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EFFECTS OF THE COULOMB INTERACTION IN NUCLEI: BINDING ENERGIES AND ISOSPIN MIXING*

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The effects of Coulomb and other charge-dependent interactions on the structure of nuclei are discussed. First, a method to predict absolute binding energies for proton-rich nuclei that allows a mapping of the proton drip-line up to $A = 70$ is presented. Then the effects of isospin-symmetry breaking on tests of the standard model for the weak interaction are examined. The principal areas discussed are (1) superallowed Fermi beta decay, which provides a test of the conserved vector current hypothesis and the unitarity of Cabibbo–Kobayashi–Maskawa matrix; and (2) parity violation in electron scattering, which offers a window into the neutral-current sector of the weak interaction.

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1. Introduction

Theoretical predictions for nuclear binding energies and the location of the proton drip-line require an understanding of the role played by both the strong and Coulomb interactions in nuclei. Of these two, the contribution due to the strong interaction is the more dominant and the more difficult to predict, leading to overall uncertainties of at least 300 keV for nuclei with $A \leq 50$, and considerably more for heavier nuclei. For many proton-rich nuclei, however, binding energies have been measured for their corresponding mirrors. Consequently, an accurate estimate of the binding energy can be made by exploiting analog (or isospin) symmetry. In particular, all that is needed is to estimate the shift in the binding energy caused by the Coulomb interaction and add this quantity to the experimental binding energy of the neutron-rich mirror.

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With predictions for binding energies, it is then possible to explore a range of phenomena ranging from the *rp*-process to exotic decay modes. One mode that is predicted by binding energy studies, but is yet to be observed, is the emission of two correlated protons, otherwise known as di-proton emmission. The number of candidate nuclei amenable to experiment, however, is severely limited because of an extreme sensitivity in the decay lifetime on the two-proton separation energies.

Atomic nuclei also represent a laboratory in which "fundamental" symmetries, such as the standard model for the weak interaction, can be tested. Two excellent examples are: (1) the *ft*-values for superallowed Fermi beta decay, which test the conserved vector current hypothesis and the unitarity of the Cabibbo–Kobayashi–Maskawa matrix; and (2) parity violation in electron scattering on $N = Z$ nuclei, which offers a probe of the neutral-current sector of the weak interaction. In each case, deviations from the predictions of the standard model might be interpreted as signals of "new" physics. The principal feature of both examples are that they were chosen so as to minimize the effects due to nuclear structure. In fact, in the limit that isospin is a good quantum number, the measured observables are essentially insensitive to nuclear structure. However, because of the Coulomb interaction and other small charge-dependent components of the strong interaction, isospin symmetry is violated and small corrections are expected.

This lecture explores the role played by the Coulomb interaction in nuclei, and is organized in the following manner. First, Section 2 gives a brief description of the isospin quantum number. The method used to predict binding energies for proton-rich nuclei and possible candidates for di-proton emission are presented in Section 3. The effects of isospin-symmetry breaking on superallowed Fermi beta decay and parity-violating electron scattering are discussed in Sections 4 and 5, respectively, and conclusions are gathered in Section 6.

2. Isospin

Shortly after the discovery of the neutron, Heisenberg proposed a new quantum number that expressed a symmetry between protons and neutrons. Heisenberg's hypothesis was to regard protons and neutrons as a charge doublet of essentially identical particles. This property can be described by introducing a new spin, called isospin, whose quantized z-component, t_z, is related to the electric charge by $q = e(1/2 + t_z)$.

Isospin introduces a mechanism for labeling the many-body states in nuclei with a particular symmetry. If the interaction between protons and neutrons were charge independent, the energy spectrum for a system of A nucleons with Z protons and N neutrons would be identical to that of the nucleus with N protons and Z neutrons. Namely, that these nuclei with $T_z = \pm|Z-N|/2$ are mirror images of one another.

