Spin and Weak Interactions in Atoms and Nuclei

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ABSTRACT

The use of spin observables to study the semi-leptonic and non-leptonic weak interaction in atoms and nuclei is surveyed. In particular, the use of semi-leptonic neutral current scattering and atomic parity violation to search for physics beyond the Standard Model is reviewed. The status of nuclear parity violation as a probe of the weak N-N interaction is surveyed. Possible atomic and nuclear signatures of parity conserving, time-reversal violating interactions are also discussed.

The use of atomic and nuclear processes to elucidate the structure of the weak interaction has a long and illustrious history. With the advent of very high-precision, high-energy studies at LEP, SLC, and the Tevatron, it is natural to ask what role, if any, low-energy weak interaction studies might continue to play in uncovering new aspects of electroweak physics. In this talk, I wish to focus on three areas in which such a role can be envisioned: (a) constraining possible extensions of the Standard Model in the neutral current (NC) sector; (b) probing the strangeness-conserving non-leptonic weak interaction; (c) searching for signatures of interactions which conserve parity invariance but violate time-reversal invariance. In each case, I will emphasize the insight which might be derived from the analysis of spin-observables.

Neutral Current Studies

Although $Z$-pole observables from LEP and the SLC are placing increasingly tight constraints on possible extensions of the Standard Model (SM), there still exists a window of opportunity for low-energy observables. To illustrate, one may consider three different types of “new physics” which may appear in NC interactions: (a) additional neutral gauge bosons, (b) effective interactions arising from lepton and quark compositeness, and (c) additional heavy physics which modifies the SM vector boson propagators. While extensions of types (a) and (b) – known as “direct”
contribute at tree level, those of type (c) arise via loops and are correspondingly referred to as “oblique”. Insofar as new direct interactions are associated with mass scales differing from $M_Z$, the high-energy $e^+e^-$ accelerators will be rather transparent to their presence. In contrast, $Z$-pole studies place non-trivial constraints on oblique corrections, since the latter modify the corresponding observables.

The presence of additional, neutral gauge bosons is expected within the context of a variety of grand unified theories in which some group $G$ associated with an un-broken gauge symmetry at a high mass scale spontaneously breaks down to the $SU(3)_c \times SU(2)_L \times U(1)_Y$ symmetry of the weak scale. A particularly useful scheme for considering the generation of additional neutral gauge bosons is associated with the group $E_6$, which arises naturally within the context of heterotic string theory [1]. The various scenarios for $E_6$ breakdown can be characterized by assuming that at least one relatively low-mass neutral boson $Z'$ generated in the process. Moreover, one may decompose this $Z'$ as

$$Z' = \cos \phi Z_\psi + \sin \phi Z_\chi$$

where $\psi$ and $\chi$ denote the U(1) groups appearing in the breakdowns $E_6 \rightarrow SO(10) \times U(1)_\psi$ and $SO(10) \rightarrow SU(5) \times U(1)_\chi$. Different scenarios for the breakdown of $E_6$ are reflected in different values of the angle $\phi$.

From the standpoint of low-energy NC observables, the $Z_\chi$ is the most interesting. The reason is that the $Z_\psi$ has only axial vector couplings to the known leptons and quarks. In the limit that the $Z - Z'$ mixing angle vanishes, it therefore cannot contribute to atomic PV or PV electron scattering, which are my focus here. I consider only the case where the $Z'$ does not mix with the SM $Z$, since results from $Z$-pole measurements place stringent constraints on $M_{Z'}$ for mixing angle $\theta_{ZZ'}$ differing non-negligibly from zero. For $\theta_{ZZ'} \approx 0$, the $Z$-pole observables are relatively insensitive to the presence of a low-mass $Z'$. In contrast, low-energy observables offer the possibility of constraining the mass of an un-mixed $Z'$. Direct searches for such a $Z'$ have been reported by the CDF collaboration, yielding a lower bound on $M_{Z'}$ of between 500 and 585 GeV for the $Z_\chi$ [2]. High-precision low-energy semi-leptonic PV measurements could increase this lower bound by nearly a factor of two.

