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Appendix

NUCLEAR STRUCTURE PUZZLES

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Starting from a nuclear Hamiltonian based on realistic interactions, we can in principle calculate all of the properties of nuclei. But sometimes the theoretical expectation is contradicted by the empirical evidence. We collect below a number of such puzzles. Several of these have already survived a decade of scrutiny and theoretical effort to understand them.

7.1 The Coulomb Energy Problem

In 1969 NOLEN and SCHIFFER [1,2] surprised us by noting that there was about an 8% discrepancy in the calculation of the mass difference of mirror nuclei, e.g. ^{41}Ca - ^{41}Sc, or of the analog-parent differences, e.g. ^{208}Bi - ^{208}Pb. The experimental mass difference of ^{41}Sc and ^{41}Ca is 7.28 MeV. A careful Hartree-Fock calculation by NEGELE [3] yielded only 6.70 MeV, leaving a .58 MeV discrepancy. This theoretical value included direct Coulomb (6.97 keV), exchange Coulomb, finite proton size, spin orbit and vacuum polarization. Since then there has been considerable controversy about the origin of this anomaly, with numerous mechanisms advocated for the resolution of the problem. These include:

1) Smaller Valence Orbits. NOLEN and SCHIFFER [1,2] suggested that the radii of the valence orbits (or neutron excess) should be smaller than the values obtained with the usual Wood-Saxon or Hartree-Fock one-body potentials. Recent elastic magnetic scattering experiments by SICK and others [5] seem to support this conclusion. However, the inclusion of exchange currents in the magnetic scattering calculations gives the same effect on the cross sections as using smaller valence orbits [6]. On the theoretical side, attempts have been made to justify smaller orbits using the velocity- and energy-dependence of the one-body potential, but these have not been successful.

2) Core Polarization. AUERBACH, KAHANA and WENESER [4] suggested that the valence neutron polarize the core, increasing the proton core radius relative to the core neutron radius. This would have exactly the same effect on the Coulomb energy difference as a smaller valence orbit radius. Hartree-Fock theories based on density-dependent inter-

actions give an effect of the correct sign. The magnitude of the effect depends on the interaction that is used. Using the usual range of Skyrme interactions, one can account for 0-250 keV of the 580 keV discrepancy in $^{41}Ca - {}^{41}Sc$.

3) Charge-Symmetry Violating Interaction. It has been proposed that there is a short range charge symmetry breaking interaction [7,8]. This is supported by the fact that the mass difference of $^{3}He - {}^{3}H$ cannot be accounted for by the Coulomb interaction, indeed the discrepancy is about 100 keV. This suggestion is probably correct and would help to resolve the discrepancy in heavy nuclei. However, the microscopic origin of this interaction has not yet been found, within conventional meson-exchange physics. Indeed, the origin of the interaction may be buried in whatever hadronic processes that give the neutron-proton mass difference.

It would help resolve this discussion to measure the difference in proton and neutron radii. For the nucleus ^{48}Ca, the charge radius has been compared with the matter radius deduced from 1 GeV elastic proton scattering [9-11]. The errors are large enough to make a firm conclusion impossible, but it seems that the Hartree-Fock theory gives a reasonable account of the radius difference between neutrons and protons. We should also mention here a pion scattering experiment [12], which indicates a smaller radius difference than Hartree-Fock theory predicts.

A related question to the Nolen-Schiffer anomaly is the isospin mixing between nearby nuclear states [13-17], for example the two 1^{+} states in ^{12}C at 12.71 MeV and 15.11 MeV. The observed isospin mixing is too large to be explained with a Coulomb interaction within the p^{n} shell configurations. However, the levels are nearly unbound and the Coulomb distortion of the single-particle wavefunctions can produce enough isospin mixing to explain the data [18-20].

