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ON THE PROPERTIES OF THE $N = 82$ EVEN–EVEN NUCLEI

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The work presents systematic theoretical investigation of even–even $N = 82$ isotones from ^{132}Sn up to ^{146}Gd, based on the random phase approximation with finite range two-body effective interaction identical in the particle–particle, particle–hole and pairing channels. Comparison with the experiment concerning the values of energies of levels, probabilities of electromagnetic transitions and electromagnetic moments available by now is offered.

На основе метода хаотической фазы с эффективным взаимодействием конечного радиуса, единым в частично-частичном, частично-дырочном и в спаривательном каналах, произведено систематическое теоретическое исследование четно-четных ядер с $N = 82$, от ^{132}Sn до ^{146}Gd. Представлено сравнение с экспериментом, касающееся спектров уровней, вероятностей электромагнитных переходов и ядерных моментов, в которое включены имеющиеся к настоящему времени экспериментальные данные.

INTRODUCTION

By the present time there are accumulated a large set of experimental data [1–10] concerning the properties of even–even nuclei with $N = 82$. The investigated region of the mentioned isotones covers the interval from the neutron excess double magic nuclei ^{132}Sn up to the neutron deficient nuclide ^{146}Gd, in which the proton configuration with $Z = 64$ manifests some properties of the closed shell.

It should be mentioned that the nuclei considered here were the subject of intent attention in many theoretical papers. So, still in the works [11, 12] the set of proton single-particle energies for the $N = 82$ nuclei was established by using the inverse gap equations yielding as a result the evidence of an energy gap between the d5/2, g7/2 and the h11/2, d3/2, s1/2 levels for nuclei close to ^{146}Gd. In works [13–16] the properties of the mentioned nuclei were considered using the TD approximation accounting the pairing correlations and particle number projection. In [17] the $N = 82$ isotones were considered in the framework of generalized seniority scheme. Shell model calculations of level energies for even and odd nuclei with $N = 82$ were carried out in [18], while in [15, 16, 19–21] the calculations of selected electromagnetic properties were performed demonstrating that some of them are very sensitive to the values of spectroscopical factors, i.e., to the values of single-particle energies.

At the same time numerous recent experimental data available by now need a new trying to review the situation as a whole. Due to this it seems quite reasonable to perform a theoretical analysis of the even $N = 82$ nuclei in the framework of common theoretical scheme using

for all of them identical parameters that define residual interaction and the mean field. As all the mentioned nuclides are close to the filled shells they may be considered as having the spherical form. As such a theoretical tool we use here the random phase approximation considering the pairing correlations and smoothing of the Fermi step (QPRPA method).

The mentioned method but without pairing correlations was successfully used by us in calculations of magic nuclei or nuclei having the structure «magic nuclei plus two (quasi)particles» [22–27]. Here we use the more general scheme that covers a larger variety of nuclei and automatically transforms to the method used by us earlier when the numbers of nucleons approach to that ones forming the closed shells. Everywhere if possible we make a comparison with the experiment.

1. BASIC RELATIONS DEFINING THE SPECTRA OF LEVELS

Here and below we shall proceed from the assumption that the nuclei are described by the Hamiltonian $\hat{H}$ with an effective two-body interaction $\hat{\vartheta}(x_1, x_2)$, where $\hat{\vartheta}$ is a scalar relatively to rotations in coordinate-spin and isospin spaces. In the case of two interacting protons we shall also include in $\hat{\vartheta}$ the pairing Coulomb interaction that destroys isotopic invariance:

$$\hat{H} = \sum_{\alpha\beta} \langle\beta|\hat{T}|\alpha\rangle a^+_\alpha a_\beta + \frac{1}{4} \sum_{\alpha\beta\gamma\delta} {}_a\langle\alpha\beta|\hat{\vartheta}|\gamma\delta\rangle_a a^+_\alpha a^+_\beta a_\delta a_\gamma. \tag{1}$$

Here $\hat{T}$ is a kinetic energy operator, ${}_a\langle\alpha\beta|\hat{\vartheta}|\gamma\delta\rangle_a$ is an antisymmetrized two-body matrix element of the effective interaction $\hat{\vartheta}$, $|\alpha\rangle$ are single-particle eigenfunctions of an arbitrary single-particle set while a^+_α and a_β are creation and annihilation nucleon operators.