Just as $Z$-pole observables are transparent to the presence of an un-mixed $Z'$, they are also fairly insensitive to the existence of effective four fermion interactions arising from the assumption that leptons and quarks are composite particles. These effective interactions have the form [3]

$$\mathcal{L}_{\text{composite}} = \frac{4\pi}{\Lambda^2} \tilde{f}_1 \Gamma f_2 \tilde{f}_3 \Gamma f_4,$$  

where the $f_i$ are fermion fields and $\Lambda$ is a mass scale associated with compositeness. Current bounds from atomic PV suggest that the corresponding distance scale $\Lambda^{-1}$ is less than 0.01 of the compton wavelength of the $Z$. Future improvements in the precision of atomic PV or PV electron scattering could improve this bound by more
than 50%. These expectations compare favorably with prospective limits attainable at the Tevatron [4].

The oblique corrections arising from modifications of the $Z$ and $W$ propagators are conveniently characterized in terms of two parameters, $S$ and $T$ [5]. Physically, the former is associated with the presence of degenerate heavy physics, such as an additional generation of fermions in which the members of an iso-doublet have the same mass. The parameter $T$ signals the presence of weak isospin-breaking heavy physics, such as a non-degenerate pair of new heavy fermions. Most significantly, the present global constraints on $S$ favor a central value which is slightly negative, in contrast with the prediction of standard technicolor theories [6]. Both atomic PV and PV electron scattering from a $(J^\pi, I) = (0^+, 0)$ nucleus are essentially sensitive to $S$ and manifest only a slight sensitivity to $T$. In this respect they contrast with most other electroweak observables, from the $Z$-pole on down. The present constraints from atomic PV are consistent with negative values for $S$. Although, by themselves, these results do not significantly affect the 68% or 95% CL contours of the global fits, future measurements with improved precision could impact the location of the central values for $S$ and $T$.

Atomic PV and PV electron scattering are sensitive to these three examples of SM extensions via their dependence on the so-called “weak charge” of the nucleus, $Q_W$. The weak charge can be decomposed as follows:

$$Q_W = Q_W^{(\text{SM})} + \Delta Q_W^{(\text{new})} + \Delta Q_W^{(\text{had})} \ ,$$

where “SM” denotes the contribution arising within the framework of the SM, “new” represents corrections arising from SM extensions as outlined above, and “had” indicates corrections that depend on hadronic and nuclear structure. In the case of atomic PV, $Q_W$ enters the PV amplitude $A_{PV}$ arising from the electron’s axial vector NC interacting with the vector NC of the nucleus. The most precise limits on $Q_W$ have been obtained for atomic cesium, in which $A_{PV}$ is extracted from a PV $6s \rightarrow 7s$ transition in the presence of Stark-induced level mixing. The PV amplitude for this transition can be written as [7]

$$A_{PV} = Q_W \vec{\epsilon} \cdot \langle F'M_F | \vec{\sigma} | FM_F \rangle \ ,$$

where $FM_F$ etc. denote atomic hyperfine levels and $\vec{\sigma}$ is the spin of a valence electron. From the PV transition between states having the same hyperfine quantum numbers, one extracts the ratio

$$|A_{PV}/A_{\text{Stark}}| = \xi Q_W \ ,$$

where $\xi$ as an atomic structure-dependent quantity. As the value of this quantity $\xi$ requires input from atomic theory, one encounters a theoretical, as well as experimental, uncertainty in the corresponding value of $Q_W$. It may well be that the primary challenge for improved constraints on $Q_W$ from atomic PV is the theoretical uncertainty.
In the case of PV electron scattering, the relevant observable is the helicity-difference “left-right” asymmetry $A_{LR}$, which may be expressed as [8]

$$A_{LR} = \frac{N_+ - N_-}{N_+ + N_-} = a_0 |Q^2| \left\{ Q_W + F(Q^2) \right\}, \quad (6)$$

