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Studying Nuclear Responses to Electroweak Probes in Effective Field Theory

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Abstract.
Effective field theory is expected to provide a powerful framework for describing low-energy nuclear phenomena in a model-independent manner. I shall give here a brief account of some of the latest developments in this domain of nuclear physics. I wish to emphasize the importance of what may be called a *symbiotic* relation between this new approach and the standard nuclear physics approach. This feature is particularly important in providing reliable nuclear physics inputs to the latest Sudbury Neutrino Observatory experiments that have given strong evidence for neutrino oscillations.

It is well known that the phenomenological potential picture has been extremely successful in describing a great variety of nuclear properties. In this picture, the responses of nuclei to external electroweak probes are given by one-body impulse-approximation terms and exchange-current terms acting on nuclear wave functions obtained by solving the Shroedinger equation which is governed by a potential; furthermore, the exchange currents (usually taken to be two-body operators) are derived from one-boson exchange diagrams. The vertices appearing in the relevant diagrams are determined to satisfy the low-energy theorems and current algebra [1]. We shall refer to formalisms based on this picture as the *standard nuclear physics approach* (SNPA). SNPA has been used extensively to describe nuclear electroweak processes, and the well-documented agreement between theory and experiment [2] gives a strong indication that SNPA captures essential physics. At the more fundamental level, however, it is a great challenge in contemporary nuclear physics to establish a link between SNPA and the fundamental quantum chromodynamics (QCD). Effective field theory (EFT) is expected to provide a natural and insightful framework for this purpose. In my talk I shall first give a brief account of EFT as applied to nuclear physics. I then proceed to discuss interplay between EFT and SNPA; my main goal is to illustrate what may be called a "symbiotic" relation that exists between EFT and SNPA. As an example, I consider the neutrino-deuteron reactions, which are of great current importance in connection with the celebrated SNO experiments that have provided clear evidence of neutrino oscillations.

The main idea of EFT is in fact very simple. In describing phenomena that belong to a typical energy-momentum scale Q, it is natural to expect that we can exclude from our Lagrangian those degrees of freedom that belong to energy-momentum scales much higher than Q. Motivated by this expectation we introduce a cut-off scale Λ that is sufficiently large compared with Q and classify our fields (to be generically represented by Φ) into two groups: high-energy fields $\Phi_{_H}$ and low-energy fields $\Phi_{_L}$. After *integrating out* $\Phi_{_H}$ we arrive at an *effective* Lagrangian which involves only $\Phi_{_L}$ as explicit dynamical variables. Thus the effective Lagrangian $\mathcal{L}_{\rm eff}$ is related to the original Lagrangian $\mathcal{L}$ as

$$\int [d\Phi] \mathrm{e}^{\mathrm{i}\int d^4x \mathcal{L}(\Phi)} = \int [d\Phi_{_L}] \mathrm{e}^{\mathrm{i}\int d^4x \mathcal{L}_{\rm eff}(\Phi_{_L})} \tag{1}$$

$\mathcal{L}_{\rm eff}$ defined this way is guaranteed to inherit all the symmetries (and the patterns of symmetry breaking, if there are any) of the original Lagrangian $\mathcal{L}$. According to this prescription, $\mathcal{L}_{\rm eff}$ should be the sum of all possible monomials of $\Phi_{_L}$ and their derivatives that are consistent with the symmetry requirements imposed by $\mathcal{L}$. Since a term that involves n derivatives scales like $(Q/\Lambda)^n$, the terms in $\mathcal{L}_{\rm eff}$ can be organized into a perturbative series with Q/Λ serving as an expansion parameter. The coefficients of the monomials appearing in this expansion scheme are called the low-energy constants (LEC). Provided all the LEC's up to a specified order n are known either from theory or from fitting to the experimental values of the relevant observables, we can use $\mathcal{L}_{\rm eff}$ as a complete (and hence model-independent) Lagrangian to a given order of expansion.

