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UNPAIRED NUCLEONS AS PROBES OF CORE COLLECTIVE FIELDS

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ABSTRACT

Field theory techniques are applied to the study of interacting boson systems. A dynamical quasiparticle theory derived by the equations of motion method is compared with IBFM as a means of coupling an odd particle to IBM cores. For a single j-shell, the two models are found to be roughly equivalent on a phenomenological level. This equivalence requires that the IBFM parameter Λ_o is a function of the core, SU(5) vs. SU(3) or O(6), but not an explicit function of the Fermi level. IBFM is found to 'Coriolis attenuate' in strong-coupled bands - partly decoupled cases must be investigated further. The dynamical quasiparticle method has advantages with respect to (i) the physical significance of the pairing Δ as opposed to the IBFM $\Lambda_{jj'}$'s, (ii) particle transfer, at least in principle, and (iii) easy generalizability. In an application of static mean field theory, the α-cluster interpretation of the SU(4) model for the Ra isotopes is tested. It is found that a larger cluster would be required to account for experimental odd-nucleon decoupling factors.

1. Introduction

Boson methods and field theory are complementary approaches to the many-body problem. The basic point of this contribution is that boson models may lead to the concept of fields and can thus be tested by the well-developed techniques of field theory. The conception of fields can occur either in attempts to visualize intuitively the meaning of the boson symmetry groups, or in a more rigorous fashion. Section 2 gives an example of the former, while Section 3 describes a theory for the latter. In both cases we focus on the odd particle as a probe of nuclear structure.

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2. Probing the Mean Fields: Do Ra Nuclei Resemble Diatomic Molecules?

The doubly even nuclei in the Ra region have unusually low-lying negative-parity states. There is a long tradition of applying molecular-type group theory to account for such spectra[1,2], and at this meeting Daley has discussed a specific decomposition of the $U(6) \otimes U(4)$ symmetry group[3,4]. The p-boson of that model has been intuitively associated with an axisymmetric α-cluster degree of freedom[5,6] partly because the U(4) group has previously been applied to di-molecules[7] and partly because configuration mixing between the two 0^+ states of the model accounts for the experimental systematics of their relative α-decay widths[3,4]. Equally compatible with the symmetry would be for example a core - ^{14}C di-molecule, made plausible by the recent observation of spontaneous ^{14}C emission from ^{223}Ra[8,9], or the pear shapes obtained as static equilibrium configurations from mean field theory[10].

2.1 Characteristic effects in odd-nucleon orbits. We shall see that odd-nucleon spectra can distinguish between these different realizations of the intrinsic symmetry. To understand how the odd particle can respond in a characteristic way to an asymmetric cluster, it is useful to start from the spherical shell model picture shown schematically in Fig. 1. Dashed and solid lines represent j-shells of opposite parity. Each major shell has one intruder j-shell of high j and unnatural parity. Reflection symmetric deformations (e.g. β_2) do not mix the intruder wave functions very much due to the parity selection rule, and this allows phenomena like backbending and strongly decoupled bands to occur in deformed nuclei. However, even a small asymmetric (e.g. β_3) deformation spreads the high-j strength over the valence shell. Thus the quenching of otherwise large $\hat{j}_+$ matrix elements can be a quantitative measure of reflection asymmetry.

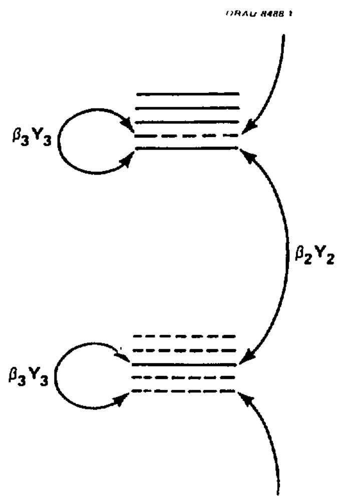


Fig. 1. Schematic picture of spherical shell structure which illustrates why the high-j intruder sub-shells are more sensitive to octupole than quadrupole deformation.

Experimental evidence for this effect on $\hat{j}_+$ matrix elements has been found in decoupling factors[11], notably the ones in ^{225}Ra, also in the Coriolis coupling between K=3/2 and 5/2 bands of ^{225}Ac[12] and in the absence of backbending in ^{222}Th[13]. Here we shall focus on the decoupling factors in ^{225}Ra, since decoupling factors are normally accurately described by mean field theory[14]. Five different intrinsic shapes will be tested: (i) a spheroid, (ii) an α particle half emerged from the spheroidal core, (iii) an α particle completely outside the core but connected by a thick neck, (iv) an A = 14 cluster likewise outside the core, and (v) a smoother pear shape corresponding to the Strutinsky equilibrium.