where $N_+$ ($N_-$) denotes the number of scattered electrons for an incident beam of positive (negative) helicity electrons, $a_0$ is a constant depending on the Fermi constant and EM fine structure constant, and $F(Q^2)$ is a term dependent on nuclear form factors. Note that the term containing $Q_W$ is nominally independent of hadron or nuclear structure. In fact, there do exist structure-dependent corrections, contained in $\Delta Q_W$(had), associated with two-boson exchange dispersion corrections. The scale of these corrections has yet to be evaluated reliably by theorists. The strategy for extracting $Q_W$ from $A_{LR}$ is to perform a kinematic separation of the two-terms in Eq. (6) by exploiting the $Q^2$-dependence of the second term. One therefore requires sufficiently reliable knowledge of the form factors in order to successfully carry out this program.

How well might future $Q_W$ determinations from either of these processes do in constraining new physics? To illustrate, I will use cesium atomic PV and PV electron scattering from a $(0^+, 0)$ nucleus, which appear to be the best cases from a variety of standpoints. The present results for $Q_W$(Cs), for which the combined experimental and theoretical error is about 2.5%, constrain the mass of the $Z_\chi$ to be greater than about 0.5 TeV, a limit roughly comparable with the bounds from the Tevatron. A future 1% determination would push this bound to 0.8 TeV. Similarly, a 1% measurement of $Q_W(0^+, 0)$ from PV electron scattering would yield a bound of 0.9 TeV (the difference between the two processes follows from the different $u$- and $d$-quark content of the respective nuclei). In terms of compositeness, the present cesium results require $\Lambda > 10$ TeV. A 1% determination of $Q_W$(Cs) or $Q_W(0^+, 0)$ would place this lower bound at about 16 TeV. Finally, a factor of four improvement in the cesium precision would shift the global central value of $S$ by a factor of four or so, assuming the same central value of $Q_W$ is obtained in a future experiment [6]. A similar statement applies to extractions of $Q_W$ from PV electron scattering. In short, it is apparent that pushing for improved precision in these low-energy processes could yield significant constraints on SM extensions which complement those obtained from other NC observables.

The prospects for achieving such improved precision are promising. In the case of cesium atomic PV, the systematic error is already at the 0.5% level, and one anticipates achieving a total experimental error of 0.5% in the future. The present atomic theory error is 1.2%, and it remains to be seen whether atomic theorists can push this error below one percent in the future. Previous experience with the MIT-Bates $^{12}$C PV electron scattering experiment indicates that achieving systematic error on the order of 1% is within reason, while the high luminosity available at TJNAF implies that obtaining a similar level of statistical error is possible with an experiment of realistic running time. Moreover, the present PV program at
MIT-Bates, TJNAF, and MAMI should yield sufficient information on the NC form factors $F(Q^2)$ appearing in Eq. (6) to render them a negligible source of uncertainty for a $Q_W$ extraction. The primary theoretical challenge appears to be obtaining a realistic evaluation of the uncertainty in $\Delta Q_W(\text{had})$ for this process.

**Nuclear Parity Violation**

From the standpoint of electroweak theory, nuclear PV observables are of interest as a window on the $\Delta S = 0$ hadronic weak interaction. Within the framework of the SM, this interaction is composed of $I = 0, 1, 2$ components. One may correspondingly write down a two-body nuclear Hamiltonian having the same isospin content:

$$\hat{H}_{NN}^{\text{PV}} = \sum_i h_i \hat{O}_i(2) \quad ,$$

(7)

where the $h_i$ are constants dependent on the hadronic weak interaction, the $\hat{O}_i(2)$ are two-body nuclear operators containing various spin, isospin, and momentum structures, and the index $i$ runs over the possible channels containing $I = 0, 1, 2$. Given the hard core of the strong $NN$ potential, it is unlikely that the weak $NN$ force is mediated by the exchange of a weak vector boson between two nucleons. Rather, one expects the exchange of mesons to dominate the weak two-nucleon potential. Under this ansatz, the $h_i$ are given as products of strong and weak PV meson-nucleon couplings:

$$h_i = g_{NNM}^i h_{NNM}^i \quad ,$$

(8)

where the $h_{NNM}^i$ ($g_{NNM}^i$) are weak PV (strong) meson ($M$) nucleon ($N$) couplings. In order to obtain all of the isospin components required by the structure of the $\Delta S = 0$ quark-quark interaction, one must include the exchange of the $\pi, \rho, \omega$ mesons at a minimum.