Having sketched the general idea of EFT, we move to a more specific discussion of EFT as applied to nuclear physics. The original Lagrangian $\mathcal{L}$ in this case is the QCD Lagrangian but we note that, for the energy-momentum scale $Q \ll \Lambda_\chi \sim 1$ GeV, which is a typical regime of interest in nuclear physics, the effective degrees of freedom that should feature in $\mathcal{L}_{\rm eff}$ are hadrons rather than the quarks and gluons.

Moreover, if we are looking at low-energy phenomena whose typical energy-momentum scale is $Q \lesssim m_\pi$, it should be a good approximation to retain only the pions and nucleons as explicit degrees of freedom that appear in $\mathcal{L}_{\rm eff}$. It is important to ensure that chiral symmetry of QCD (and the pattern of its small violation due to the finite quark masses) be inherited by our $\mathcal{L}_{\rm eff}$. An EFT obtained in this picture is called chiral perturbation theory (χPT) [4, 6]. It has been used in the meson sector with great success.

A hurdle we encounter in incorporating the nucleons in χPT is that, since the nucleon mass $m_{\rm N}$ is comparable to the cut-off scale Λ_χ, a straightforward application of expansion in Q/Λ does not work. A solution to this problem is provided by heavy-baryon chiral perturbation theory (HBχPT), which essentially consists in moving the reference point of the nucleon energy from 0 to $m_{\rm N}$ and in integrating out the small component of the nucleon field along with the anti-nucleonic degrees of freedom. HBχPT therefore involves as effective degrees of freedom the pions and the large components of the redefined nucleon field (denoted by B).

The effective Lagrangian $\mathcal{L}_{\rm ch}^{\rm HB}$ in HBχPT has as its expansion parameters Q/Λ_χ, m_π/Λ_χ and $Q/m_{\rm N}$. Since $m_{\rm N} \approx \Lambda_\chi$, it is customary to combine chiral and heavy-baryon expansions and introduce the chiral index $\bar{\nu}$ defined by $\bar{\nu} = d + (n/2) - 2$. Here n is the number of fermion lines belonging to a given vertex, while d is the number of derivatives (with m_π counted as one derivative). The leading-order ($\bar{\nu} = 0$) term in this expansion looks like [5]

$$\mathcal{L}^{(0)} = \frac{f_\pi^2}{4}{\rm Tr}[\partial_\mu U^\dagger \partial^\mu U + m_\pi^2(U^\dagger + U - 2)] \\ + \bar{B}(iv \cdot D + g_A S \cdot u)B - \frac{1}{2}\sum_A C_A(\bar{B}\Gamma_A B)^2$$

Here $U(x) \equiv \exp(2i\sum \tau_a \pi_a/f_\pi)$ is a 2×2 matrix field which involves the pion fields π_a ($a = 1, 2, 3$), and which gives a non-linear representation of chiral symmetry; $\xi \equiv \sqrt{U}$; $u_\mu \equiv i(\xi^\dagger \partial_\mu \xi - \xi \partial_\mu \xi^\dagger)$, $S_\mu = i\gamma_5 \sigma_{\mu\nu} v^\nu/2$, and D_μ is the covariant derivative acting on the nucleon.

We can introduce a similar power counting scheme for Feynman diagrams as well. Here is a brief explanation of Weinberg's counting scheme [6], which we follow here. A Feynman diagram that contains N_A nucleons, N_E external fields, L loops and N_C disjoint parts can be shown to be proportional to $\mathcal{O}(Q^\nu)$, where the chiral order ν is defined by

