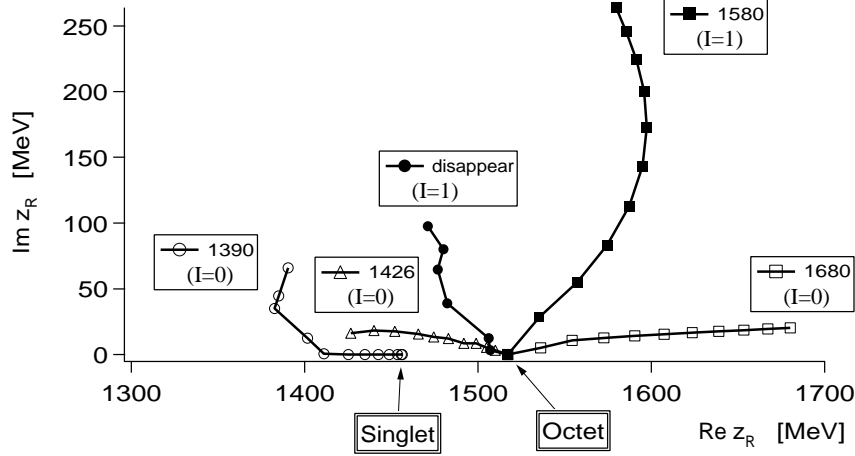


Fig. 1. Trajectories of the poles in the scattering amplitudes obtained by changing the SU(3) breaking parameter x gradually. At the SU(3) symmetric limit ($x = 0$), only two poles appear, one is for the singlet and the other for the octet. The symbols correspond to the step size $\delta x = 0.1$. The results are from ⁸.

$$\begin{aligned} m_i^2(x) &= m_0^2 + x(m_i^2 - m_0^2), \\ a_i(x) &= a_0 + x(a_i - a_0), \end{aligned} \quad (4)$$

where $0 \leq x \leq 1$. In the calculation of Fig. 1, the values $M_0 = 1151$ MeV, $m_0 = 368$ MeV and $a_0 = -2.148$ are used.

The complex poles, z_R , appear in unphysical sheets. In the present search we follow the strategy of changing the sign of the momentum q_l in the $G_l(z)$ loop function for the channels which are open at an energy equal to $\text{Re}(z)$.

The splitting of the two $I = 0$ octet states is very interesting. One moves to higher energies to merge with the $\Lambda(1670)$ resonance and the other one moves to lower energies to create a pole, quite well identified below the $\bar{K}N$ threshold, with a narrow width. On the other hand, the singlet also evolves to produce a pole at low energies with a quite large width.

We note that the singlet and the $I = 0$ octet states appear nearby in energy and what experiments actually see is a combination of the effect of these two resonances.

Similarly as for the $I = 0$ octet states, we can see that one branch of the $I = 1$ states moves to higher energies while another moves to lower energies. The branch moving to higher energies finishes at what would correspond to the $\Sigma(1620)$ resonance when the physical masses are reached. The branch moving to lower energies fades away after a while when getting close to the $\bar{K}N$ threshold.

The model of Ref. ² reproduces qualitatively the same results, but the $I = 1$ pole also stays for $x = 1$. Nevertheless, in both approaches there is an $I = 1$ amplitude with an enhanced strength around the $\bar{K}N$ threshold. This amplitude

has non negligible consequences for reactions producing $\pi\Sigma$ pairs in that region. This has been illustrated for instance in Ref. ⁹, where the photoproduction of the $\Lambda(1405)$ via the reaction $\gamma p \rightarrow K^+\Lambda(1405)$ was studied. It was shown there that the different sign in the $I = 1$ component of the $|\pi^+\Sigma^- \rangle, |\pi^-\Sigma^+ \rangle$ states leads, through interference between the $I = 1$ and the dominant $I = 0$ amplitudes, to different cross sections in the various charge channels, a fact that has been confirmed experimentally very recently ¹⁰.

