NUCLEAR STRUCTURE -- THEORY

SIZES AND REACTION CROSS SECTIONS OF LIGHT NUCLEI

G.F. Bertsch and B.A. Brown

Reaction cross sections can now be measured for unstable nuclei using high energy heavy ion beams.[1-3] A study of p-shell isotopes found ^{11}Li to have a much larger cross section than expected from systematics.[1]

We are trying to understand this behavior theoretically, determining the cross sections from Hartree-Fock densities and the eikonal or Glauber model for the reaction. We calculated the Hartree-Fock densities using the SGII interaction of Sagawa and van Giai[4] as described in Ref. 5. The predicted r.m.s. matter radii of Be isotopes are shown in Table I. The reaction

Table I. Calculated r.m.s. matter radii for the Be isotopes. Point nucleon densities are folded with a finite nucleon size of 0.64 fm^2.

A	J^π	r.m.s. radius (fm)
7	$3/2^-$	2.37
8	0^+	2.40
9	$3/2^-$	2.43
10	0^+	2.47
11	$1/2^-$	2.58
	$1/2^+$	2.94

cross section is calculated from the formula

$$\sigma_R = 2\pi \int bdb[1 - \exp(\sigma \int d^2r\ \rho_z(r)\rho_z(|\vec{b}+\vec{r}|)) \quad (1)$$

Here σ=40 mb is the nucleon-nucleon cross section. The densities ρ_z are the mass densities of the target and projectile integrated over the z coordinate,

$$\rho_z(r) = \int dz\ \rho\ (\sqrt{r^2+z^2})$$

Eq. (1) is an approximation to the Glauber model, but it is entirely adequate for our purposes.

The predicted reaction cross section for Be isotopes on a ^{12}C target is shown in Fig. 1. The cross section varies very smoothly with mass, assuming pure p-shell configurations, as shown by the line connecting the points. This agrees with the findings of Ref. 6. When one nucleon is excited to the sd shell there is an increase in the cross section as shown by the arrow. For ^{11}Be this is in fact the ground state configuration. The radial wave functions of the sd-shell orbitals were constrained to reproduce the experimental neutron separation energy of 0.5 MeV. The trend of the calculated cross sections for the Be isotopes is close to that found experimentally.

We have not carried out specific calculations for Li isotopes, but the large cross section implies that at least two neutrons must be promoted to the sd shell.

Another possible explanation is that the ^{11}Li nucleus has a low-lying unstable state which is easily excited by the Coulomb field of the target.[7] However, this possibility is readily dismissed. According to Ref. 8, the Coulomb excitation of the giant dipole resonance in a light nucleus by the Coulomb field of a heavy target is about 40 mb. One gains a factor of the logarithm of the relative excitation energies when exciting a low-lying state, but this is more than offset by the dependence on the target charge, which is quadratic. Coulomb excitation by a carbon target should be completely negligible. However, it could be tested by varying the target Z.

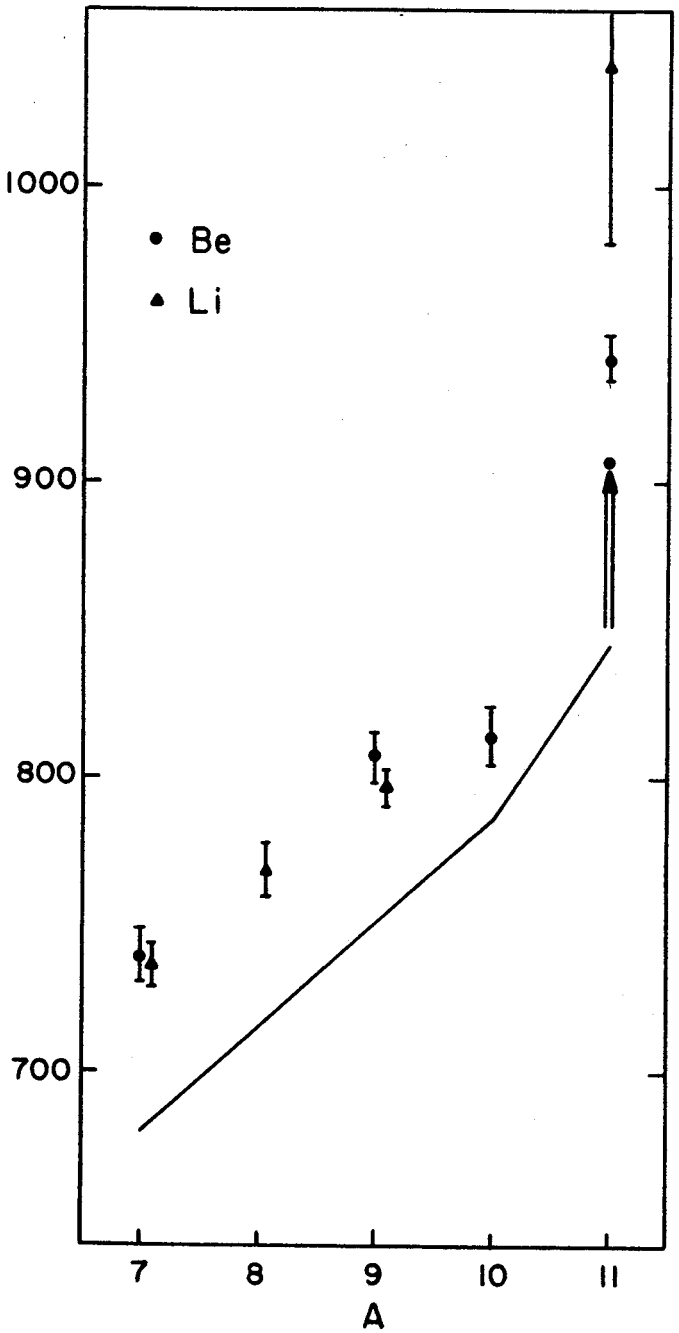

Fig. 1 Reaction cross sections for light nuclei on a carbon target. The solid line shows the theoretical cross sections for Be isotopes. The arrow shows the increase in cross section for ^{11}Be assuming a configuration with one neutron promoted to the sd shell. The data points are from Refs. 1 and 2.

