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TIME-ODD COMPONENTS IN THE ROTATING MEAN FIELD AND IDENTICAL BANDS*,**

J. Dobaczewski

Institute of Theoretical Physics, Warsaw University
Hoża 69, PL-00-681 Warsaw, Poland

AND

J. Dudek

Centre de Recherches Nucléaires
IN_2P_3–CNRS/Université Louis Pasteur
F-67037 Strasbourg Cedex 2, France

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A systematic construction of the energy-density functional within the local density approximation is presented. The Hartree–Fock equations corresponding to such a functional are solved in case of rotating superdeformed nuclei. The identical bands in ^{152}Dy, ^{151}Tb, and ^{150}Gd are investigated and the time-odd components in the rotating mean field are analyzed.

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1. Introduction

The nuclear rotation is an example of a collective motion for which time-dependent linear combinations of stationary states of a given spin can be identified with a rotating wave packet. A description of such states can

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be performed within a mean-field theory in which properties of the system are determined by the one-body density matrix. Because the time-reversal symmetry is broken for the rotating states, the density matrix, as well as the resulting mean field have both time-even and time-odd components.

Properties of nuclear time-even mean fields are known rather well, because they are reflected in multiple static phenomena which can be studied experimentally, *cf.* review [1]. On the other hand, very little is known about properties of the time-odd mean fields. The present study is devoted to an attempt to analyse properties of these time-odd components in rapidly rotating superdeformed nuclei. In particular, we aim at studying the phenomenon of identical bands, which may provide a unique information on poorly known sector of the nuclear effective interaction.

In Sect. 2 we present a construction of the energy-density functional based on a few simple assumptions concerning its structure. Section 3 presents results of selfconsistent cranking method applied to nuclei in the $A\simeq150$ region; it supplements the results of similar calculations presented in Ref. [2].

2. Local density approximation

The density of nuclear matter in the interior of atomic nuclei has a well defined value, called the saturation density, which is independent of the nuclear size. This property of matter composed of strongly interacting nucleons is a basic feature of the nucleon-nucleon interaction which (depending also on the intrinsic symmetries of the nucleonic states) gives the strongest attraction at a certain distance and becomes weaker or repulsive at other distances. Therefore, as a first approximation one can consider a state of the matter in a given point of the nucleus as given by that of the infinite matter at the saturation density. In other words, properties of nuclear matter in a given point weakly depend on other points of the nucleus.

Such an observation underlies the local density approximation (LDA) in nuclear physics [3]. It has been adopted from similar ideas developed in atomic physics [4] and reflects fundamental properties of many-fermion systems. In such systems, the Pauli correlations play a very important role in defining properties of states near its ground state. To some extent, details of interactions are less important, and in particular, the local density approximation can be valid irrespectively of the interaction range.

Because of the finite nuclear size, the simplest version of the LDA has to be corrected by taking into account not only the density itself but also its local derivatives. This amounts to extending the Thomas-Fermi approximation beyond its extreme version, or to applying the systematic Wigner–Kirkwood $\hbar$ expansion [3]. In a phenomenological approach one may use the ideas of the LDA to construct the energy-density functionals, and in what follows we proceed along these lines.

Principal assumptions of such a LDA can be formulated as follows:

- The total energy of the nucleus is given by the integral of the local energy density $\mathcal{H}(\boldsymbol{r})$,
$$\mathcal{E} = \int d^3\boldsymbol{r}\mathcal{H}(\boldsymbol{r}) , \tag{2.1}$$
- The energy density depends on the nuclear one-body density matrix and its derivatives up to the second order.

Neglecting for simplicity the isospin degree of freedom, *i.e.*, considering one type of nucleons only, the density matrix in spatial coordinates can be defined by
$$\rho(\boldsymbol{r}\sigma, \boldsymbol{r}'\sigma') = \langle\Phi|a^+(\boldsymbol{r}'\sigma')a(\boldsymbol{r}\sigma)|\Phi\rangle , \tag{2.2}$$
where $\boldsymbol{r}$ is the position and σ is the spin of a nucleon, while $|\Phi\rangle$ is a many-body nuclear wave function.

