I review the status of chiral perturbation theory in the one–nucleon sector and give some predictions to be tested. I then discuss various methods to go to higher energies (inclusion of the \( \Delta(1232) \), dispersion relations) and also to higher precision (virtual photons and isospin violation).

1 Effective field theory of QCD

In the sector of the three light quarks \( u, d \) and \( s \), QCD admits a global chiral symmetry softly broken by the quark mass term,

\[
H_{\text{QCD}} = H_{\text{QCD}}^0 + \sum_{i=u,d,s} m_i \bar{q}_i q_i ,
\]

where "light" means that the current quark mass at a renormalization scale of \( \mu = 1 \text{ GeV} \) can be treated as small compared to the typical scale of chiral symmetry breaking, \( \Lambda \chi \simeq 4\pi F_\pi \simeq 1.2 \text{ GeV} \), with \( F_\pi \simeq 93 \text{ MeV} \) the pion decay constant. The tool to investigate these issues is chiral perturbation theory (CHPT). In CHPT, the basic degrees of freedom are the Goldstone boson fields coupled to external sources and matter fields, like e.g. the nucleons. QCD is mapped onto an effective hadronic Lagrangian formulated in terms of these asymptotically observed fields. Any matrix element involving nucleons, pions, photons and so on can be classified according to its chiral dimension, which counts the number of external momenta, quark mass insertions and inverse powers of heavy mass fields. Denoting these small parameters collectively as \( p \), CHPT allows for a systematic perturbative expansion in powers of \( p \), with the inclusion of loop graphs and local terms of higher dimension. The latter are accompanied by a priori unknown coupling constants, the so–called low–energy constants (LECs). This is the so-called chiral expansion, which is nothing but an energy expansion reminiscent of the ancient Euler–Heisenberg treatment of light–by–light scattering in QED at photon energies much smaller than the electron mass. In QCD, the equivalent heavy mass scale is essentially set by the first non–Goldstone resonances, i.e. the \( \rho, \omega \) mesons. This dual expansion in small momenta and quark masses can be mapped one–to–one

\cite{Meissner98}
onto an expansion in powers of Goldstone boson loops, where an $N$–loop graph is suppressed by powers of $p^{2N}$. The leading terms are in general tree graphs with lowest order insertions leading to the celebrated current algebra (CA) results. While in the meson sector one only has terms with even powers of $p$, in the baryon sector terms with odd and even powers in $p$ are allowed, so that for example a complete one–loop calculation includes terms of order $p^3$ and $p^4$. In what follows, I will mostly be concerned with the two–flavor sector, i.e. the pion–nucleon system. For a detailed review, I refer to ref.

2 A short status report

I first consider a variety of predictions in comparison to the existing data and then discuss some predictions to be tested. Pion–nucleon scattering and pion induced pion production are the classic testing grounds for baryon CHPT and precise data have become available over the last years. In fig.1 (left panel), the CHPT predictions for the S–wave $\pi N$ scattering lengths are shown in comparison to the beautiful measurements of the level shifts of pionic hydrogen and deuterium performed at PSI. Most important is the observation that the shift from the CA point to the CHPT square is largely a pion loop effect, i.e. one sensitively probes the chiral structure of QCD. In the right panel of fig.1,

![Figure 1](image_url)

Figure 1: Left panel: CHPT prediction for the S–wave scattering lengths in comparison to the pionic hydrogen (H–band) and deuterium (D–band) data from PSI. Right panel: Prediction for $\pi^- p \rightarrow \pi^+ \pi^- n$ at leading (dashed) and next–to–leading (solid line) order. The new TRIUMF data (at $T_{\pi} = 0.18, 0.184, 0.19, 0.20 \text{GeV}$) were obtained after the CHPT prediction.