In addition to mirror symmetry, we also expect to find analogs of the states in the $T_z = \pm|Z-N|/2$ nuclei in all nuclei with the same nucleon number and $|T_z| < |Z-N|/2$. If isospin is a good quantum number, then the wave functions for these analog states may be obtained by the repeated operation of the isospin lowering operator on the state with maximum T_z. An illustrative example is given by the two nucleon system, where the $T_z = \pm 1$ states have total isospin $T = 1$, with the wave functions given by $|T = 1T_z = 1\rangle = |\uparrow\uparrow\rangle$ and $|1-1\rangle = |\downarrow\downarrow\rangle$, respectively. For $T_z = 0$, there are two states with $T = 1$ and 0 and wave functions

$$|00\rangle = \frac{1}{\sqrt{2}}(|\uparrow\downarrow\rangle - |\downarrow\uparrow\rangle)\,, \tag{1}$$

$$|10\rangle = \frac{1}{\sqrt{2}}(|\uparrow\downarrow\rangle + |\downarrow\uparrow\rangle). \tag{2}$$

It is easy to verify that Eq. (2) is the analog of the $T_z = \pm 1$ states.

The degree to which isospin may be used is determined by the charge symmetry in the nucleon-nucleon interaction. For the most part, the interaction between nucleons is predominantly charge independent, and isospin may be thought of as an approximate quantum number. Consequently, isospin remains a powerful spectroscopic tool. On the other hand, it must be pointed out that the isospin-symmetry breaking has important consequences for some processes, such as superallowed Fermi beta decay and parity violation in electron scattering as described below.

3. Coulomb energy differences

If the nuclear Hamiltonian is composed of at most two-body parts, it may be separated into three components. The dominant part, which is also responsible for most of the nuclear binding, is isoscalar and is due to the strong interaction. The other two components are due to the Coulomb interaction and charge non-symmetric parts of the nucleon-nucleon interaction, and are isovector and isotensor in character. If the isovector and isotensor components are weak, then the binding energies for the members of an isospin multiplet may be described using the isobaric mass multiplet equation [2]

$$BE(A,T,T_z,i) = a(A,T,i) + b(A,T,i)T_z + c(A,T,i)T_z^2. \tag{3}$$

The coefficients a, b, and c separately depend on the isoscalar, isovector, and isotensor components of the nuclear Hamiltonian, respectively.

From Eq. (3), the binding energy difference between iosbaric analogs with $T_z = \pm T$ is given by

$$BE(A, T, T_z = T, i) - BE(A, T, T_z = -T, i) = 2b(A, T, i)T. \qquad (4)$$

Therefore, an accurate way to predict binding energies for proton-rich nuclei whose analog has an experimentally measured mass is to compute the b-coefficient and add $2bT$ to the experimental binding energy of the neutron-rich analog.

Shell-model calculations for the Coulomb energy difference have been carried out using empirical "Coulomb" interactions [3] for 37 nuclei in the mass region $36 \leq A \leq 48$ with active particles in the $0d_{3/2}$ and $0f_{7/2}$ orbits [4] and for 75 nuclei with $46 \leq A \leq 70$ using the $0f_{7/2}$, $0f_{5/2}$, $1p_{3/2}$, and $1p_{1/2}$ orbits [5]. The empirical "Coulomb" interactions were found to reproduce experimental b-coefficients at level of 30-45 keV [3], and, consequently, the predicted binding energies have an accuracy of approximately $40|Z - N|$ keV.

3.1. Di-proton emission

Because of the pairing interaction, a nucleus with an even number of protons (Z, N) is generally more tightly bound than a $(Z - 1, N)$ nucleus, but, because of the symmetry energy and Coulomb repulsion, it may be unbound relative to the $(Z - 2, N)$ system. In this case, the parent nucleus may decay by the emission of a di-proton, *i.e.*, a correlated proton pair. With the predicted binding energies, candidates for di-proton emission may be identified. The most important observation is that the number of candidates for which experimental detection is feasible is sharply limited by the two-proton separation energy. This is in part due to the fact that β^+ emission is a competing decay mechanism with lifetimes of the order 1–100 ms. In addition, a lower limit of approximately 1 ns is often imposed by the experimental apparatus. On the other hand, the decay rate for diproton emission is determined by the probability to penetrate through the Coulomb barrier, which is exponentially dependent on the two-proton separation energy. Because of this, the number of candidates for which the observation of diproton decay is practical, is limited to nuclei with two-proton separation energies between 0.9 and 1.4 MeV.