2.2 Parametrization of the mean field. For numerical calculations, the mean field experienced by the odd nucleon will be described by a folded Yukawa potential plus spin-orbit and Coulomb terms[15]. The parameters of such single-particle potentials were determined by the Los Alamos-Lund group[16] in 1974 using data from deformed actinide and rare earth nuclei. These parameters have been found to extrapolate accurately to other regions of nuclei[17-19].

Shapes like molecular cluster configurations can be obtained with the three-quadratic-surface parametrization[20] illustrated in Fig. 2.

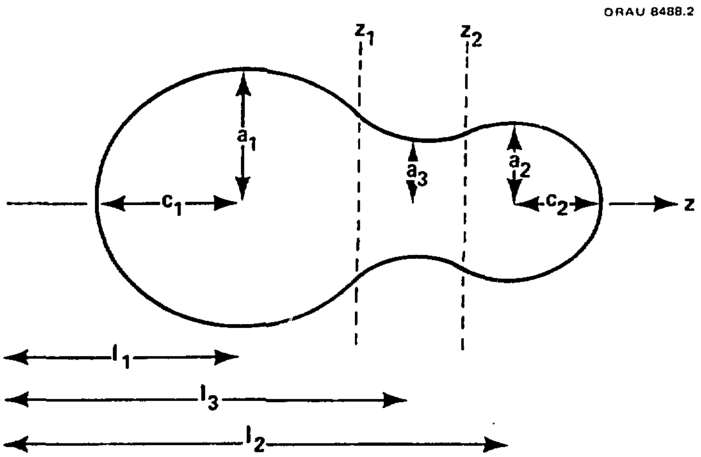


Fig. 2. Parametrization of the intrinsic shape.

There are nine shape coordinates: l_1, l_2, l_3, a_1, a_2, a_3, c_1, c_2, and c_3. Three constraints are given by volume conservation for given mass, A, and by the requirement of smooth connections at $z = z_1$ and z_2 where the surfaces intersect. The remaining six constraints must be chosen to give the shapes listed above. The following six requirements were imposed: (i) The semi-axes a_2, c_2 are scaled to the size of a cluster of mass A'=4 or 14. (ii) The center of mass is fixed by $(A-A')\, l_1 + A' l_2 = 0$. (iii) The separation of the centers, $l_2 - l_1$, is c_1 for a half emerged cluster, and $c_1 + c_2$ for clusters outside the core. (iv) The neck region, $z_2 - z_1$, is c_2 for a half emerged cluster and $2c_2$ for clusters outside the core. (v) The eccentricities of the two clusters is equal, $a_1/c_1 = a_2/c_2$. (vi) The cluster eccentricities correspond to $\varepsilon_2 = 0.14$ with the α cluster, and $\varepsilon_2 = 0.10$ with the ^{14}C cluster. Since no self-consistency criteria are being applied, the choice of these eccentricities is arbitrary. However, the odd-nucleon wave functions are not very sensitive to the eccentricity[11]. (Note that for molecular configurations the eccentricity of each cluster would be enhanced by the Coulomb field of the other[21].) The pear shape from Strutinsky theory is described by Nilsson's ε_2, ε_3,... parametrization[10].

2.3 Evidence from decoupling factors in ^{225}Ra. Decoupling factors in ^{225}Ra were recently determined in γ-ray spectroscopy[22]. The spin I = 1/2 of the ^{225}Ra ground state has also been confirmed by a laser measurement of hyperfine splitting[23]. The decoupling factor, a, extracted from experimental energy levels is related to the matrix element of $\hat{j}_+$ between $\hat{\pi}$ and $\hat{R}_1$ conjugates of the intrinsic single-particle state by[11]

$$a = - p \langle \hat{\pi}\Psi_{1/2} \mid \hat{j}_+ \mid \hat{R}_1\Psi_{1/2}\rangle \tag{1}$$

where p is the total parity. Both parities were observed in experiment, and the values extracted for the j_+ matrix element are listed at the bottom of Table I. In the strong coupling limit these values would be identical, and the observed values do indeed lie near an average of -2. (According to Ref. [11] the spreading arises, in the strong coupling representation, due to off-diagonal matrix elements of $\hat{\pi}$ analogous to the off-diagonal Coriolis matrix elements of $\hat{j}_+$).

