We discuss the formation of s- and d-wave baryon resonances as predicted by the chiral SU(3) symmetry of QCD. Based on the leading order term of the chiral Lagrangian a rich spectrum of molecules is generated, owing to final-state interactions. The spectrum of s- and d-wave baryon states with zero and non-zero charm is remarkably consistent with the empirical pattern. In particular, the recently announced $\Sigma_c(2800)$ of the BELLE collaboration is reproduced with realistic mass and width parameters. Similarly, the d-wave states $\Lambda_c(2625)$ and $\Xi_c(2815)$ are explained naturally to be chiral excitations of $J^P = \frac{3}{2}^+$ states. In the open-charm sector exotic multiplet structures are predicted. These findings support a radical conjecture: meson and baryon resonances that do not belong to the large-$N_c$ ground state of QCD should be viewed as hadronic molecular states.

1. Introduction

The task of constructing a systematic effective field theory for resonance physics of QCD is one of the key challenges in hadronic physics. Considerable progress was achieved in the last few years though there are certainly still some loose ends to be pondered over at this stage [12, 13, 14, 15, 16, 17, 18, 19]. Astounding results have been worked out demonstrating that many meson and baryon resonance states can be easily understood as being formed due to final state interaction. The computations of the authors [2, 3, 4, 8, 9, 10, 11] were driven by a hypothesis put forward some years ago: meson and baryon resonances not belonging to the large-$N_c$ ground states are generated dynamically by coupled-channel interactions [2, 3, 4]. Upon selecting a few ‘fundamental’ hadronic degrees of freedom the zoo of resonances is conjectured to be a result of the interactions of the latter. The identification of the proper set was guided by properties of QCD in the large-$N_c$ limit. We do not support for instance the approach of Chew and Low [12], who viewed the isobar resonance, a member of the large-$N_c$ baryon ground state of QCD, as being generated by pion-nucleon interactions. According to [2, 3] the decuplet states should be considered on an equal footing with the baryon octet states. Similarly, contrary to the work by Zachariasen and Zemach [13], we consider the $\rho$-meson to be ‘fundamental’, being a member of the large-$N_c$ meson ground states.

Most spectacular are results where the driving term of the final state interaction is
predicted unambiguously by the chiral symmetry of QCD. This is the case for the s-wave interaction of the Goldstone bosons with any of the 'fundamental' hadrons. Striking predictions for resonances with \( J^P = 0^+, 1^+, \frac{1}{2}^+, \frac{3}{2}^- \) quantum numbers are the consequence. They are interpreted to be chiral excitations of the large-\( N_c \) ground states with \( J^P = 0^-, 1^-, \frac{1}{2}^+, \frac{3}{2}^- \) quantum numbers. Though the study of \( 0^+ \) and \( \frac{1}{2}^- \) molecules [14-17] has a long history, the claim that \( 1^+ \) and \( \frac{3}{2}^- \) resonances are molecules too is novel and may be provocative to the community. In this talk we review the formation of baryon resonance with charm quantum numbers zero and one. New unpublished results on d-wave resonances with open charm are included. An interpretation of the newly discovered \( \Sigma_c(2800) \) as a chiral molecule is given. Similarly, the previously established d-wave states \( \Lambda_c(2625) \) and \( \Xi_c(2815) \) are reproduced naturally within chiral coupled-channel dynamics.

2. Chiral coupled-channel dynamics

The basis of an effective field theory approach to the resonance spectrum of QCD is the chiral Lagrangian [19]. The latter manifests the chiral properties of QCD in a systematic way, at the price of an infinite hierarchy of interaction terms. Many of the vertices are predicted in terms of a chiral order parameter \( f \simeq 90 \text{ MeV} \), which parameterizes the weak decay processes of the Goldstone bosons. Still, an infinite number of interaction terms is unknown and has to be determined by either experiment or additional dynamical information from QCD, like large-\( N_c \) sum rules. Predictive power arises as a consequence of counting rules that estimate the importance of a given contribution based on the naturalness assumption. The important property is that at a given order of accuracy only a finite number of parameters enter the calculation.

Most striking is the leading-order prediction for the s-wave interaction of the Goldstone bosons, \( \Phi_a(x) \), with the baryon octet, \( B_a(x) \), and decuplet fields, \( B_{\mu a c}(x) \). Given the kinetic term of the baryon field the interaction follows by a chiral rotation unambiguously in terms of the chiral order parameter:

\[
\mathcal{L} = \frac{i}{8 f^2} \text{tr} \left( (\bar{B} \gamma^\mu B) \cdot [\Phi, (\partial_\mu \Phi)]_- \right) + \frac{3i}{8 f^2} \text{tr} \left( (\bar{B}_\nu \gamma^\mu B^\nu) \cdot [\Phi, (\partial_\mu \Phi)]_- \right). \tag{1}
\]

The octet and decuplet fields, \( \Phi, B \) and \( B_\mu \), posses an appropriate matrix structure according to their SU(3) tensor representation.

