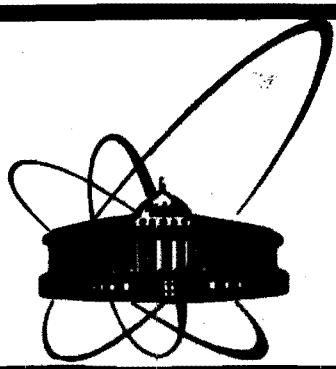

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# GIANT MONOPOLE TRANSITION DENSITIES WITHIN THE LOCAL-SCALE ATDHF APPROACH

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## I. Introduction

In recent years considerable experimental and theoretical efforts have been devoted to the study of nuclear giant monopole resonances (GMR). At present, the isoscalar (T=0) giant monopole resonance (ISMR), the nuclear breathing mode, has been established experimentally in a wide mass range of nuclei [1]. New experimental developments [2] have also opened the possibility of studying the isovector (T=1) giant monopole resonance (IVMR).

The experimental study of the GMR and more generally of any monopole transition with a projectile like $\alpha$, d, $^3$He, and light heavy ions, requires measurements at very small momentum transfer. The identification of the resonance and its physical characteristics is provided with comparison between the measured angular disrtibutions and those calculated using the Distorted-Wave Born Approximation (DWBA). The basic element of the nuclear system relevant to any DWBA analysis [3] of the GMR experimental data is the nuclear transition density $\rho_{tr}(\vec{r})$ which should incorporate nuclear structure effects that play a role in such monopole modes.

The transition densities for GMR excitations can be obtained within the microscopic framework using for example the Random Phase Approximation (RPA) [3-5], the Generator Coordinate Method (GCM) [6,7], the Time-Dependent Hartree-Fock Method (TDHF) [8]. However, these transition densities are difficult to be used within the DWBA analysis. They are not of closed analytical form and depend sensitively on the amount of np-nh configuration mixing in ground and excited $0^+$ states.

Usually, simple models of the monopole vibrations are widely used within the DWBA description of the GMR experimental data.

For example, the well-known Tassie transition density [9] follows from a simple radial scaling $\vec{r} \to \alpha \vec{r}$ of the ground state density distribution $\rho(\vec{r})$. The collective parameter $\alpha$ is one-to-one related to the nuclear mean squared radius and the resulting transition density is of the form:

$$\rho_{SM}(r)=A(3\rho(r)+r\partial\rho(r)/\partial r) \quad . \qquad (1)$$

The Tassie transition densities, eq. (1), are known to be in quantitative agreement to those of RPA calculations for heavy nuclei [4]. In the practical applications, the ground state density $\rho(r)$ in eq. (1) is usually taken in the Fermi-type form, $\rho(r) \equiv \rho_F(r) = \rho_F(r;R,b)$, considering A, R and b as a set of phenomenological parameters.

Another examlpe is the transition density $\rho_{TB_3}(r)$ following from the collective model [10] and used within the DWBA method [11]:

$$\rho_{TB_3}(r)=A\,\partial\rho(r)/\partial R+B\,\partial\rho(r)/\partial b \quad , \qquad (2)$$

where the nuclear ground state density $\rho(r)$ is presented as a Fermi-type distribution $\rho_F(r)$ and the set of the phenomenological parameters involves A, B, R, and b. In eq. (2) a surface monopole mode (diffuseness or b-type oscillations) is seen together with a bulk monopole mode (half-radius or R-type oscillations). The transition densities of the $\rho_{TB_3}(r)$ form have shown a better reproducing of the experimental ISMR cross section data [11] than $\rho_{SM}(r)$.

The main disadvantages of the phenomenological transition densities however is that there is no connection between their phenomenological parameters and the effective NN-interaction used within the microscopic approaches.

Recently, we have reported transition densities of the type of eq. (2), providing a better understanding of GMR properties. They



are of a simple analytical form and can be successfully applied to the reproduction of the experimental results using the DWBA method, as has been shown in [10,11]. In contrast with the collective model [10], however, our transition densities have been obtained within an Adiabatic Time-Dependent Hartree-Fock (ATDHF) approach [12]. In this approach both ISMR and IVMR have been described in a unified way as dynamically coupled bulk and surface vibrations of the half-radius R and the diffuseness parameter b of the local density distribution $\rho(r)$ considered as a symmetrized Fermi-type (SF) density distribution $\rho_{SF}(r) \equiv \rho_{SF}(r;R,b)$. Therefore, there is no ajustable phenomenological parameters in our transition densities which are completely obtained on the base of the used Skyrme-type effective forces [13,14].

