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NUCLEAR TRANSITIONS FROM ALIGNED STATES: EXAMPLE OF OCTUPOLE BANDS IN ACTINIDES

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I - INTRODUCTION

The negative parity states in deformed nuclei may be divided schematically into the collective states generated by octupole vibrations of the nuclear surface and those belonging to the two-quasiparticle bands. The large collectivity of the $K = 0^-, 1^-, 2^-$ excitation was shown through RPA calculations [1-4]. The Coriolis coupling of these excitations influences the spectrum of negative parity states already at moderate spin values (I). The alignment of vibrational angular momentum (j) in such states in actinides is felt at lower values of I as the effect of rotational activation of the two-quasiparticle degrees of freedom[4]. The existing experimental data[5-10] on the negative parity states contain appreciable information, both on the spectra and on the electromagnetic transitions from these states at $I \gg 1$, which stimulates the search for the general properties of nuclear transitions from the aligned nuclear states. In this paper an attempt is made to give a qualitative description of nuclear transitions from the aligned states taking as example octupole states at $I \gg 1$. The latter are treated here within a very simple theoretical model which, however, allows to reproduce many experimental results concerning the negative parity states in actinides.

II - THE MODEL

The following study of the structure of octupole states influenced by the strong Coriolis force is based on the Hamiltonian

$$H = \frac{1}{2J} \sum_{i=1}^{3} (\hat{J}_i - \hat{j}_i)^2 + \sum_{|K|\leq 3} \omega_{|K|}\, b_K^+ b_K , \qquad (1)$$

where the first term describes the kinetic energy of rotation of the nuclear core while the second term gives the energy of octupole oscillations. The states generated by the operators b_K^+ are supposed to have fixed values of the quantum numbers of intrinsic angular momentum ($j = 3$) and of its projection (K) on the axial symmetry axis ($-3 \leq K \leq 3$). The symmetry axis is taken to be the third of the axes of the intrinsic frame of reference.

When $I \gg 1$ the Coriolis term in (1) may be approximated by

$$\frac{1}{J} \sum_{i=1}^{3} \hat{J}_i\, \hat{j}_i \approx \frac{1}{J}\, I \cdot \hat{j}_1 \quad . \qquad (2)$$

With this approximation the study of (1) is simplified by going from the representation where j_3 is diagonal (K-representation) to the one where j_1 is diagonal (τ-representation). In the latter

$$(H)_{\tau\tau'} = \left(\frac{1}{2J} I(I+1) - \frac{\tau}{J}\right) \delta_{\tau\tau'} + \sum_{K\geq 0} d^3_{K\tau}(\tfrac{\pi}{2})\, d^3_{K\tau'}(\tfrac{\pi}{2}) \frac{\omega_K (1+(-1)^{\tau+\tau'})}{(1+\delta_{K,o})} \qquad (3)$$

$d^\lambda_{K\tau}(\theta)$ being the middle part of the Wigner function. As is seen, the Hamiltonian matrix (3) splits into two blocks each operating among the states with the same parity of τ. Even (odd) values of τ generate the states with positive (negative) signature and, consequently, the nuclear states with even (odd) values of I.

Much could be learned considering the special case when

$$\omega_1 = \omega_2 = \omega_3 = \omega_o + \Delta . \qquad (4)$$

The following analysis is restricted to the situation where equation (4) is satisfied.

The positive signature states under the condition (4) are characterized by the even numbers $j_x = 0, \pm 2$ which are equal to the aligned vibrational angular momentum

$$j_x = \langle \psi^{(\pm)} | \hat{j}_1 | \psi^{(\pm)} \rangle . \qquad (5)$$

In (5), $\psi^{(\pm)}$ is the eigenvector of (3) corresponding to the positive (negative) signature. For positive signature, $\hat{j}_I \psi^{(+)} = j_x \psi^{(+)}$ and the energy is :

$$E^+_{j_x}(I) = \frac{1}{2J}\left[(I - j_x)(I - j_x + 1)\right] + \omega_o + \Delta - j_x (j_x - 1)/2J . \qquad (6)$$

The last term in (6) is in fact an error introduced in the exact expression for the energy by the approximation in eq.(2). As is seen, this approximation does not affect the energy differences between the states with the same j_x but shifts the band-heads. The condition (4) is too specific to apply the model for the analysis of the positive signature states and we turn to the other case.