a prediction based on a next–to–leading (NLO) calculation for $\pi^- p \rightarrow \pi^+ \pi^- n$ is shown in comparison to the two older data points (large error bars) and the new ones from TRIUMF (small error bars) obtained after the prediction. Note that the agreement shows the expected pattern: The closer one is to
threshold, the better the agreement. More details on these topics are given by
Martin Sevior. Much experimental and theoretical activity has been focused
on photo–nucleon reactions. In particular, neutral pion production off the pro-
on and off deuterium has been measured very precisely. A typical result is
shown in the left panel of fig.2, based on the calculation of ref. Clearly, one
observes the unitary cusp at the \( \pi^+n \) threshold and the small value of the
electric dipole amplitude at the \( \pi^0p \) threshold is mostly due to the so–called
triangle diagram. Furthermore, these data clearly rule out the “NOLET” by
many standard deviations. Note, however, that the convergence of the chiral
expansion in this case is not very rapid and better tests are indeed given in
the P–waves, as discussed by Thomas Walcher. To my opinion, a particularly

![Figure 2: Left panel: CHPT prediction (solid line) for the electric dipole amplitude in \( \gamma p \rightarrow \pi^0p \) in comparison to the data from MAMI and SAL. The dashed line gives the prediction of the ancient “NOLET”. Right panel: Reduced total cross section for \( \gamma d \rightarrow \pi^0d \). The data from SAL are depicted by the boxes, the CHPT threshold prediction is the star on the the dotted line (indicating the threshold photon energy).](image)

nice example of chiral dynamics at work is the prediction of the electric dipole
amplitude for neutral pion production off deuterium which is very sensitive
to the elementary neutron amplitude. The chiral prediction for the deuteron is
\( E_d = (-1.8 \pm 0.2) \cdot 10^{-3}/M_\pi \), only 20% above the empirical value from
SAL \( E_{\text{emp}} = (-1.45 \pm 0.09) \cdot 10^{-3}/M_\pi \) and overlaps within 1.5 \( \sigma \). The
corresponding reduced total cross section (after subtraction of the breakup
channel) measured at SAL is shown in fig.2 (right panel) together with the
CHPT prediction. This shows that the elementary neutron amplitude is in
fact as large as predicted by CHPT and at total variance with the “NOLET”
prediction, \( E_{\sigma^+n} \approx 0 \). One can also extend these calculations to SU(3). While
there are many problems due to the large kaon mass and (subthreshold) reso-
nances, in some cases like the magnetic moments or the recoil polarization in
\( \gamma p \rightarrow K\Lambda \) the conventional framework is useful, see e.g. ref. I will come back
to this later on. Of course, there are also problems. For example, the recent
TRIUMF measurement of radiative muon capture can be used to infer the
induced pseudoscalar coupling constant. The extracted value is about 50% larger than the very precise CHPT prediction. At present, this discrepancy
has not been resolved. A more detailed status report of chiral dynamics (in-
cluding also the meson sector) can be found in refs. To end this section,
I turn to some predictions which have not been tested. The importance of
looking at spin–dependent Compton scattering was first pointed out in ref.
and detailed predictions for the so–called spin–polarizabilities can be found in ref. From these four structure constants, $\gamma_1$ and $\gamma_3$ are particularly sensitive
to the chiral pion loops. In virtual Compton scattering, one can measure the
so–called generalized polarizabilities, as discussed by Nicole d’Hose. In fig.3,
I show some genuine predictions due to Hemmert et al. Of special interest is the generalized magnetic polarizability $\beta(q^2)$ since its positive slope at
$q^2 = 0$ is a unique chiral prediction. Also, the P–wave low–energy theorems

![Figure 3: Left panel: CHPT prediction for the generalized sp
in–independent polarizabilities as a function of the squared momentum transfer. Right panel: $p^3$ prediction for the four independent spin–dependent
generalized polarizabilities as a function of the squared momentum transfer.](image)

(LETs) for pion electroproduction have so far only been used in the analysis
of the data on $\gamma^*p \to \pi^0p$ measured at NIKHEF and MAMI. More complete
angular distributions and polarization observables will have to be measured to
uniquely extract these P–waves and test the LETs. These are only a few
examples of the richness of chiral predictions in the respective threshold regions
which pose stringent tests on our understanding of chiral symmetry in QCD.
3 Perspectives

So far, I have considered the effective pion–nucleon field theory. It is most predictive in threshold situations. Naturally, one would like to extend the scheme to higher energies. I discuss here a few possibilities to do that. Also, to address certain questions one still needs a more refined and precise machinery even in the $\pi N$ sector as I will argue below. First, however, let me describe some attempts to extend the range of applicability of the EFT.