The spin degrees of freedom can be separated by defining the scalar and the vector parts of the density matrix, $\rho(\boldsymbol{r}, \boldsymbol{r}')$ and $\boldsymbol{s}(\boldsymbol{r}, \boldsymbol{r}')$, respectively, *i.e.*,

$$\rho(\boldsymbol{r}, \boldsymbol{r}') = \sum_{\sigma} \rho(\boldsymbol{r}\sigma, \boldsymbol{r}'\sigma) \tag{2.3a}$$
$$\boldsymbol{s}(\boldsymbol{r}, \boldsymbol{r}') = \sum_{\sigma\sigma'} \rho(\boldsymbol{r}\sigma, \boldsymbol{r}'\sigma')\langle\sigma'|\boldsymbol{\sigma}|\sigma\rangle . \tag{2.3b}$$

Since the density matrix is Hermitian, its scalar and vector parts have the following properties with respect to exchanging spatial arguments:

$$\rho^*(\boldsymbol{r}, \boldsymbol{r}') = \rho(\boldsymbol{r}', \boldsymbol{r}) , \tag{2.4a}$$
$$\boldsymbol{s}^*(\boldsymbol{r}, \boldsymbol{r}') = \boldsymbol{s}(\boldsymbol{r}', \boldsymbol{r}) . \tag{2.4b}$$

On the other hand, the density matrices corresponding to time-reversed states read:

$$\rho^T(\boldsymbol{r}, \boldsymbol{r}') = \rho^*(\boldsymbol{r}, \boldsymbol{r}') , \tag{2.5a}$$
$$\boldsymbol{s}^T(\boldsymbol{r}, \boldsymbol{r}') = -\boldsymbol{s}^*(\boldsymbol{r}, \boldsymbol{r}') . \tag{2.5b}$$

These properties mean that for the time-even state the scalar density matrix is a real and symmetric function of spatial arguments, whereas the vector density matrix is imaginary and antisymmetric. Therefore, the local parts of the scalar and of the vector densities,

$$\rho(\boldsymbol{r}) = \rho(\boldsymbol{r}, \boldsymbol{r}) , \tag{2.6a}$$
$$\boldsymbol{s}(\boldsymbol{r}) = \boldsymbol{s}(\boldsymbol{r}, \boldsymbol{r}) , \tag{2.6b}$$

respectively, do not and do change the sign under the time reversal. Hence the matter density $\rho(\boldsymbol{r})$ is time-even, and the spin density $\boldsymbol{s}(\boldsymbol{r})$ is time-odd.

The local derivatives of the density matrix have been defined in Ref. [5] and are summarized in Table I. (See Ref. [6] for another possible set of definitions.) A systematic construction [5] of these derivatives up to a given order consists in acting on the scalar and vector density matrices (2.3) with the operators $(\nabla - \nabla')/2i$ and setting the arguments equal, $\boldsymbol{r}{=}\boldsymbol{r}'$, after the differentiation is completed. This amounts to calculating derivatives with respect to the relative coordinate $\boldsymbol{r}-\boldsymbol{r}'$ at the relative coordinate equal to zero. At this point, the derivatives $(\nabla+\nabla')/2$ with respect to the total coordinate $\boldsymbol{r}+\boldsymbol{r}'$ are not considered, because they amount to calculating the derivatives *after* setting $\boldsymbol{r}{=}\boldsymbol{r}'$, see below.

TABLE I

Local densities of a fermion system up to the second order in derivatives with respect to the relative coordinate

Type of density	Order	Definition
Matter:	0	$\rho(\boldsymbol{r}) = \rho(\boldsymbol{r}, \boldsymbol{r})$
Current:	1	$\boldsymbol{j}(\boldsymbol{r}) = (1/2i)[(\nabla - \nabla')\rho(\boldsymbol{r}, \boldsymbol{r}')]_{r=r'}$
Kinetic:	2	$\tau(\boldsymbol{r}) = [\nabla \cdot \nabla' \rho(\boldsymbol{r}, \boldsymbol{r}')]_{r=r'}$
Spin:	0	$\boldsymbol{s}(\boldsymbol{r}) = \boldsymbol{s}(\boldsymbol{r}, \boldsymbol{r})$
Spin current:	1	$J_{\mu\nu}(\boldsymbol{r}) = (1/2i)[(\nabla_\mu - \nabla'_\mu) s_\nu(\boldsymbol{r}, \boldsymbol{r}')]_{r=r'}$
Spin kinetic:	2	$\boldsymbol{T}(\boldsymbol{r}) = [\nabla \cdot \nabla' \boldsymbol{s}(\boldsymbol{r}, \boldsymbol{r}')]_{r=r'}$