Listed in Table I are nuclei with di-proton halflives that are predicted to be of the order 1 ms or shorter. Note that since the uncertainties in the two-proton separation energies are of the order 200 keV, the exponential sensitivity in the halflife on the separation energy leads to a fairly wide range of halflives. From Table I, the best candidates for experimental observation are ^{45}Fe, ^{48}Ni, and ^{66}Se.

TABLE I

Range of halflives for di-proton emitter candidates. Also listed are the predictions for the one- and two-proton separation energies (S_p and S_{2p}).

^{A}Z	S_p (MeV)	S_{2p} (MeV)	$t_{1/2}$ (s)	$t_{1/2}^{\min}$ (s)	$t_{1/2}^{\max}$ (s)
^{38}Ti	0.438(164)	−2.432(132)	9×10^{-16}	4×10^{-16}	2×10^{-15}
^{45}Fe	−0.010(198)	−1.279(181)	10^{-6}	10^{-8}	10^{-4}
^{48}Ni	0.505(351)	−1.290(330)	4×10^{-6}	5×10^{-9}	0.09
^{59}Ge	0.058(211)	−1.343(192)	10^{-3}	10^{-5}	0.3
^{63}Se	0.069(288)	−1.530(262)	6×10^{-5}	3×10^{-7}	5×10^{-2}
^{66}Kr	−0.001(351)	−2.832(325)	3×10^{-12}	2×10^{-13}	6×10^{-11}
^{67}Kr	0.155(288)	−1.538(262)	2×10^{-3}	10^{-5}	0.2

4. Superallowed Fermi beta decay

Superallowed Fermi β transitions in nuclei provide an excellent laboratory for precise tests of the properties of the electroweak interaction, and have been the subject of intense study for several decades. According to the conserved-vector-current (CVC) hypothesis, for pure Fermi transitions the product of the partial half-life, t, and the statistical phase-space factor, f, should be nucleus independent and given by

$$ft = \frac{K}{G_V^2 |M_{\rm F}|^2}, \tag{5}$$

where $K/(\hbar c)^6 = 2\pi^3 \ln 2\hbar/(m_e c^2)^5$, G_V is the vector coupling constant for nuclear β decay, and $M_{\rm F}$ is the Fermi matrix element, $M_{\rm F} = \langle \psi_f \mid T_\pm \mid \psi_i \rangle$. By comparing the decay rates for muon and nuclear Fermi β decay, the Cabibbo–Kobayashi–Maskawa (CKM) mixing matrix element [7] between u and d quarks (v_{ud}) can be determined and a precise test of the unitarity condition of the CKM matrix is possible.

For tests of the standard model, two nucleus-dependent corrections must be applied to experimental ft values. The first is a series of radiative corrections to the statistical phase-space factor embodied in the factors $\delta_{\rm R}$ and $\Delta_{\rm R}$, giving [8]

$$f_{\rm R} = f(1 + \delta_{\rm R} + \Delta_{\rm R}), \tag{6}$$

where $\delta_{\rm R} \sim 1.3\%$ is due to electromagnetic radiative corrections, currently evaluated to order $Z^2\alpha^3$, and $\Delta_{\rm R} \sim 2.3\%$ is what has been referred to as the "outer" radiative correction and includes vector-axial vector interference terms. The second correction, which is discussed in detail here, arises because of the presence of isospin-nonconserving (INC) forces in nuclei that lead to a renormalization of the Fermi matrix element [9, 10]. This correction is denoted by $\delta_{\rm C}$ and modifies the Fermi matrix element by $\mid M_{\rm F} \mid^2 = [T(T+1) - T_{Z_i}T_{Z_f}](1-\delta_{\rm C})$.