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7.2 Spin-Orbit Interaction

The problem of accounting for the one-body spin-orbit interaction in nuclei has been a long and outstanding one. The importance of the spin-orbit interaction was realized in 1949 by MAYER [1] and JENSEN [2] who showed that it played a key role in the shell structure of nuclei. The spin-orbit interaction was also incorporated with the optical potential description of nucleon-nucleus scattering where it gives rise to polarization phenomena. Despite its successes in phenomenological descriptions of nuclei, the origin of this interaction based on the underlying two nucleon potential embedded in a many-body framework still remains a puzzle. Some recent attempts at understanding its origin will now be given.

First, the average one-body spin-orbit interaction has a simple phenomenological form,

$$V_{LS}(r) = V_{LS}(\vec{\ell}\cdot\vec{s})\ \frac{1}{r}\frac{d}{dr}\ \rho(r). \tag{1}$$

The $\rho(r)$ is the nuclear density. The average behavior of Eq. (1) is to give rise to splitting between two different j states of a given ℓ which is $\Delta\epsilon_{\ell s} = -20(\vec{\ell}\cdot\vec{s})$ MeV$^{2/3}$, with the higher j-state, $j_+ = \ell + 1/2$, lying lower in energy than the lower j-state, $j_- = \ell - 1/2$. In the following discussions, calculations will be considered for two types of closed shell nuclei. One type is called spin-saturated and these nuclei have both j_+, j_- levels fully occupied or fully unoccupied. Examples of spin-saturated nuclei are ^{16}O and ^{40}Ca. On the other hand, a spin-unsaturated nucleus is a system with one level, j_+, fully occupied and the other, j_-, empty. Examples of spin-unsaturated nuclei are ^{48}Ca and ^{208}Pb.

The first calculations of the spin-orbit interaction based on the Brueckner theory of the effective interaction were made by WONG [3]. His spin-orbit splitting was too weak for spin-saturated nuclei, but even worse, it had the wrong sign for spin-unsaturated nuclei. More recently this has been studied by SCHEERBAUM [4]. For spin-saturated nuclei, Scheerbaum finds that both the two-body spin-orbit force and the two-body tensor force give substantial contributions, and the total agrees with experiment in mass 15. This conclusion also follows from simpler considerations based on the nucleon-nucleon scattering phase shifts directly [5]. Thus the spin-saturated nuclei seem to be well understood, except for one relatively minor point. This is the relative splitting of particle and hole orbits at the closed shell. Theoretically, the particle orbit has the larger splitting, due to the higher ℓ. Experimentally, it is the other way around near mass ^{16}O: the $p_{3/2} - p_{1/2}$ splitting is 6.3 MeV, while the $d_{5/2} - d_{3/2}$ splitting is 5.4 MeV.

In spin-unsaturated nuclei, the very serious problem persists with the spin-orbit splitting predicted to be negative [6]. The exchange parts of the central and tensor interactions give rise to this effect in Scheerbaum's calculation. This fallacious prediction of nuclear theory has been confirmed by GOODMAN and BORYSOWICZ [7], who calculate the spin-orbit splitting of the $h_{11/2} - h_{9/2}$ orbits as a function of mass in nuclei near ^{208}Pb. The filling of spin-saturated shells increases this splitting, while the filling of the spin-unsaturated $i_{13/2}$ decreses the splitting. GOODMAN [8] finds support in the experimental data for the predicted dependence on mass, but not of course for the overall magnitude of the splitting.

No obvious suggestions for resolving this discrepancy come to mind. It would be interesting to see what conclusions would follow from more modern interactions, which have weak tensor force, such as the Paris potential [9].

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7.3 Coriolis Interaction

The structure of deformed nuclei is understood in terms of independent particle motion in a rotating frame. The description of a particle(s) plus rotor requires in the Hamiltonian the Coriolis interaction,

$$H_c = -\frac{\hbar^2}{2\Theta}(j_+I_- + j_-I_+)$$

where Θ is the moment of inertia of the even-even core. This interaction makes its presence felt in the spectra, by the decoupling of the $K = \frac{1}{2}$ bands, and as was seen in Faessler's lectures, by the backbending phenomenon. Other aspects of the structure are not so easy to reproduce with the above H_c. As BOHR and MOTTELSON [1] discuss, H_c couples bands with $\Delta K = 1$, introducing to the quadrupole transition amplitudes a term proportional to the quadrupole moment of the core. However, the empirical quadrupole matrix elements are not so large; agreement can only be obtained if H_c is attenuated by a factor of two [2].