In the first order one obtains two new local densities, the current density $\boldsymbol{j}$ and the spin current density $J_{\mu\nu}(\boldsymbol{r})$, Table I. In the second order one has to act on the scalar and vector densities with the tensor operator $(\nabla_\mu - \nabla'_\mu)(\nabla_\nu - \nabla'_\nu)$. However, the resulting traceless-symmetric-tensor density is not interesting, because there is no other zero-order tensor density with which it could have been contracted to construct the corresponding term in the energy density, see below. Therefore, only the scalar part of the second-order operator should be considered. Due to the identity

$$(\nabla - \nabla')^2 = (\nabla + \nabla')^2 - 4\nabla \cdot \nabla', \tag{2.7}$$

the action of $(\nabla - \nabla')^2$ can be expressed by that of $\nabla \cdot \nabla'$, and therefore, following Ref. [5] one uses kinetic densities shown in Table I.

We may now proceed by constructing the local energy density as a sum of terms depending on the above local densities. Such a construction is based on the following rules:

- **Every term should be quadratic in local densities.**
  An *a priori* arbitrary coupling constant can be used in front of every

term. It is well known that the energy densities which are quadratic in the matter density do not lead to saturating systems, and therefore, a density dependence of the coupling constants should be allowed.

- **The energy density should be invariant with respect to parity, time reversal, and rotations.**
  This invariance should be considered in a limit of unbroken symmetries. In principle, an extension of the functional beyond this limit is arbitrary. For example, for broken parity symmetry one could use in the functional either the term $C^{\rho}\rho^2$, or two independent terms $C^{\rho^+}\rho_+^2$ and $C^{\rho^-}\rho_-^2$, where $\rho_\pm=(\rho\pm\rho^P)/2$, and where ρ^P corresponds to the parity inversed state. It is obvious that for the unbroken symmetry both choices give the same result, provided that $C^{\rho}=C^{\rho^+}$. Moreover, for both choices the functional is invariant with respect to parity. On the other hand, for the broken symmetry the first choice gives the following contribution to the total energy:

$$\int d^3\boldsymbol{r}C^{\rho}\rho^2 = \int d^3\boldsymbol{r}C^{\rho}(\rho_+^2+\rho_-^2)\,, \tag{2.8}$$

  and therefore it requires that $C^{\rho^-}=C^{\rho^+}$, while in the second choice these two coupling constants can be chosen independently. These ambiguities are even larger for density-dependent coupling constants. Up two now, mostly the first choice has been considered, *i.e.*, the functionals are constructed in the unbroken symmetry limit.
- **Terms beyond the second order are not taken into account.**
  The order of the given term is defined as a sum of orders of both densities on which this term depends. The requirement of neglecting the terms beyond the second order is consistent with not considering the densities beyond the second order, see Table I. This amounts to an expansion of the energy density up to the second order in relative coordinates.

The construction of the energy density is illustrated in Table II. The first column gives the densities presented in Table I. The spin current density $J_{\mu\nu}$ is split into the standard scalar, vector, and tensor parts ($J^{(0)}$, $\boldsymbol{J}$, and $J^{(2)}$, respectively) according to

$$J_{\mu\nu} = \tfrac{1}{3}J^{(0)}\delta_{\mu\nu} + \tfrac{1}{2}\epsilon_{\mu\nu\kappa}J_{\kappa} + J^{(2)}_{\mu\nu}\,, \tag{2.9}$$

where

$$J^{(0)} = \sum_{\mu} J_{\mu\mu}\,, \tag{2.10a}$$

$$J_\kappa = \sum_{\mu\nu} \epsilon_{\kappa\mu\nu} J_{\mu\nu} \,, \tag{2.10b}$$

$$J^{(2)}_{\mu\nu} = \tfrac{1}{2}(J_{\mu\nu} + J_{\nu\mu}) - \tfrac{1}{3}\delta_{\mu\nu} \sum_{\kappa} J_{\kappa\kappa} \,. \tag{2.10c}$$

TABLE II

Construction of all possible terms up to second order in the energy density of a fermion system (for comments see text).