It is instructive to discuss first the SU(3) limit where the scattering channels may be decomposed into invariants

\[
8 \otimes 8 = 27 \oplus 10 \oplus 8 \oplus 8 \oplus 1, \quad 8 \otimes 10 = 35 \oplus 27 \oplus 10 \oplus 8. \tag{2}
\]

According to (2) the scattering of the Goldstone bosons of the baryon octet states has 6 independent channels. The striking implication of QCD’s chiral SU(3) symmetry is that attraction is predicted in the two octet and singlet channels only, the 27plet channel is repulsive. No interaction is provided in the decuplet channels. Amazingly, in just the channels where attraction is foreseen a rich spectrum of baryon resonances is observed in experiment. Those channels coincide with the prediction of the quark model which suggest that baryon resonances are formed primarily in channels with quantum numbers...
which can be realized in terms of three quarks only. A similar analysis reveals that the interaction of Goldstone bosons with the decuplet states is strongly attractive in the octet and decuplet sectors, but repulsive in the 35plet. The 27plet sector shows weak attraction only. Relative to the octet channel its attraction is reduced by a factor 6. The attraction in the decuplet channel relative to the octet channel is reduced by the factor 2 only.

Interaction terms analogous to (1) can be written down describing the s-wave interaction of Goldstone bosons with the open-charm $\frac{1}{2}^+$ and $\frac{3}{2}^+$ ground states [20, 21, 22]. They form anti-triplet and sextet representations of the SU(3) group. The Clebsch-Gordon coefficients that describe the relative interaction strength of the various channels reflect the SU(3) multiplet structure of the hadron the Goldstone bosons are scattered off. In the SU(3) limit the decomposition follows

$$8 \otimes \bar{3} = \bar{3} \oplus 6 \oplus 15,$$

$$8 \otimes 6 = \bar{3} \oplus 6 \oplus 15 \oplus 24.$$ (3)

For $8 \otimes \bar{3}$ scattering the leading order chiral Lagrangian predicts attraction in the anti-triplet and sextet sectors, but repulsion for the anti-15plet. For $8 \otimes 6$ scattering there is attraction in the anti-triplet, sextet but also in the anti-15plet sectors. The interaction is repulsive in the 24plet sectors. Chiral SU(3) symmetry predicts a hierarchy of strengths with the strongest attraction in the triplet channels, which is five times as strong as the attraction in the anti-15plet. In the sextet sector the attraction is reduced by a factor 3/5 only. It is also rewarding to compare the amount of attraction in the anti-triplet channels as they result from the reduction of $8 \otimes 3$ versus $8 \otimes 6$. Chiral symmetry predicts stronger binding in the latter case. The amount of attraction is larger by a factor 5/3 as compared to the former case.

Given all these interesting predictions, how do we translate this information into a spectrum of resonances? Clearly, if the interaction of the Goldstone bosons with a hadron is sufficiently attractive in some channel, one would expect the formation of a molecule. Unfortunately, it is a priori ill-defined to insert the interaction vertex (1) into a coupled-channel scattering equation, like the Bethe-Salpeter or Lippmann-Schwinger equation. The quasi-local nature of the interaction gives rise to short-range divergencies. A renormalization program can be set up in a perturbative approach defining unique answers at each order in the perturbation. However, the formation of a molecule requires the summation of an infinite number of diagrams. This asks for novel tools. In a phenomenological approach one may introduce a cutoff to tame the coupled-channel equation, however, the price is that the binding energy of the such molecule depends sensitively on the precise value of the cutoff. More seriously, the introduction of a cutoff in a meson-baryon system is at odds with the counting laws that permit a truncation of the interaction in the first
place. Rather, then giving a formal discussion of these problems (see [3, 10, 23]), an intuitive and physical argument how to remedy this unsatisfactory situation is given.