3.1 Inclusion of the $\Delta$

Although the inclusion of the decuplet was originally formulated for SU(3)\(^2\), let us focus here on the simpler two–flavor case. Among all the resonances, the $\Delta(1232)$ plays a particular role for essentially two reasons. First, the $N\Delta$ mass splitting is a small number on the chiral scale of 1 GeV, $\Delta = m_\Delta - m_N = 293$ MeV $\simeq 3 F_\pi$, and second, the couplings of the $N\Delta$ system to pions and photons are very strong, much stronger than any other nucleon resonance. So one could consider $\Delta$ as a small parameter. It is, however, important to stress that in the chiral limit, $\Delta$ stays finite (like $F_\pi$ and unlike $M_\pi$). Inclusion of the spin–3/2 fields like the $\Delta(1232)$ is therefore based on phenomenological grounds but also supported by large–$N_c$ arguments since in that limit a mass degeneracy of the spin–1/2 and spin–3/2 ground state particles appears. Recently, Hemmert, Holstein and Kambo\(^25\) proposed a systematic way of including the $\Delta(1232)$ based on an effective Lagrangian of the type $L_{\text{eff}}[U, N, \Delta]$ which has a systematic “small scale expansion” in terms of three small parameters (collectively denoted as $\epsilon$). These are $\frac{E_\pi}{\Lambda}$, $\frac{M_\pi}{\Lambda}$ and $\frac{\Lambda}{\Lambda}$, with $\Lambda \in [M_\rho, m_N, 4\pi F_\pi]$. Starting from the relativistic pion–nucleon–$\Delta$ Lagrangian, one writes the nucleon ($N$) and the Rarita–Schwinger ($\Psi_\mu$) fields in terms of velocity eigenstates (the nucleon four–momentum is $p_\mu = mv_\mu + l_\mu$, with $l_\mu$ a small off–shell momentum, $v \cdot l \ll m$ and similarly for the $\Delta(1232)$, $N = e^{-imv^\mu(x)}(H_v + h_v)$, $\Psi_\mu = e^{-imv^\mu(x)}(T_\mu v + t_\mu v)$), and integrates out the “small” components $h_v$ and $t_{\mu v}$ by means of the path integral formalism developed in Ref.\(^17\). The corresponding heavy baryon effective field theory in this formalism does not only have a consistent power counting but also $1/m$ suppressed vertices with fixed coefficients that are generated correctly (which is much simpler than starting directly with the “large” components and fixing these coefficients via reparametrization invariance). Since the spin–3/2 field is heavier than the nucleon, the residual mass difference $\Delta$ remains in the spin–3/2 propagator and one therefore has to expand in powers of it to achieve a consistent chiral power counting. The technical details how to do that, in particular how to separate the spin–1/2 components from the spin–3/2 field, are given in ref.\(^25\) Let
me now consider two examples. The first one is related to the scalar nucleon form factor (for an early calculation capturing the essence of the small scale expansion, see ref.\textsuperscript{26}). The scalar form factor is given by

\[
\sigma_{\pi N}(t) = \frac{1}{2}(m_u + m_d) \langle p' | \bar{u}u + \bar{d}d | p \rangle, \quad t = (p' - p)^2. \tag{2}
\]