1	2	3	4			5	
Density	Derivative	Order	Symmetry			Energy density	
			T	P	space	prime	dependent
$\rho(\mathbf{r})$		0	+	+	scalar	ρ^2	
	$\nabla\rho$	1	+	−	vector		$(\nabla\rho)^2$
	$\Delta\rho$	2	+	+	scalar	$\rho\Delta\rho$	
	$(\nabla\times\nabla)^{(2)}\rho$	2	+	+	tensor		
$\tau(\mathbf{r})$		2	+	+	scalar	$\rho\tau$	
$J^{(0)}(\mathbf{r})$		1	+	−	scalar	$\left(J^{(0)}\right)^2$	
	$\nabla J^{(0)}$	2	+	+	vector		
$\mathbf{J}(\mathbf{r})$		1	+	−	vector	$\mathbf{J}^2$	$\nabla\rho\cdot\mathbf{J}$
	$\nabla\cdot\mathbf{J}$	2	+	+	scalar	$\rho\nabla\cdot\mathbf{J}$	
	$\nabla\times\mathbf{J}$	2	+	+	vector		
	$(\nabla\times J)^{(2)}$	2	+	+	tensor		
$J^{(2)}_{\mu\nu}(\mathbf{r})$		1	+	−	tensor	$\left(J^{(2)}\right)^2$	
$\mathbf{s}(\mathbf{r})$		0	−	+	vector	$\mathbf{s}^2$	
	$\nabla\cdot\mathbf{s}$	1	−	−	scalar	$(\nabla\cdot\mathbf{s})^2$	
	$\nabla\times\mathbf{s}$	1	−	−	vector		$(\nabla\times\mathbf{s})^2$
	$(\nabla\times s)^{(2)}$	1	−	−	tensor		$\left((\nabla\times s)^{(2)}\right)^2$
	$\Delta\mathbf{s}$	2	−	+	vector	$\mathbf{s}\cdot\Delta\mathbf{s}$	
$\mathbf{j}(\mathbf{r})$		1	−	−	vector	$\mathbf{j}^2$	$\mathbf{j}\cdot\nabla\times\mathbf{s}$
	$\nabla\cdot\mathbf{j}$	2	−	+	scalar		
	$\nabla\times\mathbf{j}$	2	−	+	vector	$\mathbf{s}\cdot\nabla\times\mathbf{j}$	
	$(\nabla\times j)^{(2)}$	2	−	+	tensor		
$\mathbf{T}(\mathbf{r})$		2	−	+	vector	$\mathbf{s}\cdot\mathbf{T}$	

The second column of Table II presents derivatives of the densities from the first column. The derivatives are constructed in the vector-coupled form. For example, from the (vector) gradient operator ∇ and the scalar density ρ one can obtain the vector density $\nabla\rho$, the scalar density $\Delta\rho$, and the tensor density $(\nabla\times\nabla)^{(2)}\rho$. Since ρ itself is a zero-order density, the density $\nabla\rho$ is of the first order and the other two densities are of the second order, as indicated in the third column of the Table II.

The derivatives are calculated only up to the second order. Moreover, the rank-3 tensors are not taken into account and the derivatives of $J^{(2)}_{\mu\nu}(\boldsymbol{r})$ are not included in the Table II, because there are no other densities with the same symmetries to construct terms in the energy density. Apart from these two omissions, the Table presents a complete set of densities and their derivatives which can be obtained within the specified limits.

The fourth column of the Table gives the symmetries of every density in the unbroken symmetry limit, *i.e.*, assuming that ρ is a time-even, parity-even scalar density, and $\boldsymbol{s}$ is a time-odd, parity-even vector density. Then the symmetries presented in the Table result from the fact that the gradient operator ∇ is a time-even, parity-odd vector operator.

Finally, the fifth column of the Table II presents the terms in the energy density which can be constructed from the densities listed in the first two columns. In each case a term is obtained by forming a scalar, time-even, and parity-even product of two densities, *i.e.*, by multiplying densities having exactly the same symmetries. Two avoid the double-counting, a given density is multiplied only by densities appearing higher in the Table II. The rule of not including terms above the second order is also respected.

For example, according to this construction scheme, the matter density ρ gives the simplest term in the energy density, ρ^2, as shown in the first line of the Table. Similarly, the derivative $\nabla\rho$ gives rise to the term $(\nabla\rho)^2$. This term is listed in the column of "dependent" terms, because for a density-independent coupling constant it can be expressed by the "prime" term $\rho\Delta\rho$ given in the third line. This can be done by integrating by parts the integral of the energy density (2.1). The corresponding expressions for five "dependent" terms are as follows:

$$(\nabla\rho)^2 = -\rho\Delta\rho\,, \tag{2.11a}$$

$$\nabla\rho\cdot\boldsymbol{J} = -\rho\nabla\cdot\boldsymbol{J}\,, \tag{2.11b}$$

