CHIRAL SYMMETRY AND PARITY–VIOLATING MESON–NUCLEON VERTICES∗ #

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ABSTRACT

In this lecture, I review progress made in the calculations of the parity-violating meson-nucleon interaction regions. The underlying framework is the topological chiral soliton model of the nucleon. Emphasis is put on the computation of theoretically and experimentally accessible nuclear parity violating observables. I stress the importance of the interplay of strong and weak interactions which makes this field interesting and challenging. I also discuss recent developments pointing towards the importance of strange quark admixtures in the proton wave function.

INTRODUCTION

Our understanding of the hadronic weak interactions has progressed considerably in the last two decades. Still, the almost unique tool to study the non-leptonic, strangeness conserving part of the weak Hamiltonian are few-nucleon systems. In general, nuclear parity-violating (pv) observables cannot be calculated reliably enough so that we could deduce stringent limits on the standard model from them. Stated in another way: We are still far away from extracting e.g. the Weinberg angle to some decent precision from nuclear parity violation. On the other hand, there are now very sophisticated parametrizations of the strong force between two nucleons available which allow us to test our understanding of the hadronic weak interactions in terms of meson exchanges. Direct W- or Z-exchange between nucleons is wiped out by the hard core of the NN-force, but there still remains a long-range component

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of the weak interactions between nucleons, which can be parametrized in terms of
pv meson-nucleon interaction vertices. One way to calculate these pv couplings is
to make use of the quark model.\textsuperscript{1} There is, however, a considerable uncertainty in
these calculations which stems from the fact that the pertinent multiquark opera-
tors have to be calculated at low energies ($E \ll M_{W,Z}$). Gluonic corrections arise,
and unavoidably one enters the non-perturbative regime where the strong coupling
constant $\alpha_s(q^2)$ becomes larger than unity. These problems are most pronounced in
the case of the pion which dominates the long-range part of the pv potential. The
Goldstone-boson-like character has always posed problems for quark-model practi-
tioners. Quite contrary, the recently popular topological soliton models of the nu-
cleon like the Skyrme model\textsuperscript{2} and generalizations thereof\textsuperscript{3} naturally incorporate the
pseudoscalar as well as the low-lying vector multiplets. Here we have reached our
starting point for a calculation of the parity-violating meson-nucleon couplings and
form factors.\textsuperscript{4,5,6,7} The soliton approach to the nucleon is far from being perfect,
but it has the conceptual advantage that it allows for a simultaneous calculation of
the strong and weak interaction regions, a point which is generally overlooked by
quark model enthusiasts. Furthermore, nuclear parity violation can also be used as
a testing ground to find out the limitations of the soliton scenario — often more can
be learned from the failures of a model than from its successes.

Another interesting aspect of nuclear parity violation is the quest for finding
few-nucleon systems which can be calculated with some reliability and where the
experimenters have a change of detecting a clear signal. Here, I will focus on two
rather different systems. In proton-proton scattering, one can observe longitudi-
nal asymmetries of the order $10^{-7}$, which appear to be awfully small. However,
progress in experimental techniques now allows for experiments with an accuracy of
$\delta A_L \simeq \pm 1.0 \cdot 10^{-8}$ and therefore a fairly sensitive test of the meson-exchange pic-
ture underlying the theoretical description of this process. A very different system is
the nucleus $^{18}$F, in which nuclear amplification takes place and the observed circular
polarization of emitted $\gamma$-rays is of the order $2 \cdot 10^{-4}$. Luckily for the theorists, the
$\beta$-decay of the daughter nucleus $^{18}$Ne allows one to gauge the rather involved shell
model calculations,\textsuperscript{8} although sceptical minds tend to look at these calculations with
a certain dose of disbelief. As we will see, too few “good” nuclear systems are con-
sidered at present and therefore the restrictions on the pv meson-nucleon couplings
are by far too soft.

Finally, during the last year, effective field theory methods have been used to gain
further insight into the strength of the pv meson–nucleon couplings.\textsuperscript{9} These results
seem to indicate a large enhancement from operators involving strange quarks to
various coupling constants. Furthermore, some couplings not considered so far (like e.g. the pv πNγ vertex) might be of importance. I will discuss these topics in the end of this lecture.

**PV MESON-NUCLEON INTERACTION REGIONS**

In the meson-exchange parametrization of the weak nuclear force, one usually only considers the exchange of charged pions and the vector mesons ρ and ω. CP invariance does not allow for the coupling of neutral scalar or pseudoscalar mesons to nucleons, eliminating the infamous scalar mesons, the η, η' and the π⁰ (the δ⁺ are considered a form factor corrections to the π⁺-exchange). Then there remains the φ(1020) — its coupling to the nucleon is generally supposed to be OZI-suppressed and not considered.¹⁰ This might, however, be a too simplistic approach in light of the discussion surrounding the admixture of strange operators into the proton’s wavefunction. At present no final conclusion can be drawn and I will make life easy on us and neglect the φ for the time being. I will pick up this theme in the final section.