Of particular interest in the analysis of \(\sigma_{\pi N}\) is the Cheng–Dashen point, \(t = 2M_\pi^2, \nu = 0\) (at this unphysical kinematics, higher order corrections in the pion mass are the smallest) and one evaluates \(\Delta\sigma_{\pi N} \equiv \sigma_{\pi N}(2M_\pi^2) - \sigma_{\pi N}(0)\). To one loop and order \(O(\epsilon^3)\), \(\Delta\sigma_{\pi N}\) is free of counter terms and just given by two simple one loop diagrams (one with an intermediate \(N\) and the other with an intermediate \(\Delta\))\textsuperscript{26,27},

\[
\Delta\sigma_{\pi N} = \frac{3g_\pi^2 M_\pi^3}{64\pi F_\pi^2} + \frac{g_\pi^2 M_{\pi^0}^2}{6\pi^2 F_\pi^2} \frac{M_\pi^4}{M_\Delta^4} \left\{ \frac{5}{18} - \frac{\pi}{24} + \frac{5}{6} \ln \frac{2\Delta}{M_\pi} + \ldots \right\} = (7.4 + 4.1) \text{ MeV} = 11.5 \text{ MeV}, \tag{3}
\]

where the first term comes from the intermediate nucleon and scales as \(M_\pi^3\) whereas the other terms come from the loop graph with the intermediate \(\Delta(1232)\) and scale as \(M_\pi^4/\Delta\), which are therefore both \(O(\epsilon^3)\). In the chiral expansion of QCD, however, the first term is \(O(p^3)\) while the second is of order \(p^4\). Therefore, the epsilon expansion can be considered as a rearrangement of the chiral expansion. The correction due to the delta goes in the right direction, but the resulting number is still about 30% below the one of the dispersion–theoretical analysis (supplemented by chiral symmetry constraints) of Gasser, Leutwyler and Sainio, \(\Delta\sigma_{\pi N} = (15 \pm 1) \text{ MeV}\)\textsuperscript{28} This points towards the importance of higher order effects, as I will discuss in more detail below. The interplay between the chiral and the small scale expansion has been studied in ref.\textsuperscript{29} There, the nucleons’ electroweak form factors are considered. These calculations also serve as a first exploratory study of renormalization and decoupling within the small scale expansion. In particular, it is argued that the pion–nucleon effective theory in the presence of the \(\Delta\) differs from the “pure” \(\pi N\) EFT, in particular, the \(\beta\)–functions related to various dimension three operators needed for the renormalization obtain additional terms in the EFT with explicit deltas. As an example, consider the electromagnetic nucleon structure as parametrized by the electroweak formfactors (ff). To third order in the chiral expansion, the isovector Pauli ff can be predicted without any LEC, whereas the expression for the Dirac ff contains one LEC. This can be adjusted to give the empirical isovector charge radius. The numerical value of this LEC differs, of course, in the chiral and in the small scale expansion. The
predictions of the chiral and the small scale expansion are consequently indistinguishable for \( F_V^1(q^2) \), cf. fig.4. For the isovector Pauli ff, the novel graphs with intermediate deltas (see right panel in fig.4) improve the prediction of the radius and lead to a better description of the ff at small and moderate momentum transfer, see. fig.4. Of course, this EFT is supposed to work best in the \( \Delta \) region, see e.g. the talk by Gellas on the \( \Delta \to N\gamma \) transition form factors and the related work on pion photoproduction in the \( \Delta \) region. Furthermore, there has been sizeable activity in computing real and virtual Compton scattering in the framework of the epsilon expansion, see e.g. What is still missing is a complete renormalization of the generating functional to third order and some explicit calculations to \( \mathcal{O}(\epsilon^4) \), to compare e.g. with the precise chiral predictions for the nucleons’ em polarizabilities.