Table I. The $\hat{j}_+$ matrix elements between $\hat{R}_1$-conjugated[14] and parity-conjugated Ω = 1/2 orbitals, calculated for the shapes described in the text and deduced from experiment[22] for ^{225}Ra.

	S_n(MeV)	orbit	$\langle\hat{\pi}\psi \mid \hat{j}_+ \mid \hat{R}_1{}'\rangle$
Spheroid π = +	-7.46 -7.08	# 63 # 65	$\begin{pmatrix} 1.42 & 4.85 \\ & -0.27 \end{pmatrix}$
Spheroid π = -	-6.05	# 69	-7.76
1/2 α	-6.45	# 67	-6.53
α	-6.75 -5.80	# 66 # 71	$\begin{pmatrix} -4.16 & -0.42 \\ & -4.65 \end{pmatrix}$
^{14}C	-6.76 -5.88	# 67 # 70	$\begin{pmatrix} 1.16 & 2.49 \\ & -2.59 \end{pmatrix}$
Octupole ε_3 = 0.12	-5.64 -4.64	# 69 # 72	$\begin{pmatrix} -1.43 & -1.83 \\ & -4.40 \end{pmatrix}$
Experiment ^{225}Ra	-4.90	# 69	$\left.\begin{matrix} -1.53 \\ -2.59 \end{matrix}\right\}$

Calculated values for the different shapes are listed in Table I. All K = 1/2 orbitals calculated to come within ±1.3 MeV of the Fermi level are included in Table I, except the orbital [501 1/2] which would easily be distinguished if it were observed experimentally; its wave function remains rather pure for all the shapes considered here. The top two orbitals in Table I illustrate a slight complication that may occur in the interpretation of the calculated numbers. Two Nilsson orbitals which would normally have large decoupling factors of opposite sign are mixed, and both the diagonal j_+ matrix elements are reduced, not due to reflection asymmetry but simply because the orbitals happen to come close in energy and therefore mix. Cases like this are easily recognized in the calculations, however, by a small energy spacing combined with a large off-diagonal matrix element of j_+. In the excitation spectrum of the nucleus, the unmixed strongly decoupled orbitals would be restored by this off-diagonal matrix element. For this reason, j_+ matrices and not just diagonal elements are shown in Table I.

From Table I it is seen that neither the spheroid nor the half emerged α cluster can provide candidates for the ^{225}Ra ground state. With the α cluster outside the core, there is some quenching of large negative j_+ values around the Fermi level, but not enough for agreement with the data. With a ^{14}C cluster, on the other hand, orbit #70 is compatible with the data considering the uncertainties of the present calculation. With an octupole shape near the self-consistent equilibrium, orbit #69 is right at the Fermi level of ^{225}Ra and accounts for the observed decoupling factor.

In summary, the odd particle signature seems to exclude both the reflection symmetric and the α cluster alternatives in ^{225}Ra.

3. Probing the Dynamical Fields of Boson Cores.

The previous section discussed an odd particle probing the mean, or static, field from the core. The present section takes up a method in field theory to couple fermions to the dynamical field of a core[24]. The dynamical field of a quantal core is defined by all the energy levels, the transition matrix elements between them and the diagonal moments. The dynamical field of an Interacting Boson Model (IBM) core is thus defined by the output of codes like PHINT or NPBOS[25], and the method can be straightforwardly applied. In connection with the IBM, this "dynamical quasiparticle" method is a distinct alternative to the Interacting Boson Fermion Model (IBFM) of Iachello and Scholten[26,27].

Section 3.1 below reviews the general method and Section 3.2 describes its application to IBM cores; further specifics can be found in Refs. [28,29]. Section 3.3 presents a numerical comparison of the dynamical quasiparticle method with IBFM, for IBM cores with SU(5), SU(3) and O(6) symmetry. A realistic application using the "mixed IBM2" mercury cores of the Tucson group[30,31] is discussed in a separate contribution to this conference[32], and applications to some other kinds of boson cores can be found in the literature[33,34].