When seeking to extract information on the hadronic weak interaction from PV nuclear observables, one must undertake several levels of analysis. From a phenomenological standpoint, the problem is to determine whether one may obtain a consistent set of $h_i$ from a global analysis of observables. In this respect, one must also rely upon nuclear theory to provide computations of nuclear matrix elements (such as $\langle A|\hat{O}_i(2)|A\rangle$) in order to extract the $h_i$ from experimental quantities. Within the framework of the meson-exchange picture, one would also like to understand how the values of the $h_{NNM}^i$ arise from the PV four-quark weak Hamiltonian, $\mathcal{H}_{\text{PV}}$. This objective presents hadron structure theorists with the problem of reliably computing weak matrix elements $\langle NM|\mathcal{H}_{\text{PV}}|N\rangle$, a non-trivial task. In fact, there has been little progress in this direction since the “benchmark” quark model calculation of Ref. [9] more than 15 years ago. Clearly, deepening our understanding of the $\Delta S = 0$ hadronic sector of the SM requires progress on a variety of fronts.

To date, most experimental information has been derived from two broad classes of observables: so-called “direct” $NN$ studies, such as $p^-+p$ scattering, and studies of
PV in light nuclei, such as the PV $\gamma$ decays of $^{18}$F and $^{19}$F. In the latter instance, the size of the PV observable is enhanced by the mixing of nearly-degenerate opposite parity states by $\hat{H}_{NN}^{PV}$. In both direct and light-nuclei studies, the observables involve some form of spin polarization. At present, one has yet to achieve a consistent set of the $h_i$ from a rather broad sample of observables. Of particular interest is the constraint from $^{18}$F PV on $h_{NN\pi}$, which is in conflict with constraints from $\vec{p} + \alpha$, $^{19}$F, and $^{21}$Ne experiments [10]. The reason for this discrepancy is not fully understood, but the possibilities include (a) an error in one of the experiments (b) significant nuclear theory uncertainty in the computation of PV nuclear matrix elements, or (c) the omission of important terms from the model of Eq. (7). A recent analysis of Brown and collaborators [11], performed in a truncated model space but including higher $nh\omega$ configurations shift the $^{21}$Ne constraints in such a manner as to bring all the bounds on $h_{NN\pi}$ into agreement. Whether this trend emerges from a computation using a complete model space remains to be seen [12].

The measurement of new observables is clearly desirable, as such measurements could yield new and, in principle, complementary constraints on the $h_i$. In this respect, two possibilities have received considerable interest recently: scattering of epithermal polarized neutrons from heavy nuclei [13] and observation of the nuclear anapole moment via atomic PV experiments [14]. In the case of the neutron scattering experiments, one measures a neutron transmission asymmetry associated with incident neutrons of opposite helicity. Such an asymmetry arises from the mixing of $s$-wave resonances into $p$-wave resonances by the interaction of Eq. (7). Due to the large level densities for nuclei such as $^{238}$U, the energy denominators associated with the mixing are small. Moreover, the $s$-wave resonances couple strongly to the continuum. As a consequence of these two features, the PV asymmetry can be significantly enhanced. From a theoretical standpoint, the presence of a large number of neighboring $s$- and $p$-wave resonances complicates the analysis, and one must resort to statistical approaches in extracting information about the PV nuclear force. To the extent that the assumptions of the statistical models are valid, one derives from the asymmetry a mean square PV nuclear matrix element. The corresponding constraints on the $h_i$ take the shape of a quadratic form in the multi-dimensional space of PV couplings. These constraints appear to be consistent with the constraints on $h_{NN\pi}$ derived from the $^{18}$F and $^{21}$Ne $\gamma$-decays but not with the constraints obtained from $^{19}$F or $\vec{p} + \alpha$ experiments. In the case of one nucleus, $^{232}$Th, the results display a deviation from the pattern expected within the conventional statistical model. The mean value of the measured $^{232}$Th asymmetries differs from zero by more than two standard deviations (one expects this average to be zero in the statistical model). Various extensions of the statistical approach used to account for this deviation yield a mean value for the PV matrix element which is two orders of magnitude larger than the scale of PV matrix elements implied by other PV measurements. While this so-called “sign problem” arises only in the case of $^{232}$Th, one has yet to obtain a satisfactory explanation.