$$(\nabla\times\boldsymbol{s})^2 = -(\nabla\cdot\boldsymbol{s})^2 - \boldsymbol{s}\cdot\Delta\boldsymbol{s}\,, \tag{2.11c}$$

$$\sum_{\mu\nu}\left((\nabla\times s)^{(2)}_{\mu\nu}\right)^2 = \tfrac{2}{3}(\nabla\cdot\boldsymbol{s})^2\,, \tag{2.11d}$$

$$\boldsymbol{j}\cdot\nabla\times\boldsymbol{s} = \boldsymbol{s}\cdot\nabla\times\boldsymbol{j}\,. \tag{2.11e}$$

Altogether, this construction procedure gives eighteen terms in the energy density, of which thirteen are "prime" and five are "dependent". Three terms correspond to the squares of the spin current density $J_{\mu\nu}(\boldsymbol{r})$ for the three different values of the intermediate-coupling angular momentum. Usually, these three independent terms are considered together in the following single term:

$$\overleftrightarrow{J}^2 = \sum_{\mu\nu} J_{\mu\nu}^2 = \tfrac{1}{3}\left(J^{(0)}\right)^2 + \tfrac{1}{2}\boldsymbol{J}^2 + \sum_{\mu\nu}\left(J_{\mu\nu}^{(2)}\right)^2 \tag{2.12}$$

with a common coupling constant.

It turns out that each density, or a derivative of the density, gives rise to at most one "prime" term in the energy density, and these terms and coupling constants can be conveniently labelled by the corresponding densities. We then obtain the following expression for the interaction energy density:

$$\mathcal{H}^{\rm int}(\boldsymbol{r}) = \mathcal{H}^{\rm even}(\boldsymbol{r}) + \mathcal{H}^{\rm odd}(\boldsymbol{r}) \,, \tag{2.13}$$

with

$$\mathcal{H}^{\rm even}(\boldsymbol{r}) = C^{\rho}\rho^2 + C^{\Delta\rho}\rho\Delta\rho + C^{\tau}\rho\tau + C^{J}\overleftrightarrow{J}^2 + C^{\nabla J}\rho\nabla\cdot\boldsymbol{J}, \tag{2.14a}$$

$$\begin{aligned}\mathcal{H}^{\rm odd}(\boldsymbol{r}) &= C^{s}\boldsymbol{s}^2 + C^{\Delta s}\boldsymbol{s}\cdot\Delta\boldsymbol{s} + C^{T}\boldsymbol{s}\cdot\boldsymbol{T} + C^{j}\boldsymbol{j}^2 \\ &\quad + C^{\nabla j}\boldsymbol{s}\cdot(\nabla\times\boldsymbol{j}) + C^{\nabla s}(\nabla\cdot\boldsymbol{s})^2 \,.\end{aligned} \tag{2.14b}$$

Each of these terms can have the isoscalar and the isovector form, and the kinetic energy density has to be added to obtain the total energy density of a nuclear system.

Apart from the last term in $\mathcal{H}^{\rm odd}(\boldsymbol{r})$, all these terms appear in the energy density corresponding to the Skyrme interaction [5, 2]. The additional term in Eq. (2.14b) can only be obtained if the tensor interaction is added to the Skyrme force, *i.e.*, when the results of Ref. [7] are generalized to the case of broken time-reversal symmetry. In the following we do not consider such a possibility, and we set $C^{\nabla s}=0$.

By varying the total energy (2.1) with respect to the single-particle wave functions one obtains the Hartree–Fock time-even and time-odd mean fields [5]. In the convention corresponding to the energy density of Eq. (2.14), the relevant expressions are given in Ref. [2] and will not be repeated here. Instead, in the following section we present the results of the Hartree–Fock cranking method applied to the description of superdeformed bands.

3. Superdeformed bands in ^{152}Dy, ^{151}Tb, and ^{150}Gd

In Fig. 1 we show the results of calculations for the yrast superdeformed band in ^{152}Dy and for the first excited band in ^{151}Tb, denoted by ^{152}Dy(1) and ^{151}Tb(2), respectively. Details of the calculation methods can be found in Ref. [2]. The coupling constants of the energy-density functional have been determined from the parameters of three Skyrme forces, SkM* [8], SkP [9], and SIII [10]. The left-hand-side panels show the results obtained for the complete functionals, while those at the right-hand-side correspond to omitted time-odd terms $\mathcal{H}^{\mathrm{odd}}(\boldsymbol{r})\equiv 0$, *i.e.*, to

$$C^s = C^{\Delta s} = C^T = C^j = C^{\nabla j} = 0\,. \tag{3.1}$$

For the ^{151}Tb(1) and ^{150}Gd(2) bands, the analogous results are shown in Fig. 2.