Once the pole positions are found, one can also determine the couplings of these resonances to the physical states by studying the amplitudes close to the pole and identifying them with

$$T_{ij} = \frac{g_i g_j}{z - z_R} . \quad (5)$$

The couplings g_i are in general complex valued numbers. In the Table we summarize the pole positions and the complex couplings g_i obtained from the model of Ref. ³ for isospin $I = 0$. The results with the model of ² are qualitatively similar.

z_R ($I = 0$)	1390 + 66i		1426 + 16i		1680 + 20i	
	g_i	$ g_i $	g_i	$ g_i $	g_i	$ g_i $
$\pi\Sigma$	-2.5 - 1.5i	2.9	0.42 - 1.4i	1.5	-0.003 - 0.27i	0.27
$\bar{K}N$	1.2 + 1.7i	2.1	-2.5 + 0.94i	2.7	0.30 + 0.71i	0.77
$\eta\Lambda$	0.010 + 0.77i	0.77	-1.4 + 0.21i	1.4	-1.1 - 0.12i	1.1
$K\Xi$	-0.45 - 0.41i	0.61	0.11 - 0.33i	0.35	3.4 + 0.14i	3.5

We observe that the second resonance with $I = 0$ couples strongly to $\bar{K}N$ channel, while the first resonance couples more strongly to $\pi\Sigma$.

3. Influence of the poles on the physical observables

In a given reaction the $\Lambda(1405)$ resonance is always seen in $\pi\Sigma$ mass distribution. However, the $\Lambda(1405)$ can be produced through any of the channels in the Table. Hence, it is clear that, should there be a reaction which forces this initial channel to be $\bar{K}N$, then this would give more weight to the second resonance, R_2 , and hence produce a distribution with a shape corresponding to an effective resonance narrower than the nominal one and at higher energy. Such a case indeed occurs in the reaction $K^-p \rightarrow \Lambda(1405)\gamma$ studied theoretically in Ref. ¹¹. It was shown there that since the K^-p system has a larger energy than the resonance, one has to lose energy emitting a photon prior to the creation of the resonance and this is effectively done by the Bremsstrahlung from the original K^- or the proton. Hence the resonance is initiated from the K^-p channel. This is also the case in the reaction $\gamma p \rightarrow K^*\Lambda(1405)$ which has also the $\Lambda(1405)$ initiated by $\bar{K}N$ through the vertex $\gamma \rightarrow K^*K$ ¹².

4. The interaction of the decuplet of baryons with the octet of mesons

Given the success of the chiral unitary approach in generating dynamically low energy resonances from the interaction of the octets of stable baryons and the pseudoscalar mesons, in ¹³ the interaction of the decuplet of $3/2^+$ with the octet of pseudoscalar mesons was studied and shown to lead to many states which were associated to experimentally well established $3/2^-$ resonances.

The lowest order chiral Lagrangian for the interaction of the baryon decuplet with the octet of pseudoscalar mesons is given by ¹⁴

$$L = i\bar{T}^\mu D_\nu \gamma^\nu T_\mu - m_T \bar{T}^\mu T_\mu \quad (6)$$

where T_{abc}^μ is the spin decuplet field and D^ν the covariant derivative given by in ¹⁴. The identification of the physical decuplet states with the T_{abc}^μ can be seen in ¹⁵, where a detailed study of this interaction and the resonances generated can be seen. The study is done along the lines of the former sections, looking for poles in the second Riemann sheet of the complex plane, the coupling of the resonances to the different channels and the stability of the results with respect to variations of the input parameter, which in our case is just the subtraction constant, a . This allows the association of the resonances found to existing states of the particle data book, and the prediction of new ones. A detail of the results obtained can be seen in Fig. 2.

Another interesting result is the generation of an exotic state of $S = 1$ and $I = 1$ which is generated by the interaction of the ΔK channels and stands as a ΔK resonance. The pole appears in a Riemann sheet below threshold when also the sign of the momentum is changed, but it leads to a ΔK amplitude which accumulates strength close to threshold and produces a broad peak in the cross section, see Fig.3, in contrast to the $I = 2$ cross section which is much smaller and very smooth.