Unavoidably I will have to define the basic couplings which parametrize the pv nuclear potential. For the pion, there is only a ΔI = 1 (isovector) coupling (to first order in the pion field)

\[
\mathcal{L}_{pv}^{\pi N} = -\frac{G_\pi(q^2)}{\sqrt{2}} \frac{E}{M_N} \chi_f^\dagger (\vec{\tau} \times \vec{\pi})_3 \chi_i
\]

with χi,f denoting nucleon spinors, q² the invariant momentum transfer squared at the πN-vertex and I have (for simplicity) given the non-relativistic reduction of this vertex. In the case of the ω, ΔI = 0 and ΔI = 1 couplings are possible,

\[
\mathcal{L}_{\omega N}^{pv} = \chi_f^\dagger \left[ h_0^\omega(q^2) \tau^a + h_1^\omega(q^2) \tau^3 \right] \left[ \frac{E}{M_N} \vec{\sigma}_T + \vec{\sigma}_L \right] \cdot \vec{\omega} \chi_i
\]

with \(E = \left(M_N^2 + \vec{q}^2/4\right)^{1/2}\) and \(\vec{\sigma}_{L,T}\) the longitudinal and transverse spin-operator, respectively. For the ρ, one has isoscalar, isovector and isotensor vertices

\[
\mathcal{L}_{\rho N}^{pv} = \chi_f^\dagger \left[ h_0^\rho(q^2) \tau^a + h_1^\rho(q^2) \delta^{a3} + \frac{h_2^\rho(q^2)}{2\sqrt{6}} (3\tau^3 \delta^{a3} - \tau^a) \right] \times \left[ \frac{E}{M_N} \vec{\sigma}_T + \vec{\sigma}_L \right] \cdot \vec{\omega} \chi_i - \frac{iE}{2M_N^2} h_1^\rho(q^2) \chi_f^\dagger \vec{\sigma} \cdot \vec{q} (\vec{\tau} \times \vec{\rho}^0)_3 \chi_i.
\]

³
Generally, the coupling $h_\rho^1$ is neglected, but I will not follow this historical path here. From Eqs. (1) – (3) it becomes obvious that the pv interaction regions are characterized by coupling constants $h_M = h_M(q^2 = 0)$ and form factors $F_M^{pv}(q^2) = h_M(q^2)/h_M$ (in case of the pion, I use $G_\pi \equiv h_\pi$).

In the topological chiral soliton model underlying the calculation of the pv vertices, nucleons arise as solitons of a non-linear meson theory. This non-linear meson theory is constructed in harmony with chiral symmetry and anomaly constraints and all its parameters are fixed from mesonic reactions like e.g. $(\rho^0 \rightarrow \pi^+\pi^-, \omega \rightarrow \pi^+\pi^-\pi^0, \omega \rightarrow \pi^0\gamma, \ldots)$. The Lagrangian and its parameters are completely determined in the meson sector, and the calculation of nucleon properties proceeds without any new parameters, i.e. no fudging is possible! That is certainly an appealing aspect of the soliton approach to the nucleon and it poses several restrictions. Of course, the model does not perfectly predict all nucleon properties.

Now: How can we calculate the pv couplings appearing in Eqs. (1) – (3)? For that, we consider the current $\times$ current form of the weak Hamiltonian with the currents being of ($V-A$)-type. To pick out the pv pieces, consider the $\Delta I = 0, 1$ or 2 components of products like $V_\mu A^\mu$ and $I_\mu A^\mu$, with $V_\mu$ the vector, $I_\mu$ the isoscalar and $A_\mu$ the axial current. These currents are already given in terms of the meson fields which make up the soliton, and their explicit expressions can be found e.g. in Ref. [5]. One then makes use of the “background-scattering” method, which amounts to an expansion of the meson fields around the soliton background. For the pion, we write

$$\vec{\pi} = \vec{\pi}_S + \vec{\pi}_f$$

and similarly for $\rho$ and $\omega$. $\vec{\pi}_S$ is the “hard” component of the pion field making up the soliton and $\vec{\pi}_f$ a small pionic fluctuation (“soft” component). Inserting the expressions (4) into the soliton currents and these into the weak Hamiltonian, all one has to do is to find the terms linear in $\vec{\pi}_f$ (or $\vec{\rho}_f$ or $\vec{\omega}_f$). Quantizing the respective operators which are given in terms of the collective variables $(A, \dot{A})$,$^2,^3$ one can immediately read off the coupling constants and form factors for the meson under consideration. In particular, one cannot only construct pv meson-nucleon vertices, but also the equivalent pv $N\Delta$-transition couplings. I will come back to this point later on. For details, the interested reader should consult Refs. [5,6]. I will not give any explicit formula here, but rather make a few comments on the results of Ref. [5]. First, the pv $\pi N$ coupling is completely dominated by the neutral current contribution.
Table: Effective weak meson-nucleon coupling constants in units of $10^{-7}$.

