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IN FEW NUCLEON SYSTEMS

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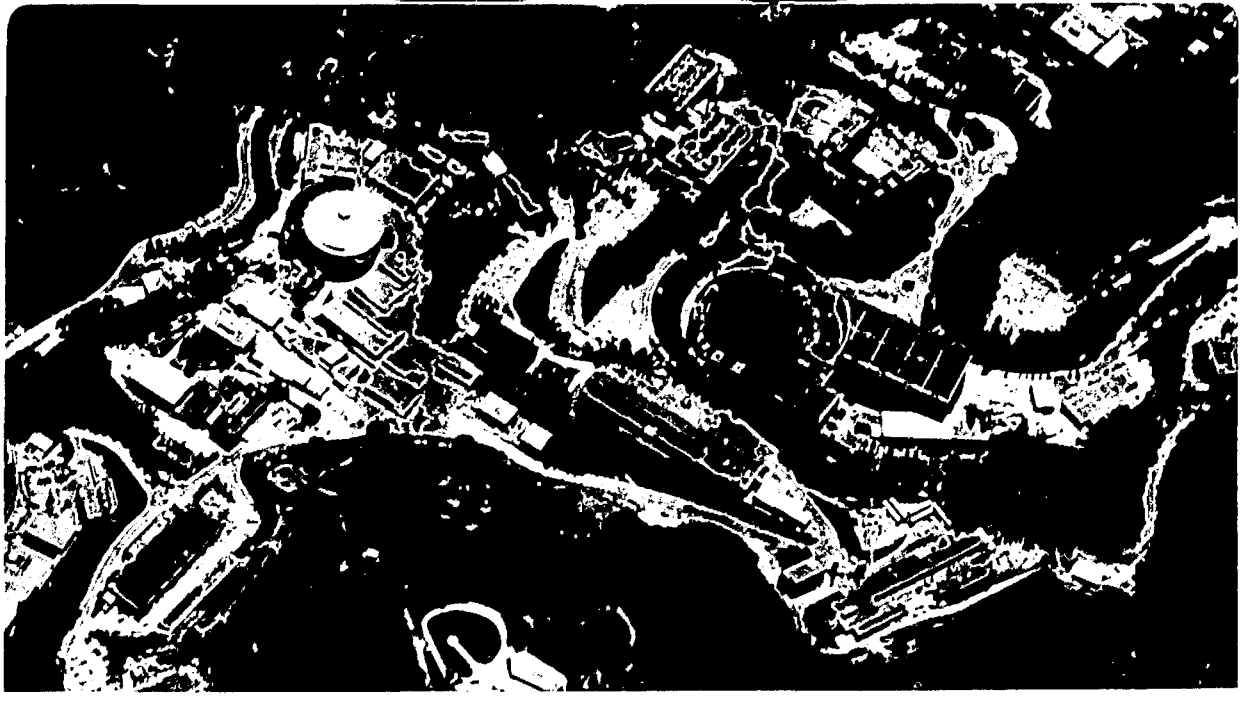

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One Pion Exchange Exchange Effects in Few Nucleon Systems[1]

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ABSTRACT

Exchange effects in ^{2}H, ^{3}He, induced by a one-pion exchange potential between quarks, are investigated in the context of a QCD-like potential model. For reasonable nucleon quark core sizes and πqq vertices such that the pion plays a non-negligible role in nucleon structure, these effects are found to be effectively repulsive and of a magnitude significant on the scale of nuclear binding.

1. Introduction

The role of one-pion exchange (OPE) in generating the long-range part of the NN interaction is well known. P-wave NN scattering lengths, the asymptotic D/S ratio, η, and quadrupole moment, Q, of the deuteron, all of which are primarily sensitive to large NN separations, are accurately pionic, sufficiently so in the case of η, Q to provide also a test of the presence of OPE at intermediate distances[1]. In addition, even in the presence of residual quark forces, a significant portion of the point-like OPE tensor force appears necessary down to $r_{NN} \sim O(1\ \mathrm{fm})$ in order to understand the deuteron binding[2]. The notion that OPE forces can be generated from an effective πNN coupling, however, is clearly valid only at large r_{NN}, where exchange overlaps are exponentially suppressed. In the language of the bag model, and, assuming the bag surface is not infinitely stiff, one expects antisymmetrization (exchange) effects to be important at internucleon distances such that two nucleon bags either overlap or have surface separation of order the typical zero point surface motion, the result being that, at such separations, the notion of a πNN coupling is ill-defined.

The problem of finding a self-consistent framework in which pion exchange can be incorporated at intermediate NN distances is easily solved by the introduction of a "fundamental" πqq coupling. Such a coupling, in fact, occurs rather naturally in both bag and potential model pictures of baryon structure. In the former it

provides a means of restoring the (approximate) chiral symmetry associated with the (approximately) massless current current u and d quarks, which is violated by the confining boundary condition of the bag[3]. There results a two-phase picture of baryon structure, in which chiral symmetry is realized in the Wigner mode inside the bag and in the Nambu-Goldstone mode outside it. Models (the so-called cloudy[4] and chiral[5] bags) based on this chiral bag picture have received extensive phenomenological application. Potential models, on the other hand, are most naturally interpreted as an effective low energy theory corresponding to a picture of the QCD vacuum in which the typical chiral symmetry breaking scale, Λ_χ, is $\sim$ 1 GeV, considerably greater than the deconfinement scale, $\Lambda_C \sim 200$ MeV[6]. Such a picture, suggested by Shuryak[7] in response to the observation that the vacuum gluon energy density extracted from QCD sum rules is an order of magnitude greater than the volume energy constant of the bag, is in agreement with both the estimate of Λ_χ obtained from an analysis of $K \to 3\pi$ using effective chiral Lagrangians[8] and the rather large value of m_π^2 when measured on the scale of the light current quark masses. Since the emergence of the pion is naturally associated with the spontaneous breakdown of chiral symmetry, the resulting effective theory involves constituent quarks (whose scale size is expected to be a few tenths of a fermi) interacting with both pions and gluons[6].