The energies of the states with negative signature are given by the formula :

$$E^{(-)}_\lambda(I) = \omega_c + \frac{1}{2J} I(I+1) + \Delta . \chi . \varepsilon_\lambda(\chi), \quad \chi = \frac{I}{2J.\Delta} = \frac{1}{2}\frac{\omega_{rot}}{\Delta}, \qquad (7)$$

where $\varepsilon_\lambda(\chi)$ satisfies the equations

$$\varepsilon = 1/\chi - Z(\chi),$$

$$Z\{ Z\{ Z\{ Z - 1/\chi\} - 40 \} + 16/\chi\} + 144 = 0 . \qquad (8)$$

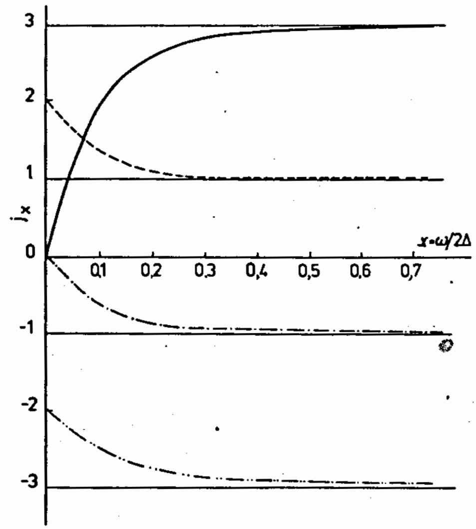


Fig. 1 - Aligned vibrational angular momentum in the octupole one-phonon negative signature states versus the rotational parameter

$$\chi = \frac{1}{2}\frac{I}{J} = \frac{\omega_{rot}}{2\Delta} .$$

The lowest of the solutions of eqs. (7), (8) describes the $K^\pi = 0^-$ band distorted by the Coriolis force.

The energy intervals inside the bands depend on the aligned angular momentum :

$$E_\lambda(I+1) - E_\lambda(I-1) \approx 2\frac{dE_\lambda(I)}{dI} = \frac{2}{J}\left[(I+1/2) - j_x\right] . \qquad (9)$$

The aligned angular momentum in the negative signature bands is shown on fig. 1 as a function of the rotational parameter χ defined in eq.(7). To the 0^- band there corresponds a solution to eqs.(7),(8) with the aligned angular momentum $j_x \to 3$ at $\chi \to \infty$. In fact the limiting values of j_x are quite close for all the bands already at $\chi = 1$.

III - ANALYSIS OF SPECTRA OF 0^- BANDS IN ACTINIDES

The changes in the nuclear structure at large angular momenta may be visualized plotting the aligned angular momentum as function of the rotational frequency of the core and comparing the above estimations with the correspondingly treated experimental data.

Using eq.(9) one obtains :

$$j^{exp}_x = (I+1/2) - J\,\omega^{0^-}_{exp}, \quad \omega^{0^-}_{exp} = \frac{1}{2}\left(E_{0^-}(I+1) - E_{0^-}(I-1)\right) . \qquad (10)$$

We use this formula adopting the definition for the moment of inertia J given in ref.(11).