An aim of the present work is to propose such obtained transition densities as an effective computational tool for research teams dealing with the experimental analysis of the reactions with excitation of the GMR. We suggest the method (part 2.), the particular form (part 3.) and all necessary coefficients of the transition densities obtained for a number of spherical even-even nuclei with the set of Skyrme forces SkM* [13] and SIII [14].

Transition densities of naturally appearing antiscaling type isoscalar and isovector monopole states [7,19] , which are not experimentally observed up to now [11], are also given.

## II. The Method

Recently elaborated Local-Scale Transformation Method (LSTM) [15,16] has allowed A-particle wave functions to be adopted for a systematical investigation of the GMR. The starting point is to

perform a local-scale point transformation (LST) on a certain model wave function $\overline{\phi} = \overline{\phi}(\vec{r}_1, \vec{r}_2, \ldots, \vec{r}_A)$. As a result, a wave function $\phi[f] \equiv \phi_f(\vec{r}_1, \vec{r}_2, \ldots, \vec{r}_A)$ is obtained which depends on the scalar LST function $f(\vec{r})$. The latter is one-to-one related to the ground state density distribution $\rho(\vec{r})$. The LST function $f(\vec{r})$ is thus presented as a functional $f_\rho(\vec{r})$ of the local density $\rho(\vec{r})$. Therefore, the local density $\rho(\vec{r})$ creates the basic states $\phi[\rho] = \phi_{f_\rho}(\vec{r}_1, \vec{r}_2, \ldots, \vec{r}_A)$ and collective parameters $u=(u_1, u_2, \ldots)$ can be naturally introduced in $\phi(u)=\phi[\rho_u]$ by means of proper characteristics of the local density distribution $\rho_u(\vec{r}) \equiv \rho(\vec{r})$.

In the ATDHF approach [12] we impose constrained $0^+$-density vibrations on the nucleus whose density profile is taken to be of SF-type [17]:

$$\rho_q(r) \equiv \rho_{SF}(r; R, b) = \rho_{0q}\,\mathrm{sh}(R_q/b_q)/(\mathrm{ch}(R_q/b_q)+\mathrm{ch}(r/b_q)). \qquad (3)$$

where $R_q$ and $b_q$ correspond to the density half-radius and surface skin-thickness parameter of the neutron (q=n) and proton (q=p) densities $\rho_n(r)$ and $\rho_p(r)$, respectively. The SF local densities (3) satisfy the normalisation conditions:

$$\int \rho_q(r)\,d^3r = A_q, \quad (q=n,p) \qquad (4)$$

through the relation:

$$\rho_{0q} = \rho_0(R_q, b_q) = (3A_q/4\pi R_q^3)/(1+(\pi b_q/R_q)^2)^{-1}, \qquad (5)$$

where $A_n$ ($A_p$) is the number of neutrons N (protons Z) and A=N+Z.

The corresponding to $\rho_{SF}$ A-particle wave function $\phi(u)=\phi[\rho_{SF}]$ built up on the base of the LSTM using a single Slater harmonic-oscillator shell model wave function [16] depends on the SF density collective variables $u=(R_n, R_p, b_n, b_p)$. Then according to refs. [7,12], the resulting wave functions $\phi(u)$ transform the mean value of the Hamiltonian H of the system in question into an algebraic function of the collective variables



$$V(u)=V(R_n, R_p, b_n, b_p)=\langle \phi(R_n, R_p, b_n, b_p)|H|\phi(R_n, R_p, b_n, b_p)\rangle \qquad (6)$$

which represents the collective potential energy $V \equiv V(u)$ in the ATDHF approximation.

Table 1.