Using the standard procedure [28, 29] we can pass to the quasi-particle basis, $a^+ \to \xi^+$:

$$a^+_\alpha = u_{|\alpha|}\xi^+_\alpha - v_{|\alpha|}\varphi_\alpha \xi_{-\alpha}; \quad u^2_{|\alpha|} + v^2_{|\alpha|} = 1. \tag{2}$$

Here φ_α is the phase of the particle–hole transformation (we use below the designations $u_{|\alpha|} = u_\alpha$ and $v_{|\alpha|} = v_\alpha$ omitting the magnetic quantum numbers that enter only in φ_α), $|-\alpha\rangle$ is a state with the opposite sign of magnetic quantum number with respect to $|\alpha\rangle$. Defining the Hartree–Fock–Bogoliubov vacuum $|\Phi_0\rangle$ of the even–even system as satisfying the stability condition, $\langle\Phi_0|\hat{H}\xi^+_i\xi^+_k|\Phi_0\rangle = 0$, we may transform the expression (1) to the form [30, 31]:

$$\hat{H} = E_0 + \sum_r E_r \xi^+_r \xi_r + H_{\rm int}, \tag{3}$$

where

$$E_0 = \sum_r v^2_r \left(\varepsilon_r - \frac{1}{r} \sum_s v^2_s {}_a\langle rs|\hat{\vartheta}|rs\rangle_a) \right) - \frac{1}{4} \sum_r \frac{\Delta^2_r}{E_r}, \tag{4}$$

$$\langle r|\hat{T}|m\rangle + \sum_n {}_a\langle rn|\hat{\vartheta}|mn\rangle_a v^2_n = \varepsilon_r \delta_{rm}, \tag{5}$$

$$v_r^2 = \frac{1}{2}\left[1 - \frac{\varepsilon_r - \lambda}{E_r}\right], \quad E_r = \sqrt{(\varepsilon_r - \lambda)^2 + \Delta_r^2}, \quad N(p,n) = \sum_r v_r^2(p,n), \tag{6}$$

$$\Delta_r = -\frac{1}{4}\sum_s \frac{{}_a\langle r - r|\hat{\vartheta}|s - s\rangle_a \varphi_r \varphi_s \Delta_s}{E_s}. \tag{7}$$

In formula (3) E_0 is the energy of vacuum in the sense of the Hartree–Fock–Bogoliubov method and all the single-particle indices entering equations (3)–(7) correspond to the self-consistent HFB orbitals that are defined by the formulas (5)–(7) while the $H_{\rm int}$ term presents the normal, with respect to the vacuum Φ_0, potential energy term from the equation (1). It has the form $H_{\rm int} = H_{22} + H_{40} + H_{04} + H_{31} + H_{13}$, where H_{ik} is an item that contains «i» operators ξ^+ and «k» operators ξ. As the v_r, the Δ_r in equations (4)–(7) also do not contain the dependence on magnetic quantum numbers.

We shall not be interested below in the total binding energies supposing also that the nuclei in question have the spherical form with the mean field generated by some phenomenological potential. In this case we should remain from the system (4)–(7) only the (6) and (7) relations, where $|r\rangle = |n_r l_r j_r(m_r)\rangle$ is an eigenfunction of a spherical mean field.

Here we use spherical harmonics with the phases of the work [32], where $Y_{l-m} = (-1)^m Y_{lm}^*$, and define φ_r as $\varphi_r = (-1)^{l_r + j_r - m_r}$. Under the mentioned conditions and using short range attraction in the pairing channel, the $\Delta_r = \Delta_{n_r l_r j_r}$ values are proved to be of the same sign while the system of equations defining the superfluid characteristics of nuclei has the form

$$\Delta_{nlj}(p,n) = -\frac{1}{2}\sum_{n'l'j'} \sqrt{\frac{2j'+1}{2j+1}} \frac{\Delta_{n'l'j'}(p,n)}{E_{n'l'j'}(p,n)} \langle (n'l'j')^2; 0^+|\hat{\vartheta}|(nlj)^2; 0^+\rangle (-1)^{l+l'}, \tag{8}$$

$$N(p,n) = \sum_{n\,l\,j}(2j+1)v_{nlj}^2(p,n) = \sum_{n\,l\,j}\frac{2j+1}{2}\left[1 - \frac{\varepsilon_{nlj}(p,n) - \lambda(p,n)}{E_{nlj}(p,n)}\right], \tag{9}$$

$$E_{nlj}(p,n) = \sqrt{[\varepsilon_{nlj}(p,n) - \lambda(p,n)]^2 + \Delta_{nlj}^2(p,n)}. \tag{10}$$

