Chiral dynamics with strange quarks

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Abstract. In the first part of the talk, I review what we know (or rather do not know) about the structure of the QCD vacuum in the presence of strange quarks. Chiral perturbation theory allows to study reactions of pions and kaons and to further sharpen our understanding of symmetry violation in QCD. I review recent progress on the description of pion-kaon scattering, in particular concerning isospin violation and the extraction of threshold and resonance parameters from Roy-Steiner equations. In the third part, it is shown how a unitary extension of chiral perturbation theory leads to novel insight into the structure of the Λ(1405).

INTRODUCTION: S QUARK MYSTERIES

The strange quark plays a special role in the QCD dynamics at the confinement scale. Here, I will discuss some open questions surrounding chiral dynamics with strange quarks, pertinent to the structure of the strong interaction vacuum as well as to the structure of light mesons and baryons. Some of these issues are: Since \( m_s \sim \Lambda_{\text{QCD}} \), is it appropriate to treat the strange quark as light or should it be considered heavy, as in the so–called heavy kaon effective field theory, see [1–3]? Why is the OZI rule so badly violated in the scalar sector with vacuum quantum numbers? One example is the reaction \( J/\Psi \to \phi \pi \pi / K K \), which is OZI suppressed to leading order, but even has an additional doubly OZI suppressed contribution. The \( \pi^+ \pi^- \) event distribution shows a clear peak at the energy of 980 MeV, which is due to the \( f_0 \) scalar meson. This lets one anticipate that the dynamics of the low-lying scalar mesons and the mechanism of OZI violation are in some way related. More generally, it is of interest to learn about the phase structure of SU\((N_c)\) gauge theory at large number of flavors \( N_f \). In QCD, we know that asymptotic freedom is lost for \( N_f \geq 17 \) but from the study of the two-loop \( \beta \) function one expects that there is a conformal window around \( N_f \approx 6 \) [4]. This lets one contemplate the question whether there is already a rich phase structure even for the transition from \( N_f = 2 \) to \( N_f = 3 \) ? Some lattice studies seem to indicate a strong flavor sensitivity when going from \( N_f = 2 \) to \( N_f = 4 \) [5,6]. As discussed by many speakers at this conference, the nature of the low–lying scalar mesons is still very much under debate (a topic I will not entertain in detail). In the baryon sector, there are also some “strange” states with non-vanishing strangeness.
More precisely, what is the nature of some strange baryons like the Λ(1405) or the $S_{11}(1535)$, are these three quarks states or meson-baryon bound states? The latter scenario was already contemplated many years ago by Dalitz and collaborators [7] and has been rejuvenated with the advent of coupled channel calculations using chiral Lagrangians to specify the driving interaction. In the following, I address some of these issues.

THE VACUUM IN THE PRESENCE OF S QUARKS

There are many phenomenological as well as theoretical indications that the chiral symmetry ($\chi_S$) of three–flavor QCD is spontaneously broken, abbreviated as $S\chi_S$. Now the question arises what are the order parameters of $S\chi_S$? Consider the current-current correlator between vector and axial currents,

$$\Pi_{ab}^{\mu\nu}(q) = i \int d^4x e^{iqx} \langle 0 | T \{ V_{\mu}^a(x)V_{\nu}^b(0) - A_{\mu}^a(x)A_{\nu}^b(0) \} | 0 \rangle .$$

(1)

In the three flavor chiral limit, it can be written in terms of meson and continuum contributions and worked out explicitly,

$$\Pi_{\mu\nu}^{ab}(0) = -\frac{1}{4} g_{\mu\nu} \delta^{ab} F^2(3) .$$

(2)

If $\Pi_{\mu\nu}^{ab}(0) \neq 0$, then we have $S\chi_S$. We have thus identified an order parameter of spontaneous chiral symmetry breaking, namely the pion decay constant in the chiral limit,

$$\lim_{m_u,m_d,m_s \to 0} F_\pi = F(3) .$$

(3)