There have been several suggestions for resolving this discrepancy. The pairing correlations will reduce the K admixture to the wavefunction. However, a careful analysis by HAMAMOTO [3] indicates that these correlations are not strong enough to reduce H_c to the empirical value.

A likely resolution of this problem was given by RING, et al. [4] who note that these difficulties with the structure do not appear in the Hartree-Fock-Bogoliubov cranking model, which treats all particles on the same footing. Good spectral fits are obtained in situations which require the attenuated H_c in the particle-rotor model [5]. The apparent fault of H_c in the usual rotor model is that the contribution of the odd particle to the angular momentum is neglected. Apparently much of the total angular momentum is due to the last particle. KREINER [6] has formulated a correction associated with this effect; the attenuation factor in H_c is shown to be

$$\left[1 - \frac{\langle j_x^2 \rangle}{\langle j_x I_x \rangle}\right].$$

The contribution of the particle to the angular momentum, $\langle j_x^2 \rangle$, turns out to be large when evaluated numerically in the cranking model.

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7.4 Missing M1 Strength in Heavy Nuclei

The M1 strength is a fundamental property of nuclei that shows, among other things, how the spin symmetry is broken in the ground state. It is not difficult to make a plausible theoretical model for the M1 strength in a nucleus such as ^{208}Pb. Using an unperturbed shell model ground state of ^{208}Pb, the giant M1 strength would reside in a linear combination of two states $[h^{\pi}_{9/2}h^{\pi}_{11/2}{}^{-1}]^{1+}$ and $[i^{\nu}_{11/2}i^{\nu}_{13/2}{}^{-1}]^{1+}$. The particle-hole splitting of these two states, as taken from experiment, are respectively 5.36 and 5.86 MeV. The residual interaction should push both states up in energy, and mix them so that the isovector combination lies higher. Most of the M1 strength would reside in the higher state because the magnetic moment of the neutron and proton are of opposite sign. The magnitude of the residual interaction is <1.5 MeV, as calculated in the Brueckner theory. Thus, all of the M1 strength should lie below 8 MeV excitation.

Experimentally, the search for the M1 strength has been a confusing process of elimination, with supposed M1 states turning out to be E1. By 1978, the only strong state left was one at 7.99 MeV, having 25% of the shell-model strength [1]. Now even the M1 character of this state is doubted [2], leaving no important M1 strength below 9 MeV [1-4]. The state would not be observed if it lies above this energy and is highly fragmented.

That the M1 state should be strongly fragmented is no surprise. It has been known for over a decade that several 2 particle-2 hole 1^+ states can come below the 1^+ isovector state in a shell model calculation [5-7]. Refs. [5] and [6] consider a related problem, the hindrance of Gamow-Teller β decays. For example the hindrance of the decay of the ground state of ^{56}Ni to the 1^+ state in ^{56}Co at 1.72 MeV is due to the fact that this lowest 1^+ $T=1$ state in ^{56}Co has very little amplitude (0.016) of the configuration $f^{\nu}_{f/2}f^{\pi}_{7/2}{}^{-1}$, but is predominantly a 2 particle-2 hole state. The work of LEE and PITTEL [7] is directly concerned with ^{208}Pb.

The two links in the theoretical argument that the strength lies below 8 MeV are the single-particle energies and the residual interaction. The residual interaction is probably not at fault; the usual ideas on the nuclear force are adequate for describing the M1 strength in ^{12}C [8], and for the related topic of the $\sigma\tau_-$ strength in ^{90}Zr [9]. The treatment of single-particle energies is called into question by BROWN and SPETH [10]. The mistake, in their view, is that the calculations were made using experimental single-particle energies, rather than Hartree-Fock energies. Since the Brueckner-Hartree-Fock theory gives an effective mass of about 0.7, and the empirical single particle energies are described with an effective mass of about 1.0, Brown and Speth argue that the single particle energies should be raised by 1/0.7. However, for the M1 states, a larger single-particle splitting would require a stronger spin-orbit potential. We have already seen that theory would like to reduce this quantity. It would be interesting in pursuing this problem, to compare the strengths of the spin-orbit potential as deduced from elastic scattering, and as deduced from the single-particle splittings. Another difficulty with this viewpoint is the necessity of reconciling all of the giant vibration data on the same footing. While the isovector excitations call for a reduced effective mass [11,12], the giant quadrupole energy is perfectly consistent with an effective mass of 1.