$$J = J_o + \omega^2 J_1 \quad ; \quad I = J . \omega . \qquad (11)$$

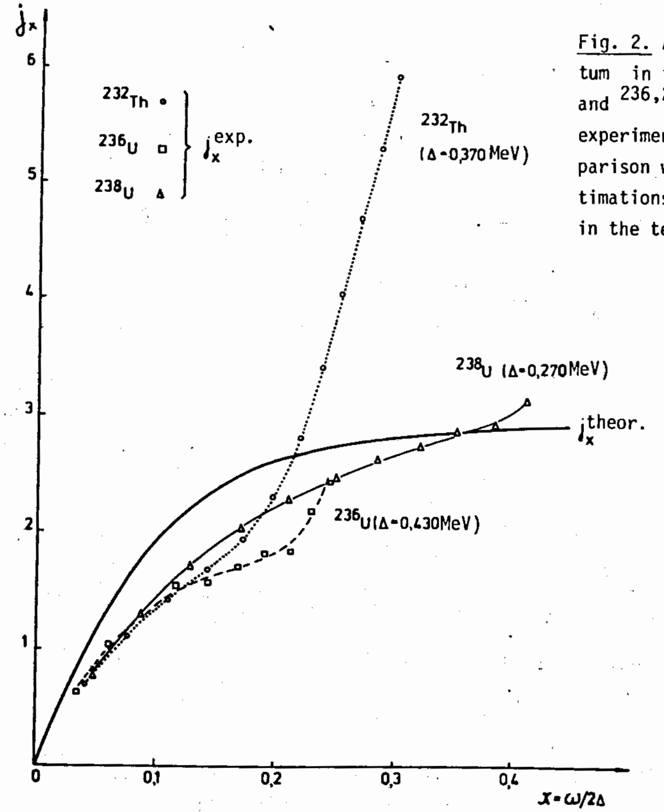


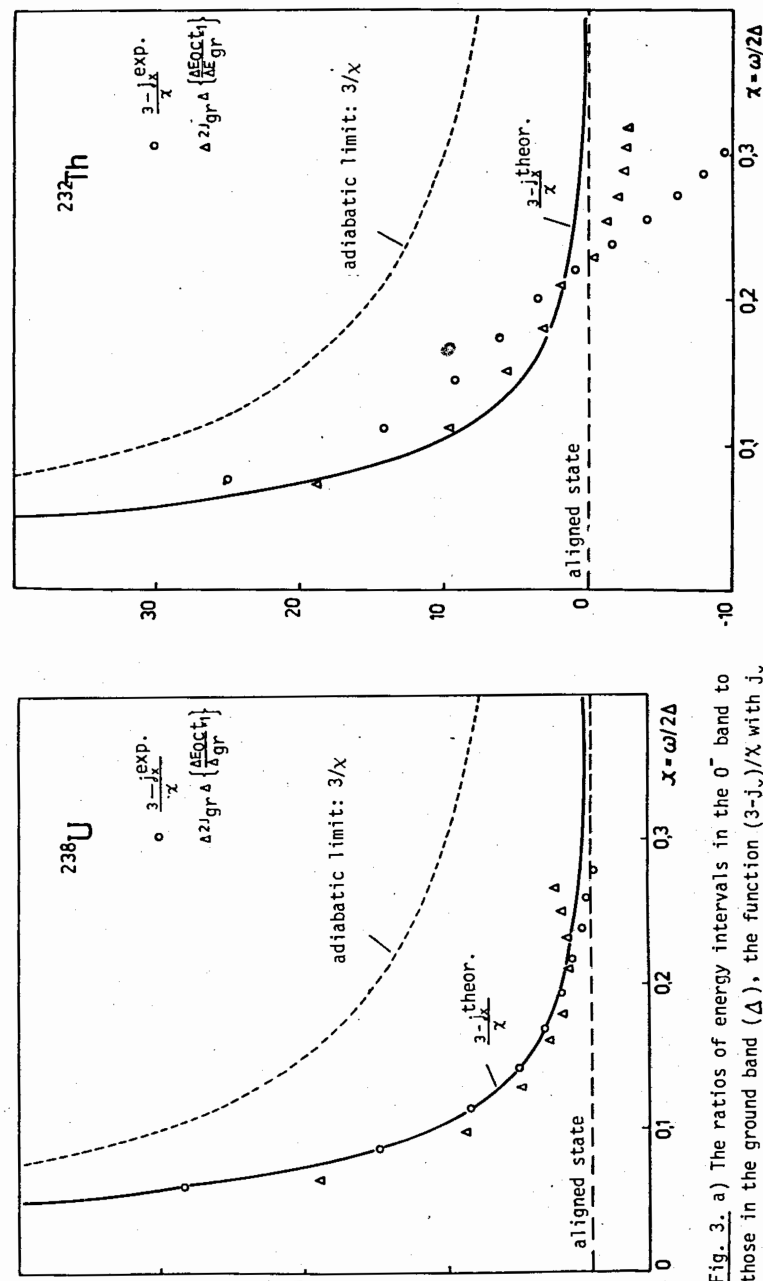
Fig. 2. Aligned angular momentum in the 0^- bands in ^{232}Th and $^{236,238}U$ deduced from the experimental data (6-9) in comparison with theoretical estimations (see explanations in the text).

The parameters J_0 and J_1 are chosen so as to reproduce the energies in the ground state band at low spins.

Fig. 2 shows the aligned angular momenta deduced in this way from the experimental data on ^{232}Th (6,7,8,10), ^{238}U (5,9) and ^{236}U(8) in comparison with the above model. The mean value of RPA calculations in(12) for the states $K^{\pi} = 1^-$, 2^- and 3^- is taken for $\omega_0 + \Delta$. The slightly different representation of the effects of alignment is given on fig. 3.
In the model, the following relation holds

$$\left\{ \frac{[E\,(I+3) - E\,(I+1)]_{oct}}{[E(I) - E\,(I-2)]_{gr}} - 1 \right\} \cdot 2\,J\,\Delta = \frac{3 - j_x(\chi)}{\chi} \qquad (12)$$

Both sides of (12) may be calculated independently using the experimental data. The moment of inertia J enters in a different way the left and right parts of eq.(12). Thus, the same set of experimental energies generates in fig. 3 two families of points.



b) the same for ^{232}Th.

Fig. 3. a) The ratios of energy intervals in the 0^- band to those in the ground band (Δ), the function $(3-j_x)/\chi$ with j_x deduced from the experimental data (O) and the theoretical estimations for the function $(3-j_x)/\chi$ (solid line). The limiting cases of complete alignment and no alignment are indicated by dashed lines. Case of ^{238}U.; b) the same for ^{232}Th.