With the above corrections, a "nucleus-independent" $\mathcal{F}t$ can be defined by

$$\mathcal{F}t = ft(1+\delta_{\rm R}+\Delta_{\rm R})(1-\delta_{\rm C}), \tag{7}$$

and the CKM matrix element v_{ud} is given by

$$\mid v_{ud} \mid^2 = \frac{\pi^3 \ln 2}{\mathcal{F}t}\frac{\hbar^7}{G_{\rm F}^2 m_e^5 c^4} = \frac{2984.38(6)\ {\rm s}}{\mathcal{F}t}, \tag{8}$$

$G_{\rm F}$ is obtained from muon β-decay. Currently, ft values for nine superallowed transitions have been measured with an experimental precision of 0.2% or better [11]. With these precise measurements, the CVC hypothesis can be confirmed by checking the constancy of the $\mathcal{F}t$ values for each nucleus, while the unitarity condition of the CKM matrix is tested by comparing the average value of v_{ud} with the values determined for $v_{us} = 0.2199(17)$ [12] and $v_{ub} < 0.0075$ [13], namely that $v_{ud}^2 + v_{us}^2 + v_{ub}^2 = 1$.

The correction $\delta_{\rm C}$ can be evaluated within the framework of the nuclear shell model [9, 10]. Due to computational limitations and uncertainties associated with determining an effective Hamiltonian, almost all shell-model calculations for nuclei with $A \geq 10$ are performed within a single major oscillator shell, *e.g.*, for ^{10}C the model space spanned by the $0p_{3/2}$ and $0p_{1/2}$ oribitals. Within this context, two types of isospin mixing must be accounted for. The first is due to the mixing between states that lie within the shell-model configuration space. For example, for $A = 10$, there are 2, 7, and 1 configurations leading to $J^\pi = 0^+$ and $T = 0$, 1, and 2, respectively. Because of its two-body nature, the INC interaction is composed of isospin operators of rank zero, one, and two, and consequently mixes together all the $J^\pi = 0^+$ states. Traditionally, the configuration mixing correction is denoted as $\delta_{\rm IM}$ and in Ref. [10] it was shown that $\delta_{\rm IM}$ is best evaluated using an INC interaction that reproduces the Coulomb energy shifts between members of isospin multiplets [3]. Of the two types of mixing, $\delta_{\rm IM}$ is the smallest with a magnitude of approximately $0.04 - 0.1$%.

In addition to the mixing between states contained within the configuration space, mixing with states that lie outside the model space must also be accounted for. In particular, the Coulomb interaction can strongly mix $1p-1h$, $2\hbar\Omega$ excitations, *e.g.*, $0p_{3/1} \rightarrow 1p_{3/2}$, into the ground state. Excitations of this type are accounted for by examining differences in the single-particle radial wave functions. Indeed, for closed-shell configurations, mixing with $1p-1h$ states is properly accounted for at the level of Hartree–Fock.

Hence, the second correction to the Fermi matrix element, denoted by $\delta_{\rm RO}$, was estimated by evaluating the mismatch in the radial overlap between the single-particle wave functions of the converted proton and the corresponding neutron. The explicit details for the calculation of $\delta_{\rm RO}$ are given in Refs. [9, 10]. Schematically, though, one finds $\delta_{\rm RO} \approx 1 - \int r^2 dr R^p R^n$, where $R^{p(n)}$ is the proton(neutron) single-particle wave function. In general, $\delta_{\rm RO}$ is the larger of the two components ($\delta_{\rm C} = \delta_{\rm RO} + \delta_{\rm IM}$), and has a magnitude of the order 0.1-0.8%.

Two methods for evaluating $\delta_{\rm RO}$ have been espoused. The first (THH) [9] uses Woods–Saxon (WS) radial wave functions, while in the second (OB) [10], Hartree–Fock (HF) wave functions are employed. Generally speaking, the two methods yield approximately the same dependence on nucleon number A, but the HF values are systematically smaller by 0.1% for the magnitude of the correction. The reason for the difference lies in the HF mean field. The principal effect of the Coulomb interaction is to push the proton wave functions out relative to the neutrons. In HF, however, the proton and neutron mean fields are coupled, and the Coulomb interaction induces an attractive isovector mean field between the protons and neutrons. In effect, the Coulomb interaction pushes the protons out, but because of the strong interaction, the protons pull the neutrons out with them, hence, reducing the magnitude of the radial overlap mismatch.