The equilibrium values of the symmetrized Fermi-type local density parameters: half-radii ($R_n, R_p$ for $\rho_n \neq \rho_p$ and R for $\rho_n \approx \rho_p$) and diffuseness parameters ($b_n, b_p$ for $\rho_n \neq \rho_p$ and b for $\rho_n \approx \rho_p$)

| Nuclei | Forces | $R_n$[fm] | $R_p$[fm] | $b_n$[fm] | $b_p$[fm] | R [fm] | b [fm] |
|---|---|---|---|---|---|---|---|
| $^{12}C$ | SIII | 2.29664 | 2.26910 | .441586 | .457215 | 2.28552 | .448469 |
| | SkM* | 2.23480 | 2.09736 | .468286 | .513233 | 2.19435 | .482203 |
| $^{16}O$ | SIII | 2.67962 | 2.68578 | .433752 | .441407 | 2.68304 | .437422 |
| | SkM* | 2.67135 | 2.66859 | .463645 | .475312 | 2.67101 | .469087 |
| $^{28}Si$ | SIII | 3.26208 | 3.28755 | .440357 | .446958 | 3.27505 | .443536 |
| | SkM* | 3.17290 | 3.19699 | .467997 | .477127 | 3.18539 | .472365 |
| $^{32}S$ | SIII | 3.35195 | 3.37515 | .489380 | .500440 | 3.36421 | .494581 |
| | SkM* | 3.21876 | 3.22784 | .525957 | .544118 | 3.22547 | .534178 |
| $^{40}Ca$ | SIII | 3.76861 | 3.80476 | .457895 | .465446 | 3.78705 | .461468 |
| | SkM* | 3.71196 | 3.74951 | .492538 | .502852 | 3.73134 | .497397 |
| $^{48}Ca$ | SIII | 4.09639 | 3.95802 | .458863 | .443752 | 4.03424 | .451689 |
| | SkM* | 4.04394 | 3.84990 | .488197 | .480486 | 3.96387 | .484735 |
| $^{56}Ni$ | SIII | 4.24725 | 4.29703 | .441477 | .448246 | 4.27259 | .444631 |
| | SkM* | 4.13986 | 4.19497 | .471185 | .479935 | 4.16808 | .475273 |
| $^{90}Zr$ | SIII | 5.11666 | 5.03803 | .462754 | .470953 | 5.08055 | .466451 |
| | SkM* | 5.03120 | 4.91373 | .496218 | .512470 | 4.98017 | .503553 |
| $^{208}Pb$ | SIII | 6.87254 | 6.81210 | .514625 | .447845 | 6.85041 | .482834 |
| | SkM* | 6.78819 | 6.67531 | .556836 | .488960 | 6.75631 | .524144 |

The equilibrium values of the density distributions (5), $u^\circ=(R_n^\circ, R_p^\circ, b_n^\circ, b_p^\circ)$, obtained by minimizing the eq. (6) with respect to the collective variables $u=(R_n, R_p, b_n, b_p)$, are given in table 1.

for a nuclear Hamiltonian with the Skyrme-type forces SkM* and SIII. In this case, $V(u^o)$ and $\phi(u^o)$ are in quite satisfactory agreement with respect to the self-consistent HF results for the total energy $E_{HF}$ and the ground state wave function $\phi_{HF}$ obtained in [13,14]. In table 1. we also list the equilibrium density parameters $R^o$ and $b^o$ obtained with a proportional SF density approximation (SFP) assuming only two variational parameters $R=R_n=R_p$ and $b=b_n=b_p$.

The collective kinetic energy in the same ATDHF approximation [12] is of the form

$$K = \frac{m}{2}\sum_{i,j=1}^{4} m_{ij}(u)\ \dot{u}_i \dot{u}_j \ , \qquad (7)$$

where $\dot{u}=(\dot{R}_n,\dot{R}_p,\dot{b}_n,\dot{b}_p)$ is the time derivative of $u=(R_n,R_p,b_n,b_p)$. The inertial tensor matrix elements $m_{ij}$ in eq. (7) are given by the equation

$$m_{ij} = \left[\frac{\varrho_{q_i}^{SF}\varrho_{q_j}^{SF}(1-\delta_{q_i q_j}) + \varrho_{q_i}(1+\beta\varrho_{q_j})\delta_{q_i q_j}}{1+\beta(\varrho_n^{SF}+\varrho_p^{SF})}\right] v_i v_j\, d^3r \ , \qquad (8)$$

where the isospin index q =n(p) for i=1,3 (2,4) is ordered to correspond to the index i=(1,2,3,4) labeling the collective variables vector $u=(u_1,u_2,u_3,u_4)=(R_n,R_p,b_n,b_p)$. The non-locality parameter $\beta=(m/2\hbar)(t_1+t_2)$ in eq. (8) comes from the exchange Skyrme-type forces and the velocity fields $v_i$, (i=1,2,3,4), measured in units $\dot{u}_i$, are expressed as

$$v_i = \left(\int_0^r \frac{\partial\varrho_{q_i}^{SF}}{\partial u_i}\, r'^2 dr'\right)/\left(\varrho_{q_i}^{SF} r\right), \quad (i=1,2,3,4). \qquad (9)$$