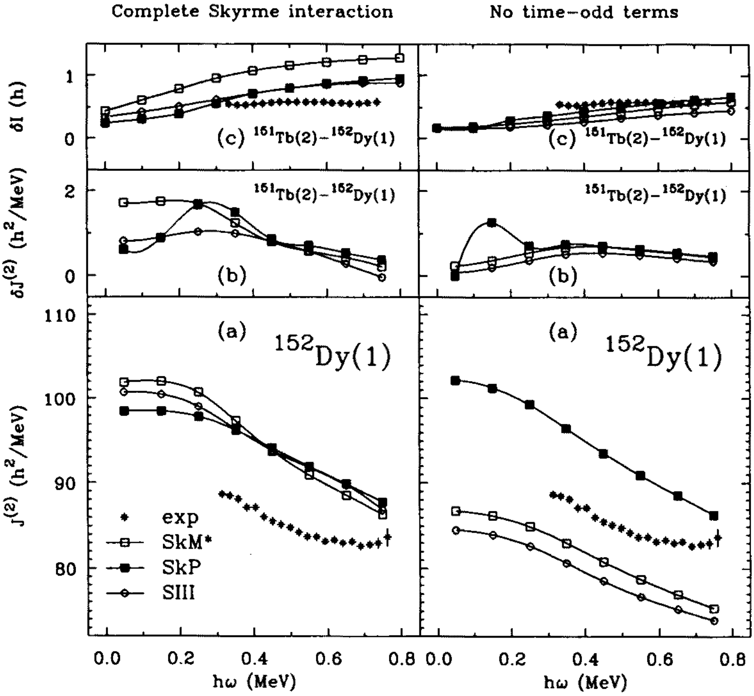

Fig. 1. Calculated dynamical moment $\mathcal{J}^{(2)}$ for the yrast band of ^{152}Dy, part (a), the relative dynamical moment $\delta\mathcal{J}^{(2)}$ calculated for the ^{151}Tb(2) and ^{152}Dy(1) bands, part (b), and the relative alignment δI between these two bands, part (c). Left-hand-side panels show the results for complete Skyrme functionals of the SkM*, SkP, and SIII interactions, and the right-hand-side panels show the results with omitted time-odd terms. The experimental points are denoted by asterisks. Note the scale in (b) expanded five times as compared to (a).

The omission of time-odd terms in the Hartree–Fock approach is analogous to the standard phenomenological rotating mean-field approach where the selfconsistent changes of the mean field due to rotations are not taken into account. Results of latter type calculations, for the same SD bands as considered in the present study, can be found in Refs [11] and [12].

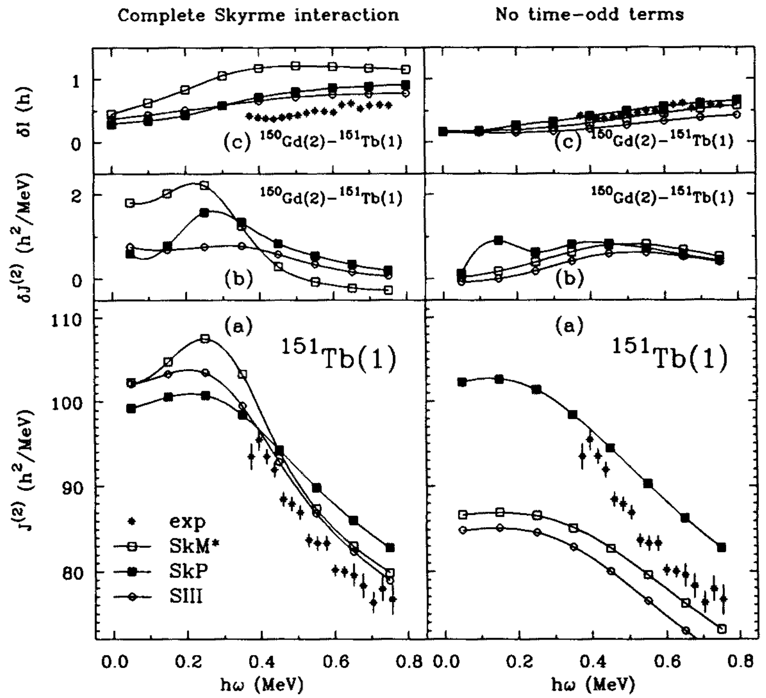

Fig. 2. Same as Fig. 1, but for the ^{151}Tb(1) band, (a), and for the differences between the ^{150}Gd(2) and ^{151}Tb(1) bands, (b) and (c).