In order to investigate this interaction we propose the study of the following reactions: 1) $pp \rightarrow \Lambda \Delta^+ K^+$, 2) $pp \rightarrow \Sigma^- \Delta^{++} K^+$, 3) $pp \rightarrow \Sigma^0 \Delta^{++} K^0$. In the first case the $\Delta^+ K^+$ state produced has necessarily $I = 1$. In the second case the $\Delta^{++} K^+$ state has $I = 2$. In the third case the $\Delta^{++} K^0$ state has mostly an $I = 1$ component. The study of the ΔK invariant mass distribution in these reactions would provide experimental information to eventually prove the prediction made here, and thus giving support to this new exotic baryonic state, which, although in the quark model would require at least five quarks and could qualify as a pentaquark, finds a much simpler interpretation as a resonant ΔK state.

5. Acknowledgments

This work is partly supported by DGICYT contract number BFM2003-00856, CI-CYT (Spain) Grant Nos. FPA2002-03265 and FPA2004-03470, the E.U. EURIDICE network contract no. HPRN-CT-2002-00311 and the Research Cooperation program of the JSPS and the CSIC.

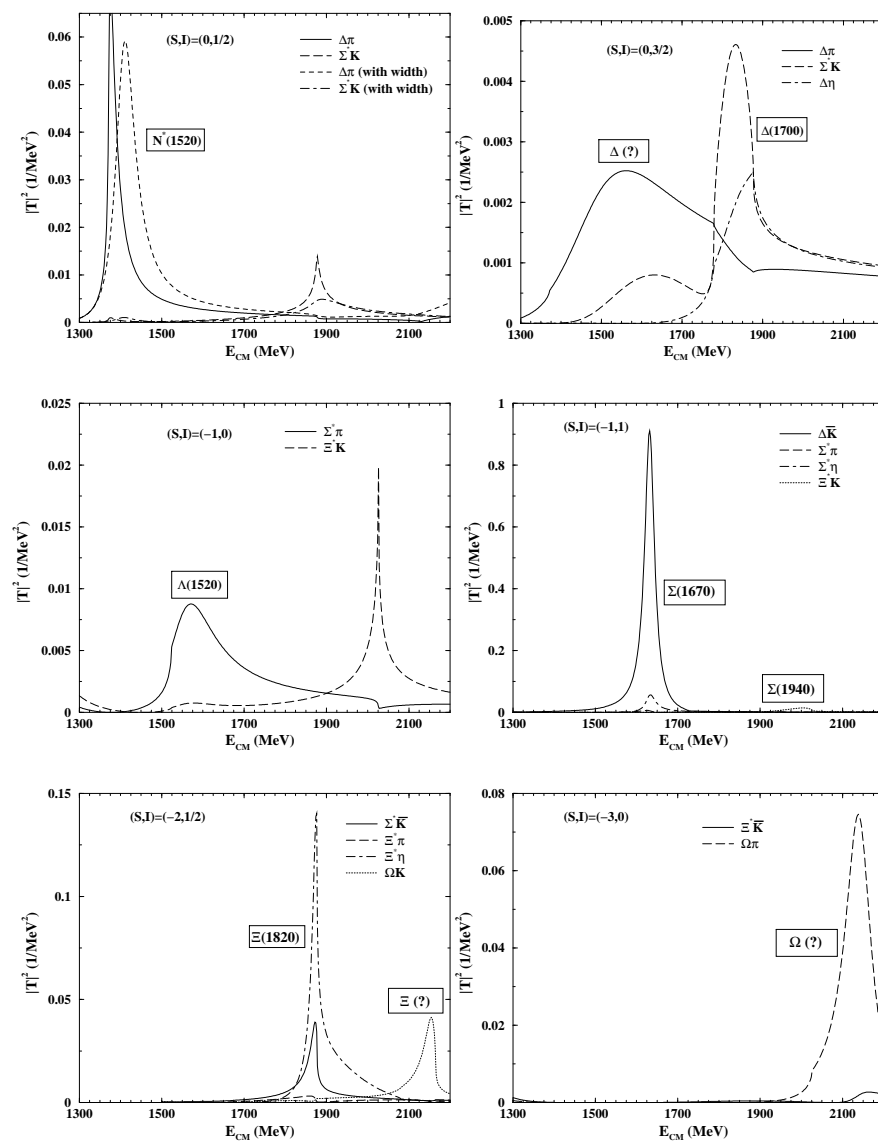
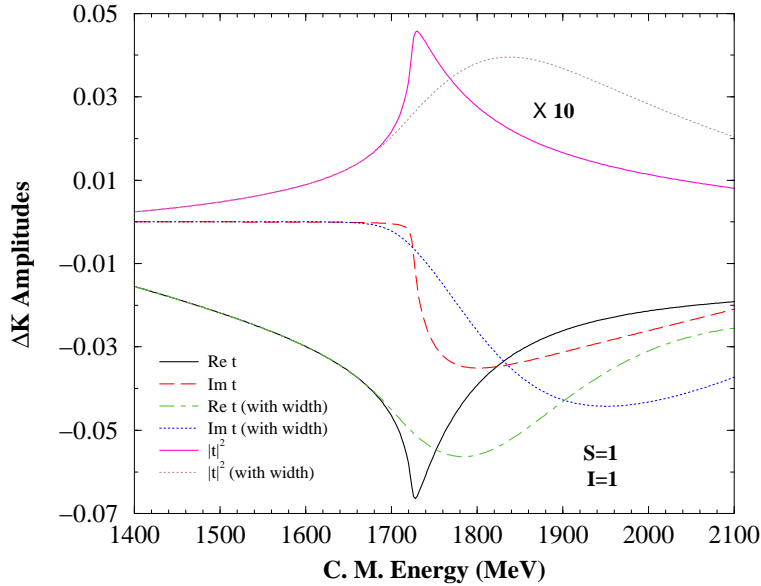