Further, we quantize the ATDHF classical Hamiltonian $\mathcal{H}_{ATDHF}=K+V$, see eq. (6) and (7), in the known harmonic approximation (HA) expanding the collective potential energy (6) around its equilibrium value $V_o = V(u^o)$ up to second order in the deviations $(u_i-u_i^o)$. Diagonalizing the resulting Hamiltonian in harmonic approximation we obtain eigen vectors $S_i^{(\alpha)}$, ($\alpha$=1,2,3,4)



which transform the original collective variables $u=(R_n,R_p,b_n,b_p)$ into a new set of normal coordinates $Q=(Q_1,Q_2,Q_3,Q_4)$ in which we have:

$$\mathcal{H}_{ATDHF} = \sum_{\alpha=1}^{4}\mathcal{H}_\alpha = (1/2)\sum_{\alpha=1}^{4}(\dot{Q}_\alpha^2 + w_\alpha Q_\alpha^2) \ . \qquad (10)$$

Therefore, in the harmonic approximation the monopole vibrations connected with the density variables $(R_n,R_p,b_n,b_p)$ are presented as four independent normal vibrations with respect to $(Q_1,Q_2,Q_3,Q_4)$. Their excitation energies $\hbar w_\alpha$, ($\alpha$=1,2,3,4) are listed in table 2.

Table 2.

The excitation energies $\hbar w_\alpha$ (in Mev) for ISMR ($\alpha$=1), IVMR ($\alpha$=2) and the antiscaling-type isoscalar (isovector) monopole vibration $\alpha$=3 ($\alpha$=4) calculated with SIII and SkM* forces.

| | SIII | | | | SkM* | | | |
|---|---|---|---|---|---|---|---|---|
| Nuclei | $\hbar w_1$ | $\hbar w_2$ | $\hbar w_3$ | $\hbar w_4$ | $\hbar w_1$ | $\hbar w_2$ | $\hbar w_3$ | $\hbar w_4$ |
| $^{12}$C | 27.03 | 32.56 | 60.83 | 67.22 | 22.23 | 27.44 | 48.24 | 58.31 |
| $^{16}$O | 28.42 | 34.68 | 61.70 | 66.30 | 23.20 | 30.57 | 50.48 | 59.00 |
| $^{28}$Si | 28.10 | 35.27 | 56.30 | 63.12 | 23.01 | 32.17 | 48.06 | 58.00 |
| $^{32}$S | 24.63 | 30.52 | 48.68 | 55.75 | 20.41 | 27.02 | 39.83 | 51.53 |
| $^{40}$Ca | 25.70 | 33.50 | 48.61 | 57.55 | 20.73 | 29.82 | 40.54 | 51.95 |
| $^{48}$Ca | 25.20 | 35.40 | 48.77 | 58.62 | 20.41 | 31.65 | 41.71 | 53.29 |
| $^{56}$Ni | 25.55 | 35.22 | 49.36 | 58.45 | 20.42 | 32.15 | 42.78 | 53.63 |
| $^{90}$Zr | 22.51 | 33.17 | 42.44 | 52.22 | 17.82 | 29.84 | 36.15 | 47.45 |
| $^{208}$Pb | 17.54 | 29.82 | 30.07 | 49.06 | 13.66 | 25.96 | 32.78 | 43.14 |

The detailed analysis of these excitation energies, the energy weighted sum rules (EWSR) and the corresponding rms radii, local and transition densities have shown that the first ($\alpha$=1) and second ($\alpha$=2) normal modes can be identified with the ISMR and IVMR, respectively. These are dynamically coupled bulk

and surface vibrations in which the density half-radius vibrates in phase (scaling-type vibrations) to the density surface. The normal mode $\lambda=3$ ($\lambda=4$) is of the isoscalar (isovector) antiscaling-type monopole vibrations in which the density half-radius vibrates out of phase to the nuclear surface.