The presence of a πqq coupling, of course, generates a pion cloud and, in consequence, meson exchange current contributions to static baryon properties. Such effects have been extensively invest-

igated in the chiral bag picture (for an excellent review, in the context of the cloudy bag, see Thomas[4]), and to a considerably lesser extent, in the "chiral potential model" picture[9]. In the present work we do not consider such effects and, as a result, will eliminate the explicit pionic degrees of freedom in favor of an effective qq interaction. Note that, in addition to the usual form factor effects generated by the finite extent of the pion sources, there will now be OPE exchange (OPEE) effects, generated by the requirement of quark antisymmetrization, which are not present in OPE NN forces produced through an effective πNN vertex. The importance of such effects in relaxing the limit on the bag radius set by η and Q has been noted previously by Guichon and Miller[10].

In what follows we investigate the consequences of the existence of OPEE effects for the binding of light nuclei. For simplicity we will consider only ^2H and ^3He and restrict ourselves to the symmetric s-wave component of the nuclear wavefunction. D-wave effects, which, owing to the possible coherence of one-gluon and one-pion exchange tensor forces, are of considerable interest, particularly in tightly bound nuclei such as ^3He, will be the subject of a future report. The calculation is performed in the potential model framework, although from the usual machinery of QCD-inspired potential models we require only the resulting nucleon wavefunction. The primary impetus for this choice is the associated flexibility in representing spatial exchange matrix elements, for which the usual boundary condition segregation of one-bag and two-bag regimes provides only a

rather crude step-function approximation. We will, however, also use Shuryak's picture of the vacuum for guidance in choosing the scale size of form factors which regulate the short distance OPE qq interaction.

2. The Model

We take, for the nucleon wavefunction,

$$\Phi_N(ijk) = \phi(ijk) A_{ijk} \left(\chi_\rho^{ijk} N_\rho^{ijk} + \chi_\lambda^{ijk} N_\lambda^{ijk} \right)/\sqrt{2} \tag{1}$$

where ijk are quark labels, A_{ijk} the usual color singlet combination of quarks i,j, and k,

$$A_{ijk} = \frac{1}{\sqrt{6}} \varepsilon_{\alpha\beta\gamma} \, | q_i^\alpha q_j^\beta q_k^\gamma \rangle \tag{2}$$

and χ_ρ, χ_λ (N_ρ, N_λ) the spin (isospin) ½ three quark wavefunctions of mixed permutational symmetry

$$\chi_\rho \uparrow = (\uparrow\downarrow - \downarrow\uparrow)\uparrow/\sqrt{2} \tag{3}$$

$$\chi_\lambda \uparrow = (\uparrow\downarrow\uparrow + \downarrow\uparrow\uparrow - 2\uparrow\uparrow\downarrow)/\sqrt{6} \tag{4}$$

(with analogous expressions for N_ρ, N_λ). The ground state spatial wavefunction, $\phi(ijk)$, in (1), is taken to be

$$\phi(ijk) = \frac{\alpha^3}{\pi^{3/2}} \exp\left(-\alpha^2 (\rho_{ijk}^2 + \lambda_{ijk}^2)/2\right) \tag{6}$$

with

$$\underset{\sim}{\rho}_{ijk} = (\underset{\sim}{r}_i - \underset{\sim}{r}_j)/\sqrt{2} \tag{7}$$

$$\underset{\sim}{\lambda}_{ijk} = (2\underset{\sim}{r}_k - \underset{\sim}{r}_i - \underset{\sim}{r}_j)/\sqrt{6} \; . \tag{8}$$

Owing to the δ-function singularity of the hyperfine piece of one-gluon exchange, the cluster parameter, α, is not usually dynamically determined. We have taken the Isgur-Karl[11] spectroscopic value of 320 MeV, which corresponds to an rms quark core radius of .6 fm, and checked, as discussed below, that the qualitative features of the results are not changed if one uses somewhat larger values such as that suggested by baryon decay analyses[12].

We denote by (123;456) the normalized but not yet fully antisymmetrized state of two nucleons coupled to I=0 S=1 in which nucleon 1 contains quarks 1, 2, and 3 and nucleon 2 quarks 4, 5, and 6. From (1) we have

$$(123;456) = \left[\Phi_N(123)\Phi_N(456)\right]_{(IS=01)}\psi(\underset{\sim}{R}_{123;456}) \tag{9}$$

where the explicit spin and isospin couplings have been suppressed, ψ is the intercluster spatial wavefunction and $\underset{\sim}{R}_{123;456}$ the relative intercluster coordinate

$$\underset{\sim}{R}_{123;456} = \frac{1}{3}(\underset{\sim}{r}_1+\underset{\sim}{r}_2+\underset{\sim}{r}_3-\underset{\sim}{r}_4-\underset{\sim}{r}_5-\underset{\sim}{r}_6). \tag{10}$$

As mentioned above we restrict ourselves to symmetric s-wave components in ψ. A normalized but not yet fully antisymmetrized state of three nucleons coupled to $I=\frac{1}{2}$ $S=\frac{1}{2}$ is similarly defined by

$$(123;456;789) = \left[\Phi_N(123)\Phi_N(456)\Phi_N(789)\right]_{(IS=\frac{1}{2}\frac{1}{2})} \cdot \psi(\underset{\sim}{R}_{123;456;789}, \underset{\sim}{L}_{123;456;789}) \tag{11}$$

with explicit spin and isospin couplings similarly suppressed and $\underset{\sim}{R}_{123;456;789}$, $\underset{\sim}{L}_{123;456;789}$ the usual nuclear Jacobi coordinates

$$\begin{aligned}\underset{\sim}{R}_{123;456;789} &= ((\underset{\sim}{r}_1+\underset{\sim}{r}_2+\underset{\sim}{r}_3)/3-(\underset{\sim}{r}_4+\underset{\sim}{r}_5+\underset{\sim}{r}_6))/\sqrt{2} \\ \underset{\sim}{L}_{123;456;789} &= (2(\underset{\sim}{r}_7+\underset{\sim}{r}_8+\underset{\sim}{r}_9)/3-(\underset{\sim}{r}_1+\underset{\sim}{r}_2+\underset{\sim}{r}_3)/3 \\ &\qquad -(\underset{\sim}{r}_4+\underset{\sim}{r}_5+\underset{\sim}{r}_6)/3)/\sqrt{6}\ .\end{aligned} \tag{12}$$