The problem is illustrated in Fig. 1, which indicates that the physics of a partial-wave scattering amplitude is perturbative in a small subthreshold energy domain only. The amplitude is characterized by its s-channel and u-channel unitarity branch points. Additional branch points implied by t-channel exchanges are suppressed in the discussion for simplicity. At sufficiently large and small energies the amplitude is necessarily non-perturbative with possible resonance signals. If one feeds an interaction kernel into a coupled-channel scattering equation the s-channel unitarity cut is constructed explicitly, i.e. the scattering amplitude satisfies s-channel coupled-channel unitarity. However, unless an infinite sum of diagrams is considered the so constructed scattering amplitude does not satisfy u-channel unitarity, i.e. the approximation to the u-channel cut is necessarily inadequate. Thus, this amplitude can be trusted only at energies to right of $\mu$ in Fig. 1. In practical applications this is not really an issue. The u-channel cut can be reconstructed uniquely by the s-channel cuts of the crossed reactions. Since the scattering amplitude has a small window where it can be evaluated in perturbation theory, an accurate scattering amplitude, applicable to the left and right of $\mu$ in Fig. 1, can be constructed by gluing together s- and u-channel unitarized amplitudes at some matching point $\mu$. As a result the full amplitude is crossing symmetric exactly at energies accessible to experiments. The glued amplitude should be smooth and not jump around at the matching point. This is not necessarily realized in phenomenological coupled-channel computations, in which a cutoff or a form factor may severely destroy the perturbative nature of the subthreshold amplitudes. In [3] we suggested to take the smooth matching as a constraint to be imposed on the unitarization scheme. This sifts out a major part of the ambiguities inherent in unitarization schemes. The renormalization program is carried out in the small and perturbative subthreshold window. Any residual scheme dependence is systematically eliminated upon the inclusion of higher order contributions to the interaction kernel.

In Fig. 2 we demonstrate the quality of the proposed matching procedure as applied for typical forward scattering amplitudes [10]. The figure clearly illustrates a successful matching of s-channel and u-channel iterated amplitudes at subthreshold energies.

Figure 2. Typical cases of forward scattering amplitudes. The solid (dashed) line shows the s-channel (u-channel) unitarized scattering amplitude. The dotted lines represent the amplitude evaluated at tree-level.
Figure 3. Baryon resonances as seen in the s-wave scattering of Goldstone bosons off the baryon octet \( N(939), \Lambda(1115), \Sigma(1195), \Xi(1315) \). Shown are real and imaginary parts of partial-wave scattering amplitudes with \( J^P = \frac{1}{2}^- \) and \( (I, S) = (\frac{1}{2}, -2), (0, -1), (1, -1), (1/2, 0) \).

3. S- and d-wave baryon resonances with zero charm

We discuss the chiral excitations of the baryon octet and decuplet ground states of QCD. The partial-wave scattering amplitudes of Fig. 3 show evidence for the formation of the \( \Xi(1690), \Lambda(1405), \Lambda(1670) \) and \( N(1535) \) resonances close to their empirical masses. An additional \( (I, S) = (0, -1) \) state, mainly a SU(3) singlet [7, 8], can be found as a complex pole in the scattering amplitude close to the pole implied by the \( \Lambda(1405) \) resonance. There is no convincing signal for a \( (1, -1) \) resonance at this leading order calculation. However, chiral corrections lead to a pronounced resonance signal in this sector [3, 8] which should be identified with the \( \Sigma(1750) \) resonance, the only well established s-wave resonance in this sector. It couples strongly to the \( K \Xi \) and \( \eta \Sigma \) states. The fact that a second resonance with \( (\frac{1}{2}, 0) \) is not seen in Fig. 3 even though the 'heavy' SU(3) limit suggests its existence [8, 9], we take as a confirmation of the phenomenological observation that the \( N(1650) \) resonance couples strongly to the \( \omega N \) channel not considered here [4]. For a more detailed discussion of the spectrum we refer to [8].