We present the result of the soliton model calculation of Kaiser and Meißner (KM) [5,6] together with the quark model results of Desplanques et al. (DDH) [1] as well as Dubovik and Zenkin (DZ) [16]. The value for $h^{'1}_\rho$ in the column DDH is taken from Holstein’s calculation in ref. [11]. The “reasonable ranges” (RR) defined by DDH are also given. The column AH gives the best fit values of Adelberger and Haxton.\(^{17}\)

|       | KM  | DDH | DZ  | AH   | RR   |
|-------|-----|-----|-----|------|------|
| $\tilde{F}_\pi$ | 0.6 | 10.8| 3.1 | 5.0  | 0.0→27.1 |
| $F_0$  | 5.9 | 15.9| 11.5| 8.0  | -15.9→43.0 |
| $F_1$  | 0.2 | 0.3 | -0.5| 0.3  | -0.1→0.6  |
| $F_2$  | 5.3 | 13.3| 9.3 | 9.8  | 10.6→15.3 |
| $H_1$  | 1.7 | 0.5 | 0.0 | 0.0  | —     |
| $G_0$  | 24.5| 8.0 | 16.3| 27.0 | -23.9→43.1 |
| $G_1$  | 4.1 | 4.8 | 9.2 | 10.0 | -3.3→8.0  |

The charged current contribution can be estimated in the factorization approximation, $G^{CC}_\pi = \cos^2 \theta_c < \pi |A_\mu|0 < p|V_\mu|n > = G_F \cos^2 \theta_c f_\pi (M_n - M_p)$. The electromagnetic mass difference of the neutron and the proton is well-reproduced in the model,\(^4\) whereas a strong part of $M_n - M_p$ is somewhat underestimated.\(^{14}\)

Taking as an upper limit the empirical value $M_n - M_p \simeq 1.3$ MeV, we find $A = G^{NC}_\pi / G^{CC}_\pi \sim 13.5$, consistent with previous estimates\(^{15}\) and the quark model calculations of Ref. [1] ($A \simeq 24$). The numerical value for the effective pion-exchange coupling, $\tilde{F}_\pi = g_{\pi NN} G_\pi / \sqrt{32}$ with $g_{\pi NN}$ the strong $\pi N$ coupling, is considerably smaller in the soliton model than in the quark model, we find $\tilde{F}_\pi^{\text{sol}} = 0.6$ vs rsus $\tilde{F}_\pi^q = 10.8,^8$ or $\tilde{F}_\pi^q = 3.1^{16}$ (in units of $10^{-7}$).

For the vector meson couplings, the results are less different. Using the standard definitions $F_i = -g_{\rho NN} h_\rho^i / 2$, and $G_i = -g_{\omega NN} h_\omega^i / 2$ and $H_1 = -g_{\rho NN} h_\rho^1 / 4$, we find
that the soliton and the quark model predict the following pattern for the $\rho$-couplings: $F_0 \gtrsim F_2 \gg F_1$. The absolute values of the constant $s F_i$, are, however, reduced in the soliton approach. For $H_1$, the soliton model predicts a value three times as large as the quark model. In the case of the $\omega$, all calculations give $G_0 > G_1$, but $G_0$ is considerably enhanced in the soliton approach and close to the “best-fit” estimate of Adelberger and Haxton. These results are summarized in the table.

As already stated, the calculation of $\text{pv} \, N\Delta M$ ($M = \pi, \rho, \omega$) transition vertices proceeds along the same lines and only differs at the step when one quantizes the collective coordinates. In our model, the $\text{pv} \, \pi N\Delta$ coupling has $\Delta I = 0$ and 1 components and reads non-relativistically

$$L^{\text{pv}}_{N\pi} = \left( \frac{E}{4M_N^2} \right) \left[ \delta_{ab} h_1^0(q^2) + e^{3ab} h_1^1(q^2) \right] \chi^\dagger_{\Delta} \bar{S} \cdot \bar{q} \sigma \cdot \bar{q} T^a \chi_{N \pi} + \text{h.c.}$$

with $\bar{S}$ and $\bar{T}$ the conventional $N\Delta$ transition spin and isospin operators, respectively. It is easy to convince oneself that $h_1^0(q^2) = 0$ in this model, naively, a non-vanishing isoscalar $\pi\Delta N$-vertex would lead to a non-zero CP-violating $\pi N$-vertex (a more foolproof argument is given in Ref. [6]). The isovector vertex does not vanish, and for the “minimal” model we find

$$h_1^1(0)/G_\pi = 1.10 \, .$$

The presence of $\pi\rho\omega$-correlations in the effective action tends, however to decrease this ratio. For the $V\Delta N$-couplings, we find ($V = \rho$ or $\omega$):

$$h_{\rho\Delta N}^i(q^2) = \frac{3}{\sqrt{2}} h_{\rho}^i(q^2) \quad (i = 0, '1, 2) , \quad h_{\omega\Delta N}^1(q^2) = \frac{3}{\sqrt{2}} h_{\omega}^1(q^2)$$

and $h_{\rho\Delta N}^1(q^2) = h_{\omega\Delta N}^0(q^2) \equiv 0$. These predictions are insofar interesting since in the seventies it was argued that e.g. the $\rho\Delta N$-couplings are negligible — quite in contrast to our results. A recently performed quark model calculation by Feldman et al. along the lines of DDH gives results rather different from the soliton model predictions. The source of these discrepancies is not yet understood. In a similar fashion, one easily derive the corresponding $M\Delta \Delta$ vertices,

$$G_{\pi\Delta \Delta}(q^2) = G_\pi(q^2)$$

$$h_{\rho\Delta \Delta}^i(q^2) = \frac{1}{5} h_{\rho}^i(q^2) \quad (i = 0, '1, 2) ; \quad h_{\rho\Delta \Delta}^1(q^2) = h_{\rho}^1(q^2)$$

$$h_{\omega\Delta \Delta}^0(q^2) = h_{\omega}^0(q^2) ; \quad h_{\omega\Delta \Delta}^1(q^2) = \frac{1}{5} h_{\omega}^1(q^2)$$