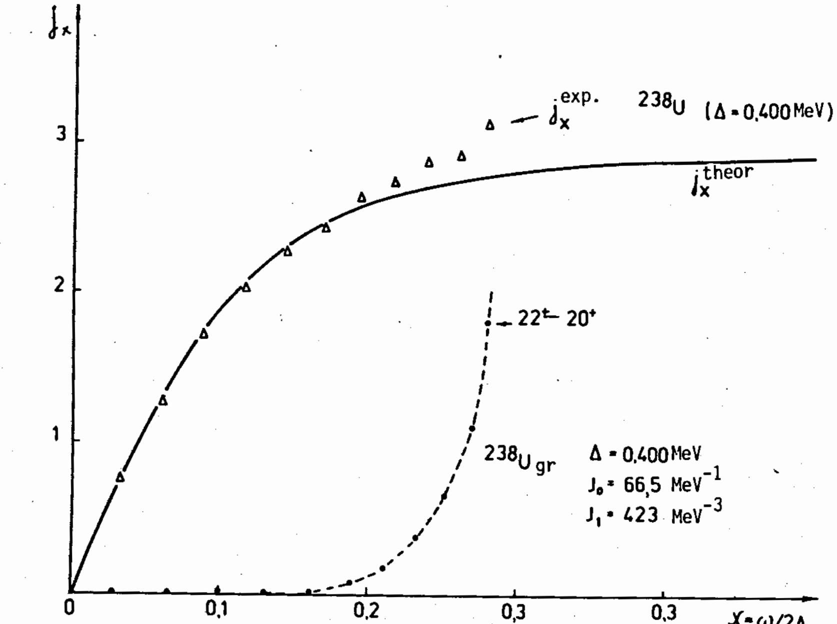


Fig. 4 - The same as in fig. 2 at a different value of Δ

At I $\lesssim$ 15 the schematic model is in reasonable agreement with the experimental data : at small I (small χ) the aligned angular momentum tends to zero. Correspondingly, the energy intervals are close to the adiabatic limit. At large spins (I $\gtrsim$ 9), j_x^{theor} is close to 3 and the numerator in (12) tends to zero. The experimental data show that at such spins the alignment is important. The agreement between the theory and experimental points depends on the choice of parameters Δ and J. Fig. 4 shows the aligned angular momentum in ^{238}U calculated using the same techniques as described before at a slightly different choice of Δ, as suggested in[13] on the basis of more detailed calculations of a semi-microscopic nature. In the quoted paper, the known experimental data on the negative parity states in ^{238}U are well reproduced with the parametrization of the bandheads corresponding to such choice of Δ. Here the agreement with the theoretical estimations at spins I $\lesssim$ 13 is quite good. At larger values of I the quasiparticle degrees of freedom become important which is indicated by the steep rise of the curve showing the aligned angular momentum in the ground band of ^{238}U (see fig. 4). The RPA calculations[12] allow to make some guesses about the differences between the three nuclei (^{232}Th, ^{236}U, ^{238}U) represented in fig. 2, considering the main two-quasiparticle components in the RPA bosons with negative parity. These include j15/2 neutron states with small K values. Due to the blocking

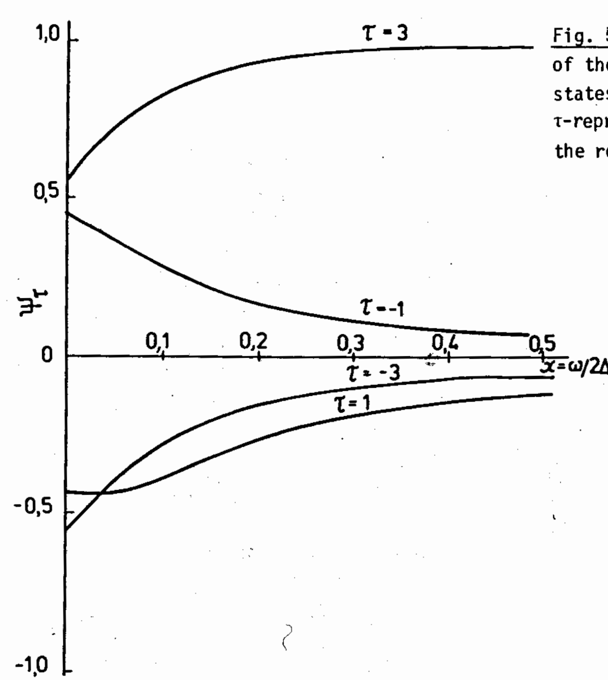


Fig. 5 - The structure of the wave function of the states of 0^- band in the τ-representation versus the rotational parameter χ.

effect, they cannot contribute to the observed extraneous aligned angular momentum associated with the alignment of a proton pair. This is consistent with the conclusions of the set of experiments on odd and even actinides[14-17].

IV - STRUCTURE OF ALIGNED STATES AND THE BRANCHING RATIOS IN ELECTROMAGNETIC TRANSITIONS.

The model described in Sec. 2 allows to study the changes in the structure of nuclear states with the increase of the aligned angular momentum The eigenvectors of the Hamiltonian matrix (3) determine weights with which the octupole bosons carrying a fixed amount of the first projection of the vibrational angular momentum.

$$b_\tau^+ = \sum_K d^3_{K\tau}\left(\frac{\pi}{2}\right) b_K^+ \qquad (13)$$

enter the wave function of a nuclear stationary state. The expression

$$|\,\psi_\lambda^{(\pm)}\rangle = \sum_\tau \psi_\tau^{(\pm)}(\lambda)\, b_\tau^+\,|\text{core}, I\rangle, \quad (b_\tau\,|\text{core}, \; I\;\rangle = 0) \qquad (14)$$

has the meaning of the intrinsic wave-function of a rotating nucleus [18,19]. As said before, in the case when eq.(4) holds $\psi_\tau^+ = \delta_{j_x,\tau}$ $(j_x = 0, \pm 2)$ and the states $|\,\psi^{(+)}\rangle$ do not change with I. The structure of the intrinsic wave - function of the lowest negative-signature band (0^- band) in the τ-

representation is shown in fig. 5. The above arguments show that the aligned angular momentum at large I determines the leading component of the wave-function. In the case of the 0^- band the leading component becomes $\psi^{(-)}_{\tau=3}$. Comparing figs. 2-4 and fig. 5, one notices that in the actinide nuclei the effects of rotation are well pronounced in the structure of the wave-function at already very moderate values of I. This must lead to certain consequences concerning the electromagnetic transitions from such states, and we start the discussion of the possible effects.