## III. Transition Densities

The transition densities corresponding to the four normal monopole vibrations mentioned in the previous section can be obtained in the following way. First we transform the original coordinates $u=(R_n, R_p, b_n, b_p)$ in the SF density distributions, eq. (3), into normal coordinates $(Q_1, Q_2, Q_3, Q_4)$ using the canonical transformation vectors $S_i^{(\lambda)}$. Then, expanding the obtained expression up to second order in $S_i^{(\lambda)}$ (the harmonic approximation), we substitute $\{Q_\alpha\}$ by the corresponding operators $Q_\alpha=(\hbar/2w_\alpha)^{1/2}(a_\alpha^+ + a_\alpha)$. Finally, we take the expectation values between ground and excited one phonon states of the corresponding $\lambda$-type.

This procedure leads to the following expression for the isoscalar (T=0) transition density:

$$\rho_\lambda^{(T=0)}(r)=A_1^\lambda(\partial\rho^{SF}/\partial R_n)+A_2^\lambda(\partial\rho^{SF}/\partial R_p)+A_3^\lambda(\partial\rho^{SF}/\partial b_n)+A_4^\lambda(\partial\rho^{SF}/\partial b_p), \qquad (11)$$

where $\rho^{SF}(r)=\rho_n^{SF}(r)+\rho_p^{SF}(r)$, eqs. (3) and (5), and its derivatives with respect to the SF density parameters ($R_n, R_p, b_n, b_p$) are taken at the equilibrium values (see table 1.). The amplitude coeffitients $A_i^\lambda$, $(i,\lambda=1,2,3,4)$ in eq. (11) defined according to the equation

$$A_i^\lambda=(\hbar/2w_\lambda)^{1/2}S_i^{(\lambda)}, \quad (i,\lambda=1,2,3,4) \qquad (12)$$

are given in tables 3 and 4 for the SkM* and SIII effective forces, respectively.



Table 3.

Transition density coeffitients obtained with Skyrme forces SkM*.