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The calculation of the associated weak form factors proceeds in a straightforward way. In Fig. 1 we show the weak $\pi N$ form factor $G_\pi(q^2)$ in comparison with the equivalent strong form factor $G_{\pi NN}(q^2)$ as well as the monopole with cut-off $\Lambda = 1$ GeV. As it turns out, all form factors can be fitted by monopoles at low $\vec{q}^2$, 

$$F^M(\vec{q}^2) = h_M \frac{\Lambda^2}{\Lambda^2 + \vec{q}^2}$$  \hspace{1cm} (9)$$

and are very similar to the respective strong form factors. This is the first time that such a calculation has been performed and its result can be understood as follows: The intrinsic scale of the meson-nucleon interaction regions is set by the topological baryon charge radius, $r_B \simeq 0.5$ fm. From that, one can deduce a cut-off scale $\Lambda \simeq \sqrt{6}/r_B \simeq 1$ GeV. It is, however, not that simple because the dynamical treatment of the vector mesons modifies this result. Defining by $R_M$ the ratio of the (averaged) weak to strong $MN$ cut-offs (all form factors of monopole type), we find

$$R_\pi = 1.15 \ , \quad R_\rho = 0.91 \ , \quad R_\omega = 0.77$$  \hspace{1cm} (10)$$
which justifies within the accuracy of the model the assumption of taking the same form factors for the strong and weak vertices as it was done e.g. by Driscoll and Miller\textsuperscript{20} in their study of the pv pp-interaction.

Another topic which can be discussed in the framework of the chiral soliton model are the corrections of pv two–pion exchange. One motivation to do this is that correlated $2\pi$–exchange gives rise to the intermediate range attraction of the parity–conserving $NN$ interaction. Furthermore, recent investigations point towards the importance of pv $2\pi$ exchange even below production threshold.\textsuperscript{31} Using the soliton model, Norbert Kaiser and I have shown that the inclusion of pion loops gives the intermediate range attraction with just the right strength as compared to the Paris potential.\textsuperscript{32} Similarly, we have worked out corrections to the various pv $\rho N$ couplings.\textsuperscript{33} The effects of irreducible two–pion corrections are generally small, of the order of $10\ldots20\%$. This is in agreement with older dispersion–theoretical investigations.\textsuperscript{34} So we finally have all the tools at hand to make contact to experiment.

PARITY-VIOLATION IN PROTON-PROTON SCATTERING

The simplest system in which one can probe certain components of the weak pv inter-nucleon force is the two nucleon system. By scattering polarized protons off a hydrogen target, parity violation shows itself in a non-vanishing longitudinal asymmetry,

$$A_L = \frac{\sigma^+ - \sigma^-}{\sigma^+ + \sigma^-}$$

assuming a 100\% longitudinal polarization of the beam and having taken care of the Coulomb-corrections $\sigma^{\pm}$ are the cross-sections for scattering positive/negative helicity protons from an unpolarized target. The calculation of this process in the DWBA as pioneered by Brown \textit{et al.}\textsuperscript{21} goes as follows. One splits the total scattering amplitude $F_{ss'}$ into a strong and weak part

$$F_{ss'} = F_{ss'} + f_{ss'}$$

for total spins $s$ and $s'$. Now it is of utmost importance to take into account the strong distortions, \textit{i.e.} calculating the weak scattering amplitude with distorted waves $\psi_{s}^{(-)}$ and $\psi_{s}^{(+)}$, \textit{i.e.} $f_{ss'} = \langle \psi_{s}^{(-)}|V_{pv}|\psi_{s}^{(+)} \rangle$ with $V_{pv}$ the pv one-meson-exchange potential. It should be pointed out that the strong distortions govern the energy-dependence of the analyzing powers $A_L$. Recently, Driscoll and Miller\textsuperscript{20} have done the most complete calculation based on the Bonn-Potential\textsuperscript{21} for the strong
force and an equivalently constructed weak potential with the pv couplings taken from the quark model\(^1\) and using the same vertex functions for the weak and strong form factors. I should point out here that for obvious reasons there is no pion contribution to this process and one essentially tests the vector-meson couplings \(h^\rho_{pp} = h_0^\rho + h_1^\rho + h_2^\rho / \sqrt{6}\) and \(h^\omega_{pp} = h_0^\omega + h_1^\omega\).