In (11) the nuclear wavefunction ψ is taken to be normalized with respect to the measure $d\tau = d^3R_{123;456;789}d^3L_{123;456;789}$. The fully antisymmetrized states corresponding to (9) and (11) are then formed by applying to the right hand sides the six and nine quark antisymmetrizers, respectively. Bearing in mind the already existing antisymmetries with respect to interchange of quarks within the same nucleon cluster, one can easily show that the expectation value of any symmetric operator, V, between such states, is given by[13,14]

$$\langle(123;456)|V|[|(123;456)\rangle - 9|(126;453)\rangle]/N_2^{\,2} \tag{13}$$

and

$$\begin{aligned}\langle(123;456;789)|V|[|(123;456;789)\rangle &- 27|(126;453;789)\rangle \\ &+ 54|(129;453;786)\rangle + 162|(169;452;783)\rangle \\ &- 36|(483;159;726)\rangle]/N_3^{\,2}\end{aligned} \tag{14}$$

where

$$N_2^{\,2} = 1 - 9\langle(123;456)|(126;453)\rangle \tag{15}$$

and

$$N_3^{\,2} = C^{(1)} - 27C^{(2)} + 54C^{(3)} + 162C^{(4)} - 36C^{(5)} \tag{16}$$

with the quantities $C^{(i)}$ given by

$$\begin{aligned} C^{(1)} &= \langle(123;456;789)|(123;456;789)\rangle = 1 \\ C^{(2)} &= \langle(123;456;789)|(126;453;789)\rangle \\ C^{(3)} &= \langle(123;456;789)|(129;453;786)\rangle \\ C^{(4)} &= \langle(123;456;789)|(169;452;783)\rangle \\ C^{(5)} &= \langle(123;456;789)|(483;159;726)\rangle. \end{aligned} \tag{17}$$

For our purposes V is given by

$$V = \sum_{i<j} V_{ij}^{\pi}(r_{ij}) \tag{18}$$

where V_{ij}^{π} is the OPE potential between quarks i and j. The first terms in (13), (14) contain the pionic contributions to the masses of the constituent nucleons plus the regulated OPE interactions between them; the second, the induced non-local two-nucleon interactions plus the effects of delocalization on interactions involving quarks in the remaining "spectator" nucleon[14]. The last three terms in (14) correspond to induced non-local three-nucleon forces, the contributions to binding of which will be considerably less than that of the single-exchange induced two-body force[14], and so neglected in what follows.

Given the form of V_{ij}^{π}, to be discussed below, one may readily evaluate the necessary spin, color and isospin matrix elements. The only difficulties lie in the spatial sector. Since we are interested in exchange effects it is the short distance behavior of the nuclear

wavefunction, ψ, which most concerns us. Unfortunately the behavior of ψ in this regime, being especially susceptible to various quark exchange and mixing effects, must be considered to be rather ill-determined. We proceed as follows. For the deuteron we take the following wavefunction obtained in Ref. 2), which, though resulting from a somewhat crude implementation of OPE, both correctly incorporates the dynamics of the short-range exchange hyperfine interaction and provides a reasonable description of the deuteron properties:

$$\psi(\underset{\sim}{R}) = \frac{1}{N}\sum_i \xi_i \exp(-\beta_i^2 R^2/2) \tag{19}$$

where

$$\begin{aligned} \underset{\sim}{\beta} &= (40,90,320)\ \text{MeV} \\ \underset{\sim}{\xi} &= (1,2.8,-3.4) \end{aligned} \tag{20}$$

and N is a normalization constant. For ^{3}He no equivalent dynamical calculation exists so we instead resort to wavefunctions of the form originally proposed by Fearing[15] and Khanna[16] and adapted to ^{3}He by Aerts and Dover[17]

$$\psi(\underset{\sim}{R},\underset{\sim}{L}) = \frac{1}{N}\exp\left(-\beta_0^2(R^2+L^2)/2\right)\prod_{i<j}\left(1 - C\exp(-\beta_1^2 R_{ij}^2/2)\right) \tag{21}$$

where $\underset{\sim}{R}$, $\underset{\sim}{L}$ are the nuclear Jacobi coordinates, as in (12), $\underset{\sim}{R}_{ij}$ is the separation of nucleons i and j, and N is a normalization constant. The parameters β_0, β_1, C are adjusted to give as good a fit as possible to the ^{3}He charge form factor. One should bear in mind here that, in making this fit, quark mixing and exchange effects have been ignored,

as have contributions from components other than the symmetric s-wave in the nuclear wavefunction. The form (21) should, therefore, be taken only as a reasonable first attempt at incorporating the effects of the short-range two-nucleon repulsion, and results following from it should not be trusted to any great numerical accuracy. We employ the following two parameter sets, cited as models I and II in Ref. 17), differences in the numerical results associated with each serving to display the sensitivity of the calculations to details of the short-distance behavior of ψ:

$$\begin{aligned} &\text{model I:} \quad \beta_0 = 117 \text{ MeV} \quad \beta_1 = 626 \text{ MeV} \quad C = 1.0 \\ &\text{model II:} \quad \beta_0 = 154 \text{ MeV} \quad \beta_1 = 328 \text{ MeV} \quad C = .925 \ . \end{aligned} \tag{22}$$