When the angular momenta are large, one may use the following expression, for the reduced matrix element of transition between the states $|\alpha_1\ I_1\ M_1>$ and $|\alpha_2\ I_2\ M_2>$ taken from the multipole operator $\hat{M}(\lambda,\mu)$ ($\lambda \ll I_1,\ I_2$) :

$$<\alpha_2\ I_2 || \hat{M}(\lambda) ||\ \alpha_1\ I_1> \ = \ \sqrt{2I}\ <\psi_{\alpha_2 I_2} | \hat{M}'\ (\lambda, \tau = I_2 - I_1) |\ \psi_{\alpha_1 I_1}>, \quad (15)$$

where $|\psi_{\alpha I}>$ stand for the intrinsic states (14) and $\hat{M}'$ is the intrinsic part of the multipole operator in the τ-representation [18,19].

For the transitions between the one-phonon and ground bands, eq.(15) becomes :

$$<gr\ I_2 || \hat{M}\ (\lambda) || \text{oct}\ I_1> = \sqrt{2I}\ \sum_{\tau'} \psi_{\tau'} <core\ I_2 |\ [M'(\lambda, \tau = I_2 - I_1), b^+_{\tau'}]\, | core\ I_1> . \quad (16)$$

In the limit when $\psi_{\tau'} \to \delta_{\tau' j}$ (the "alignment limit") one has:

$$<gr\ I_2 ||\ \hat{M}(\lambda) ||\ \text{oct}\ I_1> \to \sqrt{2I}\ <\text{core}\ I_2 | [\hat{M}'(\lambda, \tau = I_2 - I_1),\ b^+_j] | core\ I_1>. \quad (17)$$

Thus, the transitions from the band marked with the symbol j to the ground band, with different transferred angular momenta $\Delta I = I_1 - I_2$, are generated by different components of intrinsic multipole moments $M'(\lambda,\tau)$. Consequently, one expects different behaviour of such transitions. If the core function $|core, I>$ and the amplitudes ψ_τ for the states belonging to one particular band change smoothly with I, the transitions from this band at a fixed value of ΔI must also smoothly depend on I.

The emerging pattern of transition rates is more transparent for octupole electric transitions generated by the field

$$\hat{M}(E3,\mu) = m_o \sum_\tau \mathcal{D}^3_{\mu\tau}\ (\phi,\theta,\psi)\ (b^+_\tau - (-1)^\tau\ b_{-\tau}) + \quad \text{terms} \quad (18)$$

absent in the transitions from octupole bands to ground band.

In this case :

$$\hat{M}'_\tau = m_o\ \ (b^+_\tau - (-1)^\tau\ \ b_{-\tau}\ ) \quad (19)$$

and eqs.(16),(17) lead to the following expressions :

$$B(E3,\ \text{oct}\ I_1 \to \text{gr}\ I_2) = m_o^2\ |\psi_{\tau = I_2 - I_1}|^2 \to m_o^2\ \ \delta_{I_2 - I_1\ j}\ . \quad (20)$$

The octupole transitions $I_1 \to I_2$ carry away ΔI units of aligned angular momentum, and the probability of such transitions is proportional to the weight of the corresponding phonon state. In the alignment limit the octupole transitions from the band j obey the selection rule

$$I_2 - I_1 \ = \ j\ .$$

In particular, E3-transitions form the lowest band (j=3) at large I are stretched transitions : in this case $I_1 - I_2 = 3 = \max(\Delta I)$.

The probability of octupole transitions at an energy around 1 MeV is extremely small, and the states of octupole bands decay emitting dipole γ-rays. The description of this process in the framework of the model needs further assumption concerning the nature of the intrinsic dipole moments $\hat{M}'(E1)$. The simple expressions like the one in eq.(19) are not applicable here because of the different multipolarity of $\hat{M}(E1,\mu)$ and of the octupole bosons. The dipole transitions may be treated introducing the new degrees of freedom into the intrinsic Hamiltonian which describe excitations of dipole type and also the term for the coupling between the octupole and dipole modes :

$$H_{intr} = \sum_{K=-3}^{3} \omega_{|K|}\, b^+_K\, b_K + \sum_{K'=-1}^{1} \omega^{(d)}_{|K'|}\ d^+_{K'}\ d_{K'} \ + \ H_{oct,\ dip}\ . \quad (21)$$

The part of the coupling term $H_{oct,dip}$, essential for the present study of decay properties of the one-phonon octupole states, may be written in the form:

$$H_{oct,dip} = \sum_{\lambda=2,4} h_\lambda \sum_{\mu\mu'\mu''} (-1)^{\mu+\mu''}\ C^{\lambda\mu''}_{1\mu' 3\mu} (d^+_{\mu'} b_{-\mu} + (-1)^{\mu-\mu'} b^+_\mu\ d_{-\mu'})\ \hat{Q}_{\lambda\mu}\ . \quad (22)$$