Recently, Doug Driscoll and I have repeated this calculation\(^2\) by including the soliton model predictions \(h^\rho_{pp} = -5.15 \cdot 10^{-7}\) and \(h^\omega_{pp} = -8.20 \cdot 10^{-7}\). The resulting curve for \(A_L\) is shown in Fig. 2, for the quark\(^1,16\) and the soliton model.\(^2\) The shape of the curve as predicted by the soliton model follows closer the empirical trend suggested by the low-energy data.\(^23\) In fact, a \(\chi^2\) calculation for the three curves shown in fig.2 gives \(\chi^2 = 34/3\) (DDH), \(\chi^2 = 26/3\) (DZ) and \(\chi^2 = 8/3\) (KM) as discussed in ref.26 (at that time, the Bonn result was not available). Also, the maximum at \(p_{lab} = 0.95\) GeV/c is flatter than in the calculation using the quark model parameter. Furthermore, the energy at which the asymmetry changes sign is larger than the quark model predicts, which can be traced back to the fact that in the soliton model \(h^\omega_{pp} > h^\rho_{pp}\), in contrast to the quark model with \(h^\omega_{pp} < h^\rho_{pp}\). Of particular interest is the value of \(A_L\) at 222 MeV. This is the energy selected for an upcoming \(pp\) parity violation measurement at TRIUMF because \(\delta(1^S_0) + \delta(3^P_0) = 0\) at this energy and the \(j = 0\) contribution to the analyzing power consequently vanishes. The measurement of the dominant \(j = 2\) contribution gives a different combination of \(h^\rho_{pp}\) and \(h^\omega_{pp}\) than the \(j = 0\) contribution to \(A_L\), which is already measured at 15 and 45 MeV.\(^23\). The predictions using the quark\(^1\) and soliton model\(^5\) weak parameters, respectively, differ by \(\Delta A_L = 4.6 \cdot 10^{-8}\). To be more precise, the various predictions are:

\[
A_L(\text{DDH}) = 5.0 \cdot 10^{-8}, \quad A_L(\text{DZ}) = 2.6 \cdot 10^{-8}, \quad A_L(\text{KM}) = 3.7 \cdot 10^{-9} \quad (13)
\]

The projected long-term accuracy of the upcoming TRIUMF experiment is \((\delta A_L)_{\text{stat}} \simeq \pm 1 \cdot 10^{-8}\),\(^24\) which should be sufficient to discriminate between these two predictions. Notice that a similar experiment is also planned at COSY.\(^30\) This experiment should set rather stringent limits on some combinations of the pv \(\rho N\) and \(\omega M\) couplings. To stress it again, the \(pp\) system is a particularly good example of the interplay of weak and strong interactions and it is therefore mandatory to treat both of them consistently (for further discussion, see Refs. [17,20]). A possible loophole to all of this will be discussed in the last section.
Fig. 2: Parity-violating asymmetry in pp-scattering. The solid line gives the prediction based on the weak couplings as given by the soliton model,\textsuperscript{22} whereas the dashed and dashed-dotted lines are based on the quark model calculations of refs.\textsuperscript{1} and 16, respectively.

PARITY VIOLATION IN $^{18}$F AND DEUTERON PHOTODISINTEGRATION

The nucleus $^{18}$F is what I called a “good system” before. It exhibits “nuclear amplification” in that it has two close-by levels of opposite parity which are separated by only 39 keV (the next level which could mix with these is approximately 2 MeV away) and the dominant E1-transition from the level at 1.081 MeV to the ground state is suppressed, which leads to $|M_1/E2| \approx 112$. The M1-transition is, of course, only possible because of the mixing of the opposite parity-levels. Altogether, this amounts to an amplification of approximately $(2/0.039) \times 112 \approx 6 \times 10^3$ (for further details, see Ref. \textsuperscript{17}). Theoretically, one can calibrate the shell-model calculation to extract the pv circular polarization from the $\beta$-decay of $^{18}$Ne, because the pion-exchange of this $\beta$-decay up to an overall isospin rotation,\textsuperscript{25} and therefore calculation
and measurement of \(^{18}\text{Ne} \ (0^+1) \rightarrow ^{18}\text{F} \ (0^−0)\) \(\beta\)-decay serves as a “gauge” for the accuracy and amounts effectively to a large model-independent limit on the weak pion decay constant. The latter dominates completely this \(\Delta I = 1\) pv observable, and one can deduce a limit on \(\tilde{F}_\pi\), \(\tilde{F}_\pi \leq 3.4 \cdot 10^{-7}\). Here, we have used the experimental circular polarization, \(|P_{\gamma}(^{18}\text{F})| = (0.17 \pm 0.58) \cdot 10^{-3}\). The quark model prediction of Ref. [1], \(\tilde{F}_\pi = 10.8 \cdot 10^{-7}\), is clearly in contradiction to this result.