3.2 Marriage with dispersion relations

Another possibility of extending the range of applicability of CHPT (or to sharpen the chiral predictions at low energies) is the use of dispersion relations (DR), for some early work in the meson sector see e.g. (for a more pedagogical introduction, see e.g. ref). Let me first make some general remarks. A generic amplitude for some process can be written in terms of an \( n \)-times subtracted dispersion relation (where the number of subtractions depends on
the convergence properties and generally is small)
\[ A(t) = \sum_{i=0}^{n-1} a_i t^i + \frac{t^n}{\pi} \int_{t_0}^{\infty} \frac{dt'}{t^{n+1}} \frac{\text{Im} A(t)}{t - t' - i\epsilon}. \]  

(4)

These subtraction constants \( a_i \) play a role similar to the LECs in the corresponding chiral calculation and thus chiral symmetry can be used to constrain them. Similar remarks hold for the spectral function \( \text{Im} A(t) \) at low \( t \). Of course, there is no free lunch - precise cross section data at intermediate and large momenta are needed to make DRs a viable and accurate tool (note that sometimes the large momentum behaviour is modeled giving rise to some unwanted model–dependence). Note also that the analytic structure in the two approaches is the same (as demanded by analyticity). In the dispersive approach, one essentially sums up classes of loop diagrams (the so–called unitarity corrections) allowing one to obtain a precise representation at low and higher energies. What the dispersive treatment can in general not provide is linking Green functions of different processes (as it is the case in CHPT) and only under special circumstances the enhancement of certain LECs due to IR logarithms can be unraveled. Clearly, the combination of dispersion theory with CHPT constraints is a powerful tool which has not yet been explored in big detail in the nucleon sector. One exception to this is the work on the scalar form factor in ref. Since the nucleon is a composite object, \( \sigma_{\pi N}(t) \) must fulfill a once–subtracted DR. The corresponding spectral function reads
\[ \text{Im} \sigma_{\pi N}(t) = \frac{3}{2} \frac{\Gamma_{\pi}^+(t)}{4m^2 - t} \sqrt{1 - \frac{4M^2}{t}}, \]  

(5)

with \( \Gamma_{\pi} \) the scalar form factor of the pion, \( \langle \pi^a(p')|\bar{u}(\bar{u}u + \bar{d}d)|\pi^b(p)\rangle \sim \Gamma_{\pi}(t) \) and \( f_{0}^{+}(t) \) the \( I = J = 0 \) \( \pi N \) partial wave (analytically continued from the existing data). The evaluation of the dispersive integral leads to the abovementioned result for \( \Delta\sigma_{\pi N} \). The enhancement compared to the one loop chiral (or small scale) calculation can be understood as follows. First, in the scalar \( \Gamma_{\pi} \) one observes sizeable two (and higher) loop corrections even at moderate \( t \). This is related to the strong pionic final–state interactions (FSI) in the isospin zero S–wave. Similarly, the \( \pi N \) partial wave \( f_{0}^{+}(t) \) exhibits strong \( \pi N \) FSI which can not be accounted for by a one–loop calculation. Consequently, the corresponding spectral function shows a strong enhancement at the two–pion threshold. This is at the heart of the large shift of the scalar form factor. It is also worth pointing out that this does not mean that CHPT breaks down. In fact, we are dealing with a small quantity which has a slowly converging chiral expansion. Therefore, dispersion theory properly constrained by chiral
symmetry is the better tool to analyze this quantity (for details, see ref. 28). Let me now turn to another process. The Mainz group has investigated pion photoproduction in a dispersion theoretical framework. While they did not enforce chiral constraints and did not include the data very close to threshold in their analysis, the results for charged and neutral pion production off the proton are very (even surprisingly) similar to the CHPT predictions. In case of $\pi^0$ production off the neutron, they find a sizeably smaller electric dipole amplitude than the CHPT prediction. This might, however, be an artefact of the not so precise data for this channel. It certainly would be very interesting to combine their machinery with chiral symmetry constraints. Clearly, more work in this direction should be performed. I also would like to point out that for mesonic processes, the combination of DR with CHPT has become a standard tool and it has been applied successfully to many reactions, like e.g. $\pi\pi \rightarrow \pi\pi$, $\gamma\gamma \rightarrow \pi^0\pi^0$, $K^0 \rightarrow \pi^0\gamma\gamma$, $\eta \rightarrow 3\pi$ and so on.