Its non-vanishing is a sufficient and necessary condition for $S\chi_S$,

$$\Pi_{\mu\nu}^{ab}(0) \neq 0 \iff F(3) \neq 0 \iff S\chi_S .$$

(4)

Naturally, there are many other possible order parameters. Often considered is the light quark condensate,

$$\langle 0 | \bar{q}q|0 \rangle = \langle 0 | \bar{u}u|0 \rangle^{(3)} = \langle 0 | \bar{d}d|0 \rangle^{(3)} = \langle 0 | \bar{s}s|0 \rangle^{(3)} = \Sigma(3) ,$$

(5)

because the scalar-isoscalar operator $\bar{q}q$ mixes right- and left-handed quark fields. As will be discussed below, the quark condensate plays a different role than the pion decay constant. Other possible color-neutral order parameters of higher dimension are e.g. the mixed quark–gluon condensate $\langle 0 | \bar{q}^i \sigma^{\mu\nu} C^a_{\mu\nu} T^a_i q_j | 0 \rangle^{(3)}$ or certain four-quark condensates $\langle 0 | (\bar{q} \Gamma_1 q)(\bar{q} \Gamma_2 q) | 0 \rangle$ with $\Gamma_i$ some Dirac operator. It goes without saying that the spontaneously and explicitly broken chiral symmetry can be systematically analyzed in terms of an effective field theory - chiral perturbation theory (CHPT) (or some variant thereof). We now turn to the flavor dependence of these
various order parameters. For this, consider QCD on a torus, or in an Euclidean $(t \rightarrow -ix^0)$ box of size $L \times L \times L \times L$, understanding of course that we have to take the infinite volume limit at its appropriate place. Quark and gluon fields are then subject to certain boundary conditions, which are anti-periodic and periodic, in order. Analyzing the spectrum of the QCD Dirac operator, one arrives e.g. at the Banks–Casher relation. Also, the order parameters $F^2$ and $\Sigma$ are dominated by the IR end of the Dirac spectrum [8], and one therefore expects a paramagnetic effect,

$$\Sigma(N_f + 1) < \Sigma(N_f) \sim 1/L^4, \quad F^2(N_f + 1) < F^2(N_f) \sim 1/L^2,$$

indicating a suppression of the chiral order parameters with increasing number of flavors. We note that the condensate is most IR sensitive. These results are exact, the question is now how strong this flavor dependence is or how it can be tested or extracted from some observables.

In the standard scenario of $S\chi$SB, terms quadratic in the quark masses are small, as has been recently confirmed for the two flavor case from the analysis of the BNL E865 $K_{e4}$ data [9]. If these terms are also small in the three flavor case, the so-called Gell-Mann–Oakes–Renner ratio $X(3)$ stays close to one,

$$X(3) \equiv \frac{2\hat{m}\Sigma(3)}{F^2 \pi M^2} \sim 1,$$

with \(\hat{m} = (m_u + m_d)/2\) the average light quark mass. There are many successes supporting this scenario, as one example I will discuss pion–kaon scattering in the next section. However, there is also some information pointing towards a more complicated phase structure (suppression of $\Sigma(3)$), as discussed next.

![Figure 1. Flavor dependence of chiral symmetry breaking order parameters. A speculative scenario with a strongly suppressed three flavor condensate is depicted. In the standard scenario, the two lines for the pion decay constant and the condensate would be very close to each other.](image)