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CALCULATIONS OF THE PARITY VIOLATING ALPHA WIDTH OF THE 0^- STATE AT 6.88 MeV IN ^{18}O

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A measurement of the parity nonconserving (PNC) alpha-decay width of the 0^- state in ^{18}O at 6.88 MeV is being planned, in order to study PNC aspects of nonleptonic weak interactions. Current models predict a PNC affect in ^{18}F which is several times larger than the current experimental limit.[1] The present case, may shed light on this failure because in both cases the 0^- states involved have a fairly pure $1s_{1/2}$-$0p_{1/2}$ structure. The present case also will supplement the information obtained in ^{19}F.[1]

The PNC effects in the ^{18}F and ^{19}F cases are enhanced and relatively simple to interpret due to the near degeneracy of the states involved. In contrast, the PNC alpha-decay of the 2^- state in ^{16}O[2] is more difficult to interpret because the mixing presumably comes from several 2^+ states whose structures are not well established. The present case in ^{18}O is intermediate to these two situations.

We have initiated a study of the nuclear structure aspects of the ^{18}O PNC alpha decay, and in this report we summarize the present status of the results. The calculations employ the methods described in Ref. 3. In particular, the calculations are based on two-body PNC matrix elements obtained from the PNC Hamiltonian, H(PNC), of DDH,[4] together with the short range correlation function described in Ref. 3. The PNC reduced width is obtained in first-order perturbation theory from a summation over the allowed widths:

$$\gamma(0^-,\mathrm{PNC}) = \sum_n \langle 0_n^+|H(\mathrm{PNC})|0^-(6.88)\rangle\, \gamma(0_n^+)/\Delta E_n$$

where ΔE_n is the energy denominator $Ex(0_n^+)$-6.88 MeV. The allowed alpha widths were calculated by taking the overlap of the 0^+ states in ^{18}O with the ground state of ^{14}C together with a four-particle wave function in which the four particles are assumed to be in alpha-cluster configuration. The total width was then obtained by considering the ratio to the second 4^+ state at 7.12 MeV excitation:

$$\Gamma(0^-,\mathrm{PNC}) = [P_0(6.88)/P_4(7.12)] \times \times [\gamma(0^-,\mathrm{PNC})/\gamma(4^+,7.12)]\ \Gamma(4^+,7.12)$$

We use the experimental value of $\Gamma(4^+,7.12)$ = 0.138 eV. This relation is used because the ratios of the penetration factors P_0 and P_4 and the ratio of the reduced widths should be less sensitive to details of the calculations than the individual quantities themselves. The ratio of the penetration factors is 37 with a channel radius of 5.4 fm.

The initial calculations have been carried out in the full $0p_{1/2}$-$0d_{5/2}$-$1s_{1/2}$ (ZBM) model space. Relative to a closed shell for ^{16}O, this model space includes some subset of the 1 and 3 ħw excitations for the 0^- states and the 0, 2 and 4 ħw excitations for the 0^+ states. The wave functions were obtained from the strong interactions labelled "F" and "Z" in Ref. 3. The alpha-cluster was obtained by using an SU3 conserving interaction projected onto the sd part of the ZBM model space.

The PNC alpha width is dominated by a constructive addition of the terms in the above summation from the 0^+(gs) ground state (predominantly 2p in structure) and the $0^+(3.63)$ excited state (predominantly 4p-2h in structure). Other states are less important but tend to give a small destructive contribution. The results are similar for the two interactions especially for the lower two states. The PNC width obtained from the above procedure is 0.4×10^{-11} eV with the F interaction and 0.8×10^{-11}

eV with the Z interaction. In terms of the coupling constants defined in Ref. 1, the amplitude of the PNC alpha decay is proportional to the linear combination $1.3\ F_{\pi} + 0.6\ F_{0}$. This is essentially the same "proton-type" linear combination which enters into the ^{19}F case. We note that the calculated PNC alpha widths for ^{18}O are about an order of magnitude smaller than the measured width of $(1.03 \pm 0.28)\times10^{-10}$ in ^{16}O.[2] The ^{18}O experiment is thus more difficult but we hope still feasible.

It will be important to consider the effects of expanding the model space to include the $0p_{3/2}$ and $0d_{3/2}$ orbitals. From Ref. 3 it is clear that these will probably have the effect of reducing the PNC matrix elements, perhaps by as much as a factor of 2-3. They are also important in the calculation of the allowed 0^{+} alpha widths. We should also take into account the $(p)^{2}(sd)^{2}$ component of the alpha-cluster transfer in addition to the $(sd)^{4}$ component considered in the above results. Also a calculation of the M1 width of the 0^{-} state is important. It will be important to test the calculations against ancillary experimental results in A=18, such as the allowed alpha widths of the low-lying states from direct transfer reactions, in particular the ratio of the reduced widths for the $4^{+}(7.12)$ and the $0^{+}(3.63)$ states, and the M1 strengths for gamma transitions between negative parity states.