3.3 Higher precision: Virtual photons and isospin violation

Let me start with some general remarks. Since a large body of elastic scattering and charge exchange data exists, one has the possibility of deducing bounds on isospin violation from simple triangle relations, which link e.g. the processes $\pi^\pm p \rightarrow \pi^\pm p$ and $\pi^- p \rightarrow \pi^0 n$. Care has, however, to be taken since there are two sources leading to isospin violation. One is the “trivial” fact that electromagnetism does not conserve I–spin, since it couples to the charge. The other one is a strong effect, linked to the difference of the light quark masses, $m_d - m_u$. This is essentially the quantity one is after. In terms of the symmetry breaking part of the QCD Hamiltonian, i.e. the quark mass term, we have $H_{\text{QCD}}^{ab} = m_u \bar{u}u + m_d \bar{d}d = \frac{1}{2}(m_u + m_d)(\bar{u}u + \bar{d}d) + \frac{1}{2}(m_u - m_d)(\bar{u}u - \bar{d}d)$, so that the strong I–spin violation is entirely due to the isovector term. This observation lead Weinberg to address the question of I-spin violation in the pion and the pion–nucleon sector, with the remarkable conclusion that in neutral pion scattering off nucleons one should expect gross violations of this symmetry, as large as 30%. Only recently an experimental proposal to measure the $\pi^0p$ scattering length in neutral pion photoproduction off protons below the secondary $\pi^+n$ threshold has been presented and we are still far away from a determination of this elusive quantity. At present, there exist two phenomenological analyses which indicate isospin breaking as large as 7% in the S–waves (and smaller in the P–waves). Both of these analyses employ approaches for the strong interactions, which allow well to fit the existing data but can not easily be extended to the threshold or into the unphysical region. What is, however, most disturbing is that the electromagnetic corrections have been
calculated using some prescriptions not necessarily consistent with the strong interaction models used. One might therefore entertain the possibility that some of the observed I–spin violation is caused by the mismatch between the treatment of the em and strong contributions. Even if that is not the case, both models do not offer any insight into the origin of the strong isospin violation, but rather parametrize them. In CHPT, these principle problems can be circumvented by constructing the most general effective Lagrangian with pions, nucleons and virtual photons,

\[ \mathcal{L}_{\pi N} = \mathcal{L}^{(1)}_{\pi N} + \mathcal{L}^{(2)}_{\pi N, \text{str}} + \mathcal{L}^{(2)}_{\pi N, \text{em}} + \mathcal{L}^{(3)}_{\pi N, \text{str}} + \mathcal{L}^{(3)}_{\pi N, \text{em}} + \ldots \]  

where the superscript gives the chiral dimension. It is important to note that the electric charge counts as a small momentum, based on the observation that \( e^2/4\pi \sim M^2_\pi/(4\pi F_\pi)^2 \sim 1/100 \). Since a virtual photon can never leave a diagram, the local contact terms only start at dimension two. For the construction of this effective Lagrangian and a more detailed discussion of the various terms, see ref.\textsuperscript{40} Let me just review the results obtained so far (for more details, see the talk by Steininger\textsuperscript{41}). First, Weinberg’s finding concerning the scattering length difference \( a(\pi^0 p) - a(\pi^0 n) \), which is entirely given by a dimension two term \( \sim m_u - m_d \), could be confirmed. This is not surprising because for the case of neutral pions there is no contribution from the em Lagrangian of order two and three. This will change at fourth order. Second, it was noted in\textsuperscript{40} that the I–spin breaking terms in \( a(\pi^0 p) \) can be as large as the I–spin conserving ones. It is also important to note that the em effects are entirely due to the pion mass difference. Another quantity of interest is the pion–nucleon \( \sigma \)–term. Of course, one now has to differentiate between the one for the proton and the one for the neutron, whose values to third order differ by the strong neutron–proton mass difference. For the proton one finds \( \sigma_p(0) = \sigma^{IC}_p(0) + \sigma^{IV}_p(0) = 47.2 \text{ MeV} - 3.9 \text{ MeV} = 43.3 \text{ MeV} \), which means that the isospin–violating terms reduce the proton \( \sigma \)–term by \( \sim 8\% \). The electromagnetic effects are again dominating the isospin violation since the strong contribution is just half of the strong proton–neutron mass difference, 1 MeV. Furthermore, one gets \( \sigma_p(2M^2_\pi) - \sigma_p(0) = 7.5 \text{ MeV} \), which differs from the result in the isospin limit at \( \mathcal{O}(p^3) \) by a few percent. This small difference is well within the uncertainties related to the so–called remainder at the Cheng–Dashen point\textsuperscript{42}. Again, these corrections should be considered indicative since a fourth order calculation is called for. As for the channels involving charged pions, we have obtained some preliminary results. Considering only the pion and nucleon mass differences (but neglecting all other virtual photon effects), the CA prediction for the triangle–deviation (normalized to the charge–exchange
amplitude)