3.1 <u>Quasiparticles in dynamical fields.</u> First we derive general field equations for states $|I\rangle$ of the odd-A nucleus, following the equation of motion method worked out by Klein and others[35,36]. Since the final goal is a theory for coupling odd nucleons to cores with known dynamical fields, the quantities entering the field equations will be expressed in terms of the states $|R\ A-1\rangle$ and $|R'\ A+1\rangle$ of the two neighboring doubly even cores. We start by writing down the obvious identity

$$|I\rangle = \frac{1}{\Omega}\sum_j (a_j^+ a_j + \bar{a}_j \bar{a}_j^+)|I\rangle \tag{2}$$

which follows from the anticommutation relation for single-fermion creation and annihilation operators a_j^+ and a_j. Here j denotes the spherical single-particle quantum numbers (nljm), $\Omega = \Sigma(2j+1)$ and $\bar{a}_{jm} = (-)^{j+m} a_{j-m}$. Next, the completeness relations for the two cores are inserted between the Fermion operators, which gives

$$|I\rangle = \frac{1}{\Omega}\sum_{jRR'} a_j^+|RA-1\rangle\langle RA-1|a_j|I\rangle + \bar{a}_j|R'A+1\rangle\langle R'A+1|\bar{a}_j^+|I\rangle \tag{3}$$

The c-number coefficients in this expression are unknowns to be determined from the field equations. Denoting them u_j and v_j,

$$|I\rangle = \sum_{jRR'} u_j a_j^+|RA-1\rangle + v_j \bar{a}_j|R'A+1\rangle\ , \tag{4}$$

it is seen that the states $|I\rangle$ of the odd-A system are quasiparticle states obtained by a Bogolyubov transformation from particle states $a_j^+|RA-1\rangle$ and hole states $a_j|R'A+1\rangle$. These states $|I\rangle$ have good angular momentum and good particle number.

Equations of motion are obtained in the usual way, by forming commutators between the Hamiltonian and the Fermion operators and taking matrix elements between appropriate states

$$\langle I|[h,a_j^+]|R\ A-1\rangle = (E_I - E_R^{A-1})\langle I|a_j^+|R\ A-1\rangle \tag{5}$$

$$\langle I|[h,\bar{a}_j]|R'A+1\rangle = (E_I - E_{R'}^{A+1})\langle I|\bar{a}_j|R'A+1\rangle \tag{6}$$

The expressions on the right are obtained because $|I\rangle$, $|R\ A-1\rangle$ and $|R'A+1\rangle$ are eigenstates of h; they contain the energy eigenvalues, E, and coefficients identical with the ones previously denoted u_j and v_j.

Up to this point the formalism is exact. To proceed it is necessary to specify the Hamiltonian, so we assume a two-body interaction between single-particle states of the most general form that can be written as a multipole expansion

$$h = h_{sp} - \frac{1}{4}\sum_{\lambda\mu} \kappa_\lambda Q_\mu^\lambda Q_{\bar{\mu}}^\lambda - \frac{1}{4}\sum_{\lambda\mu} g_\lambda P_\mu^{\lambda +} P_\mu^\lambda \tag{7}$$

$$h_{sp} = \sum_j \epsilon_j \, a_j^+ a_j \quad \text{(spherical s.p. Hamiltonian)} \tag{8}$$

$$Q_\mu^\lambda = \sum_{jj'} \langle j \,|\, r^\lambda Y_\mu^\lambda \,|\, j' \rangle \, a_j^+ a_{j'} \quad \text{(p.h. multipole operator)} \tag{9}$$

$$P_\mu^{\lambda +} = \sum_{jj'} \langle j \,|\, r^\lambda Y_\mu^\lambda \,|\, j' \rangle \, a_j^+ \bar{a}_{j'}^+ \quad \text{(multipole pair operator)} \tag{10}$$

Inserting this Hamiltonian into the equations of motion gives the field equations

$$\begin{bmatrix} \epsilon + E_R^{A-1} + \Gamma^{A-1} & -\Delta^+ \\ -\Delta & -\epsilon + E_{R'}^{A+1} - \Gamma^{A+1} \end{bmatrix} \begin{bmatrix} u \\ v \end{bmatrix}_I = E_I \begin{bmatrix} u \\ v \end{bmatrix}_I \tag{11}$$

where

$$\Gamma^{A\pm 1} = -\frac{1}{2} \sum_\lambda \kappa_\lambda \, \langle j \,||\, r^\lambda Y^\lambda \,||\, j' \rangle \, \langle R \;\, A\pm 1 \,||\, Q^\lambda \,||\, R' \;\, A\pm 1 \rangle \, W(j'\lambda IR;jR') \tag{12}$$

$$\Delta = \frac{1}{2} \sum_\lambda g_\lambda \, \langle j \,||\, r^\lambda Y^\lambda \,||\, j' \rangle \, \langle R \;\, A-1 \,||\, P^\lambda \,||\, R' \;\, A+1 \rangle \, W(j'\lambda IR;jR') \tag{13}$$