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ISOSPIN MIXING CORRECTIONS TO SUPERALLOWED BETA DECAY

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Superallowed Fermi beta-decay has been the subject of intense study for several decades. The principal interest in these transitions is that, according to the conserved vector-current (CVC) hypothesis, their ft values should be constant for all nuclei. Two classes of nucleus-dependent corrections, however, must considered. The first is radiative corrections to the statistical rate function denoted by δ_R. The second is corrections to the nuclear matrix element due to the presence of isospin-nonconserving (INC) forces in nuclei, and is denoted by δ_C. With these correction the "nucleus-independent" ft value for Fermi transitions is defined by $Ft = (1+\delta_R)(1-\delta_C)\,ft$. From these Ft values, it is then possible to extract empirical values for the weak interaction vector coupling constant G_V. By comparing the vector coupling constant for nucleon beta-decay to that of muon beta-decay the Kobayashi-Maskawa mixing angle between u and d quarks v_{ud} can be determined[1] and a test of the 3-generation standard model of the electro-weak interaction is possible.[2]

Until recently, the calculated corrections δ_R and δ_C for the eight most accurately measured ft-values failed to yield constant Ft-values.[3] Part of this discrepancy has been removed by a recent reevaluation of δ_R.[4] We have recently re-evaluated the nuclear overlap correction δ_C and have made a number of improvements in various aspects of the calculations. In Ref. 5 we reported on new calculations for overlap between the proton and neutron involved in the transition. These calculations involved a calculation for the isospin mixing between nearby states (giving a correction denoted by δ_{IM}) and a calculation for the radial overlap (giving a correction denoted by δ_{RO}). Here we report on a new calculation of the overlap between the spectator protons and neutrons (giving a correction denoted by δ_S), present a summary of the results for the total correction ($\delta_C = \delta_{IM}+\delta_{RO}+\delta_S$), and discuss the implications.

The radial wave functions were obtained from a Hartree-Fock calculation utilizing Skyrme-type interactions. Five different Skyrme forces were used to evaluate δ_{RO} and δ_S in order to investigate the sensitivity to the parameterization of the mean field. Shown in Fig. 1 are the results obtained with the SGII interaction.[6] The values obtained with the SGII are about equal to the average of those obtained with all five Skyrme forces. The variation due

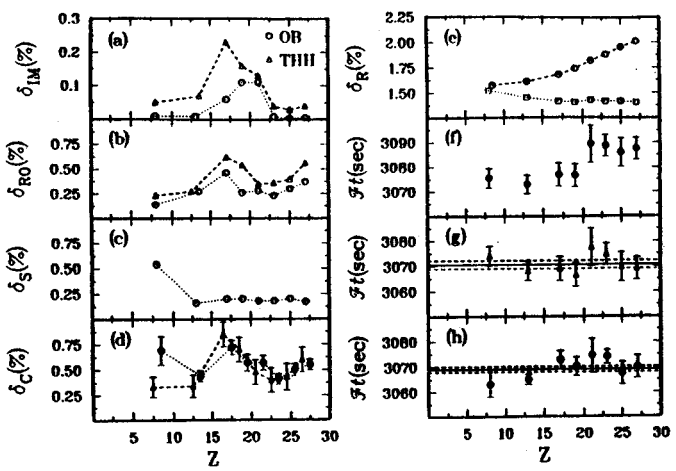

Fig. 1 (a-d) Comparison between THH (triangles) and present (open circles) corrections to the Fermi matrix element: (a) δ_{IM}, (b) δ_{RO}, (c) δ_S, (d) the total correction δ_C including theoretical uncertainties. (e-h) Also shown are the new[10] (squares) and older[7] (circles) values for the radiative corrections δ_R (e), Ft-values obtained with the older δ_R and the THH δ_C (f), new δ_R and the THH δ_C (g), and the new δ_R and the present values for δ_C (h).

to the dependence on the Skyrme parameterization (about 0.07%) is included in the uncertainty of the total correction δ_C shown in the figure.

Our results are compared in the figure to the previous work of Towner, Hardy and Harvey (THH).[7] THH assume a more conservative theoretical error of about 0.12% for δ_C. As can be seen from the figure our values for δ_{RO} and δ_{IM} are smaller than the THH values. However, when the correction δ_S is included which was ignored by THH, the two sets of δ_C are nearly identical. For the purpose of comparison, we have also evaluated δ_S using Woods-Saxon wave functions, and found these values to be much larger than their HF counterparts; δ_S(WS) = 2.4% and 0.42% for ^{14}O and ^{54}Co, respectively.

The corresponding "nucleus-independent" Ft-values obtained from our calculations are compared with the THH results on the right side of the figure. Both sets yield fairly constant Ft-values, with the THH set being slightly better; the THH values (without the WS δ_S correction) give $\langle Ft \rangle = 3070.6 \pm 1.6$ sec compared to $\langle Ft \rangle = 3069.0 \pm 1.1$ sec from our new calculations. If we increase our theoretical uncertainty to 0.12% we obtain $\langle Ft \rangle = 3069.3 \pm 1.6$ sec. With this value the Kobayashi-Maskawa mixing matrix element v_{ud} is then 0.9746 ± 0.00010. Taking $v_{us} = 0.220 \pm 0.002$[8] and $v_{ub} < 0.0075$ (90% confidence level)[9] the unitarity condition gives $v_{ud}^2 + v_{us}^2 + v_{ub}^2 = 0.9985 \pm 0.0026$, which is in good agreement with the 3-generation standard model.