\[ \Delta_{\pi N} = \frac{(T^{\text{EL}+} - T^{\text{EL}-}) - \sqrt{2} T^{\text{SCX}}}{\sqrt{2} T^{\text{SCX}}}, \quad (7) \]

with \( T^{\text{EL}+} \) the elastic scattering amplitude \( \pi^\pm p \to \pi^\pm p \) and \( T^{\text{SCX}} \) denotes the single charge exchange \( \pi^- p \to \pi^0 n \), is 0.8% (at threshold), and to second order we obtain \( \Delta_{\pi N} = 2.3\% \). This number is smaller than the one obtained in the phenomenological analysis [39, 38] but larger than expected by many. Clearly, the one loop calculation has to be completed to draw definite conclusions from this. Such an investigation is underway and the details will be soon available.

It is worth to emphasize again that we have a consistent machinery at hand to simultaneously calculate the strong and the em isospin violating effects.

### 4 Other topics

Space does not allow to cover all the interesting developments in this field since the last Baryons Conference in 1995. Still, I wish to briefly discuss two topics, which have received lots of attention over the last few years. Consider first the extension to three flavors. While that is technically straightforward, the larger kaon mass and the appearance of (subthreshold) resonances in certain channels clearly limit the usefulness of baryon CHPT. However, for particular reactions and processes, stringent predictions can be made. As an example, I mention the recent bounds on the strange magnetic form factor of the nucleon, which will be of use in analyzing the data from SAMPLE and CEBAF to constrain the strange magnetic moment [44].

Recently, a cut–off scheme was proposed to improve the convergence properties of SU(3) baryon CHPT [45]. It is argued that the natural extension of the hadrons suppresses the unwanted short–distance (non–chiral) physics. Technically, this is implemented by a cut–off which essentially decreases the large kaon and eta loop contributions. To my opinion, in its present formulation this is not a viable alternative since one shuffles the large loop contributions into a string of large higher dimension terms generated through the cut–off procedure. Certainly, this deserves more study. Thirdly, there has also been made considerable progress in setting up non–perturbative resummation schemes to deal with the abovementioned resonances, see e.g. [46, 47]. With a few parameters, one is able to describe a wealth of data for a large energy range. It should be noted, however, that such approaches seem to be more useful for studying the nature of some resonances (or quasi bound–states) than for directly testing the consequences of chiral symmetry breaking.

Another subject of much discussion is the extension of EFT approaches to few–nucleon systems. This is reviewed by David Kaplan. I only wish to add that one still has to study in more detail alternative schemes like e.g. the
hybrid approach proposed by Weinberg in which nuclear matrix elements are calculated using realistic wave functions and applying CHPT constraints to the interaction kernel. This approach has been applied quite successfully to neutral pion photoproduction off deuterium and to $\pi - d$ scattering.

Chiral nucleon dynamics has matured considerably over the last few years and many more precise data are becoming available. Furthermore, there is progress in extending the framework to deal with higher energies, the three flavor case and few–nucleon systems.

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