In order to proceed we need to specify the form of V^{π}_{ij}. The leading chiral quark model πqq coupling is pseudovector[6], equivalent, in the weak binding limit, in which the quarks are taken to be approximately on mass shell, to a pseudoscalar form. We, therefore, take V^{π} to be the usual static pseudoscalar OPE potential

$$V^{\pi}_{ij}(r) = \frac{1}{3} f_q^{\,2} \, \underset{\sim}{\tau}_i \cdot \underset{\sim}{\tau}_j \, \underset{\sim}{\sigma}_i \cdot \underset{\sim}{\sigma}_j \, \frac{\exp(-m_\pi r)}{r} \, F(r) + V^{\pi}_t(r) \tag{23}$$

where the factor $F(r)$ arises as a result of form factors at the πqq vertices and $\to 1$ as $r \to \infty$, and V^{π}_t is the tensor component of V^{π} which, in view of our restriction to s-waves, will be neglected in what follows. In (23)

$$\begin{aligned} f_q^{\,2} &= (m_\pi / 2m_q)^2 g_{\pi qq}^{\,2} / 4\pi \\ &= 9 f_N^{\,2} / 25 \end{aligned} \tag{24}$$

where

$$f_N^{\,2} = (m_\pi / 2m_N)^2 g_{\pi NN}^{\,2} / 4 \cong .08 \tag{25}$$

and the second line in (24) follows from the requirement that $g_{\pi qq}$ be chosen so as to reproduce the asymptotic OPE NN potential. OPE interactions of the form (23) have been employed previously in non-relativistic potential models[9,18-21] and provide a reasonable starting point for investigating pionic effects in the baryon sector. Note, however, that (23), which arises from a

$$v(\underset{\sim}{q}) = \underset{\sim}{\sigma} \cdot \underset{\sim}{q} \, \underset{\sim}{\tau} f(q^2) \tag{26}$$

form for the pion emission vertex (with $f(q^2)$ the vertex form factor), is undoubtedly an oversimplification. Not only does the choice (24) produce baryon decay widths which are generally too small[22,23], but detailed decay analyses require, in addition to (26), a vertex recoil term of roughly equal magnitude and opposite sign in order to reproduce the relative phases of different partial waves in certain N^* decays[12]. These objections are themselves subject to the caveat that pion cloud effects have been ignored in arriving at them.

The form of $F(r)$, in (23), depends on the structure of the vertex form factor, $f(q^2)$. For $f(q^2) = 1$ (point-like coupling)

$$F(r) = 1 - \frac{4\pi}{m^2} \, r \exp(m_\pi r) \, \delta^3(r) \quad , \tag{27}$$

the δ-function term in (27) being spread out over finite range of r for non-point-like $f(q^2)$. From (23) one can show that, as in the

cloudy bag, OPE produces attractive, but unequal, contributions to the N and Δ masses, thus generating a portion of the Δ-N splitting, and, as a result, allowing a smaller phenomenological value of α_s. The presence of non-point-like $f(q^2)$ is a natural consequence of the chiral quark model picture. As an effective low energy theory, such a model requires a cutoff, in this case of order $\Lambda_\chi \sim 1$ GeV, which may be thought of as arising from non-localities associated with the current quark substructure of pions and constituent quarks. If one imagines a constituent quark as a small region (of order $r \sim 1/\Lambda_\chi$) of chirally symmetric vacuum from which the pion field is, therefore, to be excluded, one conlcudes that the OPE force between quarks should be suppressed at distances below of order a few times $1/\Lambda_\chi$. Since, however, the effective low energy Lagrangian, in general, contains higher order non-minimal terms which also scale as $1/\Lambda_\chi$, one cannot be too precise about this statement, and, given the given the qualitative nature of the argument, little firm guidance is to be found regarding the sharpness of the suppression of the effective pion field.

Let us begin by using a monopole vertex form factor

$$f(q^2) = \Lambda^2/(\Lambda^2 + q^2) \ . \tag{28}$$

The form (28) represents a rather soft suppression of the pion field, and, in consequence, produces a rather larger effect in the baryon spectrum than would a sharper cutoff of equivalent range. As we will see below, this observation, coupled with the lack of quantitative criteria for choosing the sharpness of the field suppression, some-

what degrades our ability to correlate the importance of pionic effects in the spectrum with the size of possible OPEE shifts in few nucleon systems.

In closing this section, it is worth noting that the underlying chiral quark picture to which we have been appealing for qualitative guidance, generates, in a quite natural way, reasonable πNN form factors. The parameter, Λ, in (28), is related to the range of suppression in the usual manner

$$r = 6/\Lambda \tag{29}$$

r being, in our case, of order .2 to .3 fm on naive geometrical grounds. The induced πNN form factor is then given by

$$f_{\pi NN}(q^2) = \frac{\Lambda^2}{\Lambda^2 + q^2} \exp(-q^2/6\alpha^2) \tag{30}$$

where α is the nucleon cluster parameter of (6). For low q^2 this reduces to an effective monopole form factor with

$$\Lambda_{\pi NN} = \sqrt{\frac{6\alpha^2\Lambda^2}{6\alpha^2+\Lambda^2}} \tag{31}$$

which, for α between 320 and 400 MeV and r as above, lies in the range 725 to 925 MeV, in good agreement with nucleon monopole form factor parameters extracted from phenomenological analyses*. While

*See Ericson and Rosa-Clot[24] for the appropriate references, as well as a critical discussion of the difficulties presented by such values for a purely pion-plus-nucleon picture, in the absence of constituent effects.

this agreement is, to a certain extent, no doubt, fortuitous, it is of interest to note that significantly smaller values of $\Lambda_{\pi NN}$ can only be produced by grossly overestimating the size of the substructure of the constituent quarks.