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COMPARISON OF GAMOW-TELLER STRENGTH FUNCTIONS IN THE SHELL MODEL AND IN QRPA

G.F. Bertsch, B.A. Brown, and B. Lauritzen

The Gamow-Teller transitions to states of higher isospin, notably the β^+ transitions in N>Z nuclei, are of interest in their own right and also because of their role as virtual transitions in double beta decay. While the giant Gamow-Teller resonance in the β^- strength function is well described by RPA theory, the strength for β^+ decay is understood much more poorly. The transitions in this direction are severely suppressed with respect to the single-particle values from the independent-particle shell model. Only a portion of that suppression can be explained by the presence of the giant Gamow-Teller resonance. Away from closed shells, when the ground state is taken as the BCS-state, the Gamow-Teller strength is obtained in the quasi-particle random phase approximation (QRPA), but it too failed to reproduce the needed suppression when particle interaction parameters were fixed by external considerations.[1,2]

However, the suppression can apparently be enhanced by using a different particle interaction.[3] On the other hand, calculations in the sd shell successfully reproduce the empirical suppression without any modifications to the interaction.[4] Finally, a recent QRPA study claims that sufficient suppression can also be obtained in heavy nuclei using a realistic interaction.[5] These findings raise the following questions.

(1) To what extent does QRPA reproduce the Gamow-Teller strengths found in shell model calculations?

(2) How sensitive are the shell-model predictions to the details of the interaction?

(3) Can shortcomings of the QRPA be overcome by changing the interactions?

In this report we focus on item (2) above; the other questions are currently under study.

In the shell model the particle-particle (pp) and particle-hole (ph) matrix elements are related by angular momentum recoupling, while in QRPA, pp and ph interactions are renormalized independently. The ph strength determines the location of the giant Gamow-Teller resonance, while the β^+ strength to low lying states is sensitive to the pp strength. It has been argued that pp strength should be renormalized due to the schematic forces applied and due the truncated bases.[1]

In the sd-shell there are 63 independent two-body matrix elements of a rotational invariant and charge independent interaction. By carrying out a least square fit to ~450 experimental binding energies in the shell the energies are reproduced with an error of ~150 keV. When the number of independent variables is decreased, fixing the remaining variables by G matrix elements or setting them equal to zero, the error is only increased slightly; using a 14 parameter model interaction the rms deviation from the data set is ~250 keV.[6] Evaluation of the β^+ decay matrix elements may be accomplished within this model. Following the results of QRPA calculations, the GT strength to low lying states in the shell model calculation should be sensitive to the J=1 two body matrix elements.

We therefore wish to map the dependence of the B(GT) strength on the shell model interaction, especially on the strength in the various ST channels and investigate whether the required renormalization of pp, ph strengths in QRPA can be understood in the shell model picture.

Examples of the B(GT) strength obtained in shell model calculations are shown in Figs. 1-2. In Fig. 1 the integrated Gamow-Teller β^+ strength $^{24}Ne \rightarrow {}^{24}F$ is shown for three different parameterizations of the interaction. In all cases the interaction has a central, a spin-

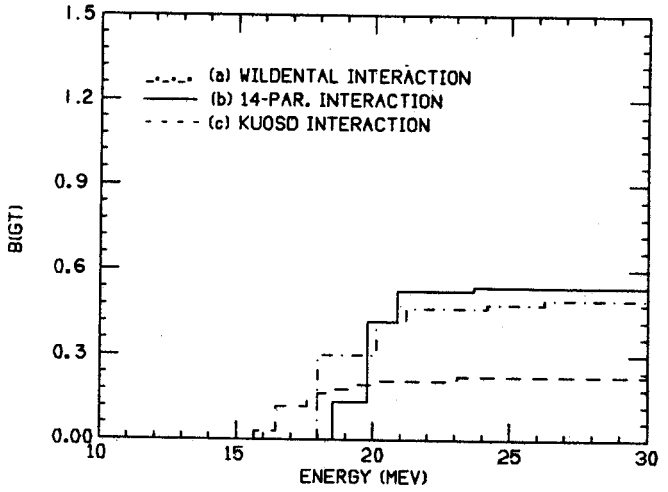

Fig. 1 Integrated B(GT) strength, $^{24}Ne \rightarrow {}^{24}F$

orbit and a tensor part. The curve (a) represents an interaction-model independent parameterization, (b) the 14-parameter interaction and (c) a G-matrix calculation based on the Hamada-Johnston interaction.[7] The total β^+ strength is reduced below the independent particle model by a factor ~6 for interactions (a) and (b), and a factor ~12 for (c).

In Fig. 2. the parameters of the central, S=1, T=0 component of the 14-parameter interaction is varied. The interaction has a monopole (M) and a medium and long range meson exchange term (OBEP). It is seen that the nuclear structure Gamow-Teller matrix elements are sensitive to changes ±20% in the OBEP but insensitive to similar changes in the monopole term. A combined change in M and in OBEP can lead to very different β^+ strengths without changing the excitation energies.

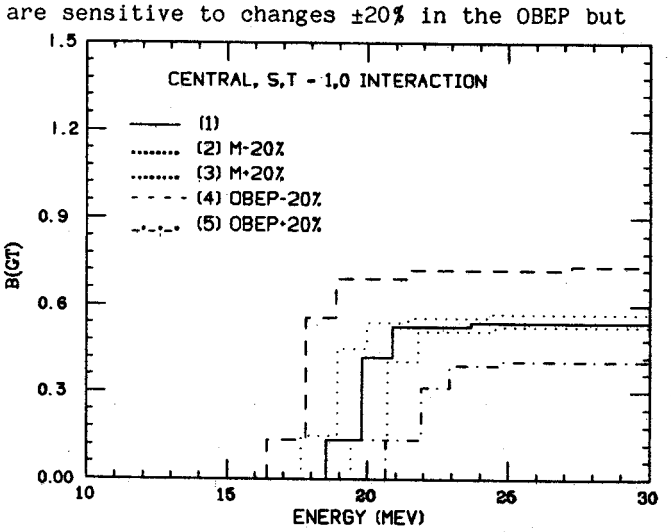

Fig. 2 Integrated B(GT) strength, $^{24}Ne \rightarrow {}^{24}F$

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Parity Splitting in Rotational Bands