The set of solutions $|I\rangle$ to the field equations in the particle and hole basis is overcomplete by a factor of two, so a density matrix must be constructed and used as a projector to identify the physical subset. This could lead to a self-consistent microscopic theory of the nucleus, which is not desirable if the core field is to be input. Therefore self-consistency is bypassed through projection with the approximate density matrix that is obtained from the "adiabatic" part of the field Hamiltonian only[24,28], in other words the part that is even under the conjugation $a_j^+ \longleftrightarrow \bar{a}_j$.

3.2 Models using IBM cores. In practical calculations the core space must be truncated to a set of states connected by large P^λ and Q^λ matrix elements, and boson cores are clearly suitable. The most complex form of the boson-fermion Hamiltonian actually used so far[29] includes a quadrupole-quadrupole interaction and quadrupole pairing,

$$H = H_{IBA}^{A\pm 1} + h_{sp} - \frac{\kappa}{2} \sum_\mu Q_\mu^2 \, q_\mu^2 - \frac{g_o}{2}(P_o^+ p_o + P_o p_o^+) - \frac{g_2}{2} \sum_\mu (P_\mu^{2+} p_\mu^2 + P_\mu^2 \, p_\mu^{2+}) \tag{14}$$

where the single-particle operators are

$$q^2_\mu = \sum_{jj'} \langle j \,|\, r^2 Y^2_\mu \,|\, j' \rangle \; a_j{}^+ a_{j'} \tag{15}$$

$$p^{\lambda +}_\mu = \sum_{jj'} \langle j \,|\, r^2 Y^\lambda_\mu \,|\, j' \rangle \; a_j{}^+ a_{j'}{}^+ \tag{16}$$

and the collective pair operators are

$$P_\circ{}^+ = \frac{\Delta_\circ}{g_\circ} s^+ \;;\; P^{2\,+}_\mu = \frac{\Delta_2}{g_2} d^+_\mu \tag{17}$$

A model to compare with IBFM is obtained with

$$H^{A+1}_{IBM} = H^{A-1}_{IBM} = H_{IBM1} \tag{18}$$

$$Q^2_\mu = e_2 (s^+ \bar{d}_\mu + s d^+_\mu + \chi \; (d^+ \bar{d})^2_\mu) \tag{19}$$

$$\frac{g_\circ}{2} \langle R' A+1 \,|\, P_\circ{}^+ \,|\, R A-1 \rangle = \Delta \delta_{RR'} \tag{20}$$

$$g_2 = 0 \tag{21}$$

Then the only difference between the models lies in the handling of exchange effects between the fermion and the core. In IBFM, exchange effects enter through the spherical occupation probability v^2 and through an explicit exchange term i.e.,

$$\begin{array}{ll} \Gamma^{jj'}_\circ (1-2v^2) \times & \text{quadrupole coupling and} \\ \Lambda^{jj'}_\circ \sqrt{v^2(1-v^2)} \times & \text{exchange term,} \end{array} \tag{22}$$

while in the dynamical quasiparticle method they arise through the Bogolyubov transformation.

Both models have the same number of parameters. Both the v^2 of IBFM, and the Fermi level λ used as a reference for the single-particle energy ε in the dynamical field equations, are determined by the particle number. The $\Gamma_\circ$ of IBFM equals κ_2 above to within a constant factor such that the models are identical for completely filled or empty j-shells. There remains one parameter for each model: $\Lambda_\circ$ and Δ, respectively.

3.3 Numerical comparison with IBFM. The dynamical quasiparticle method and IBFM will be compared in numerical calculations for a j = 9/2 shell coupled to an SU(5), SU(3) or O(6) core. For a given core, particle number and quadrupole coupling strength, the energy levels most likely to be observed in experiment will be plotted versus the parameter Δ or Λ_0, respectively. The core and the quadrupole coupling parameter are identical with those used in Refs. [25,27], with the minor exception of taking χ in Eq. (19) equal to zero for the SU(5) core, which gives vanishing diagonal core quadrupole matrix elements.