3. Results and Discussion

From the permutational symmetries of the states defined in §2 one can show that the matrix elements in (13) are given by

$$\langle(123;456)|V|(123;456)\rangle = \langle(123;456)|6V^{\pi}_{12}+9V^{\pi}_{36}|(123;456)\rangle \quad (32)$$

and

$$\langle(123;456)|V|(126;453)\rangle = \langle(123;456)|2V^{\pi}_{12}+8V^{\pi}_{13}+4V^{\pi}_{14} +V^{\pi}_{36}|(126;453)\rangle \ . \quad (33)$$

Similarly, in (14),

$$\langle(123;456;789)|V|(123;456;789)\rangle = \langle(123;456;789)|9V^{\pi}_{12} +27V^{\pi}_{36}|(123;456;789)\rangle \quad (34)$$

and

$$\langle(123;456;789)|V|(126;453;789)\rangle = \langle(123;456;789)|2V^{\pi}_{12} +8V^{\pi}_{13}+4V^{\pi}_{14}+V^{\pi}_{36}+3V^{\pi}_{78}+12V^{\pi}_{17}+6V^{\pi}_{37}|(126;453;789)\rangle \ . \quad (35)$$

Neglecting three-body terms in (14), (16) and evaluating spin, color and isospin factors the normalization factors $N_2^{\ 2}$ and $N_3^{\ 2}$ of (15),(16) become

$$N_2^{\ 2} = 1 + B^{(2)}/9 \quad (36)$$

$$N_3^{\ 2} = 1 + B^{(3)}/3 \quad (37)$$

where the $B^{(i)}$ are the single exchange spatial overlap integrals

$$B^{(2)} = \int d\tau \left[\phi(123)\phi(456)\psi(123;456)\right]^* \left[\phi(126)\phi(453)\psi(126;453)\right] \quad (38)$$

and

$$B^{(3)} = \int d\tau \left[\phi(123)\phi(456)\phi(789)\psi(123;456;789)\right]^* \left[\phi(126)\phi(453)\phi(789)\psi(126;453;789)\right] \ . \quad (39)$$

Because of the assumed structure of the nuclear states the space, color and spin-isospin sectors of (32)-(35) are decoupled. V^{π} is color independent so that the color matrix elements are simply the overlap factors 1, 1/3, 1, and 1/3 respectively. The required spin-isospin matrix elements are presented in Table 1.

An understanding of the implications of the results of (32)-(35) for few nucleon systems is greatly facilitated by the following characterization of the terms therin. Since, in (32), (34), we may specify which quarks "belong" to which nucleons, it is clear that the V_{12}, V_{36} terms generate the pionic mass shifts of the constituent nucleons and the NN OPE potentials, respectively. These are standard effects, present even in the absence of exchange contributions, are in which we are, therefore, not interested <u>per se</u>. Note that for πqq interactions which produce 10% or more of the Δ-N splitting from OPE forces the former is much greater than the latter; as a result we restrict ourselves to this limit for the sake of the ensuing discussion. The

matrix elements (32), (34) enter (13), (14) not as is, but multiplied by the factor $(1+\nu B)^{-1}$ ($=1-\nu B$ for small B), where $\nu = 1/9, 1/3$ for ^2H, ^3He, respectively. Since the pionic mass shift of the nucleon is negative, the presence of the normalization factors in (13), (14), therefore, tends to unbind the nuclear system. This is a simple consequence of the delocalization necessitated by quark antisymmetry. Similarly, the V_{78} term of (35), which depends only on the exchange overlap factor for the particular nuclear state and the expectation of V^{π}_{ij} for quarks i and j in an isolated nucleon, results purely from delocalization. We combine these two contributions into a single exchange delocalization effect. Given the pionic contribution to the nucleon mass such effects depend only on the exchange overlap. The V_{17}, V_{37} terms of (35) are less unambiguously interpreted. One might plausibly consider them as either delocalization effects on the NN OPE force or true exchange interactions. In practice these terms turn out to be small and are grouped together with the remaining terms of (35), V_{12}, V_{13}, V_{14} and V_{36} which, as in (33), generate the non-local two-nucleon force. We label such terms, generically, exchange interaction effects, and note that they are more sensitive to details of the qq OPE force than are the corresponding delocalization contributions. Note that while the three-body terms which we are neglecting in (14) are certainly small compared to the single exchange terms, (35), they may nonetheless be of order a few hundred keV in magnitude[14] and play some role in understanding the three-nucleon binding discrepancy.

Let us now turn to the actual results. We begin with the deuteron, where, as expected, the effects are relatively small due to the diffuse structure of the nuclear wavefunction. Figure 1 displays the exchange interaction, exchange delocalization and total exchange contributions for the nuclear wavefunction (19), α=320 MeV, and a range of monopole form factors. The % of the Δ-N splitting generated by OPE, as a function of Λ, is given along the upper edge of the Figure and serves as a measure of the importance of pionic effects in the baryon spectrum. The range of Λ corresponding to the qualitative physical picture discussed above is 1670 to 2500 MeV; due to the rapid convergence toward point-like values we have omitted the upper portion of this range from the Figures. Note that, for such values of Λ, the full exchange effect is dominated by exchange delocalization. For smaller values this is no longer true; the net exchange contribution, however, remains quite insensitive to Λ and small, though not negligible on the scale of the deuteron binding. Being effectively repulsive, it may leave room for contributions to the intermediate range NN attraction from two pion exchange, in addition to color polarization, without upsetting the (albeit somewhat semi-quantitative) understanding of the deuteron binding obtained in Ref. 2.

As one might guess from its considerably smaller size, exchange effects in ^3He are much enhanced over those in the deuteron. Figure 2 displays the results for the model wavefunctions I and II of (21) and (22) over a range of monopole form factors, with α=320 MeV. Similar results are presented in Figure 3 for α=400 MeV. Note that models I and II, having similar spatial exchange overlaps, $B^{(3)}$, have also

similar exchange delocalization effects. This is not the case for the exchange interaction contributions, which, as noted above, are sensitive to short-range features of the nuclear wavefunction. Note also that the results do not vary greatly with Λ over the range of physical interest discussed above. While the qualitative features of the model- and Λ-dependence are not affected by cluster size, one sees that the magnitude of the total exchange contribution is. This is not surprising since smaller cluster sizes (larger α values) lead to reduced spatial exchange matrix elements, an effect clearly seen in the Figures, and one which is, in fact, even somewhat masked because the smaller nucleon clusters have much larger pionic mass shifts, resulting, therefore, in delocalization contributions considerably larger than one would have expected on the basis of the spatial overlap factors, $B^{(3)}$, alone.