G.F. Bertsch

The light radium nuclei have low-lying odd parity bands[1] which may be interpreted as members of a parity-deformed ground state band. We have recently calculated the parity splitting of the ground-state rotational band in the nucleus Ra^{222} [2]

Hartree-Fock and other mean field calculations predict that the ground configuration of Ra^{222} has an intrinsic octupole deformation,[3,4] with the deformation parameter ε_3 in the range 0.07 to 0.10. There are two degenerate minima having opposite signs for ε_3. The band splitting is simply related to the mixing of these configurations: the linear combinations with plus and minus phases generate even and odd angular momenta, respectively.

In principle the mixing can be calculated in many ways, but we favor a simple approach based on the dominance of the pairing interaction.[5] We set up a small dimension Hamiltonian matrix, counting the number of configurations necessary to go from one minimum to the other. The off-diagonal interaction between configurations is expressed in terms of the pairing gap Δ and the pairing interaction strength G as,[6]

$$v = \frac{\Delta_n^2}{2G_n} + \frac{\Delta_p^2}{2G_p} \approx 2.9 \text{ MeV}.$$

For the case of Ra^{222}, there are 6 configuration changes, and the energy barrier for intermediate configurations is about 300 keV. The predicted splitting of the band is 500 keV, to be compared to the experimental splitting of about 220 keV. This disagreement could be lessened by decreasing the pairing gap or by increasing the potential barrier.

The strong dependence of the tunneling on the splitting on the pairing gap leads to the following prediction: at high angular momentum, the gap should be reduced and the two bands should become nearly degenerate.

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NEW EFFECTIVE INTERACTIONS

B. A. Brown, W. A. Richter,[a] R. E. Julies,[b] and L. Zhao

A study of semi-empirical effective interactions for the sd-shell nuclei has recently been completed.[1] This study considered a wide variety of forms and parameterizations which included simple schematic potentials, multipole interactions, Talmi integrals, and one-boson exchange potentials (OBEP). For the potential models the effects of density dependence were investigated. One of the optimal formulations found was a 10-parameter density-dependent OBEP added onto the standard one-pion exchange potential (OPEP) and modified by the addition of four monopole terms for the central component. This 14-parameter potential model (plus three single-particle energies) reproduced 447 sd-shell binding-energy data to within an rms deviation of about 260 keV, which compares favorably to Wildenthal's 47-parameter "model-indepedent" result of 185 keV.[1,2]

Work has been started to extend this study to other mass regions. In this report we briefly summarize results for the p shell, the fp shell, and the combined sd and fp shells. All of these analyses make use of the techniques developed in Ref. 1. In particular, we have investigated the importance of the antisymmetric spin-orbit component, the density dependence, and the mass dependence of the two-body matrix elements.

The p-shell work is an update on the classic work of Cohen and Kurath.[3] The p-shell interactions were optimized to fit 77 known binding-energy data for A=5-16. The optimal fits gave rms deviations of about 600 keV. An independent fit to the data for the He isotopes was used to predict that ^{10}He should be unbound to two-neutron emission by 1.18±0.14 MeV and marginally unbound to one-neutron emission by 0.04±0.14 MeV. This is result is consistent with a recent experimental.[4]

New two-body interactions are derived for nuclei in the lower part of the 0f-1p shell by fitting semi-empirical two-body potential forms as well as two-body matrix elements to 61 binding-energy data for A=41-49. The shell-model calculations assumed a ^{40}Ca core and valence nucleons distributed over the full fp space. Care was taken to exclude intruder states from the selected data set. Rms deviations between fitted and experimental energies of about 150 keV have been achieved with a variation of only 8 (density-dependent) OBEP parameters added onto the standard OPEP plus four monopole terms for the central component. The results obtained with this 12-parameter model are comparable in quality to those obtained in the sd shell.

The success of the particular semi-empirical interaction form used both for the fp and sd shells allowed a simultaneous fit to 447 sd-shell and 61 fp-shell binding-energy data. An overall rms of 200-250 keV for both shells was obtained in a 17 parameter fit where the 10-parameter OBEP was constrained to be the same for both shells and the four monopole terms were shell dependent (eight terms altogether). The relationship between the parameters of the interactions in the two different model spaces is being investigated.

The new fp-shell interactions will be tested in a variety of ways. Results obtained by L. Zhao for the ^{48}Ca-^{48}Sc Gamow-Teller strength distribution are shown in Fig. 1. The individual lines have been broadened by 0.32 MeV in order to simulate the experimental resolution obtained in Ref. 5. Spectrum (a) was obtained by restricting the initial state to have a 0p0h configuration and the final 1^+ states to have 0p0h-1p1h configurations relative to a closed $f_{7/2}$ shell for ^{48}Ca. Spectrum (b) was obtained