We turn to the d-wave resonance spectrum, which can be read off Fig. 4. Amplitudes, that describe the s-wave scattering of Goldstone bosons off the decuplet states are shown.
Figure 4. Baryon resonances with $J^P = \frac{3}{2}^-$ as formed in the s-wave scattering of Goldstone bosons off the baryon decuplet $\Delta(1232), \Sigma(1385), \Xi(1530), \Omega(1672)$. Shown are real and imaginary parts of partial-wave scattering amplitudes. All sectors are included for which chiral symmetry predicts an attractive interaction. One may expect that chiral dynamics does not make firm predictions for d-wave resonances since the meson baryon-octet interaction in the relevant channels probes a set of counter terms presently unknown. However, this is not necessarily so. Since a d-wave baryon resonance couples to s-wave meson baryon-decuplet states chiral symmetry is quite predictive for such resonances, under the assumption that the latter channels are dominant. This is in full analogy to an analysis of the s-wave resonances \[14, 15, 18, 5, 6, 7, 8\] that neglects the effect of the contribution of d-wave meson baryon-decuplet states. The empirical observation that the d-wave resonances $N(1520), N(1700)$ and $\Delta(1700)$ have large branching fractions ($> 50\%$) into the inelastic $N\pi\pi$ channel, even though the elastic $\pi N$ channel is favored by phase space, supports this assumption. It is a stunning success of chiral coupled-channel dynamics that it recovers the four star hyperon resonances.
Ξ(1820), Λ(1520), Σ(1670) with masses quite close to the empirical values. The nucleon
and isobar resonances $N(1520)$ and $Δ(1700)$ also present in Fig. 4 are predicted with
less accuracy. The important result here is the fact that these resonances are generated at
all. Fully realistic results should not be expected already in this leading order calculation.
For instance, chiral correction terms provide a d-wave $πΔ$ component of the $N(1520)$.
Moreover, from phenomenological analysis one anticipates the importance of additional
inelastic channels involving vector mesons (see e.g. [4]). We continue with the bound
state signal in the $(0, -3)$ amplitudes at mass 1950 MeV. Such a state has so far not been
observed but is associated with a decuplet resonance [24]. Further states belonging to
the decuplet are seen in the $(\frac{1}{2}, -2)$ and $(1, -1)$ amplitudes at masses 2100 MeV and 1920
MeV. The latter state can be identified with the three star $Ξ(1940)$ resonance. Finally
we point out that the $(0, -1)$ amplitudes show signals of two resonance states consistent
with the existence of the four star resonance $Λ(1520)$ and $Λ(1690)$. It appears that the
SU(3) symmetry breaking pattern generates the ‘missing’ state in this particular sector
by promoting the weak attraction of the 27plet. For completeness we include also the
exotic sectors $(1, -3)$, $(3/2, -2)$, $(2, -1)$, $(1, 1)$ part of the 27plet. Since chiral dynamics
predicts weak attraction only, no clear resonance signals are seen in the amplitudes. It
remains an open issue whether correction terms may conspire to increase the amount of
attraction in some selected channels leading to unambiguous exotic signals. Subsequently
our results have been confirmed in [25].

4. S- and d-wave baryon resonances with open charm

We turn to the chiral excitations of open-charm baryons. At present there is little
known empirically about the open-charm baryon resonance spectrum. Only three s-
wave resonances $Λ_c(2593), Λ_c(2880)$ and $Ξ_c(2790)$ were discovered so far [26]. Two d-
wave resonances $Λ_c(2625)$ and $Ξ_c(2815)$ are acknowledged by the particle data group [26]. An additional isospin triplet state $Ξ_c(2800)$ was announced recently by the BELLE
collaboration [27], which most likely carries $J^P = \frac{3}{2}^-$ quantum numbers also.

We confront the empirical spectrum with the spectrum of chiral molecules. Let us first
discuss the chiral excitations of the $J^P = \frac{1}{2}^+$ states, which form an anti-triplet and sextet.
The relevant scattering amplitudes are collected in Fig. 5. Amplitudes with identical
quantum numbers that describe the scattering of Goldstone bosons off the anti-triplet
and sextet states are grouped on top of each other. This is instructive since at subleading
order in the chiral expansion such amplitudes start to couple.

Consider the chiral excitations of the triplet states manifest in the first and third row
of Fig. 5. According to our discussion those states form a strongly bound anti-triplet
and a weakly bound sextet. The $(0, 0)$ resonance is identified with the $Λ_c(2880)$ being a
member of the anti-triplet. It is in fact a bound state at this leading order computation.
The isospin doublet of that multiplet, a $(1/2, -1)$ state, is a much broader object so far
unobserved. The chiral sextet excitations of the anti-triplet is visible in the $(1, 0), (1/2, -1)$
and $(0, -2)$ sectors. Most spectacular would be the exotic and bound $K_0 \Xi_c(2470)$ system,
if it survives chiral correction terms.

An even richer spectrum is formed by the chiral excitations of the sextet states as shown
in the second and fourth row of Fig. 5. Chiral SU(3) symmetry predicts attraction with
Figure 5. Open-charm baryon resonances as seen in the scattering of Goldstone bosons off anti-triplet $\Lambda_c(2284), \Xi_c(2470)$ and sextet $\Sigma_c(2453), \Xi'_c(2580), \Omega_c(2697)$ baryons. Shown are real and imaginary parts of partial-wave scattering amplitudes with $J^P = \frac{1}{2}^-$ and $(I, S) = (0, 0), (1/2, -1), (1, 0), (0, -2)$. Decreasing strength in the triplet, sextet and anti-15plet. The triplet states are identified with the $\Lambda_c(2593)$ and $\Xi_c(2790)$ resonances. The sextet manifests itself most clearly as a bound $K \Xi'_c(2580)$ system. The existence of the anti-15plet is clearly seen in the $(0, 0), (1, 0)$ and $(1/2, -1)$ sectors where it leads to narrow resonance structures around 3 GeV. Additional signals of the anti-15plet in the exotic sectors $(1/2, 1), (3/2, -1)$ and $(1, -2)$ are weaker and are not shown in Fig. 5. They can, however, be inferred from Fig. 6.