The second new approach to placing new constraints on the $h_i$ is to measure the
nuclear anapole moment (AM). The AM is an axial vector coupling of the photon to the nucleus induced by parity-mixing in the nucleus. Technically, it is an elastic matrix element of the transverse electric multipole operator, $\hat{T}_E^{J=1}$ – a matrix element which must vanish in the absence of PV. This matrix element goes like $Q^2$ for small momentum transfer. Consequently, the AM couples only to virtual photons, such as those exchanged between the nucleus and atomic electrons. Moreover, because the leading $Q^2$ of the AM cancels the $1/Q^2$ of the photon propagator, the corresponding interaction is contact-like in co-ordinate space. In this respect, the contribution made by the AM to PV observables is indistinguishable from the $V(e) \times A(N)$ NC interaction. Indeed, both induce a nuclear spin-dependent (NSD) atomic interaction of the form

$$\mathcal{H}_{\text{ATOM}}^{\text{PV}}(\text{NSD}) = \frac{G_F}{\sqrt{2}} \bar{k} \psi_e^\dagger(0)\bar{\alpha}\psi_e(0) \cdot \vec{I},$$  \hspace{1cm} (9)$$

where $\psi_e$ is the electron field, $\bar{\alpha}$ is the vector of Dirac matrices, $\vec{I}$ is the nuclear spin, and the constant $\bar{k}$ can be decomposed into a sum of NC and AM contributions: $\bar{k}_{NC} + \bar{k}_{AM}$.

Although $\bar{k}_{AM}$ is nominally suppressed with respect to $\bar{k}_{NC}$ by a factor of $\alpha$, one may nevertheless expect it to make an observable contribution for the following two reasons. First, the NC contribution is suppressed since $\bar{k}_{NC} \propto g_V^e = -1 + 4 \sin^2 \theta_W \approx -0.1$. Second, the scale of the AM grows with the square of the nuclear radius, and thus as $A^{2/3}$, whereas NC contribution receives no coherence enhancement. Hence, for heavy nuclei, $\bar{k}_{AM}/\bar{k}_{NC}$ can be as large as three or more, according to a variety of calculations [15]. Although recent results from the cesium atomic PV experiment is not conclusive, it is nevertheless consistent with theoretical expectations [16]. The result from the atomic thallium experiment [16] differs from theoretical predictions by about $2\sigma$ and has the opposite sign. There is undoubtedly room for improvement on the theoretical side as well as on the part of experiment. Ideally, future measurements of NSD atomic PV observables will achieve significantly better precision and, coupled with theoretical progress, yield new constraints on the $h_i$.

**Parity Conserving Time-reversal Violation**

Finally, I wish to touch briefly on a subject which has received renewed interest recently – searches for parity conserving time-reversal violating (PCTV) interactions. Traditionally, such searches have relied on three classes of studies: detailed balance in nuclear reactions, nuclear $\gamma$-decays, and $\beta$-decay [17]. More recently, constraints on PCTV physics have been derived from two other observables: the five-fold correlation in epithermal neutron scattering from heavy nuclei and the permanent electric dipole moments (EDM’s) of atoms and nuclei. In the case of neutron scattering, one looks for a scattering phase shift $\delta_{\text{PCTV}}$ proportional to

$$\vec{s} \cdot (\vec{I} \times \vec{p})(\vec{I} \cdot \vec{p})$$ \hspace{1cm} (10)$$

where $\vec{s}$ is the spin of an incident neutron of momentum $\vec{p}$ and $\vec{I}$ is the nuclear spin. Typically, searches for EDM’s try to detect a frequency shift $\Delta\nu \sim d \vec{J} \cdot \vec{E}$, where $d$
is the EDM, $\vec{E}$ is a static, applied electric field, and $\vec{J}$ is the spin of the quantum system of interest.