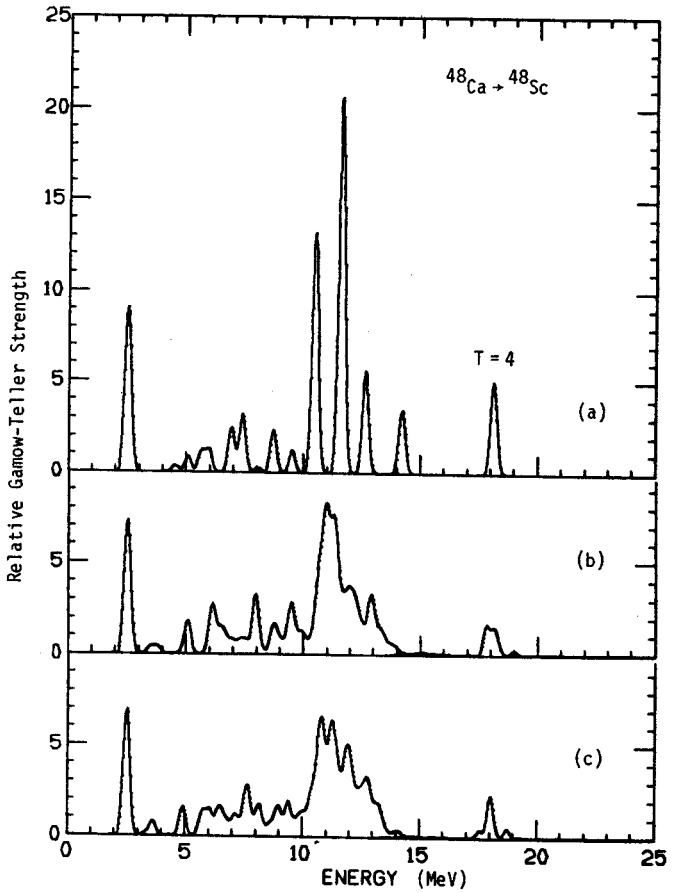

Fig. 1 (see text)

from a 0p0h-2p2h initial state and a 0p0h-3p3h final state. Spectrum (c) includes a selected set of the 0p0h-4p4h configurations for the initial state and 0p0h-5p5h configurations for the final state. Spectrum (c) had a total dimension of 3681 and contains about 500 1^+ states below 25 MeV excitation. The calculated strength distribution of spectrum (c) is in reasonable agreement with the shape of the experimental (p,n) spectrum (Ref. 5). An absolute comparison gives an excitation energy independent quenching factor of 0.6-0.7. We plan to calculate the ^{48}Ti-^{48}Sc Gamow-Teller strength and the double-beta decay matrix elements.

Further extensions of this work are being planned. These include: (1) improvements in the code to allow more nuclei in the fp shell to be considered, (2) fits to the total binding energies (as opposed, for example, to the binding energies relative to ^{16}O considered above for the sd shell) - this would allow the single-particle energies to be treated on the same footing as the interaction and (3) extensions to heavy nuclei up to the few particle and hole states around ^{208}Pb. This last extension may lead to a new universal parameterization of the effective interaction.

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STATIC PATH APPROXIMATION FOR THE NUCLEAR PARTITION FUNCTION

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The study of nuclei at high excitation energy has become an important area of research in heavy ion physics. Aspects of the behavior of hot nuclei, such as the damping of the giant resonances and multiparticle fragmentation, are strongly influenced by fluctuation effects. However, a theory adequate to deal with fluctuations still needs to be formulated.

We have begun to construct a theory capable of dealing with fluctuations, based on the thermodynamic partition function. For a many-body system with separable attractive interactions the partition function may be written as a path integral in the Hubbard-Stratonovich representation,[1]

$$Z = \int D[\sigma]\, e^{S_{eff}[\sigma]}$$

in terms of the effective action $S_{eff}[\sigma]$. The standard stationary phase approximation leads to the Hartree mean field solution. We propose to approximate the many-body partition function by summing static paths, $[\sigma]$=const., in the integral.[2] Since each static path represents a mean field configuration we are summing over certain fluctuations in the mean field. The integration may be done by a Monte Carlo sampling method where the effort involved is equivalent to one Hartree-Fock iteration for each sampled path.

We have applied the static path approximation to a simple two-level model similar to the Lipkin-Meshov-Glick model. The Hamiltonian is written,

$$H = -4\lambda(J_x)^2 - 2dJ_z.$$

where J_x, J_z are the components of the Lipkin quasi-spin operator. The total angular momentum J^2 is a constant of motion. The solutions are readily obtained in a semiclassical approximation from the intersection of H with the sphere, J^2 =const., and resembles the motion of a single nucleon in a cranked oblate deformed nucleus.[3] While the ground state is obtained in the Hartree approximation, the high temperature expansion of the energy in the Hartree approximation is very different from the exact result. The Hartree-Fock solution lacks an energy expansion in the inverse temperature and so is useless for the high temperature behavior of the model. The model contains a phase transition and the standard RPA corrections will break down near the phase transition.

On the other hand, the static path approximation reproduces the exact expansion of the energy to first order in the inverse temperature β. Corrections to the approximation to second order in β may be evaluated by Fourier expanding the integration paths σ. This is readily done in the model and is seen to give only a small contribution to the energy. Also, the static path approximation remains useful in the phase transition region.

The superiority of the static path approximation to ordinary finite temperature Hartree-Fock theory was also noted in Ref. 4.

References

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The charge-exchange reactions (p,n)[1], and (^{3}He,t)[2], and (d,pp)[3] at relativistic energies are seen to excite the Δ resonance in the target. There is a large shift in the energy of the Δ peak: for example, in (^{3}He,t), it is lowered from 325 MeV to 255 MeV when target is changed from ^{1}H to the heavier nuclei studied. The shift is rather surprising in view of the fact that the (e,e') reaction shows no target dependence in the peak position.

We report here a calculation on the effect of Δ excitation in the projectile, depicted in Fig 1b. The amplitude for this reaction in the closure approximation depends on the form factor of the projectile at the momentum transfer of the entire reaction. There will be a peak in the cross section arising simply from the kinematic cutoff associated with the form factor.