| Nuclei | $\lambda$ | $A_1$ | $A_2$ | $A_3$ | $A_4$ | $A^{(T=0)}$ | $B^{(T=0)}$ | $A^{(T=1)}$ | $B^{(T=1)}$ |
|---|---|---|---|---|---|---|---|---|---|
| $^{12}$C | 4 | .9246 | -.1595 | -.1470 | -.0138 | .3826 | -.0666 | .5420 | -.0804 |
| | 3 | .0980 | .9551 | -.0336 | -.1358 | .5266 | -.0847 | -.0429 | .0511 |
| | 2 | -.1696 | .3604 | .1368 | -.1154 | .0954 | .0107 | -.2650 | .1261 |
| | 1 | .0360 | -.3050 | .0587 | .1695 | -.1344 | .1141 | .1705 | -.0554 |
| $^{16}$O | 4 | .6083 | -.4267 | -.1186 | .0750 | .0908 | -.0207 | .5184 | -.0967 |
| | 3 | .3851 | .6168 | -.0925 | -.1250 | .5005 | -.1090 | -.1074 | .0170 |
| | 2 | -.0349 | .1314 | .0971 | -.0995 | .0309 | -.0050 | -.0821 | .0987 |
| | 1 | .0712 | .0223 | .0661 | .0949 | .0520 | .0798 | .0149 | -.0048 |
| $^{28}$Si | 4 | .4098 | -.3459 | -.0921 | .0735 | .0367 | -.0103 | .3772 | -.0826 |
| | 3 | .3131 | .4009 | -.0892 | -.1008 | .3568 | -.0950 | -.0471 | .0079 |
| | 2 | .0116 | .0510 | .0683 | -.0763 | .0182 | -.0051 | -.0196 | .0725 |
| | 1 | .1035 | .0907 | .0457 | .0581 | .0987 | .0516 | .0063 | -.0011 |
| $^{32}$S | 4 | .4314 | -.3549 | -.0834 | .0628 | .0476 | -.0117 | .3916 | -.0729 |
| | 3 | .3470 | .4448 | -.0864 | -.0894 | .3951 | -.0878 | -.0538 | .0047 |
| | 2 | -.0362 | .1560 | .0810 | -.0851 | .0468 | -.0074 | -.0967 | .0844 |
| | 1 | .0335 | -.0283 | .0522 | .0880 | .0103 | .0692 | .0205 | -.0074 |
| $^{40}$Ca | 4 | .3338 | -.3004 | -.0804 | .0694 | .0228 | -.0073 | .3164 | -.0748 |
| | 3 | .2832 | .3309 | -.0900 | -.0910 | .3074 | -.0905 | -.0302 | .0037 |
| | 2 | .0186 | .0541 | .0604 | -.0735 | .0225 | -.0069 | -.0178 | .0672 |
| | 1 | .0991 | .0905 | .0400 | .0513 | .0965 | .0453 | .0058 | -.0010 |
| $^{48}$Ca | 4 | .2199 | -.3415 | -.0572 | .0887 | -.0388 | .0142 | .2604 | -.0683 |
| | 3 | .2729 | .2398 | -.0856 | -.0855 | .2574 | -.0845 | .0993 | -.0244 |
| | 2 | -.0290 | -.0484 | .0536 | -.0598 | -.0245 | .0096 | .0085 | .0568 |
| | 1 | .1110 | .1131 | .0427 | .0260 | .1152 | .0349 | .0049 | .0076 |
| $^{56}$Ni | 4 | .2601 | -.2339 | -.0746 | .0645 | .0174 | -.0066 | .2465 | -.0694 |
| | 3 | .2029 | .2462 | -.0797 | -.0822 | .2251 | -.0811 | -.0259 | .0047 |
| | 2 | .0504 | .0050 | .0427 | -.0578 | .0124 | -.0053 | .0230 | .0503 |
| | 1 | .1224 | .1208 | .0251 | .0311 | .1225 | .0277 | .0037 | .0000 |
| $^{90}$Zr | 4 | .1805 | -.2158 | -.0576 | .0642 | -.0107 | .0051 | .1931 | -.0598 |
| | 3 | .1667 | .1938 | -.0729 | -.0739 | .1784 | -.0733 | .0382 | -.0116 |
| | 2 | .0293 | -.0161 | .0347 | -.0523 | -.0098 | .0040 | .0240 | .0428 |
| | 1 | .1145 | .1143 | .0216 | .1884 | .1145 | .0204 | .0072 | .0040 |
| $^{208}$Pb | 4 | .0706 | -.1540 | -.0268 | .0739 | .0092 | .0054 | .1141 | -.0466 |
| | 3 | .1274 | .0102 | -.0454 | -.0452 | .1029 | -.0575 | .0392 | -.0164 |
| | 2 | -.0483 | -.0955 | .0574 | .0233 | -.0049 | .0049 | .0372 | .0267 |
| | 1 | .1014 | .1003 | .0109 | .0609 | .1035 | .0079 | .0199 | .0002 |

In the proportional density approximation for the equilibrium density distributions SFP when $R_n=R_p=R$ and $b_n=b_p=b$ (see the previous section) the isoscalar transition density is of the form:

$$\rho_\alpha^{(T=0)}(r)=A^\alpha(\partial\rho^{SF}/\partial R)+B^\alpha(\partial\rho^{SF}/\partial b), \qquad (13)$$

where the equilibrium SFP density parameters $R^\circ$, $b^\circ$ have been listed in table 1.

Table 4.

Transition density coeffitients obtained with Skyrme forces SIII.