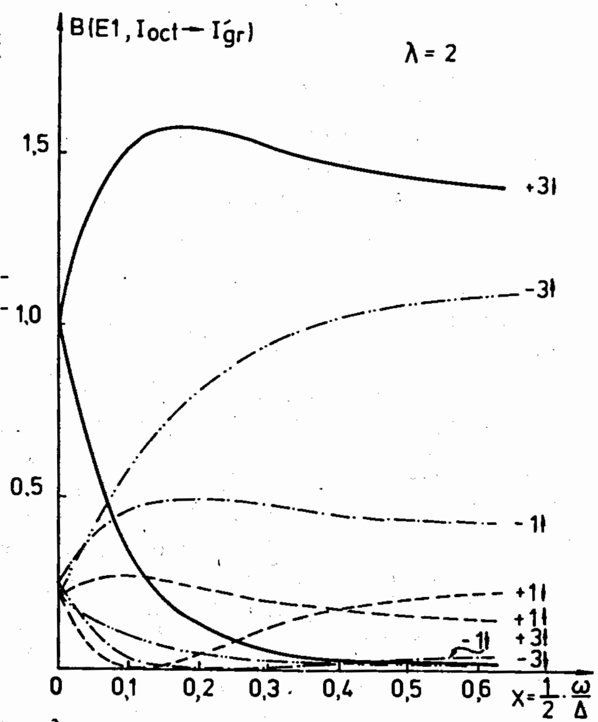
In (21),(22) d^+_μ stand for the creation operators of dipole excitations, and $Q_{\lambda\mu}$ is the operator with rotational quantum numbers $\lambda\mu$ operating on the core function. Eq.(22) is a direct consequence of a spherical symmetry of the total Hamiltonian. In the following expressions enter only the moments

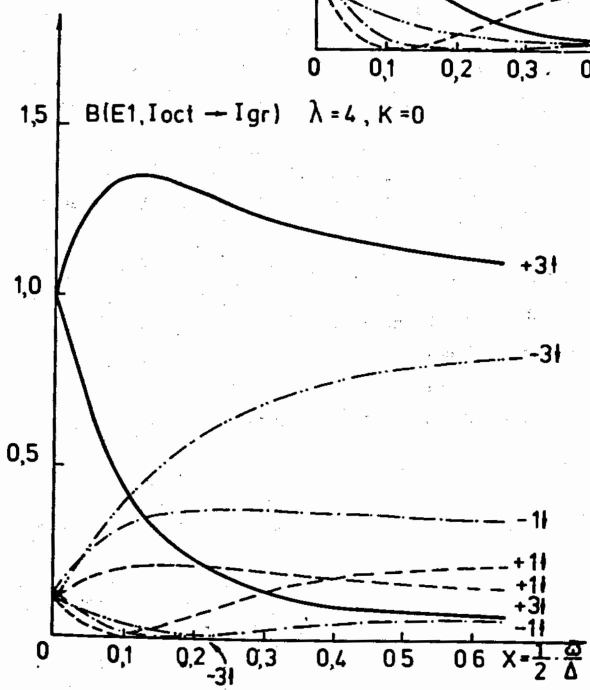
$$q_{\lambda\tau} = \ <\text{core}\ I' |\ \hat{Q}'_{\lambda\tau}\ | \text{core}\ I> \ , \quad (23)$$

which are the matrix elements of the intrinsic parts of $\hat{Q}_{\lambda\mu}$ in the τ-representation.

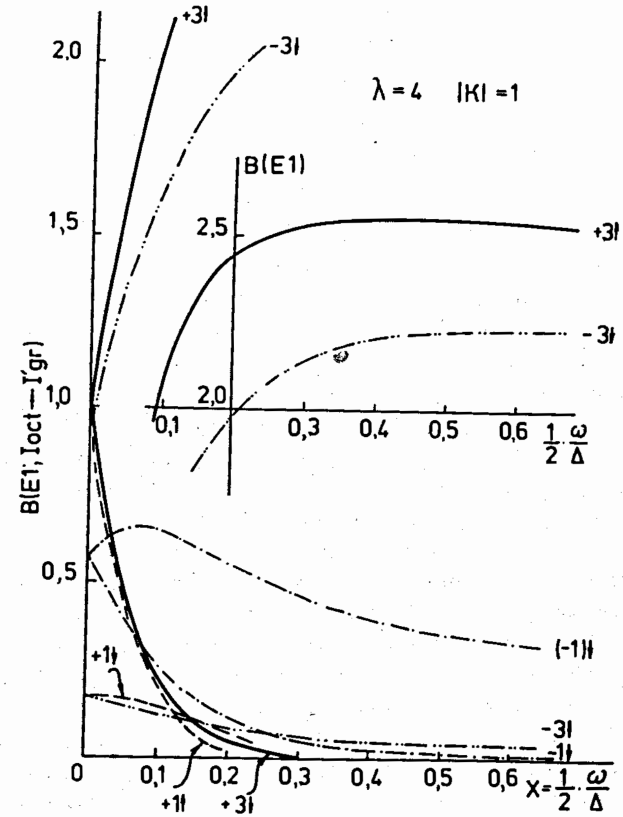
Experimentally the B(E1) factors for the transitions from octupole bands are found to be small fractions of a single-particle unit. This suggests a perturbative treatment of the coupling term in eqs.(21),(22). Writing :