We use a model of pion coupling to the nucleon and Δ to calculate the diagrams in Fig. 1. The model is similar to that of E. Oset et al.[4] except for that the momentum of the pion is evaluated in a frame intermediate to that of the initial and final baryons. The amplitude of diagram (a) for a spin and isospin $\frac{1}{2}$ projectile is

$$T_a = (-i)C\frac{\sqrt{2}f_\pi}{9m_\pi}\left(\frac{f^*_\pi}{m_\pi}\right)^2\left(\frac{\Lambda^2-m_\pi^2}{\Lambda^2-\omega^2+q^2}\right)^2 2\Sigma_p\sqrt{\frac{3}{2}-2\Sigma_p\Sigma'_p}$$

$$(-1)^{\Sigma_p-\Sigma'_p}(q_p)_{\Sigma_p-\Sigma'_p}\sum_\Sigma\left\{\sqrt{\tfrac{3}{2}+2\Sigma'_t\Sigma}\sqrt{\tfrac{3}{2}+2\Sigma_t\Sigma}(-1)^{\Sigma-\Sigma_t}\right.$$

$$\left.(q_t)_{\Sigma_t-\Sigma}(k_t)_{\Sigma-\Sigma'_t}\right\}\frac{1}{\omega^2-\omega_\pi(q)^2}\frac{1}{\omega-E_\Delta(q)+i\Gamma(k_\Delta)/2}$$

$$\times\frac{1}{\sqrt{2\omega_\pi(k)}}$$

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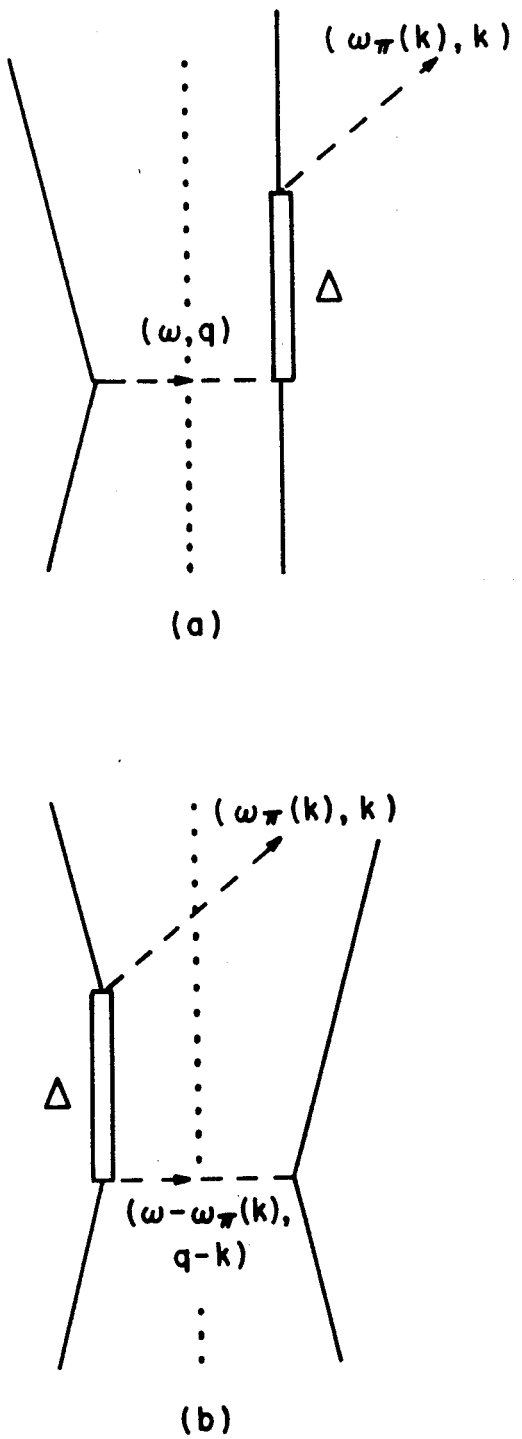

Fig. 1 Feynman diagrams of the T-matrix. The full drawn lines to the left of the dotted line represent the projectile/ejectile while those on the right represent a target nucleon. The dashed lines represent pions.

where Σ_p and Σ_t are the initial spin components of the projectile and target nucleon respectively, Σ_p' and Σ_t' refer to the final spin components, ω and q denote the total energy and momentum transfer, and k is the momentum of the final pion. The spherical tensor component m of a vector q is denoted $(q)_m$. The index p or t on a pion momenta denote that it is Lorentz-transformed to a frame intermediate to that of the baryons appearing in the vertex. The quantities $E_\Delta(q)$, and $\Gamma(k_\Delta)$ and $\omega_\pi(k)$ are defined

$$E_\Delta(q) = \sqrt{M_\Delta^2+q^2}-M_N, \quad \Gamma(k_\Delta)=\Gamma_o\left(\frac{k_\Delta}{k_b}\right)^3\left(\frac{\Lambda_\Gamma^2+k_o^2}{\Lambda_\Gamma^2+k_\Delta^2}\right)$$

and $\omega_\pi(k)=\sqrt{m_\pi^2+k^2}$.

where k_Δ is the momentum of the particle in the frame of the excited nucleon and k_o its value when $\omega=E_\Delta(q)$. The equation for the amplitude of diagram b, T_b is gotten from that of diagram a by the following replacements: $\omega \to \omega-\omega_\pi(k)$; $q \to q-k$; interchange indices p and t everywhere; and change the definition of E_Δ to

$$E_\Delta(P) = \sqrt{(M_p+\sqrt{M_\Delta^2+q'^2}-M_N)^2-q'^2}-M_p$$

where M_p is the projectile mass and q' is the momentum transfer to the projectile nucleon evaluated in the frame of the projectile.