| Nuclei | $\alpha$ | $A_1$ | $A_2$ | $A_3$ | $A_4$ | $A^{(T=0)}$ | $B^{(T=0)}$ | $A^{(T=1)}$ | $B^{(T=1)}$ |
|---|---|---|---|---|---|---|---|---|---|
| $^{12}$C | 4 | .7649 | .1622 | -.1398 | -.0319 | .4737 | -.0868 | .2943 | -.0505 |
| | 3 | .2679 | -.8443 | -.0399 | .1353 | -.2672 | .0464 | .5566 | -.0901 |
| | 2 | -.1075 | .1506 | .1298 | -.0808 | .0199 | .0079 | -.1364 | .1107 |
| | 1 | -.0068 | -.2013 | .0589 | .1512 | -.0883 | .1059 | .0553 | -.0223 |
| $^{16}$O | 4 | .5756 | -.3667 | -.1198 | .0702 | .1065 | -.0247 | .4714 | -.0948 |
| | 3 | .3033 | .5576 | -.0731 | -.1192 | .4297 | -.0962 | -.1249 | .0237 |
| | 2 | -.0235 | .0631 | .0945 | -.0849 | .0137 | -.0012 | -.0439 | .0904 |
| | 1 | .0243 | -.0087 | .0698 | .0964 | .0099 | .0831 | .0119 | -.0046 |
| $^{28}$Si | 4 | .3772 | -.3198 | -.0923 | .0752 | .0359 | -.0104 | .3476 | -.0835 |
| | 3 | .2801 | .3557 | -.0803 | -.0942 | .3172 | -.0871 | -.0449 | .0095 |
| | 2 | .0247 | .0103 | .0674 | -.0660 | .0083 | -.0024 | .0080 | .0671 |
| | 1 | .0582 | .0491 | .0511 | .0656 | .0548 | .0582 | .0061 | -.0016 |
| $^{32}$S | 4 | .3907 | -.3306 | -.0843 | .0679 | .0434 | -.0112 | .3587 | -.0757 |
| | 3 | .2986 | .3760 | -.0748 | -.0850 | .3358 | -.0795 | -.0514 | .0087 |
| | 2 | -.0297 | .0683 | .0813 | -.0705 | .0167 | -.0015 | -.0503 | .0771 |
| | 1 | .0030 | -.0300 | .0575 | .0864 | -.0113 | .0721 | .0136 | -.0058 |
| $^{40}$Ca | 4 | .3054 | -.2776 | -.0803 | .0708 | .0214 | -.0071 | .2907 | -.0753 |
| | 3 | .2516 | .2915 | -.0805 | -.0849 | .2713 | -.0825 | -.0281 | .0055 |
| | 2 | .0274 | .0144 | .0597 | -.0624 | .0114 | -.0034 | .0063 | .0614 |
| | 1 | .0626 | .0564 | .0442 | .0570 | .0607 | .0504 | .0056 | -.0015 |
| $^{48}$Ca | 4 | .1805 | -.3286 | -.0504 | .0961 | -.0502 | .0177 | .2338 | -.0668 |
| | 3 | .2684 | .1786 | -.0862 | -.0688 | .2287 | -.0779 | .1046 | -.0311 |
| | 2 | -.0108 | -.0663 | .0423 | -.0595 | -.0084 | .0006 | .0296 | .0504 |
| | 1 | .0811 | .0692 | .0512 | .0263 | .0805 | .0401 | .0098 | .0104 |
| $^{56}$Ni | 4 | .2382 | -.2157 | -.0750 | .0657 | .0165 | -.0066 | .2264 | -.0701 |
| | 3 | .1856 | .2224 | -.0740 | -.0787 | .2040 | -.0764 | -.0236 | .0057 |
| | 2 | .0585 | -.0197 | .0417 | -.0498 | .0069 | -.0032 | .0392 | .0459 |
| | 1 | .0904 | .0896 | .0293 | .0368 | .0909 | .0328 | .0038 | -.0004 |
| $^{90}$Zr | 4 | .1605 | -.2012 | -.0561 | .0677 | -.0139 | .0062 | .1754 | -.0603 |
| | 3 | .1571 | .1648 | -.0696 | -.0687 | .1599 | -.0690 | .0374 | -.0143 |
| | 2 | .0340 | -.0415 | .0361 | -.0457 | -.0029 | -.0006 | .0387 | .0377 |
| | 1 | .0930 | .0849 | .0261 | .0216 | .0894 | .0241 | .0100 | .0048 |
| $^{208}$Pb | 4 | .0519 | -.1387 | -.0217 | .0761 | -.0108 | .0058 | .1026 | -.0466 |
| | 3 | .1268 | -.0032 | -.0520 | -.0345 | .0917 | -.0548 | .0338 | .0182 |
| | 2 | -.0299 | -.0943 | .0456 | -.0260 | -.0007 | -.0001 | .0452 | .0221 |
| | 1 | .0868 | .0762 | .0143 | .0713 | .0862 | .0097 | .0218 | .0010 |

The amplitude coeffitients $A^\alpha = A_1^\alpha + A_2^\alpha$ and $B^\alpha = A_3^\alpha + A_4^\alpha$,



defined by the amplitudes $A_i^\alpha$, eq. (12) calculated in SFP approximation, are also given in tables 3 and 4.