What does the soliton model give? Of course, \(\tilde{F}_\pi\) is considerably reduced, so we expect a smaller asymmetry. The vector meson contribution is enhanced, and taking nuclear structure calculation from Ref. [17], we predict \(P_{\gamma}(^{18}\text{F}) = 2.2 \cdot 10^{-4}\), not far from the central value of the experiment. We should, however, not put too much emphasis on this closeness of the experimental and theoretical number, but rather state that the strength of the pv \(\pi N\) coupling should still be considered as the main theoretical puzzle. I am sure that \(\tilde{F}_\pi\) should come out smaller than in Ref. [1], but whether it is as small as predicted by the soliton model can only be checked if more theoretical and experimental information on the \(\Delta I = 1\) part of the pv nuclear force are available. One particularly interesting candidate to study in more detail would be the reaction \(\bar{n} + p \rightarrow d + \gamma\) or the inverse process \(\gamma + d \rightarrow n + p\). A calculation of the circular asymmetry as a function of the photon energy has been performed some time ago by Oka.\(^{27}\) I have used this calculation in Ref. 28 to investigate the sensitivity of the circular asymmetry \(A_L\) to the various pv couplings. Considering photon energies below 30 MeV, \(A_L(\omega)\) increases linearly with energy when one uses the quark model couplings of DDH or DZ, with the slope determined by the strength of the pv \(\pi N\) coupling. For the DDH-case the pion contribution is completely dominant for all energies, whereas for the DZ-parameters the reduced \(\pi NN\) strength leads to an overall decrease of \(A_L(\omega)\). For the soliton model, however, things are significantly different. First, between 1 and 20 MeV, \(A_L(\omega)\) shows a flat minimum at about \(\omega_L \approx 12\) MeV and only after \(\omega_L \geq 20\) MeV a gradual rise in \(A_L(\omega)\) sets in. Also, the overall magnitude of the effect is an order of magnitude smaller for the weak parameters predicted by the soliton model. It would be worthwhile to measure the asymmetry say at 10 and 20 MeV incident energy, although the effect is small, the tremendously different slope of \(A_L(\omega)\) should be detectable in a dedicated experiment. Of course, as already mentioned, a more thorough theoretical study has also to be done. First, a more consistent calculation employing e.g. the Bonn-potential and the equivalently constructed weak potential should be performed. Second, the effects of meson-exchange currents, which play an important role in the accurate description of the deuteron properties have to be included. Therefore, these results should only be considered as a guide, but the trends exhibited will certainly
not be wiped out by a more elaborate calculation. A more detailed discussion is given in Ref.28.

In the last section, I will discuss some medium renormalization effects which might come to the rescue of the large value for the pv pion–nucleon coupling as predicted by DDH. However, for the deuteron photodisintegration process just discussed, such a renormalization cannot be operative since the deuteron is essentially an ensemble of two free nucleons.

THE NUCLEON ANAPOLE MOMENT

Apart from the electric dipole moment, there is one other pv coupling of the photon to nucleons (spin-1/2 fields), the so–called anapole moment. It has recently attracted new interest since its contribution might be enhanced considerably in nuclei, similar to the case of $^{18}$F just discussed. For on–shell nucleons, current conservation and Lorentz invariance require that pv corrections to matrix elements of the electromagnetic current take the form

$$<N(p')|J_{\mu, pv}^\text{em}(0)|N(p) >= \frac{a(q^2)}{M_N^2} \bar{u}(p')[\gamma_\nu q^\nu q_\mu - q^2 \gamma_\mu] \gamma^5 u(p)$$

(14)

with $q^2 = (p' - p)^2$. In the Breit frame, where the photon transfers no energy, this matrix element reads

$$<N(\vec{q}/2)|J_{\mu, pv}^\text{em}|N(-\vec{q}/2) >= \frac{E\vec{q}^2}{M_N^3} a(\vec{q}^2) \chi_i \vec{\sigma} T \chi_i$$

(15)

with $\vec{\sigma} T = \vec{\sigma} - \hat{q}\vec{\sigma} \cdot \hat{q}$ the transverse spin operator and $a(q^2)$ the nucleon anapole form factor. The anapole moment has isoscalar and isovector components,

$$a(0) = a_S(0) + a_V(0) \tau_3$$

(16)

In the soliton model, one can easily calculate the anaple moment and form factor. For that, one identifies the matrix element in (15) with the Fourier transform of the pv electromagnetic current. For the usual hedgehog ans"atze, its has the general form

$$\vec{J}_{\rho \nu}^\text{pv}(\vec{r}) = \Gamma_1(\vec{r}) \vec{\sigma} + \Gamma_2(\vec{r}) \hat{r} \vec{\sigma} \cdot \hat{r}$$

(17)

with $\Gamma_{1,2}(r)$ functions of the various meson profiles whose explicit form we do not need here. However, one immediately encounters a difficulty. Current conservation demands $\Gamma_1'(r) + \Gamma_2'(r) + 2\Gamma_2(r)/r = 0$, where the prime denotes differentiation.
with respect to $r$. This condition is not met. Interestingly, if one switches off the $\rho$–meson fields and considers the so–called $\omega$–stabilized Skyrmion\textsuperscript{37}, then the divergence condition is fulfilled. This peculiar behaviour might be traced back to the fact that in the isoscalar channel one has exact vector meson dominance (VMD) but not in the isovector one (compare the discussion of Hwang and Nigoyi\textsuperscript{38} of VMD and gauge invariance for pv photon–nucleon couplings). To get an idea of the size of the anapole moment, let me crudely restore gauge invariance by subtracting the pieces which violate current conservation. In that case, the ”minimal” model gives $a_s(0) = 4 \cdot 10^{-8}$, and the extension of the pv $\gamma N$ vertex is given by a mean square radius of about 0.4 fm corresponding to a monopole form factor with a cut–off of $\Lambda = 1.23$ GeV. At present, I can not offer a solution to the problem concerning the violation of current conservation, but I suspect that is it related to the rather crude quantisation procedure used (which is known to do harm to e.g. the chiral algebra of the charges\textsuperscript{38}).