$$\nu = 2L + 2(N_C - 1) + 2 - (N_A + N_E) + \sum_i \bar{\nu}_i \,, \tag{3}$$

with the summation running over all the vertices in the Feynman diagram. In fact, although this counting scheme works well for one-nucleon cases, it cannot be immediately applied to systems with more than one nucleon for the following reason. As is well known, nuclei involve very low-lying excited states, and the existence of this small energy scale spoils the original counting rule [6]. The situation can be compared to a well-known problem in ordinary quantum mechanics that perturbation expansion breaks down in the presence of an intermediate state that gives a very small energy denominators. A method proposed by Weinberg to avoid this difficulty is as follows. Classify Feynman diagrams into two categories. Diagrams in which every intermediate state has at least one meson in flight are classified as irreducible diagrams, and all others are called reducible diagrams. Apply the above-mentioned chiral counting rules only to irreducible diagrams. Treat the total contribution of all the irreducible diagrams (up to a specified chiral order) as an effective operator acting on nuclear wave functions. By summing up an infinite series of irreducible diagrams (solving either the Schrödinger equation or the Lippman-Schwinger equation), we can include the contributions of reducible diagrams [6]. For convenience we refer to this two-step procedure as *nuclear* χPT.

I have briefly summarized above a *nuclear* χPT in the Weinberg scheme. It should be mentioned that an alternative version of nuclear EFT, called the power divergence subtraction (PDS) scheme, has been developed with much success. The PDS scheme, first studied in seminal papers by Kaplan, Savage and Wise [7], employs a counting scheme different from the above-mentioned Weinberg counting scheme. An advantage of the PDS scheme is that it allows us to maintain formal chiral invariance, whereas the Weinberg scheme loses manifest chiral invariance. We note, however, that in many practical applications this formal problem is not worrisome up to the chiral order of our interest, *i.e.*, the chiral order up to which we evaluate our irreducible diagrams. For a recent discussion of the relation between PDS and the Weinberg scheme, I refer you to Ref. [8].

Nuclear χPT is one of the most active research frontiers in contemporary nuclear physics. Many important results have been reported in both the Weinberg scheme and PDS; for recent reviews, see *e.g.* Ref. [8–10]. I choose here to concentrate on our own work based on the Weinberg scheme, but this of course does not mean I am minimizing the exciting achievements made in PDS; for a comprehensive survey of PDS, see Beane *et al.* [10].

In applying nuclear χPT to a process that involves (an) external current(s), we derive a nuclear transition operator $\mathcal{T}$ by evaluating the complete set of all the irreducible diagrams (up to a given chiral order ν) with the relevant external current(s) duly attached. In order to preserve consistency in chiral counting, the nuclear matrix element of $\mathcal{T}$ must be calculated with the use of nuclear wave functions which are governed by the nuclear interactions that subsume all irreducible A-nucleon diagrams up to ν-th order. If this program is faithfully carried out, we would have an *ab initio* calculation. At present, however, it is not always easy to generate needed nuclear wave functions strictly within the EFT framework and therefore we often reply on wave functions obtained with the use of phenomenological nuclear interactions. For convenience, we refer to this hybrid approach as EFT*. This eclectic method has been applied to many cases with great success. One of the best known examples is an EFT* calculation by Park *et al.* [11] of the isovector M1 transition amplitude that governs the $n(\mathrm{thermal}) + p \rightarrow d + \gamma$ reaction.

Although EFT* is known to have been successful for many observables, it is desirable to provide its formal justification. Fortunately, at least for the two-nucleon systems, its partial justification has been given by Park, Min, Rho and myself (PKMR) [12]. PKMR carried out an *ab initio* calculation up to next-to-leading-order (NLO) of various observables in the A=2 systems. Here is a brief summary of Ref. [12]. Since no loop diagrams appear up to NLO, we can work with a potential, and its generic form can be parameterized as

$$\mathcal{V}(\mathbf{q}) = -\tau_1 \cdot \tau_2 \frac{g_A^2}{4f_\pi^2} \frac{\sigma_1 \cdot \mathbf{q}\, \sigma_2 \cdot \mathbf{q}}{\mathbf{q}^2 + m_\pi^2} + \frac{4\pi}{m_{_N}} \left[C_0 + (C_2 \delta^{ij} + D_2 \sigma^{ij}) q^i q^j \right] \quad (4)$$