For the complete Skyrme functionals, the dynamic moments $\mathcal{J}^{(2)}$ of the ^{152}Dy(1) band obtained for SkM*, SkP, and SIII interactions are very similar, see part (a) in the left-hand-side panel of Figs 1 and 2. This is not the case when the time-odd terms are omitted, as show in the corresponding right-hand-side panels. For SkM* and SIII, the omission of the time-odd terms significantly decreases the values of $\mathcal{J}^{(2)}$, while for SkP the changes of $\mathcal{J}^{(2)}$ are much smaller. This feature is related to the values of the isoscalar nuclear-matter effective mass, which for SkP is equal to the free nucleon mass, $m^*=m$, while for SkM* and SIII is smaller by the factor of 0.79 and 0.76, respectively. As a consequence, for SkP the isoscalar coupling constant C_0^j is equal to zero [2], and the modifications introduced by omitting the time-odd terms, Eq. (3.1), do not influence $\mathcal{J}^{(2)}$ in a very strong way. On the other hand, whenever this coupling constant is relatively large, as is for SkM* and SIII, the corresponding changes in $\mathcal{J}^{(2)}$ are well pronounced.

These results illustrate the importance for the rotational properties of the nucleus of the $\boldsymbol{j}^2$ term in the energy density. The cranking term mainly induces the nonzero flow of nuclear matter, as given by the current density $\boldsymbol{j}$, which influences the time-odd nuclear mean field provided the corresponding coupling constant C^j is nonzero.

In spite of the strong influence on the values of $\mathcal{J}^{(2)}$, the omission of the time-odd terms has a very small effect on the sameness of $\mathcal{J}^{(2)}$ in pairs of identical bands. In parts (b) of Figs 1 and 2 we show the relative dynamical moments $\delta\mathcal{J}^{(2)}$, *i.e.*, the differences between the values of $\mathcal{J}^{(2)}$ calculated for two pairs of identical bands ^{151}Tb(2)–^{152}Dy(1) and ^{150}Gd(2)–^{151}Tb(1). One can see that the omission of the time-odd terms influences the values of $\mathcal{J}^{(2)}$ in a very similar way for both members of each pair. Even if the values may change by as much as 15 $\hbar^2$/MeV, the relative values are always well below 2 $\hbar^2$/MeV. This shows that the sameness of $\mathcal{J}^{(2)}$ is not governed by the time-odd terms in the mean-fields, but rather can be attributed to general geometric properties of the underlying orbital $\pi[301]1/2(r{=}{+}i)$.

A different patterns is obtained for relative alignments shown in parts (c) of the figures. Here the omission of the time-odd terms leads to important changes of the relative alignments, and also different relative alignments are obtained of the three Skyrme forces studied.

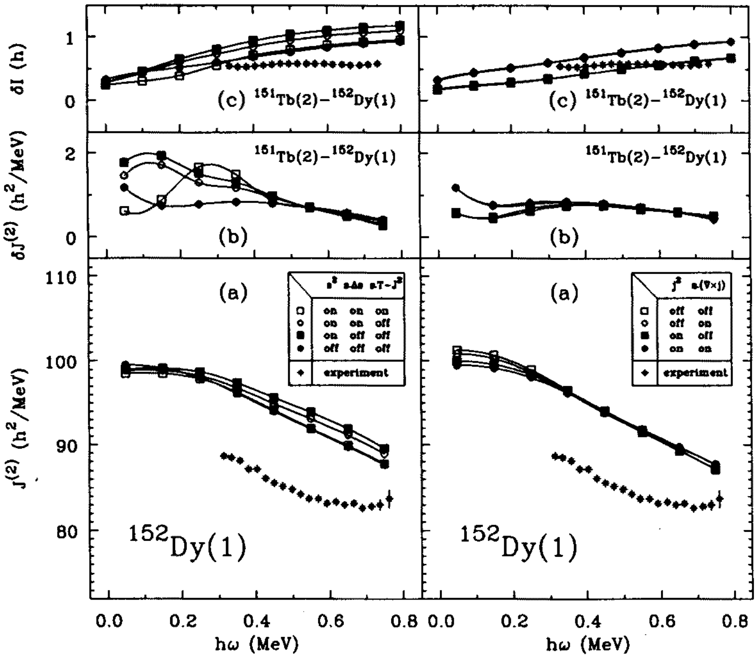

Fig. 3. Similar as in Fig. 1, but for different time-odd terms omitted in the energy density, see text.