With the *known* corrections applied to the experimental data, the $\mathcal{F}t$ values are found to be essentially constant, with a $\chi^2/\nu \sim 0.7$, but the unitarity limit is violated at the level of approximately 0.4(1)% or 0.3(1)% for the OB and THH corrections, respectively. This discrepancy from unitarity is difficult to reconcile, as one implication is another quark generation. However, this implies a larger CKM matrix element than presently found for v_{ub}. In addition, this would break the three-flavor symmetry that is at least partially confirmed by the fact that the number of light neutrino species is limited to three from both cosmology and high-energy experiments. Alternatively, perhaps a small correction due to nuclear structure is still unaccounted for. Indeed, the separation between $\delta_{\rm IM}$ and $\delta_{\rm RO}$, while presently necessary, is somewhat unsatisfying. This issue has recently been addressed for ^{10}C, where the isospin-mixing corrections have been evaluated [14] within the framework of a no-core, large-basis shell-model calculation (up to $4\hbar\Omega$) using an effective interaction derived from a realistic nucleon-nucleon interaction including Coulomb and charge-dependent terms. A value of $\delta_{\rm C} \approx 0.1\%$ was obtained, which is in excellent agreement with the previous estimates. Two other possibilities, which warrant further study, are the CKM matrix element v_{us} and better evaluation of the radiative correction involving vector-axial vector interference.

5. Parity-violating electron scattering

In addition to the charged vector current, nuclei can also be used to test the neutral-current sector of the weak interaction. For this, experiments involving the elastic scattering of electrons on even-even, $N = Z$ nuclei are planned, with the observable of interest being the parity-violating electron-scattering asymmetry [15, 16]

$$\mathcal{A} = \frac{\mathrm{d}\sigma^+ - \mathrm{d}\sigma^-}{\mathrm{d}\sigma^+ + \mathrm{d}\sigma^-} = -\left(\frac{G_\mathrm{F} q^2}{2\pi\alpha\sqrt{2}}\right)\frac{\tilde{F}_\mathrm{C}(q)}{F_\mathrm{C}(q)}, \tag{9}$$

where $q = |\boldsymbol{q}|$ is the magnitude of the three momentum transfer. The dependence on nuclear structure is embodied in the Coulomb monopole form factors $F_\mathrm{C}(q)$ and $\tilde{F}_\mathrm{C}(q)$ for the electromagnetic and neutral currents, respectively. In the shell model, both form factors have the same form, and are given by

$$F_\mathrm{C}(q) = \sum_{\mu}^{\text{protons}} n_\mu^p q_E^p f_\mu^p(q) + \sum_{\mu}^{\text{neutrons}} n_\mu^n q_E^n f_\mu^n(q), \tag{10}$$

where μ denotes the labels for each single-particle orbit, $n_\mu^{p(n)}$ the number of protons(neutrons) occupying each single-particle orbit, and $f_\mu^{p(n)}(q)$ is the form factor for the individual single-particle states given by

$$f_\mu^{p(n)}(q) = \frac{1}{\sqrt{4\pi}}\int_0^\infty r^2\mathrm{d}r(R_\mu^{p(n)}(r))^2 j_0(qr), \tag{11}$$

where $R_\mu^{p(n)}(r)$ is the proton(neutron) radial wave function for the single-particle orbit μ. For the Coulomb part, the charges are given by $q_E^p = 1$ and $q_E^n = 0$, while for the neutral current, the corresponding weak charges are $q_\mathrm{W}^p = (1 - 4\sin^2\theta_\mathrm{W})/2$ and $q_\mathrm{W}^n = -1/2$, where θ_W is the Weinberg angle.