where $\sigma^{ij} = 3/\sqrt{8}[(\sigma_1^i \sigma_2^j + \sigma_1^j \sigma_2^i)/2 - (\delta^{ij}/3)\sigma_1 \cdot \sigma_2]$, with $\mathbf{q}$ the momentum transfer between the two nucleons. The first term arises from Goldstone-boson (pion) exchange, and hence its structure is uniquely dictated by χPT. The contact terms in the square brackets, whose structure is also dictated by HBχPT, encode the effects of the high-energy degrees of freedom that have been "integrated out" in our EFT. We remark that, if we push the idea of EFT to an extreme, we could even work with a very low energy-momentum regime in which the pion is treated as a *heavy* particle to be "integrated away". Needless to say, if the pion is eliminated as a massive particle, we cannot any longer talk about chiral symmetry but it is legitimate and informative to consider a "pionless" (or "nucleon-only") EFT for describing very low-energy nuclear phenomena. In [12] both "pionful" and "pionless" cases were studied.

When we solve the Lippman-Schwinger equation corresponding to the potential in Eq.(4) to obtain a *fully distorted* two-nucleon wave functions, there will appear an infinite series of divergences because of the δ-function-type potential. To regularize this singular behavior, we may introduce a momentum-cutoff parameter Λ. For a specified value of Λ, we can determine the LEC's (C's and D) in Eq.(4) by relating them (after renormalization) to the low-energy observables in the two-nucleon systems: the scattering length and the effective range for each of the scattering channels, and a selected set of the deuteron observables. Once the LEC's are fixed, we can make predictions for the N-N scattering phase shifts and the deuteron properties (other than those treated as input). It is also possible to make the *parameter-free* calculation for electroweak observables in the two-nucleon systems. An *ab initio* calculation in [12] has shown that (a) all the calculated quantities show good agreement with the corresponding experimental values; (b) the numerical results of the calculation are very stable against changes in the cut-off parameter Λ, so long as Λ stays within a physically reasonable range. This range has been found to be Λ =100 - 300 MeV for the "pionless" case ("nucleon-only EFT"), and Λ =200 - 500 MeV for the "pion-ful" case; these results are in conformity with what we generally expect in the EFT picture. (c) Since the *ab initio* calculation essentially reproduces the results of EFT*, the use of the latter is justified.

Yet, a legitimate question you could raise here is: "Did the use of EFT in the above example really give us anything new beyond SNPA ?" In particular, can we see any real difference between EFT (as used here) and the familiar effective-range expansion formula ? As a matter of fact, in the "pionless' treatment, the two LEC's (C_0 and C_2) in Eq.(4) play much the same role as the scattering length

a and the effective range r_e. Such a trivial correspondence, however. disappears when the pion is included in $\mathcal{L}_{\rm EFT}$, and the χPT calculation in [12] contains more physics than the well-known effective-range expansion method. We should also mention that χPT offers a unified and systematic expansion scheme for both N-N scattering and electroweak transition processes.

We have discussed above a case in which observables of our interest are governed by the lowest chiral-order contributions. The usefulness of χPT, however, is not confined to such cases. To give an example for which higher order terms play a crucial role, let us consider the spin-dependent observables in the $\vec{n} + \vec{p} \rightarrow d + \gamma$ reaction for the thermal neutron. The photon circular polarization P_γ and the photon anisotropy in this case are known to be sensitive to the small isoscalar M1 and E2 transition matrix elements, and these matrix elements arise from higher chiral order terms. A χPT calculation for these spin observables has been done in [13] and in [14]. Park *et al.* [13] adopted the Weinberg scheme, while Chen *et al.* [14] used PDS, and these two calculations give values of P_γ that are compatible with each other and that agree with the existing experimental value. One might in general expect that higher chiral-order calculations would involve too many LEC's and can quickly get out of control. The above example, however, shows that in some favorable cases the physical amplitudes depend on a sufficiently limited number of combinations of LEC's, allowing χPT to maintain its predictive power.