Results for an SU(5) core are shown in Fig. 3a. The case which is shown has $v^2 = 0.2$ and exchange effects play a significant role. Clearly, similar spectra can be obtained by varying the parameters Δ and Λ_0 of the two models, which suggests that the two different ways of handling exchange effects are in fact roughly equivalent on the phenomenological level. For the O(6) core and $v^2 = 0.3$ in Fig. 3b, changes of Λ_0 have a similar but somewhat stronger effect on the spectrum compared to changes of Δ.

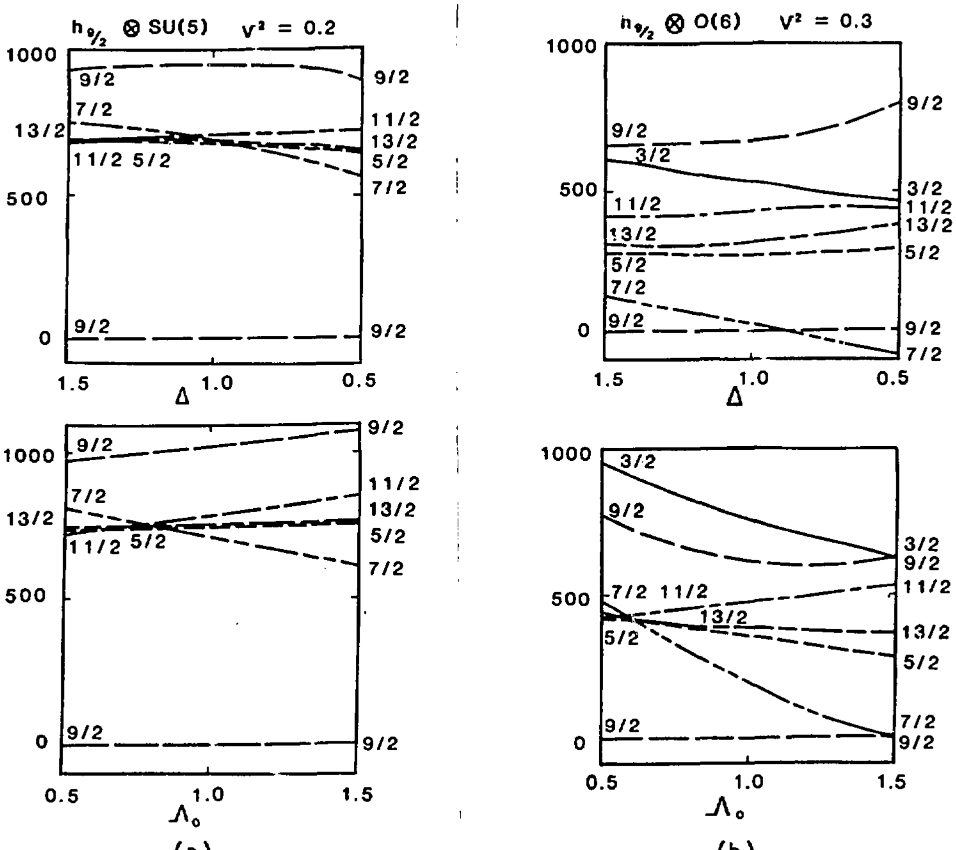
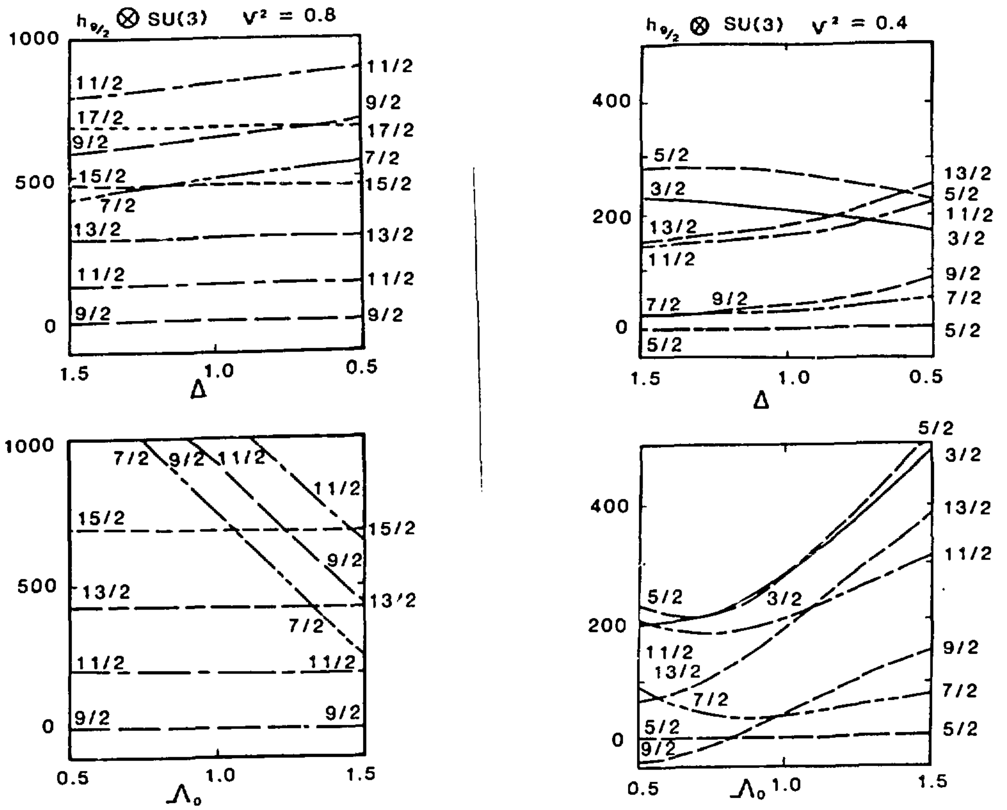