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7.5 Nuclear Level Densities

Beyond a few MeV excitation, we cannot hope to understand the structure of individual states, unless some quantum number, such as isospin or angular momentum singles them out from their neighbors. In fact, we have seen in the lectures of French that we should not expect such states to show individuality. Thus, the only property left to describe is their numbers. The theory of the level density of a Fermi gas gives the following formula [1,2]

$$\rho(N,Z,E) = \frac{6^{\frac{1}{4}}}{12} \frac{g_0}{(g_0 E)^{5/4}} \exp[2(aE)^{\frac{1}{2}}]$$

Here E is the excitation energy, and the single-particle level density is

$$g_0 \approx \frac{2}{\pi^2} mk_F V \approx \frac{3}{2}\frac{A}{\varepsilon_F}$$

This is related to the conventional level density parameter a by $a = \frac{\pi^2}{6} g_0$. The theoretical a is

$$a \approx \frac{3}{2}\frac{A}{\varepsilon_F}\frac{\pi^2}{6} \approx \frac{A}{15}.$$

Experimentally, the level density is higher. The neutron resonance data is best fit with

$$a \approx \frac{A}{8},$$

twice as large as theory. This factor of two discrepancy is the problem to be resolved.

BOHR and MOTTELSON [3] suggest that the harmonic oscillator level density be used for g_0. Since the valence particles occupy a relatively larger volume in the harmonic oscillator model, this would increase a. However, the concomitant prediction of the oscillator model, that core nucleons occupy a smaller volume, is easily seen to be incorrect by examining density distributions.

Other possible effects which could influence a are:

--effective mass corrections to the single-particle spectrum. This is certainly important in the description of another Fermi system, liquid ^{3}He. But as we noted previously, the empirical effective mass for nucleons is close to 1 at the Fermi surface.

--the contribution of collective states to level densities. SOLOVIEV, et al. [4] have computed level densities in a model assuming that particle-hole phonons make up the elementary excitations. The result is an increase in level density in certain regions. In particular, the presence of 2^+ and 3^- phonon states can increase the level density by a factor of 10. One difficulty of this description is that it does not respect the Pauli Principle, since the Fermion degrees of freedom are replaced by bosons. We therefore expect that it would overestimate the density of levels.

The problem of level densities appears in a particularly acute form at closed shell nuclei. Let us examine that favorite example, ^{208}Pb. Experimentally, the density of 1^- states at 7-8 MeV excitation is 10/MeV [5,6]. The empirical particle-hole energy gap in ^{208}Pb is 4.2 MeV for protons and 3.4 MeV for neutrons. Thus, in the independent particle description, states at 7-8 MeV excitation will be no more complicated than 2 particle-2 hole, and in fact none of them will have quantum numbers 1^-. To include the residual interactions via the phonon description, we note that the lowest phonon in ^{208}Pb is at 2.7 MeV. It is just possible to make odd parity states in the 7-8 MeV range with two or three phonons. However, the number predicted is an order of magnitude lower than observed.

This problem is also seen at lower excitations, with unanticipated intruder states mingling with the expected shell model states. These are typically strongly deformed states, and it is possible that the increased level density at higher energy is associated with deformation degrees of freedom. The presence of rotational bands in nuclei whose ground state is deformed of course will increase the level density [7]. More interesting, the level density of spherical nuclei can be substantially increased by the presence of deformed states at high excitation [8].