There is an overall numerical factor C taking into account isospin couplings which depends on whether the target is a proton or neutron and in the case of a neutron whether the final pion is a π^o or π^+ :

target nucleon	final pion	C diag. a	C diag. b
proton	π^+	3	-1
neutron	π^o	$-\sqrt{2}$	$\sqrt{2}$
neutron	π^+	1	1

The outgoing final pion is assumed to be on-shell although it might be absorbed in the target. The T-matrix of the reaction is taken as the sum of the contributions from the two diagrams with the factors given above.

We used the following parameters in our calculation: Λ=700 MeV; f_π=1; f_π^*=2; M_Δ=1236 MeV; Λ_Γ=300 MeV; Γ_o=140 MeV. Fig. 2 shows the model for p(p,n), proving a good agreement with data from this parameter set.

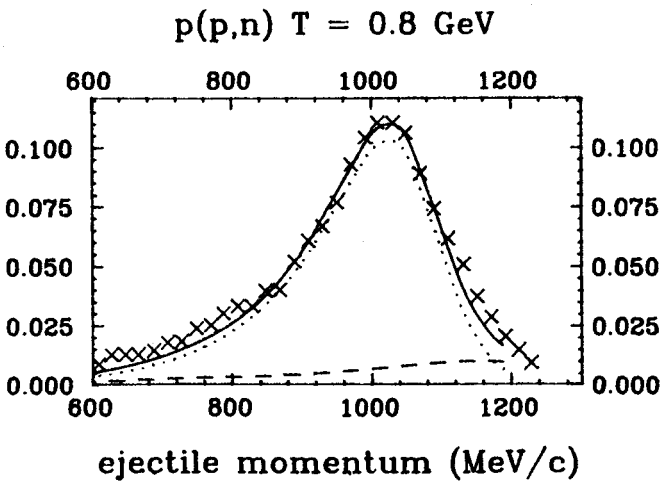

Fig. 2 $d^2\sigma/(d\Omega dp_n)$ in mbarn/(ster·MeV/c) for p(p,n) as a function of the momentum of the outgoing neutron for forward scattering. Crosses mark experimental data, the dotted line the contribution from the Δ excitation of the target nucleon, the dashed lines the excitation of the projectile and the solid curve the sum of them.

Figure 3 shows the result for the p(^{3}He,t) reaction. The peak is well described except at small energy transfer where the cross section is underestimated. The peak is dominated by the diagram (a). For the n(^{3}He,t) reaction this is not the case and the peak is moved towards lower energies due to the delta excitation of the projectile/ejectile system as is seen in Fig. 4. This shows that on heavy nuclei where the projectile will see a preponderance of neutrons due to the surface character of the reaction the cross section should be closer to the solid curve in Fig. 4 than that of Fig. 3.

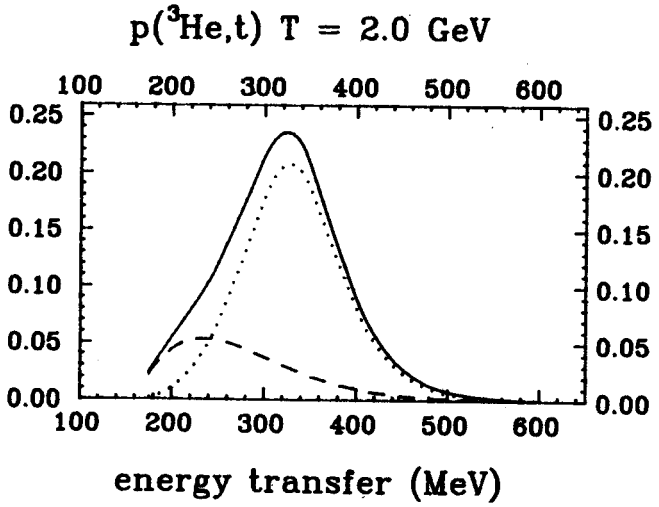

Fig. 3 $d^2\sigma/(d\Omega d\omega)$ for $p(^3He,t)$ forward scattering as function of the energy transfer ω. The beam energy is 2 GeV. The curves has the same meaning as in Fig. 1.

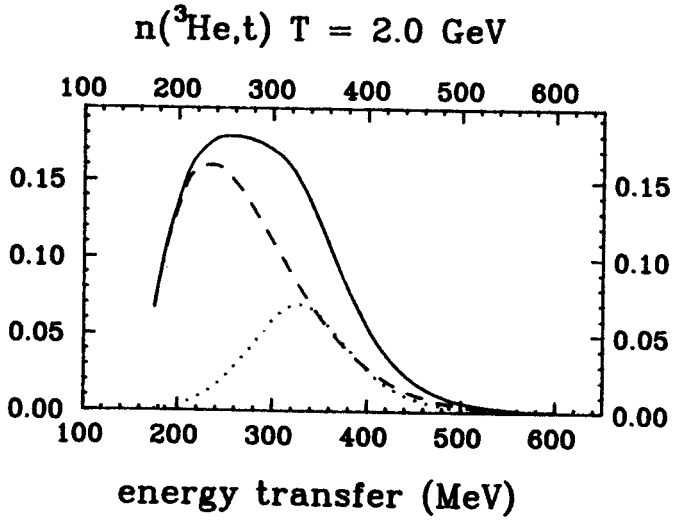

Fig. 4 The same as Fig. 3 but for $n(^3He,t)$

We propose that our explanation of the apparent shift of the delta resonance in nuclei be checked by means of the $d(^3He,t)$ reaction.

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