Fig. 2. Resonances obtained from the interaction of the octet of mesons with the decuplet of baryons.

Fig. 3. Amplitudes for $\Delta K \rightarrow \Delta K$ for $I = 1$

References

1. G. Ecker, Prog. Part. Nucl. Phys. **35** (1995) 1; V. Bernard, N. Kaiser and U.-G. Meißner, Int. J. Mod. Phys. E **4** (1995) 193.
2. J. A. Oller and U.-G. Meißner, Phys. Lett. B **500** (2001) 263.
3. E. Oset, A. Ramos and C. Bennhold, Phys. Lett. B **527** (2002) 99 [Erratum-ibid. B **530** (2002) 260].
4. E. Oset and A. Ramos, Nucl. Phys. A **635** (1998) 99.
5. A. Ramos, E. Oset and C. Bennhold, Phys. Rev. Lett. **89** (2002) 252001.
6. N. Kaiser, P. B. Siegel and W. Weise, Nucl. Phys. A **594** (1995) 325
7. T. Inoue, E. Oset and M. J. Vicente Vacas, Phys. Rev. C **65** (2002) 035204.
8. D. Jido, J.A. Oller, E. Oset, A. Ramos, U.G. Meissner, Nucl. Phys. A **725** (2003) 181.
9. J. C. Nacher, E. Oset, H. Toki and A. Ramos, Phys. Lett. B **455**, 55 (1999).
10. J. K. Ahn et al., Proceedings of the XVI Particles and Nuclei International Conference (PANIC02), Eds. K. Imai, T. Kishimoto and H. Toki, Nucl. Phys. A721 (2003) 715-718.
11. J. C. Nacher, E. Oset, H. Toki and A. Ramos, Phys. Lett. B **461** (1999) 299.
12. T. Hyodo, A. Hosaka, M.J. Vicente Vacas and E. Oset, Phys. Lett. B **593** (2004) 75-81.
13. E. E. Kolomeitsev and M. F. M. Lutz, Phys. Lett. B **585** (2004) 243.
14. E. Jenkins and A. V. Manohar, Phys. Lett. B **259** (1991) 353.
15. S. Sarkar, E. Oset and M. J. Vicente Vacas, arXiv:nucl-th/0407025.