My next topic is what I call a symbiotic relation between EFT and SNPA; I discuss as an example the weak-interaction processes in the two-nucleon systems. As far as SNPA is concerned, the latest calculations can be found in Ref. [15] for μ-d capture, in Refs. [15–17] for ν-d reactions, and in Refs. [12, 18, 19] for pp-fusion. For the μ-d capture and ν-d reactions, experimental data are available although with somewhat large ($\sim$10 %) experimental errors, and the SNPA calculations reproduce these data very well. How about EFT calculations for the weak-interaction processes in the two-nucleon systems ? I compare here the results of the SNPA and EFT calculations for the ν-d reactions. These particular processes are chosen because they are of great current importance in connection with the highly consequential solar neutrino experiments at the Sudbury Neutrino Observatory (SNO) [20, 21].

At SNO a 1-kiloton heavy water Cerenkov counter is used to detect, among other things, the solar neutrinos. SNO can monitor the following reactions: $\nu_e + d \rightarrow e^- + p + p$, $\nu_x + d \rightarrow \nu_x + p + n$, $\bar{\nu}_e + d \rightarrow e^+ + n + n$, $\bar{\nu}_x + d \rightarrow \bar{\nu}_x + p + n$, and $\nu_x + e^- \rightarrow \nu_x + e^-$. Here x stands for a neutrino of any flavor (e, μ or τ). An important aspect of SNO is that it can detect the charged-current (CC) and neutral-current (NC) reactions simultaneously but separately. With this unique feature, SNO is an ideal facility for studying solar neutrino oscillations. The recent SNO experiments [20, 21] have established that the total solar neutrino flux (summed over all flavors) agrees with the prediction of the standard solar model [22], whereas the electron neutrino flux from the sun is significantly smaller than the total solar neutrino flux. The amount of deficit in the electron neutrino flux is consistent with what used to be known as the solar neutrino problem. These results of the SNO experiments have given “smoking-gun” evidence for the transmutation of electron neutrinos into neutrinos of other flavors.

It is obvious that a precise knowledge of the ν-d reaction cross sections is of primary importance in interpreting the existing and future SNO data. A detailed calculation of the ν-d cross sections based on SNPA was carried out by Nakamura, Sato, Gudkov and myself (NSGK) [16], and this calculation has recently been updated by Nakamura *et al.* (NETAL) [17]. As pointed out in Ref. [18], the exchange currents in SNPA are known to be dominated by the Δ-particle excitation diagram [18], and the reliability of estimation of this diagram hinges on the precision with which the coupling constant $g_{\pi N\Delta}$ is known. NSGK studied two representative values for $g_{\pi N\Delta}$, one consistent with the $np \to \gamma d$ data and the second value arrived at by fitting the experimental value of Γ_β^t, the tritium β-decay rate. The calculated ν-d cross sections were found to differ by ∼3 % between these two cases, and this 3 % difference was taken as a measure of theoretical uncertainty. Meanwhile, Butler, Cheng and Kong (BCK) [23] were the first to carry out an EFT calculation of the ν-d cross sections. The results obtained in the PDS scheme agree with those of SNPA in the following sense. We first note that, in EFT calculations, there may occur some LEC's that cannot be fixed from the symmetry requirements alone and hence need to be determined from the experimental data. In the case of BCK's calculation [23], the coefficient (denoted by $L_{1\mathrm{A}}$) of a four-nucleon axial-current coupling term appears as such an unknown parameter, even though a *naturalness* argument (based on a dimensional analysis) gives an order-of-magnitude estimate, $|L_{1\mathrm{A}}| \leq 6\,\mathrm{fm}^3$. BCK therefore determined $L_{1\mathrm{A}}$ by imposing the requirement that the cross sections of NSGK be reproduced by their EFT calculation. With the use of the value of $L_{1\mathrm{A}}$ fine-tuned this way, a perfect agreement was seen between the cross sections obtained in Ref. [23] and those of NSGK for all the four reactions (CC and NC channels for ν and $\bar{\nu}$) for the entire energy range of the solar neutrinos, $E_\nu \leq 20$ MeV. Moreover, the best fit value, $L_{1\mathrm{A}} = 5.6\,\mathrm{fm}^3$, obtained by BCK [23] is consistent with the value expected from the naturalness argument. The fact that an EFT calculation (with one parameter fine-tuned) reproduces the results of SNPA very well strongly suggests that SNPA captures the basic physics involved.