Moussallam [10,11] investigated a sum rule for the OZI violating correlator $\Pi_z \sim \langle \bar{u}u(x)\bar{s}s(0) \rangle_c$, which has the form

$$\Pi_z(m_s) = \frac{1}{\pi} \int_0^\infty ds \frac{1}{s} \sigma(s),$$

and allows to relate $\Sigma(3)$ with $\Sigma(2)$. This sum rule is super-convergent. In the approximation that the spectral function can be saturated by two–particle intermediate $\pi\pi$ and $K\bar{K}$ states, it can be expressed entirely in terms of the (non)strange
scalar form factors of the pion and the kaon. These form factors can be calculated within CHPT, but are needed at higher energies here (for one particular calculation in a unitarized version of CHPT, see e.g. [12]). This gives the spectral function for energies below $\simeq 1.6$ GeV. One can use various T-matrices for the $\pi\pi \rightarrow \pi\pi/K\bar{K}$ system to get an idea of the uncertainty in this energy domain. Above that energy, one can use pQCD. Putting all the various pieces together, one obtains

$$\Sigma(3) = \Sigma(2) [1 - 0.54 \pm 0.27] ,$$

(9)

where the central value indicates a large suppression of the three flavor condensate but the uncertainties are large enough to give marginal consistency with the standard scenario. For a discussion of the stability of this result against some higher order corrections, see [11]. For more investigations of such a scenario see [13,14]. Clearly, more work is needed to further quantify such results and to reduce the uncertainties.

**PION-KAON SCATTERING**

Pion–kaon scattering is the simplest scattering process involving strange quarks. Furthermore, since to one–loop accuracy all low–energy constants (LECs) are known from other processes, one can predict e.g. the S–wave scattering lengths (given here in the basis of total isospin $1/2$ and $3/2$). This has been done long time ago [15,16] (in units of $M_\pi^{-1}$),

$$a_0^{1/2} = 0.18 \pm 0.03 [0.22 \pm 0.02] , \quad a_0^{3/2} = -0.05 \pm 0.02 [-0.045 \pm 0.008] .$$

(10)

The CHPT predictions are compared to the then existing data/Roy equation analysis in Fig. 2. Obviously, no firm conclusion could be drawn (the dark hatched ellipse comes in later).

![Figure 2. S-wave scattering lengths for πK scattering. The CHPT (CA) predictions are shown by the cross (black dot). The older data/Roy equation analysis can be traced back from [15,16]. The dark hatched ellipse refers to the new dispersive analysis of [17].](image)

In the light of more recent and more precise data from the eighties, that never were analyzed using dispersive methods, a novel evaluation of the Roy-Steiner equations
was called for. This was recently achieved by Büttiker et al. [17]. They solved the Roy–Steiner equations for the S– and P–waves using all available input from πK → πK and ππ → K̅K and employing Regge theory for large energies. The outcome of this analysis are the S– and P–wave phase shifts below the matching energy of 1 GeV and the amplitude in the interior of the Mandelstam plane, in particular (sub)threshold parameters. It turns out that the resulting phase shifts are mostly in poor agreement with the existing low energy data, e.g. the mass of the spin–1 K∗ mesons from the crossing of the P-wave isospin 1/2 phase through π/2 happens at \(905 \pm 3\) MeV, visibly different from the PDG value of \(891.7 \pm 0.3\) MeV. This needs further investigation. The resulting S–wave scattering lengths are given in the square brackets in Eq. (10), they come out consistent with the CHPT predictions, pointing toward the validity of the standard scenario. Similarly, the LECs extracted in [17] agree well with earlier determinations based on \(X(3) \simeq 1\).

Another method of extracting the S-wave scattering lengths is the precise measurements of the characteristics of pion-kaon bound states, so-called \(\pi K\) atoms. In order to relate the lifetime and the energy shift to the scattering lengths, one has to make use of modified Deser formulae that include NLO effects in isospin breaking, 

\[
\Gamma_{\pi^0 K^0} \propto \left( a_0^{3/2} - a_0^{1/2} + \epsilon \right)^2 (1 + \kappa),
\]

\[
\Delta E_{2S-2P}^{str} \propto \left( a_0^{3/2} + 2a_0^{1/2} + \epsilon' \right)^2 (1 + \kappa'),
\]