Fig. 3. Energy levels for a core plus a j = 9/2 particle in a partly filled shell. The energies relative to the I = 9/2 level obtained by the dynamical quasiparticle method are plotted as functions of the pairing gap parameter Δ (above), and by the IBFM as functions of the exchange term parameter Λ_0 (below). Case (a) has an SU(5) core and the occupation of the j-shell $\nu^2 = 0.2$, (b) an O(6) core and $\nu^2 = 0.3$.



(a) (b)

Fig. 4. Same as Fig. 3, but for an SU(3) core and (a) $\nu^2 = 0.8$, (b) $\nu^2 = 0.4$.

However, Δ and Λ_o are not in general equivalent parameters. From the way that v^2 enters the IBFM Hamiltonian, expressions (22) above, it is clear that at mid-shell ($v^2 = 0.5$), a variation of Λ_o is exactly equivalent to a variation of v^2, which would correspond to a variation of the Fermi level λ in the dynamical quasiparticle formalism. A connection between Λ_o and λ is also obvious in Fig. 4a, which has the results for an SU(3) core at $v^2 = 0.8$. The rotational model analog would be a prolate rotor with the Fermi level λ just below the K = 9/2 Nilsson orbital. An increase of Δ gives only a slight lowering of the K = 7/2 band head, related to the familiar compression of BCS quasiparticle levels, whereas an increase of Λ_o has an effect similar to moving the Fermi level further down into the shell. Results for the SU(3) core and $v^2 = 0.4$ are shown in Fig. 4b. Clearly more variation of the spectrum can be achieved with the parameter Λ_o than with Δ, however, the same range of spectra could have been achieved by adjusting both Δ and λ.

An interesting question is whether the IBFM introduces any "Coriolis attenuation". The level ordering cannot be used as a prime indicator in view of the discussion above. However, the levels appear to be consistently more spread out in energy in the IBFM which does indicate some kind of attenuation. For example, the K = 9/2 band in

Fig. 4a from the dynamical quasiparticle theory has a moment of inertia 1.38 times that of the core, ostensibly due to Coriolis coupling to the K = 7/2 band, whereas the K = 9/2 band from IBFM has a moment of inertia actually smaller than that of the core.

To conclude this section we shall look for an empirical connection between the two models by focussing on parameter values which give similar spectra. As a reference we take the dynamical quasiparticle spectrum at Δ = 1.0 MeV, which is a physically reasonable value. There is a range of Λ_o values for which IBFM gives a similar level spectrum, see e.g. Figs. 3 and 4. An estimate of the optimal Λ_o is indicated in Fig. 5 for each core and several values of v^2. Keeping κ_2 and Δ fixed and varying λ through the shell in the dynamical quasiparticle approach is seen to be roughly equivalent to keeping Γ_o and Λ_o fixed and varying v^2 through the shell in IBFM. The value of Λ_o corresponding to Δ = 1.0 MeV is about 0.8 MeV for the SU(5) core and 1.2 MeV for the SU(3) and O(6) cores. The value of Λ_o can be expected to scale with Γ_o. However, since the same Γ_o was used here with all three cores[27], it is not obvious why Λ_o comes out smaller in the SU(5) case.