RECENT DEVELOPMENTS

There are some recent developments (partly outside the soliton model) which indicate some interesting new effects and might lead to a reconsideration of some topics discussed so far. The first one is due to a calculation of Dai, Savage, Liu and Springer.\textsuperscript{9} They calculate an effective Hamiltonian for $\Delta I = 1$ nuclear parity violation, including the effects of the heavy quarks $s, c$ and $b$. At the scale of the W-boson mass, the pv $\Delta I = 1$ Hamiltonian is, of course, well known and given in terms of eight four–quark operators with known Wilson coefficients. Integrating out the $b$ and the $c$ quark successively, one has a tower of effective theories. For each of these, the anomalous dimension matrix is calculated to one loop in the QCD corrections and the effective field theories are matched. By this procedure, one can finally go down to the hadronic scale of $\Lambda_\chi = 1$ GeV and compare the Wilson coefficients $C_i(\Lambda_\chi)$ with the original ones, $C_i(M_W)$. The important observation made in ref.9 is that the operators involving strange quarks are substantially larger than the ones involving only the up and down quarks, approximately

$$\frac{C_i^{\text{strange}}(\Lambda)}{C_i^{\text{non–strange}}(\Lambda)} \sim \frac{1}{\sin^2 \Theta_W} \sim 5$$

(18)

The authors of ref.9 did not compute hadronic matrix elements at the scale $\Lambda_\chi$, but resorted to the meson–exchange picture and large $N_c$ arguments. In that case ($N_c \rightarrow \infty$), factorization can be justified and one finds for the $\rho^0N$ pv matrix element

$$<\rho^0N|H_{\rho^0\Delta I = 1}|N> = \frac{1}{3} G_F \sin^2 \Theta_W f_{\rho^*\rho} \left\{ -0.95 <N|\bar{u}\gamma_\mu \gamma_5 u + \bar{d}\gamma_\mu \gamma_5 d|N> + 13.4 <N|\bar{s}\gamma_\mu \gamma_5 s|N> \right\}$$

(19)
Consider first the case were the strange matrix element vanishes (like in DDH). In that case, one finds $h_{\rho}^{(1)} = -1.9 \cdot 10^{-8}$, quite consistent with the DDH value. However, the large relative factor in front of the new, un–colored strange contribution can easily alter this result by an order of magnitude. Combining the EMC–data and hyperon decay rates, one has $< N|\bar{u}\gamma_{\mu}\gamma_{5}u + \bar{d}\gamma_{\mu}\gamma_{5}d|N > \simeq - < N|\bar{s}\gamma_{\mu}\gamma_{5}s|N > \simeq -(0.2 \pm 0.1) S_{\mu}$ with $S_{\mu}$ the nucleon spin vector. In this case, $h_{\rho}^{(1)} = -2.9 \cdot 10^{-7}$, which is an enhancement of a factor 15. If that were true, all previous estimates of pv meson–nucleon couplings can be off the mark by large factors. However, we should not forget that in last years many of the matrix elements which indicated a large contribution of the strange quark sea to various nucleon properties have been tamed, the prime example being the famous $\pi N \Sigma$ term.

In a similar fashion, Kaplan and Savage\textsuperscript{40} have recently reanalyzed the pv–nucleon couplings making use of baryon chiral perturbation theory. They have derived the most general pv and CP–conserving effective pion–nucleon–photon Lagrangian to first order in derivatives and first order in the photon field and to all orders in the pion field. This effective Lagrangian is parametrized by a few coupling constants, which are labelled $h_{V}^{0,1,2}, h_{A}^{1,2}$ and $h_{\pi NN}^{1} = G_{\pi}$. Apart from the standard pv pion–nucleon coupling (discussed before), the authors of ref.40 mainly concentrate on the novel pv $\gamma \pi NN$ and the pv $\pi \pi NN$ vertices (the latter one has been already been considered by nuclear theorists in the seventies). Three different methods are used in ref.40 to estimate the strength of these coupling constants, namely factorization, dimensional analysis and relations to $\Delta S = 1$ hyperon decay matrix elements. From these methods, the dimensional analysis is considered most reliable. The most interesting results of this are 1) a large contribution of the strange quarks to $G_{\pi}$ (together with a large value for this coupling), 2) a sizeable strangeness enhancement for the pv $\pi \pi NN$ coupling $h_{A}^{1}$ and 3) a large value for the strength of the pv $\gamma \pi NN$ coupling. Taking these estimates face value, drastic consequences would arise. First, in the case of the $^{18}$F experiment, interference between the one–pion exchange (considered so far) and the novel $\gamma \pi NN$ vertex might complicate the analysis of the data and ultimately relax the bound on $G_{\pi}$. Similarly, for the planned TRIUMF and COSY experiments measuring parity violation in pp scattering at 230 MeV, one would have to consider two–pion exchange, not only the conventional one arising from e.g. intermediate $\Delta$ resonances, but also the one due to the large pv $\pi \pi NN$ coupling. However, before jumping too far, one should not forget that the results of ref.40 should be considered indicative – more elaborate calculations of the hadronic matrix elements are necessary (using e.g. lattice methods) and also more complex nuclear structure calculations involving these novel couplings have to be
performed before one can draw a final conclusion. For more details on these topics, please consult ref.40.

MEDIUM RENORMALIZATION OF \( G_\pi \)?