The correlation between OPE contributions to the Δ-N splitting and the range of pion field suppression, (29), displayed in Figures 2 and 3 is, of course, characteristic of the monopole form, (28). While such a form is by no means unreasonable, there is, at the same time, no compelling reason for supposing it to be correct. To explore the extent to which the results might depend on this choice, the following alternative scheme has been employed, motivated by the naive geometrical picture of suppression noted above. Recall that the central piece of the point-like OPEP is related to the coordinate space pion field, $\phi_\pi(r)=\exp(-m_\pi r)/r$, by

$$V^{\pi}_{ij}(r_{ij}) \sim \nabla^2_{r_{ij}} \phi_\pi(r_{ij}) \quad . \tag{40}$$

In order to account for the suppression of the pion field inside the region of the constituent quarks we, therefore, modify ϕ_π of (40) by a factor

$$f(r) = r^n/(r^n+r_o^{\,n}) \quad . \tag{41}$$

r_o sets the scale of the suppression and should be of order .4 fm, the point-like limit being recovered by $r_o \to 0$. Note that, while (41) might seem naively more plausible than (28), such an assessment implicitly ignores the non-local and non-minimal structure of the underlying chiral quark picture. Moreover, the derivative coupling structure of the pseudoscalar vertex tends to produce, in combination with the rather sharp cutoff (41), unphysical oscillations in V^{π}_{ij} in the vicinity of $r=r_o$. These oscillations may be removed by averaging the resulting V_{ij} over the assumed $\sim$ .2 fm extent of the constituent quarks; the effect is a decrease of .8 to 1.0 MeV in the net exchange contribution. We will not comment further on this modification. The results of the scheme (41) are displayed in Tables 2 and 3, corresponding to the choices α=320, 400 MeV, respectively. The case r_o=.4 fm, n=4 or 8 should be compared to the rightmost portion of Figures 2,3 to which it corresponds in range of pion field suppression. One notes a decrease in the net exchange contribution for all choices of nuclear wavefunction and cluster size, the bulk of this effect being a decrease in the delocalization contribution due to the smaller pionic mass shifts associated with (41). Note, however, that the model dependence, for fixed r_o, n and α, is also somewhat different, a reflection of the

sensitivity of the exchange interactions to short range features of, not only the nuclear wavefunction, but also V_{ij}^{π}. Thus, while (41) may not be particularly realistic as it stands, it serves to underscore the uncertainties associated with the present calculation. The chiral quark picture is simply not precise enough to specify either the short range behavior of V_{ij}^{π} or the relation between the range of field suppression and the importance of OPE effects in the baryon spectrum, as typified by the OPE contributions to the Δ-N splitting, the results being quantitatively sensitive to uncertainties in both. Such uncertainties aside, however, one sees that the results are in good qualitative agreement.

In conclusion, OPEE effects in ^{2}H, ^{3}He, and, by extension, other light nuclei, appear to be effectively repulsive and of non-trivial size on the scale of nuclear binding. While the magnitudes of these effects cannot be calculated to any great numerical accuracy, there appears no way they may be made small, short of employing nucleon quark cores considerably smaller than .5 fm or models in which the pion plays no significant role in determining nucleon structure. The presence of such effects, which cannot be accounted for by OPEP's generated from effective NN vertices, suggests the necessity of addition non-mesonic contributions to the attractive part of the NN force.

Table 1: Matrix Elements of the Operators $\underset{\sim}{\sigma}_i \cdot \underset{\sim}{\sigma}_j \; \underset{\sim}{\tau}_i \cdot \underset{\sim}{\tau}_j$*

(ij)	$\langle \underset{\sim}{\sigma}_i \cdot \underset{\sim}{\sigma}_j \; \underset{\sim}{\tau}_i \cdot \underset{\sim}{\tau}_j \rangle$
$(12)_d$	5
$(36)_d$	-25/27
$(12)_e$	-37/27
$(13)_e$	11/27
$(14)_e$	1/9
$(36)_e$	59/27
$(17)_e$	-5/243
$(37)_e$	35/243
$(78)_e$	-5/27

*The notation is as follows: the subscripts 'd' and 'e' refer to direct and exchange matrix elements, respectively. Results quoted are for ^{3}He; those for ^{2}H are identical, with the exception that the last three entries do not occur.

Table 2: Exchange Contributions Corresponding to the Ansatz (41) for Pion Field Suppression, for α=320MeV*

		Model I			Model II		
r_o,n	%Δ-N	Del	Int	Total	Del	Int	Total
.5,4	14.6	2.08	3.11	5.19	2.27	2.97	5.24
.5,8	14.5	2.07	3.54	5.61	2.26	3.45	5.71
.4,4	21.5	3.05	2.29	5.34	3.34	1.82	5.16
.4,8	22.7	3.22	2.46	5.68	3.52	1.96	5.48
.3,4	29.4	4.15	1.71	5.86	4.56	.87	5.43
.3,8	31.2	4.43	1.67	6.10	4.85	.75	5.60

*The notation is: %Δ-N = % of the Δ-N splitting due to OPE, Del=exchange delocalization contribution, Int=exchange interaction contribution. All entries in MeV.

Table 3: Exchange Contributions Corresponding to the Ansatz (41) for Pion Field Suppression, for α=400 MeV*

		Model I			Model II		
r_o,n	% -N	Del	Int	Total	Del	Int	Total
.5,4	14.8	.97	2.82	3.79	.91	2.05	3.42
.5,8	11.1	.73	3.45	4.18	.68	2.83	3.85
.4,4	29.0	1.91	2.19	4.10	1.78	1.71	3.49
.4,8	29.0	1.91	2.60	4.51	1.78	2.08	3.86
.3,4	47.2	3.11	1.43	4.54	2.90	.74	3.64
.3,8	50.3	3.31	1.48	4.79	3.08	.74	3.82

*Notation as in Table 2. All entries in MeV.

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FIGURE CAPTIONS

Figure 1: Exchange Contributions for ^2H

1=exchange delocalization contribution; 2=exchange interaction contribution; 3=total exchange contribution

Figure 2: Exchange Contributions for ^3He (α=320 MeV)

I=Model I; II=Model II. All other notation as in Figure 1.

Figure 3: Exchange Contributions for ^3He (α=400 MeV)

All notation as in Figure 2.