Fig.6. a) The B (E1, j_x I→gr I') - factors for the quadrupole dipole-to-octupole coupling. Labels ↑(↓) correspond to the transitions j_x I gr, I-1 (j_xI→gr,I+1). The limiting value of j_x in the band is indicated near the corresponding curve. The normalisation is such that B(E1 0^-I→gr I')=1 at X=0.;





b) The same for the case λ = 4, K = 0 ;



c) The same for the case λ = 4, K = 1.

$$\hat{M}'(E1,\tau) = m_d\,(d^+_\tau - (-1)^\tau\, d_{-\tau}) + \ldots \tag{24}$$

one comes to the following expression for the B(E1)- factors :

$$B(E1, oct\ I_1 \to gr\ I_2) = m_d^2 \left| \sum_{\lambda=2,4} \sqrt{\frac{2\lambda+1}{7}} \cdot \frac{h_\lambda}{\Delta\omega} \cdot \sum_{\tau'} C^{3\tau'}_{1,-\tau;\lambda\tau+\tau'}\, q_{\lambda,\tau+\tau'}\, \psi_{\tau'} \right|^2 , \tag{25}$$

where $\Delta\omega = \omega^{dip} - \omega^{oct}$ is assumed to be independent from τ.

The dependence of the B(E1) factors on the rotational parameter χ is shown on fig. 6 for different parametrizations of the moments $q_{\lambda\tau}$. Fig.6a corresponds to a pure "quadrupole coupling" ($h_4 = 0$), figs. 6b,c correspond to a pure "hexadecapole coupling" ($h_2 = 0$). In the calculations shown in fig. 6a,b the moments $q_{\lambda\tau}$ were associated with static axial deformation :

$$q_{\lambda K} = q_\lambda\, \delta_{K,0}$$

while in those on fig.6c the $q_{\lambda K}$ are parametrized according to the formula

$$(h_2 = 0), \quad q_{4K} = q'_4 (\delta_{K,1} - \delta_{K,-1}) \,. \tag{26}$$

One may think that the coupling generated by the Coriolis force must go predominantly through the components $Q_{\lambda,\, K = \pm 1}$ and thus favours the parametrization in eq.(26). The normalization in figs. 6 is such that at $\chi = 0$ the strongest transitions have the B(E1) factors equal to one.

Figures 6 clearly show the selectivity of dipole transitions in respect to the angular momentum transferred from the nucleus. This is especially clear in the case of transitions from the distorted 0^- band. But in contrast with the E3-transitions, selectivity is not universal : it depends on the type of coupling between the octupole and dipole bands. The relations between the different B(E1) factors appear at the first sight rather complicated but they have in fact a simple geometrical interpretation visualized in fig. 7 . In the lowest octupole band at I>>1 the vibrational angular momentum is parallel to the angular momentum of the core (see the left part of fig. 7). To reach such state in the field represented by the dipole operator applied to the wave functions of rotating core one has to ensure that the angular momenta involved in the transition add in a proper way. The rules for additions of angular momenta at high spins are just the rules of classical mechanics. One has to add to the collective rotational angular momentum of the ground-band state (right part of fig.7), the two components of the angular momentum of the dipole operator (middle part of fig. 7) : the intrinsic octupole angular momentum (j = 3) and the collective rotational part (λ = 2 or 4) compensating the two extra units of Planck of the vibrational angular momentum in $\hat{M}(E1,\mu)$. These two components look in the opposite directions, because the angular momentum associated with the dipole operator is equal to the algebraical difference of j and λ. Thus adding j to the core angular momentum of the ground-band state to obtain the state of aligned octupole band one has simultaneously to subtract λ units from the core angular momentum. This leads to the selection rule for the dipole transitions from the j = 3 band to the ground-state band :

$$(\text{oct } (j=3),\ I) \rightarrow (\text{gr},\ I-3+\lambda) \quad (\lambda = 2 \text{ or } 4).$$

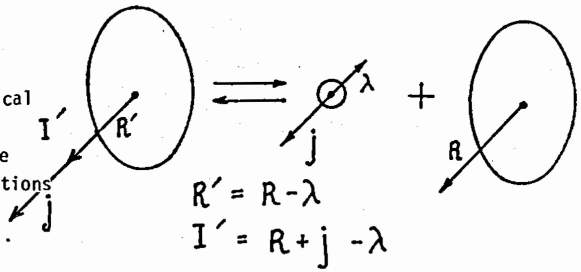


Fig. 7 - The geometrical interpretation of the selection rules for the electric dipole transitions from the aligned octupole band (j_x = 3).