where \(\epsilon\) and \(\epsilon'\) represent, respectively, the isospin violating corrections in the regular part of the scattering amplitudes \(\pi^- K^+ \to \pi^0 K^0\) and \(\pi^- K^+ \to \pi^- K^+\) at threshold, while \(\kappa\) (\(\kappa'\)) is an additional contribution only calculable within the bound state formalism. There are two sources of isospin violation, the strong contributions \(\propto m_u - m_d\) and electromagnetic contributions \(\propto \alpha = e^2/4\pi\). It is most efficient to collect these two small parameters as \(\delta \in \{m_u - m_d, \alpha\}\) and expand the corrections to order \(\delta\) in all channels. To one–loop accuracy, \(\epsilon\) and \(\epsilon'\) have been calculated in [18,19] (see also [20,21])

\[
a_0(\pi^- K^+ \to \pi^0 K^0) = -\sqrt{2}a_0 \left\{ (1. \pm 0.8\%) + (1.3 \pm 0.1\%) + (0. \pm 1.1\%) \right\},
\]

\[
a_0(\pi^- K^+ \to \pi^- K^+) = (a_0^- + a_0^+) \left\{ (1. \pm 16.1\%) + 0.2\% \right\},
\]

where we have switched to the isospin basis for the scattering amplitudes, \(T^+ = (T^{1/2} + 2T^{3/2})/3\) and \(T^- = (T^{1/2} - T^{3/2})/3\). All quoted errors for the different contributions are due to the uncertainties in the respective strong and electromagnetic LECs. Note further that to leading order, the isospin violating corrections to the elastic scattering length are entirely given in terms of the pion mass difference, and thus are of electromagnetic origin. From the above equations, we can give rather precise predictions for \(\epsilon\) and \(\epsilon'\).
\[ \epsilon = 1.3 \pm 1.2 \%, \quad \epsilon' = 1.1 \pm 3.2 \%. \] (15)

We can thus conclude that the extraction of the strong scattering amplitudes at threshold from the lifetime and level shift of the \( \pi K \) atoms is sufficiently well under control. The isospin breaking effects in both cases are only of the order of 1% with an uncertainty of 1% and 3%, respectively. What remains to be done is an equally precise calculation of the bound state corrections \( \kappa \) and \( \kappa' \) [22].

**THE NATURE OF THE \( \Lambda(1405) \)**

In this section, I will discuss some issues in the framework of SU(3) baryon chiral perturbation theory and extensions thereof. First, it is often stated that three flavor baryon CHPT does not converge due to the large kaon mass and/or unitarity corrections. While that is true in certain cases, there are many examples where indeed one can make precise predictions. As one particular example, let me consider the charge radii of the ground state baryon octet. To fourth order (complete one-loop calculation), the charge radii can be given in terms of two LECs. These parameters can be fixed from the well measured proton and neutron electric radii, so that predictions for the other members of the octet emerge. On the other hand, the radius of the \( \Sigma^- \) can be obtained by scattering a highly boosted hyperon beam off the electronic cloud of a heavy atom (elastic hadron–electron scattering). Such an experiment has been first carried out at CERN, demonstrating the feasibility of the method and later repeated with much better accuracy at FNAL. The theoretical prediction (published before the data came out) compares well with the result from the SELEX collaboration,

\[
\langle r_{\Sigma^-}^2 \rangle_{\text{th}} = 0.67 \pm 0.03 \text{ fm}^2 [23], \quad \langle r_{\Sigma^-}^2 \rangle_{\text{exp}} = 0.61 \pm 0.12 \pm 0.09 \text{ fm}^2 [24].
\] (16)