What makes $\delta_{\text{PCTV}}$ and $d$ particularly interesting is their sensitivity to PCTV flavor-conserving ($\Delta F = 0$) interactions. Within the context of renormalizable gauge theories, $\Delta F = 0$ PCTV interactions between quarks cannot arise at $O(g^2)$, where $g$ is the gauge coupling [18]. In order to generate them, one requires loops involving gauge interactions which exist beyond the framework of the SM. Alternatively, one may work with non-renormalizable effective interactions which apply below some scale $M_X$ associated with PCTV gauge interactions. In the latter instance, one finds a class of dimension seven operators which can generate a $\delta_{\text{PCTV}}$ or EDM at the quark level [19]. The list of such operators includes, for example,

$$\frac{\beta}{M_X^3} \bar{\psi} \gamma_5 D_\mu \psi \bar{\psi} \gamma_5 \gamma_\mu \psi$$

and

$$\frac{eg\tilde{\beta}}{M_X^3} \bar{\psi} \sigma_{\mu\lambda} \psi F_{(\gamma)}^\mu_{\nu} F_{(Z)}^{\lambda\nu} g_{\rho\nu}$$

where $\psi$ is a fermion field and $F_{(\gamma)}^\mu_{\nu}$ ($F_{(Z)}^{\mu\nu}$) is the field strength associated with the photon ($Z$-boson). Both of these interactions may generate contributions to the EDM of a quark. To do so, they require the presence of a PV weak interaction, since the interaction of an EDM with a static electric field is both parity and time-reversal violating. Assuming, for illustrative purposes, that the couplings $\beta$ and $\tilde{\beta}$ are of order unity, the present limits on the neutron EDM would constrain the mass scale $M_X$ to be roughly two orders of magnitude larger than $M_Z$ [20].

An alternate scheme for treating PCTV in the purely hadronic sector is to employ a meson-exchange model. In this case, the lightest allowed meson is the $\rho$, which can interact with one nucleon through the PC strong interaction and the other nucleon through a PCTV interaction. The associated PCTV coupling is conventionally denoted $g_{\rho}$. The present upper bounds on $g_{\rho}$ from detailed balance and $\delta_{\text{PCTV}}$ are about 2.5 and 22, respectively [18]. One expects the epithermal neutron bounds to improve by a factor of 100 or so with the completion of future measurements. Alternately, one may derive limits on $g_{\rho}$ from atomic EDM’s by assuming that the EDM is generated by (a) PCTV in the purely hadronic sector and (b) an additional PV weak interaction either inside the nucleus or between the nucleus and atomic electrons. In the case of the neutron EDM, all of the symmetry violating interactions are hadronic. The corresponding limits on $g_{\rho}$ obtained from the neutron EDM are roughly $< 10^{-3}$, while those obtained from atomic EDM’s are about an order of magnitude weaker [21]. With the advent of more precise atomic EDM measurements, one would anticipate deriving better limits on $g_{\rho}$. The theoretical problem of understanding how such bounds would translate into constraints on PCTV quark-quark interactions remains open.

Conclusions

In this talk, I hope to have convinced you that the use of spin-dependent atomic
and nuclear weak interaction observables have an important and on-going role to play in searching for electroweak physics beyond the Standard Model. In the case of semi-leptonic PV, for example, the prospects are good for placing bounds on the scale of compositeness and the mass of an additional neutral gauge boson which are competitive with, or better than, those one might achieve with high-energy accelerators. Similarly, the study of hadronic parity violation continues to challenge our understanding of the purely hadronic weak interaction. Finally, the analysis of spin observables such as the five-fold correlation in epithermal neutron scattering or neutron and atomic EDM’s yield constraints on the scale of PCTV interactions which may arise in certain extensions of the SM. Undoubtedly, the continuing improvement in the precision with which spin-dependent atomic and nuclear observables are measured will provide an abundant supply of grist for the electroweak theorist’s mill.

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