In the limit that isospin is a good quantum number, both $n_\mu^p = n_\mu^n$ and $R_\mu^p(r) = R_\mu^n(r)$, and Eq. (10) reduces to

$$\mathcal{A}_0 = [G_\mathrm{F} q^2/\pi\alpha\sqrt{2}]\sin^2\theta_\mathrm{W} = 3.22 \times 10^{-6} q^2. \tag{12}$$

It is the simple form of Eq. (12) that makes experiments on even-even $N = Z$ nuclei an attractive choice for testing the standard model. In particular, the experimental goal is to perform a 1% measurement of $\mathcal{A}$, and, thus, a 1% measurement of $\sin^2\theta_\mathrm{W}$. Naturally, any deviations from the simple q^2 dependence as well as the expected curvature, might be a signature of physics beyond the standard model. Isospin, however, is not a conserved quantity and the effects of this broken symmetry must be evaluated. Because isospin is an approximate quantum number, deviations from Eq. (12) are expected to be small and are quantified by the factor $\Gamma(q)$ [16] defined as

$$\mathcal{A} = \mathcal{A}_0(1 + \Gamma(q)), \tag{13}$$

where, from Eq. (9), $\Gamma(q)$ may be written as

$$\Gamma(q) = -[1 + \tilde{F}_{\rm C}(q)/2\sin^2\theta_{\rm W} F_{\rm C}(q)]. \tag{14}$$

An estimate of $\Gamma(q)$ proceeds in a similar manner as $\delta_{\rm C}$ for the Fermi matrix element. Namely, that isospin-mixing both within the configuration space and $1p-1h$ states must be accounted for. For the former, shell-model calculations in proton-neutron formalism using the INC interaction of Ref. [3] are carried out, with the principal effect being that $n^p_\mu \neq n^n_\mu$. Again, $1p-1h$ mixing is taken into account by using Hartree–Fock proton and neutron radial wave functions to evaluate the form factors in Eq. (11).

Shown in Fig. 1, is the correction factor $\Gamma(q)$ for the targets ^{4}He, ^{12}C, ^{16}O, and ^{28}Si as a function of momentum transfer up to 1.2 fm $^{-1}$. The dotted line in the figure illustrates the level at which isospin-mixing corrections exceed 1%. From the figure, it is apparent that the experiments have to be performed at the lowest of momenta transfers. From Ref. [17], the momenta transfer that maximize the figure of merit for ^{4}He and ^{12}C, which would provide a 0.7% measurement of $\mathcal{A}$, are 0.97 fm^{-1} and 0.62 fm^{-1}, respectively. From Fig. 1, the corrections due to isospin-symmetry breaking are then expected to be 0.5% and 0.7% for ^{4}He and ^{12}C, respectively.

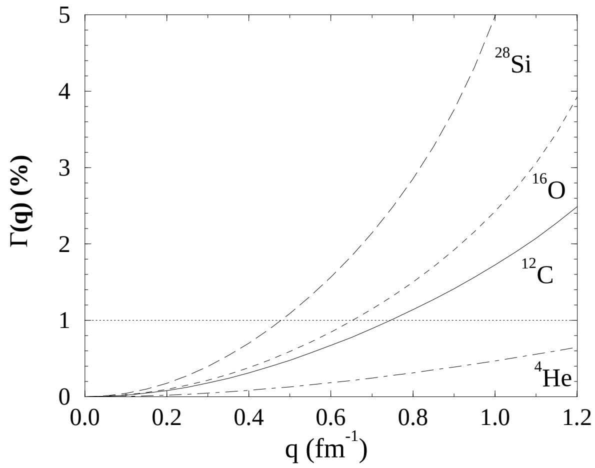

Fig. 1. Calculated values of $\Gamma(q)$ (in %) obtained for ^{4}He, ^{12}C, ^{16}O, and ^{28}Si.

6. Conclusions

Although most of the nuclear binding is provided by the charge-symmetric part of the strong interaction, the isospin-nonconserving components of the nuclear Hamiltonian, such as the Coulomb and other weaker parts of the strong interaction, play an important role determining the structure of nuclei. In addition, since these INC components are fairly weak, isospin remains an approximate quantum number and is powerful spectroscopic tool. In these lectures, it was shown that this approximate symmetry can be exploited to predict the binding energies of nuclei along the proton drip-line. The extent to which isospin-symmetry is broken can also play an important role for some observables. Two classic examples that have important ramifications regarding precise tests of the standard model for the weak interaction were illustrated in this lecture.

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