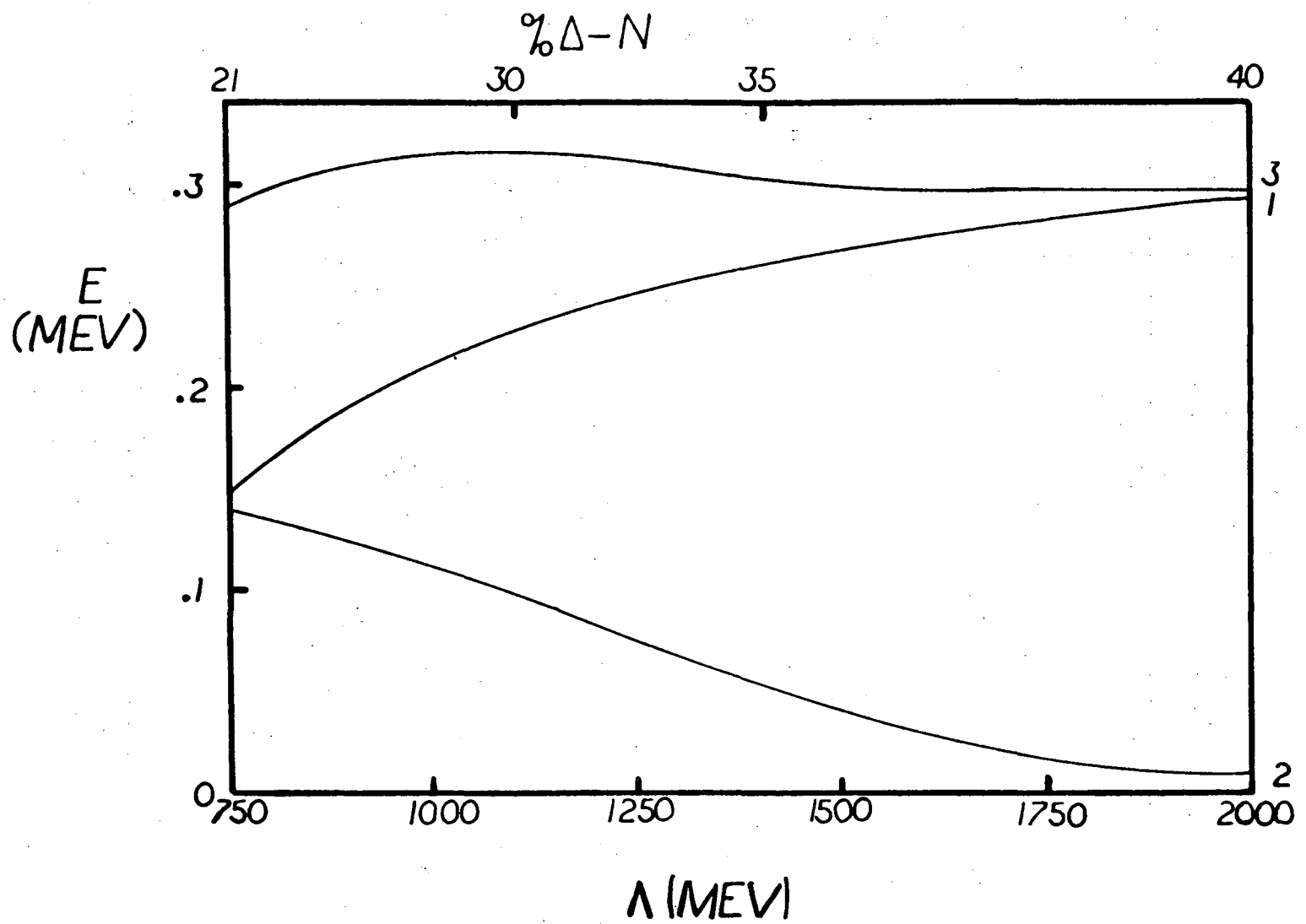

FIG. 1

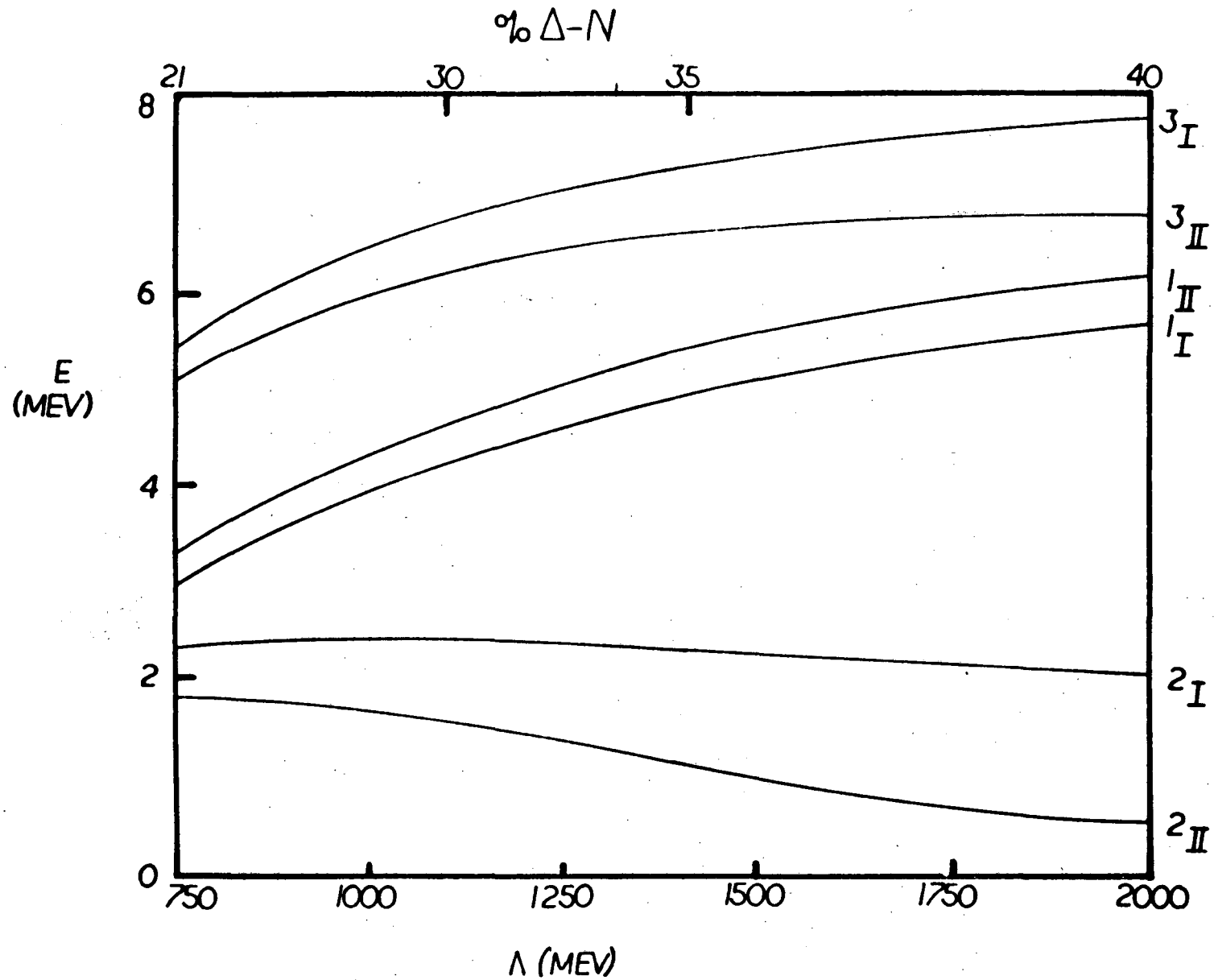
% Δ-N
21
30
35
40
8
6
4
2
0
E
(MEV)
3I
3II
1II
1I
2I
2II
750
1000
1250
1500
1750
2000
Λ (MEV)