In Figs 3 and 4 we present a detailed analysis of the time-odd terms appearing in the energy density for the SkP interaction. As shown in the

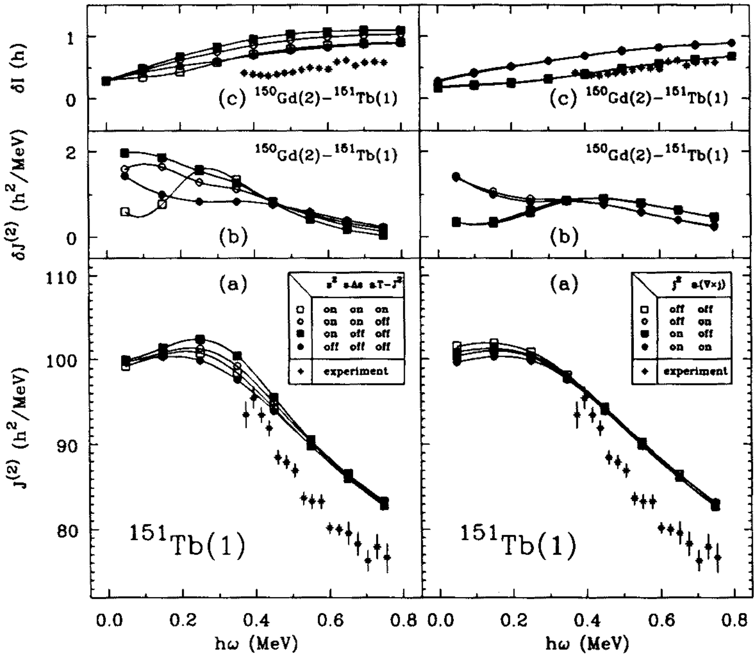

Fig. 4. Same as Fig. 3, but for the ^{151}Tb(1) band, (a), and for the differences between the ^{150}Gd(2) and ^{151}Tb(1) bands, (b) and (c).

legends, the results presented in the left-hand-side panels correspond to the following conditions: *(i)* the complete functional, *(ii)* $C_t^T = C_t^J = 0$, *(iii)* $C_t^{\Delta s} = C_t^T = C_t^J = 0$, and then *(iv)* $C_t^s = C_t^{\Delta s} = C_t^T = C_t^J = 0$. In the right-hand-panels, in addition to the last of these conditions we set to zero either one, or both of the C^j and $C^{\nabla j}$ coupling constants.

As discussed above, the value of the effective mass $m^*{=}m$ renders the dynamical moments $\mathcal{J}^{(2)}$, the relative dynamical moments, and the relative alignments δI almost independent of the C^j coupling constants. (A weak residual dependence results from a non-zero value of the isovector coupling constant C_1^j, see Ref. [2].) On the other hand, the values of the relative alignments strongly depend on the $C^{\nabla j}$ coupling constants, and are closest to the experimental data when these coupling constants are equal to zero.

4. Conclusions

In the present study we have presented a systematic construction of the energy-density functional in the frame of the nuclear local density approximation. Such a construction is based on a few simple prescriptions and leads to the functional identical to that obtained by using the Skyrme effective interaction. The only exception is the term given by the square of the divergence of the spin density, which can only be obtained if the tensor component is added to the Skyrme interaction.

Following Ref. [2], we have here presented the analysis of the time-odd components of the mean field of rotating superdeformed nuclei. Special attention has been devoted to properties of the SkP Skyrme interaction which is characterized be the effective mass $m^*{=}m$. In particular, we have shown that the sameness of the dynamical moments $\mathcal{J}^{(2)}$ *does not depend* on whether the time-odd terms are, or are not taken into account. On the other hand, the sameness of the corresponding alignments *does depend* on these time-odd terms. Therefore, one may expect that a systematic investigation of the rotational alignments may serve as a tool for establishing the properties of time-odd component of the rotating nuclear mean field.

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