For a more detailed discussion of the status of SU(3) baryon CHPT, see e.g. [25]. Next, I discuss \( K^- p \) scattering. For this process, a purely perturbative treatment is not possible (for an explicit demonstration, see [26]) due to the strong channel couplings and the appearance of a subthreshold resonance, the \( \Lambda(1405) \), which is supposed to be a meson-baryon bound state rather than a genuine 3-quark state. First speculations about its possible unconventional nature date back to [7]. Since then many (QCD-inspired) models have been considered, but the first work of supplementing coupled channel dynamics with chiral Lagrangians which allows to dynamically generate the \( \Lambda(1405) \) was reported in [27], see also [28] and the review [29]. A non-perturbative resummation scheme is mandatory to generate a bound state or a resonance. There exist many such approaches, but it is possible and mandatory to link such a scheme tightly to the chiral QCD dynamics. Such an improved approach was developed for pion–nucleon [30] and later applied to \( \bar{K}N \) scattering [31]. The starting point is the T–matrix for any partial wave, which can be represented in closed form if one neglects for the moment the crossed channel (left-hand) cuts (for more explicit details, see [30])
\[ T = \left[ \tilde{T}^{-1}(W) + g(s) \right]^{-1}, \tag{17} \]

with \( W = \sqrt{s} \) the cm energy (note that the analytical structure is much simpler when using \( W \) instead of \( s \)). \( \tilde{T} \) collects all local terms and poles (which can be most easily interpreted in the large \( N_c \) world) and \( g(s) \) is the meson-baryon loop function (the fundamental bubble) that is resummed by e.g. dispersion relations in a way to exactly recover the right-hand (unitarity) cut contributions. The function \( g(s) \) needs regularization, this can be best done in terms of a subtracted dispersion relation and using dimensional regularization. It is important to ensure that in the low-energy region, the so constructed \( T \)-matrix agrees with the one of CHPT (matching). In addition, one has to recover the contributions from the left-hand cut. This can be achieved by a hierarchy of matching conditions,

\[
\mathcal{O}(p): \tilde{T}_1(W) = T_1^\chi(W), \\
\mathcal{O}(p^2): \tilde{T}_1(W) + \tilde{T}_2(W) = T_1^\chi(W) + T_2^\chi(W), \\
\mathcal{O}(p^3): \tilde{T}_1(W) + \tilde{T}_2(W) + \tilde{T}_3(W) = T_1^\chi(W) + T_2^\chi(W) + T_3^\chi(W) \\
+ T_1^g(W) g(s) \tilde{T}_1(W), \tag{18}
\]

and so on. Here, \( T_n^\chi \) is the \( T \)-matrix calculated within CHPT to \( \mathcal{O}(p^n) \). Of course, one has to avoid double counting as soon as one includes pion loops, this is achieved by the last term in the third equation (loops only start at third order in this case). In addition, one can also include resonance fields by saturating the local contact terms in the effective Lagrangian through explicit meson and baryon resonances (for details, see \cite{30}). In particular, in this framework one can cleanly separate genuine quark resonances from dynamically generated resonance-like states. The former require the inclusion of an explicit field in the underlying Lagrangian, whereas in the latter case the fit will arrange itself so that the couplings to such an explicit field will vanish. It was observed in \cite{31} that there are indeed two poles close to the nominal \( \Lambda(1405) \) resonance, as earlier found in the cloudy bag model \cite{32}, and later confirmed in \cite{33,34}. The physics behind these two poles was recently revealed in \cite{35}. Starting from an SU(3) symmetric Lagrangian to couple the meson octet to the baryon octet (in that limit, all octet Goldstone boson masses and all octet baryon masses are equal), one could in principle generate a variety of resonances according to the SU(3) decomposition,

\[ 8 \otimes 8 = 1 \oplus 8_s \oplus 8_a \oplus 10 \oplus \overline{10} \oplus 27. \tag{19} \]