If λ = 2 the dipole transitions from the distorted 0^- band lead predominantly to the decrease of the nuclear angular momentum ; if λ = 4, the situation is opposite and the nuclear angular momentum increases by one unit of Planck after the dipole transition from the 0^- band.

The same arguments lead to the selection rule for transitions from the "antialigned" band (j = -3). Here the strongest transitions are such that :(oct, j = -3 ; I)→(gr , I + 3 - λ). The classical addition rules tell also that it is impossible to reach the octupole state in which the vibrational angular momentum looks askew to the collective rotational one without the change in the orientation of the latter in the intrinsic frame, i.e. without the excitation of a wobbling motion of the system.

This explains why the B(E1) factors for the bands $j \neq \pm 3$ are smaller compared to the B(E1) factors for transitions obeying the above selection rules.

The unadiabatic effects in the branching rules for E1-transitions from the bands distorted by the Coriolis coupling have already been noticed in literature [5]. One may expect that the octupole-to-dipole coupling is, as a rule, the result of the deformation of the nuclear average field [20]. The predominance of the quadrupole deformation of the field suggests that the quadrupole coupling dominates also in the E1-transitions. However, the hexadecapole coupling cannot be excluded especially in the case when it is generated by the Coriolis excitations of the core : the larger is the multipolarity of intrinsic excitations, the stronger are the effects of the Coriolis force.

As was mentioned before, the alignment of vibrational angular momentum in deformed nuclei is well pronounced already in the first states of negative parity bands. The above estimations indicate that in ^{232}Th, ^{236,238}U the states of the 0^- band with angular momenta I $\gtrsim$ 5 carry appreciable amount of aligned angular momentum. Thus the selectivity of transitions from such states must be well marked. This agrees with the existing experimental data. As a rule, the dipole transitions $I^- \rightarrow (I \pm 1)^+$ from the 0^- band are seen together only when I $\lesssim$ 5 ; 7. At larger spins only a part of the transitions allowed by the angular momentum conservation law is seen. Examples of the γ-ray spectrum from the Coulomb excited ^{238}U nucleus is given in references [5,9]. In these examples the transitions (0^-, I)→(gr,I - 1) are well resolved up to I = 17. The last of the experimentally registered transition(0^-, I)→(gr,I + 1) corresponds to I = 3. There is a possibility of weak transitions of this type from the states 0^-,I = 5 and I = 7, which cannot be resolved from some other transitions. This is in accordance with the picture obtained on the basis of a simple model corresponding to the quadrupole dipole-to-octupole coupling.

The above example might be called typical : the enhancement of transitions $I^- \rightarrow (I-1)^+$ is observed in $^{234,236}U$ (8).

However, a qualitatively different situation is also observed. Figure 1, in reference (7) which gives the γ-ray spectrum of ^{232}Th, shows that the transitions $I^- \rightarrow (I+1)^+$ are enhanced in a way which resembles the case of an hexadecapole coupling ($\lambda = 4$, $K = 1$). Both theory and experiment give a ratio $B(E1, 5^- \rightarrow 6^+)/B(E1 ; 5^- \rightarrow 4^+) \approx 20$. At higher spins only the transitions $I^- \rightarrow (I+1)^+$ are seen, which again agrees with the theoretical estimations.

Thus, the theoretical conclusions concerning the electromagnetic properties of aligned states also have some experimental justification. The difference of the decay properties of the states of 0^- bands in ^{232}Th and ^{238}U must be related to the details of microscopic structure of these two nuclei. Phenomenologically this difference may be ascribed to the alignment of the angular momentum of quasiparticle nature which proceeds in a very different way in the two nuclei. The difference of the RPA bosons of the two nuclei which was discussed at the end of Sec. 3 must also play an important role for the coupling between the octupole and dipole degrees of freedom.

V - CONCLUSION

In the above sections a very simple theoretical model was presented for the study of effects of alignment on the electric properties of nuclei. Being rather simple, the model turns out to be adequate for a qualitative description of the known experimental data. The mathematical simplicity of this model allows to make an attempt of generalization of the selection rules for the electromagnetic transitions which were discussed in Sec. 4.

It follows from the geometrical interpretation in fig. 7 that the transitions from the band in which the alignment of the intrinsic angular momenta j_α associated with some nuclear degree of freedom (α) is well pronounced are selective in respect to the angular momentum transferred from the nucleus. The enhanced transition $I \rightarrow I'$ must obey the classical addition rules for the angular momenta involved in the transition so that $I - I' = j_\alpha + \Delta R$, where ΔR is the change of the rotational angular momentum. This change is necessarily present when the multipolarity of the emitted γ-ray (λ) is different from j_α. In a general case one knows only that $|j - \lambda| \leqslant \Delta R \leqslant j + \lambda$. The precise value of ΔR depends on the nuclear structure and may be associated with a coupling between the different nuclear degrees of freedom. The measure of selectivity depends on the following factors :

1. The intrinsic excitations with a certain value j_α must play a predominant role in the formation of nuclear intrinsic states ;

2. The angular momentum associated with this mode must be aligned ;
3. The particular value of ΔR must dominate the transitions (i.e. the coupling of nuclear modes).

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