There exist ample evidence that suggests scale changes of fundamenral properties of nucleons in nuclei. Some pertinent examples are the first EMC effect, the quenching of the axial–vector coupling constant \( g_A \) in nuclear \( \beta \)-decay or the behaviour of the longitudinal and transverse strength functions in quasi–elastic electron scattering off nuclei. These effects are there and they are important, but their origin still remains to be explained in a consistent treatment of many–body effects and fundamental scale changes of the nucleon properties. The chiral soliton model allows to systematically investigate the constraints from chiral symmetry on such possible medium modifications.\(^{41,42}\) The basic idea is the following: In the soliton model, baryon properties are fixed once the mesonic input is determined. We know, however, that meson masses and coupling constants change in the baryon–rich environment.\(^{43}\) This immediately leads to density or temperature–dependent nucleon properties.\(^{44}\) For the meson sector, I will use here results from the Nambu–Jona-Lasinio model which have been obtained in collaboration with Véronique Bernard.\(^{42}\) For not too large densities \( \rho \), one finds for the pion decay constant and the vector and scalar meson masses (all other quantities are essentially unaffected)

\[
F^*_\pi = F_\pi(0)[1 - R_\pi \frac{\rho}{\rho_0}], \quad m^*_V = m_V(0)[1 - R_V \frac{\rho}{\rho_0}], \quad m^*_\sigma = m_\sigma(0)[1 - R_\sigma \frac{\rho}{\rho_0}] \quad (20)
\]

where the '\( * \)' denotes quantities in the medium and \( \rho_0 \) is the nuclear matter density. The range of values for \( R_{\pi,V,\sigma} \) is discussed in ref.42. For simplicity, let me take an universal and equal value, \( R_\pi = R_V = R_\sigma = R \). This is not a direct consequence of the NJL model but compatible with it. For the sake of the argument I will make here, this simplification is justified. In ref.42, which is a widely overlooked paper, I have shown that most of the pv meson–nucleon couplings are very sensitive to such medium effects, quite in contrast to their strong counterparts. In particular, the most important pion–nucleon couplings show the following medium renormalization (for \( R = 0.2 \) and at nuclear matter density)

\[
\frac{G^*_\pi(\rho_0)}{G_\pi(0)} = 0.65, \quad \frac{g^*_\pi NN(\rho_0)}{g_{\pi NN}(0)} = 0.99 \quad (21)
\]

and similar results for the vector meson couplings. One can understand this very different behaviour if one takes a closer look at the expressions for the various coupling
constants. Using the dimensionless variable \( x = gF_\pi r \), with \( g \) the universal vector-meson–pion coupling, one notices that the weak couplings depend one much higher powers of \( F_\pi \) than their strong counterparts and thus are more sensitive to medium modifications. A more detailed account of this can be found in ref.42.

Finally, let me point out some recent work by Grach and Shmatikov\textsuperscript{45} which concerns yet another mechanism to bring down the value of the pv \( \pi N \) coupling in the medium. The basic idea of their work is that the rescattering of emitted pions leads to a strong suppression of \( G_\pi \) (the basic Feynman diagrams are shown in fig.3). Using monopole form factors with a cut off \( \Lambda \simeq 7M_\pi \simeq 1 \text{ GeV} \) to regulate the diverging loop integrals, they find

\[
G_\pi^{(r)} = G_\pi (1 + \frac{g_{\pi NN}^2}{8\pi^2} I_1 + \frac{g_{\pi \Delta N}^2}{8\pi^2} \frac{14}{27} I_2) \\
= G_\pi (1 - 0.76 + 0.01)
\]

which leads to

\[
\tilde{F}_\pi^{(r)} = G_\pi^{(r)} g_{\pi NN} / \sqrt{32} \simeq 2.9 \cdot 10^{-7}
\]

which is below the bound from the\textsuperscript{18}F experiment. This is an interesting suggestion, but it definitively needs a better treatment (better regularization procedure) and should also be applied to the other pv meson–nucleon couplings. Also, one should understand the relation to the soliton model results discussed before.

\textit{Fig. 3: Strong pion rescattering in the medium.\textsuperscript{45} The solid, double and dashed lines denote nucleons, \( \Delta \)'s and pions, in order. Strong meson–nucleon vertices are depicted by open circles and weak vertices by the crossed circles.}
OPEN PROBLEMS

Instead of rephrasing what I have said so far, let me just mention the two salient problems which have to be addressed in the framework of the chiral soliton model to allow for a deeper understanding of the $pv$ meson–nucleon interaction regions.

- Realistic versions of the three flavor Skyrme model are now available. They do not indicate a large strange component in proton wave function. It would be worthwhile and necessary to extend the analysis of the $pv$ interaction region discussed here. This would also allow to address such questions like the strength of the $\phi$–couplings and the relation to the $\Delta S = 1$ hyperon decay matrix elements. Ultimately, such calculations will shed some light on the recent developments concerning the possible enhancement of various weak couplings due to the strange color–singlet operators.

- An old problem is whether the soliton model calculations should be supplemented by strong interaction enhancement factors or whether these are already contained in the non–perturbative soliton currents. This question was to some extent addressed in ref.7, where it was argued that the inclusion/omission of these factors would at most lead to uncertainties of the order of 30 per cent, i.e. lead to corrections within the accuracy of the model. To my opinion, this question is not yet settled. Its resolution will also bring about the answer to the question of including operators which are not of the canonical $V_\mu A^\mu$–type.

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