As it turns out, the leading order transition potential is attractive only in the singlet and the two octet channels, so that one a priori expects a singlet and two octets of bound states. However, the two octets come out degenerate, see Fig. 3. This has no particular dynamical origin but rather is a consequence of the actual values of the SU(3) structure constants. In the real world, there is of course SU(3) breaking of various origins. This was parameterized in \cite{35} in terms of a symmetry breaking parameter \( x \) in the expressions for the meson \( M_i \) and baryon masses \( m_i \) as well as the
subtraction constants \( a_i \) via \( M_i^2(x) = M_0^2 + x(M_i^2 - M_0^2) \), \( m_i(x) = m_0 + x(m_i - m_0) \) and \( a_i(x) = a_0 + x(a_i - a_0) \), with \( M_0 = 368 \text{ MeV}, m_0 = 1151 \text{ MeV} \) and \( a_0 = -2.148 \), where \( 0 \leq x \leq 1 \). The motion of the various poles in the complex energy plane as a function of \( x \) is shown in Fig. 3. We note that the two octets split, in particular, one moves to lower energy (\( I = 0, 1426 \text{ MeV} \)) close to the position of the singlet (\( I = 0, 1390 \text{ MeV} \)). These are the two poles which combine to give the \( \Lambda(1405) \) as it appears in various reactions.

**Figure 3.** Trajectories of the poles in the scattering amplitudes obtained by changing the SU(3) breaking parameter \( x \) gradually. In the SU(3) limit (\( x = 0 \)), only two poles appear, one corresponding to the singlet and the other to the two degenerate octets. The symbols correspond to the step size \( \delta x = 0.1 \).

The question is now how these two different poles can actually be disentangled in experiments? For that, one has to determine the couplings of these resonances to the physical states by studying the amplitudes close to the pole and identifying them with \( T_{ij} = g_ig_j/(z - z_R) \) where \( z_R \) is the pole position and the \( g_i \) are in general complex numbers. As shown in [35], in the \( I = 0 \) channel the first resonance couples more strongly to \( \pi\Sigma \) while the second one has a stronger coupling to the \( \bar{K}N \) channel. We thus conclude that there is not just one single \( \Lambda(1405) \) resonance, but two, and that what one sees in experiments is a superposition of these two states. Then, in the case that the \( \Lambda(1405) \) is produced from the \( \bar{K}N \) initial state, the peak is narrower as if it were produced from an \( \pi\Sigma \) initial state. Therefore it is clear that, should there be a reaction which forces the initial channels to be \( \bar{K}N \), then this would give more weight to the second resonance 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 \( \bar{K}^-p \to \Lambda(1405)\gamma \) studied theoretically in Ref. [36]. It was shown there that since the \( \bar{K}^-p \) 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 \( \bar{K}^- \) or the proton. Hence the resonance is initiated from the \( \bar{K}^-p \) channel and leads to a peak structure in the invariant mass distribution which is narrower and appears at higher energies than the experimental \( \Lambda(1405) \) peaks observed in hadronic experiments performed so far. Experiments of producing the \( \Lambda(1405) \) with (real or virtual) photons have been performed or are underway or will be done at SPRING-8, JLab and ELSA. Clearly, these should be able to verify (or falsify) the two pole nature of this particular baryon resonance.
In the coupled channel approach matched to CHPT, there is also an interesting enhancement of the $I = 1$ amplitudes in the vicinity of the $\Lambda(1405)$. Independently of whether this can be interpreted as a resonance or as a cusp, the fact that the strength of the $I = 1$ amplitude around the $\Lambda(1405)$ region is not negligible should have consequences for reactions producing $\pi \Sigma$ pairs in that region. This has been illustrated for instance in [37], where the photoproduction of the $\Lambda(1405)$ via the reaction $\gamma p \to K^+ \Lambda(1405)$ 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 [38].

CONCLUDING REMARKS

There are many fascinating open problems in the large field of chiral dynamics with strange quarks. I have addressed three particular recent issues here and refer the reader to [39] for a much broader exposition. Certainly, one of the most important projects in the near future is to combine the chiral coupled channel dynamics with covariant quark models such as the Bonn one (see [40] and references therein) to solve the outstanding problem of the strong decay widths in such type of models and to get a better handle on the true nature of a variety of meson and baryon resonances, which have been one of the central issues of this conference.

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