Since EFT is a general framework, it is capable of giving model-independent results, *provided* all the LEC's needed to specify $\mathcal{L}_{eff}$ up to a given order are known *a priori*. The PDS scheme used by BCK, however, does contain an unknown LEC, $L_{1\mathrm{A}}$. Meanwhile, although SNPA involves a certain degree of model dependence, its basic assumptions and the parameters contained in it have been tested extensively for many observables. Insofar as these tests are accepted as valid, SNPA has predictive power. We thus recognize that EFT and SNPA are playing complementary roles here.*

Although it is highly significant that the SNPA and EFT calculations of the ν-d cross sections agree with each other (in the sense explained above), it is desirable to perform an EFT calculation that contains no adjustable LEC. Is this type of calculation possible ? The answer is fortunately yes, if we invoke the EFT* approach mentioned earlier. Let me describe here the latest developments in this regard.

An important conceptual difference between the nuclear χPT à la Weinberg and the PDS used by Butler *et al.* [23] is that the former uses the notion of wave functions whereas the latter does not. This point may require a little more explanation. If we, following Weinberg, sum all the A-body ladder diagrams, each composed of Weinberg-irreducible vertices connected by A-nucleon propagators, we arrive at (in the language of scattering theory) the wave operator Ω, which changes the free A-nucleon wave function, $|\Psi_A^0\rangle$, into the *fully distorted* wave functions, $|\Psi_A\rangle$. This feature is analogous to what is done in SNPA. The only difference is that the kernel of the Lippman-Schwinger equation, which generates Ω, is constructed here in accordance with χPT whereas in SNPA the kernel corresponds to a phenomenological nuclear potential.

We now note that, as far as one-body operators [or impulse approximation (IA) terms] are concerned, the nuclear matrix elements calculated with the use of EFT-generated wave functions are expected to be very close to those calculated with the SNPA wave functions. Thus EFT and EFT* should give practically the same IA matrix elements. Next, we note that the ratio of the exchange current contributions to those of the one-body operators should be much less sensitive to the details of the nuclear wave functions than the absolute values of the nuclear matrix elements. Through these considerations EFT* acquires a higher status than an expedient, for we now know that it is indeed a natural thing to do to rely on χPT for deriving transition operators and evaluate their matrix elements using the realistic wave functions obtained in SNPA. EFT*, justified this way, provides the possibility of applying χPT beyond the two-nucleon systems. We note in this connection that, in A-nucleon systems (A$\geq$3), the contributions of transition operators involving three or more nucleons are intrinsically suppressed by chiral counting. This means that, up to a certain chiral order, the transition matrix element for an A-nucleon system can be evaluated using the same χPT-based 1-body and 2-body operators as used for the two-nucleon systems. Therefore, insofar as high-quality wave functions are obtainable from SNPA, one can calculate

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*SNPA is the only theoretical framework available at present for calculating the ν-d cross sections beyond the solar energy region. An EFT calculation in Ref. [23] adopts a so-called “nucleon-only” EFT, excluding all the degrees of freedom other than that of the nucleon, and therefore the validity of this approach is limited to very low incident neutrino energies (typically the solar neutrino energies). On the other hand, we encounter no immediate conceptual problems in applying SNPA to an energy region beyond the solar neutrino energy.

a transition matrix element for the A-nucleon system with precision comparable to that for the corresponding two-nucleon case. We will discuss below an important consequence of this observation

As for PDS, we emphasize that PDS, based on an expansion scheme for transition amplitudes themselves, does not employ the notion of wave functions. This feature may or may not be an advantage. A disadvantage in the present context is that we cannot readily relate the transition matrix elements in an A-nucleon and (A+1)-nucleon systems. Each nuclear system requires a separate parameterization in PDS. A consequence of this feature is that L_{1A} remained undetermined in the work of BCK [23], for no experimental data is available to fix L_{1A} within the two-nucleon systems.