In the same manner the isovector [18] transition densities $\rho_\alpha^{(T=1)}(r)$ follow from eqs. (11) and (13) by substituting $A_i^\alpha = -A_i^\alpha$ for i=2,4 and $\alpha$ =1,2,3,4 in eq. (12).

The realistic behaviour of the presented transition densities can be seen from the following points observed comparing our numerical results with the experimental data and other theoretical estimates:

1) The calculated ISMR excitation energies ($\hbar\omega_1$, table 2., SkM*) are in an excellent agreement with the experimental peak energy data [1]. At the same time, the ISMR and IVMR excitation energies ($\hbar\omega_1$ and $\hbar\omega_2$, table 2., SIII) reproduce the corresponding RPA results [5].

2) Both $\rho_\alpha^{(T=0)}$ and $\rho_\alpha^{(T=1)}$ for ISMR and IVMR are rather close to the densities obtained within the RPA method [4]. In particular, the ISMR densities, eq. (13), almost exactly reproduce the results obtained in the GCM calculations [7] for light nuclei, while for heavy nuclei this transition density is near to the usually used Tassie transition density, eq. (1) and to the scaling-type (T=0) density [19] obtained by means of the so-called Extended Thomas-Fermi method (ETF) [20].

3) Our calculations show that ISMR and IVMR exhausts almost completely the isoscalar and isovector EWSR [5] respectively.

## IV. Summary

In conclusion, it is evident that using eqs. (3), (13) and the values of the transition density coeffitients $R^\circ$, $b^\circ$ (table 1) and $A^\alpha$, $B^\alpha$ (tables 3 or 4) one can easily apply $\rho_\alpha^{(T=0)}$ to a number of particular problems dealing with the exitation of ISMR

($\alpha$=1) or IVMR ($\alpha$=2). In more precise numerical calculations one can also apply the transition densities from eqs. (3), (11) with the corresponding values of $R_n^c, R_p^c, b_n^c, b_p^c$ (table 1 ) and $A_i^\alpha$ (tables 3 and 4 ). In the latter case a small contribution of isovector (isoscalar) component appears in the ISMR (IVMR) reproduced by $\alpha$=1 ($\alpha$=2).

Moreover, the tables 1 , 3 , and 4 allow the isoscalar (isovector) antiscaling type vibrations, the mode $\alpha$=3 ($\alpha$=4) to be investigated using the eqs. (3) and (11) or (13). In this respect it is very interesting to answer the question: what kind of evidences can be found out within the available experimental data about the actual clarification of these highly-excited antiscaling-type monopole states.

In the case of another nuclei and Skyrme-type forces one can obtain the necessary transition density coeffitients using the method discussed in section 2.

Finally, we hope that the suggested in the present work transition densities obtained on the base of an effective NN-interaction in nuclei will prove to be a useful tool in the description of giant monopole isoscalar and isovector monopole experimental data.

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Переходные плотности гигантских монопольных резонансов в рамках локально-масштабного варианта адиабатического ВЗХФ метода

В рамках локально-плотностной версии адиабатического времязависящего метода Хартри-Фока с силами Скирма SM* и SIII для ряда четно-четных ядер рассчитаны переходные плотности, соответствующие ядерному гигантскому монопольному резонансу. В работе изложен подход, в котором получено конкретное выражение и рассчитаны необходимые коэффициенты переходных плотностей. Последние получены в простой форме и могут быть использованы, например, для анализа данных по неупругому ядерному рассеянию частиц методом искаженных волн, что дает возможность проверять теоретическую интерпретацию гигантских монопольных резонансов.

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Giant Monopole Transition Densities within the Local-Scale ATDHF Approach

We propose transition densities for even-even nuclei corresponding to nuclear giant monopole resonances obtained within a local-scale ATDHF approach in terms of effective Skyrme-type forces SkM* and SIII. The approach, the particular form and all necessary coefficients of these transition densities are reported. They are of a simple analytical form and may be directly used for example in DWBA analyses of inelastic scattering experiments and in such a way allowing a crucial test of the theoretical interpretation of giant monopole resonances.

The investigation has been performed at the Laboratory of Theoretical Physics, JINR.

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