We turn to the d-wave spectrum of open charm baryons. It is induced by the s-wave scattering of the Goldstone bosons off the sextet $J^P = \frac{3}{2}^+$ ground states. Here we assume a value for the $\Omega_c(2770)$ mass, so far unknown, which is suggested by quark model calculations [28] and lattice simulations [29]. The scattering amplitudes are shown in Fig. 4 for all sectors in which chiral dynamics predicts attraction. The pattern is analogous to the already discussed spectrum arising from the scattering of the Goldstone bosons off the sextet of $\frac{1}{2}^+$ ground states. This is an immediate consequence of the identical interaction strengths of the two systems. So far only three states are observed.
Figure 6. Open-charm baryon resonances with $J^P = \frac{3}{2}^-$ as seen in the scattering of Goldstone bosons off the sextet $\Sigma_c(2520), \Xi_c(2644), \Omega_c(2770)$ baryons. Shown are real and imaginary parts of partial-wave scattering amplitudes.

experimentally $\Lambda_c(2625), \Sigma_c(2800)$ and $\Xi_c(2815)$. The masses are rather well reproduced, given the fact that the amplitudes are computed at the parameter-free leading order. We underestimate the binding for the $\Lambda_c(2625)$ and $\Xi_c(2815)$ by about 75 MeV. Whereas the width of the $\Sigma_c(2800)$ is quite consistent with the recent BELLE measurement we overestimate the width of the $\Xi_c(2815)$ and $\Lambda_c(2656)$. Note that these states are somewhat underbound. Allowing for small attractive correction terms will pull down these states giving them a smaller width within reach of their measured values. Besides the anti-triplet excitations, chiral dynamics predicts again a sextet and anti-15plet of excitations, however, with reduced binding strength. Most spectacular is the prediction that the $K$ is bound at the $\Sigma_c(2520)$ and the $\bar{K}$ at the $\Xi_c(2644)$. The former is rather weakly bound, if at all, being a member of the anti-15plet. The latter enjoys a binding energy of about 40 MeV being a member of the sextet. Since it couples equally strong to the $\eta \Omega_c(2774)$ state
a derivation of its electromagnetic properties requires a coupled-channel study. Striking
evidence for the existence of the anti-15plet are narrow resonance structures in the (0, 0)
and (1, 0) sectors around 3.1 GeV. These resonances couple strongly to the $K \Xi^c(2664)$
state. A further somewhat broader anti-15plet state with $(1/2, -1)$ quantum numbers is
predicted at around 3.2 GeV. This state is dominated by its $K \Omega_c(2774)$ component.

5. Summary and outlook

In this talk we reviewed the spectrum of chiral excitations of the large-$N_c$ ground state
baryons. The leading order term of the chiral Lagrangian predicts a rich spectrum of
molecules in astounding agreement with the empirical resonance pattern. No adjustable
parameters enter the spectrum. It is determined by the well known value of the pion
decay constant. The properties of the chiral excitations depend crucially on the values
of the current quark masses of QCD. An SU(3) limit where the pion mass is as heavy
as the empirical kaon mass leads to a bound state spectrum, typically. The SU(3) limit
where the kaon mass is as light as the empirical pion mass implies the disappearance of the
resonance signals. This is a striking prediction of chiral dynamics which should be verified
by unquenched lattice QCD simulations. Besides reproducing well-established states new,
so far unobserved, exotic multiplets are foreseen. New and unpublished results concern
the d-wave spectrum of open charm baryons. The recently announced $\Sigma^c(2800)$ state of
the BELLE collaboration is interpreted as a chiral molecule. The large width of that state
is naturally explained. Similarly, the $\Lambda_c(2625)$ and $\Xi_c(2815)$ states are well understood in
terms of chiral coupled-channel dynamics. The results support the radical conjecture of
the authors that meson and baryon resonances that do not belong to the large-$N_c$ ground
state of QCD should be viewed as hadronic molecular states. To further substantiate
these findings extensive and refined coupled-channel computations are necessary that will
reveal the resonance spectrum of QCD more quantitatively.

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