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7.6 Spin Modes

It is only recently that the spin degrees of freedom of the nucleus could be studied as carefully as the charge and density. This has been brought about by better electron scattering experiments, and the availability of intermediate energy protons and pions as nuclear probes.

On the theoretical side, our ignorance is considerable, as we saw in the discussion of the M1 strength. In the words of MOTTELSON [1],

> "At the present time we are only beginning to gain some understanding of the collective modes associated with spin-dependent fields. We are in fact overwhelmed by questions which we cannot properly answer; such as, do the pair of modes $[(\sigma Y_{\lambda-1})^{\lambda}$ and $(\sigma Y_{\lambda+1})^{\lambda}]$ in fact mix strongly? What are the strengths (even the signs) of the collective potentials associated with deformations having these spin dependent structures? Are there incipient instabilities associated with some of these models? (e.g. pion condensates should appear in an analysis of $\tau = 0$, $\kappa = 1$, $\lambda = 0$ modes) etc."

The shell model has particular difficulties with spin densities at high momentum transfer. One interesting example is the elastic scattering of electrons from ^{17}O, measuring the magnetic density [2]. In the shell model this nucleus consists of a $d_{5/2}$ nucleon plus an ^{16}O core. The magnetic moment of ^{17}O is very close to the Schmidt moment and in the early days, this was used as evidence for the validity of the shell model. However, the spin density deviates significantly from the shell model for momentum transfers in the range $q = 1 - 3\ fm^{-1}$. At about $q = 1$ there is a big dip in the experimental cross section, which is called M3 suppression. This can probably be explained by core polarization [3]. Roughly speaking, if we consider the operator $Q_z\sigma_z$ (part of the M3 operator at low momentum transfer), then its expectation value in the ^{16}O core is $(Q_\uparrow - Q_\downarrow)$ where $Q_\uparrow$ is the quadrupole moment of the spin up particles etc. Thus to explain the suppression it is necessary that spin up nucleons have a different quadrupole moment than spin down particles. The M5 region is quite confusing; on the low momentum transfer side the experimental form factor is smaller than the shell model value, on the high momentum transfer side it is larger.

Shell model calculations in the $(0p)^8$ basis predict a smaller cross section at this momentum transfer by an order of magnitude [7]. This is quite surprising, since the low momentum transfer, which relates to the magnetic moment, is quite well fit by such calculations [6]. DUBACH and HAXTON [7] show that a transition density can be constructed within the $(0p)^8$ basis that fits the data. Whether this transition density can be derived from a reasonable Hamiltonian remains to be seen.

There are two 'explanations' of the second peaks which may not be as different as they sound. One is a straight core polarization calculation by SUZUKI, et al. [8], in which it is shown that it is essential to use a tensor interaction. In their calculation, exchange currents are included but do not seem to play an important role.

A more dramatic explanation is afforded by M. ERICSON [9]. She cites the experiment as possible evidence for a precursor to pion condensation. The summed strength of all 1^+ excitations in ^{12}C is equal to the ground state expectation value of the square of the M1(q) operator, and hence is a measure of the fluctuation of this operator. The ^{12}C nucleus in this picture fluctuates between the normal state and the pion condensate state in which the nucleus spins are aligned as in an anti-ferromagnetic, presumably with a characteristic distance d separating the spin up and spin down nucleus. The peak in the inelastic scattering at $q = 2\ fm^{-1}$ may mean that d is of the order of 1/q. Ericson sees a strong analogy between the secondary peak and critical opalescence.

It has been suggested by TOKI and WEISE [10] and FAYANS, et al. [11] that perhaps proton inelastic scattering is the more relevant probe of critical opalescence rather than electron scattering. The Fourier transform of the tensor interaction is proportional to $\sigma_1\cdot q\ \sigma_2\cdot q$. This interaction may be responsible for spin coherence. The electromagnetic scattering goes as $\vec{\sigma} \times \vec{q} \cdot \varepsilon$. It should be noted that a backward peak is seen in proton inelastic scattering to both the $T = 1, 1^+$ state at 15.11 MeV and the $T = 0, 1^+$ state at 12.7 MeV [12].

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