FIG. 2

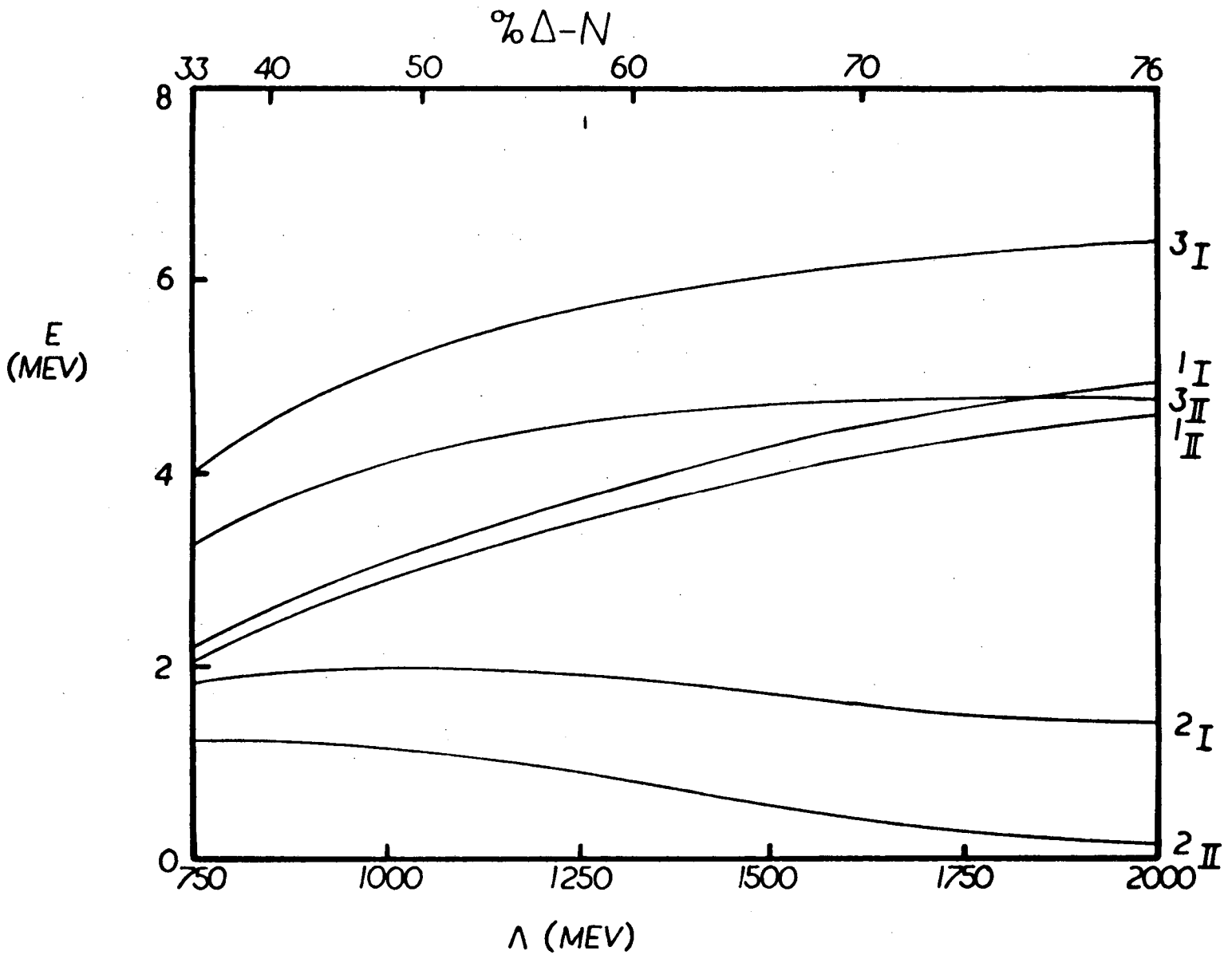
%Δ-N
33
40
50
60
70
76
8
6
4
2
0
E
(MEV)
750
1000
1250
1500
1750
2000
Λ (MEV)
3 I
1 I
3 II
1 II
2 I
2 II

FIG. 3

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