Recently, Park *et al.* [24, 25] proposed to use EFT* to get a constraint on an LEC in $\mathcal{L}_{eff}$ from the experimental information on an A=3 system. Here is a brief description of that investigation and subsequent related studies. To the lowest meaningful order in Weinberg's counting, the two-body currents for the Gamow-Teller transition involve one unknown LEC, which was denoted by $\hat{d}_R$ in [12]. Like L_{1A} featuring in the PDS scheme [23], $\hat{d}_R$ controls a strength of four-nucleon contact-term coupling to the axial current. Unfortunately, the existing experimental data involving the A=2 systems do not have sufficient precision to control the value of $\hat{d}_R$. It turns out, however, that the same parameter, $\hat{d}_R$, features also in the tritium β-decay rate Γ^t_β. The facts that the experimental value of Γ^t_β is known with high precision, and that the accurate wave functions of ^{3}H and ^{3}He are available from a well-developed variational calculation [26], enable us to determine $\hat{d}_R$ with sufficient accuracy for our purposes. After fixing the value of $\hat{d}_R$ this way, Park *et al.* [24] carried out the first parameter-free calculation of the solar pp fusion rate. Their formalism based on a momentum cutoff scheme contains a momentum cutoff parameter Λ. As a measure of the consistency of their method, Park *et al.* demonstrated that the results are stable against changes in Λ. The estimated uncertainty in the calculated pp fusion rate is $\sim$0.5%, which improves by a factor of $\sim$10 the precision of the previous calculations.*

Let us now return to our discussion of the ν-d reactions. The EFT* approach of Refs. [24,25] can be readily applied to these reactions, and thus we are in a position to carry out an EFT-motivated parameter-free calculation of the ν-d cross sections. Such a calculation has recently been carried out by Ando *et al.* [27]. The ν-d cross sections obtained in this calculation are found to agree within 1% with what Nakamura *et al.* [17] obtained in SNPA adopting the value of $g_{\pi N\Delta}$ adjusted to reproduce Γ^t_β. These latest developments show that the ν-d cross sections used in interpreting the SNO experiments [20,20] are reliable at the 1% precision level,** and hence the evidence for neutrino oscillations reported in those

We remark en passant that the same EFT method was successfully applied also to the calculation of the cross section for the hep reaction: $^3\mathrm{He} + p \rightarrow\ ^4\mathrm{He} + \nu_e + e^+$ [25].

**For radiative corrections to be applied here, see Ref. [28].

experiments is robust against nuclear physics ambiguities.

Even though the determination of $\hat{d}_R$ from Γ^t_β is good enough for all practical purposes, it seems worthwhile to examine a possibility (including a futuristic one) to fix $\hat{d}_R$ without using information from three-nucleon systems. A rather obvious candidate is $\mu^- + d \to \nu_\mu + n + n$. It is true that the rather large energy-momentum transfer involved in the disappearance of a μ^- seems to make the applicability of EFT a delicate issue, but we can show that, as far as the hadron sector is concerned, μ-d capture is in fact a "gentle" process; this is because (1) ν_μ carries away most of the energy, and (2) there is a large enhancement of the transition amplitude in a kinematic region where the relative motion of the final two nucleons is low enough to justify the use of EFT. Ando *et al.*'s recent study [29] has shown that μ - d capture can be useful for controlling $\hat{d}_R$, provided the quality of experimental data improves sufficiently.

I have given above a very limited survey of the current status of nuclear χPT. I need to repeat my disclaimer that I have left out many important topics. Among others, I did not discuss a very important attempt by Epelbaum, Glöckle and Meißner [30] to construct a formally consistent framework for applying χPT to complex nuclei (A=3,4 ...). It should be highly informative to apply this type of formalism to electroweak processes and compare the results with those of EFT*.

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