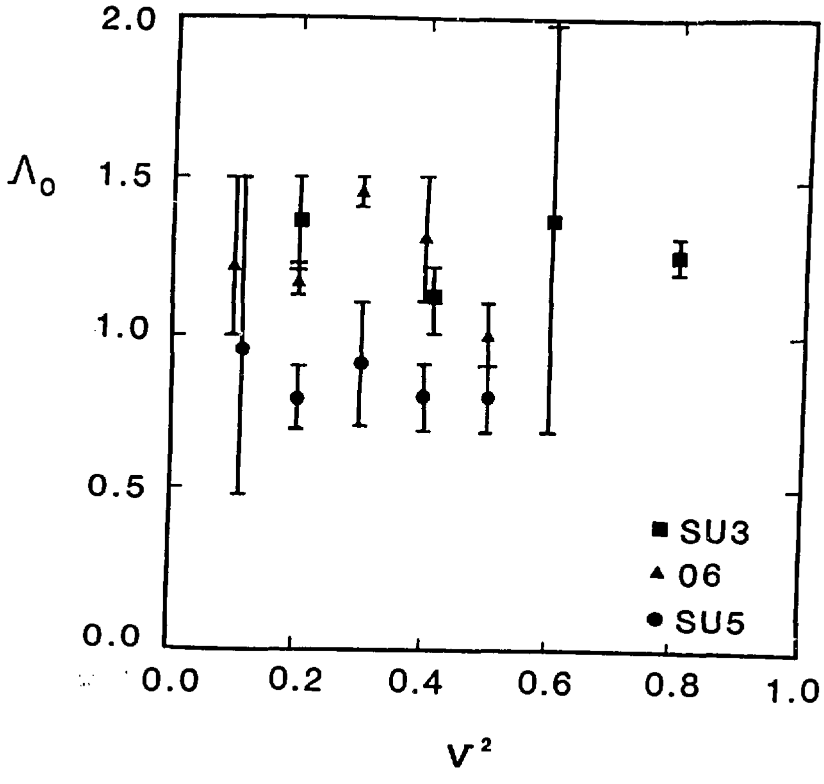


Fig. 5. The Λ_o values for which the IBFM spectrum is most similar to the spectrum obtained by the dynamical quasiparticle method with Δ = 1 MeV, for three different cores and different occupation probabilities in the j-shell, ν^2. The two methods give identical spectra at ν^2 = 0 and 1 for all values of Λ_o, and they are symmetrical around ν^2 = 0.5 for the SU(5) and O(6) cores, so those points are not shown.

4. Conclusions and Discussion.

Molecular configurations that have been speculatively inferred from boson models for the Ra nuclei were tested by studying the motion of odd nucleons in the intrinsic mean fields implied by such configurations. It was found that the interpretation of the p-boson in terms of an α cluster and a reflection symmetric core does not account for the decoupling factors - or the K = 1/2 ground state - of ^{225}Ra. It was previously known that these and other single-particle properties are well described with the octupole shapes obtained from mean field theory[11]. Experimental data for the somewhat lighter odd-A isotopes are presently scarce but would be very interesting.

The spectroscopy of odd-mass nuclei can be used more generally to test the description of collective modes in boson models. In order to calculate odd-A spectra using boson cores, the dynamical quasiparticle method is put forth as a viable alternative to IBFM. On a phenomenological level, the two models were shown to be similar with respect to easily measurable energy levels (c.f. Refs. [37,38] for a discussion of not so easily observable differences). However, the dynamical quasiparticle method has some practical advantages. It is applicable to any kind of core with equal ease. Its parameter Δ is more "physical" than the corresponding IBFM parameters $\Lambda^{jj'}$ in the sense of being more readily available from theory or experiment. Extensions to include for example the quadrupole pair field are straightforward. The problem of single-particle transfer should also be mentioned. The correct form of the operator is unknown in IBFM, as was emphasized at this meeting in the session on supersymmetries. The dynamical quasiparticle theory does not bring out supersymmetries; however the explicit presence of fermion operators and two cores in the formalism does imply that spectroscopic factors are unambiguously obtained from matrix elements of the type $\langle R\ A\pm1 | a_j + a_j^+ | I \rangle$.

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