Supposing the presence of correlations in the true ground state $|\tilde{0}\rangle$ of an even–even nuclei we define the creation operator $Q_{n,JM}^+$ of the one-phonon excited state $|\omega_n, JM\rangle$ with $|\omega_n, JM\rangle = Q_{n,JM}^+|\tilde{0}\rangle$ in the following way:

$$Q_{n,JM}^+ = \sum_{a\geq b} X_{j_a j_b}^{n,J}\left[\xi_a^+\xi_b^+\right]_{JM} - \sum_{c\geq d} Y_{j_c j_d}^{n,J}\left[\xi_c\xi_d\right]_{JM}, \tag{11}$$

where

$$\left[\xi_a^+\xi_b^+\right]_{JM} = \frac{1}{\sqrt{1+\delta_{j_a j_b}}}\sum_{m_a m_b} C_{j_a m_a j_b m_b}^{JM}\xi_{j_a m_a}^+\xi_{j_b m_b}^+, \tag{12}$$

$$\left[\xi_c^+\xi_d^+\right]_{JM} = \frac{1}{\sqrt{1+\delta_{j_cj_d}}}\sum_{m_cm_d} C^{JM}_{j_cm_cj_dm_d}\xi^+_{j_c-m_c}\xi^+_{j_d-m_d}\varphi_c\varphi_d, \tag{13}$$

$$X^{n,J}_{j_aj_b} = \langle\omega_n; JM|\left[\xi_a^+\xi_b^+\right]_{JM}|\tilde{0}\rangle,\ Y^{n,J}_{j_aj_b} = \langle\omega_n; JM|\left[\xi_a\xi_b\right]_{JM}|\tilde{0}\rangle. \tag{14}$$

Here and below $j_a \equiv n_a, l_a, j_a, (t_{z_a})$; for even–even nuclei both the indices «*a*» and «*b*» («*c*» and «*d*») in (11) simultaneously belong to protons or to neutrons: $a, b(c, d) \in p$ or $a, b(c, d) \in n$.

Using the relation $[H, Q_n^+] = \omega_n Q_n^+$, which is equivalent to the Shrödinger equation and calculating the commutators from both the sides of the mentioned relation with the pair operators $\xi_i^+\xi_k^+$ and $\xi_i\xi_k$ one may obtain the set of the QPRPA equations [30, 31] that define the amplitudes «X» and «Y» of the states $|\omega_n, JM\rangle$ and the eigenvalues ω_n. They have the form

$$\left\|\begin{array}{cc}[(E-\omega)I+A] & B\\ -B & -[(E+\omega)I+A]\end{array}\right\|\left\|\begin{array}{c}X\\Y\end{array}\right\| = 0. \tag{15}$$

In formulas (15) $E = E_{ab} = E_{j_a} + E_{j_b}$, $I_{cd,ab} = \delta_{j_aj_c}\delta_{j_bj_d}$ while the matrix elements of the submatrices A and B in the case of even–even nuclei occur to be as follows:

$$\begin{aligned}A_{cd,ab} \equiv A^J_{j_cj_d,j_aj_b} &= (u_{j_c}u_{j_d}u_{j_a}u_{j_b} + v_{j_c}v_{j_d}v_{j_a}v_{j_b})\,{}_a\langle j_cj_d; J|\hat{\vartheta}|j_aj_b; J\rangle_a +\\ &+ (u_{j_c}v_{j_d}u_{j_a}v_{j_b} + v_{j_c}u_{j_d}v_{j_a}u_{j_b})\,{}_a\langle j_c\bar{j_d}; J|\hat{\vartheta}|j_a\bar{j_b}; J\rangle_a +\\ &+(-1)^{j_a+j_b+J+1}(v_{j_c}u_{j_d}u_{j_a}v_{j_b} + u_{j_c}v_{j_d}v_{j_a}u_{j_b})\,{}_a\langle j_c\bar{j_d}; J|\hat{\vartheta}|j_b\bar{j_a}; J\rangle_a;\end{aligned} \tag{16}$$

$$\begin{aligned}B_{cd,ab} \equiv B^J_{j_cj_d,j_aj_b} &= (u_{j_c}u_{j_d}v_{j_a}v_{j_b} + v_{j_c}v_{j_d}u_{j_a}u_{j_b})\,{}_a\langle j_cj_d; J|\hat{\vartheta}|j_aj_b; J\rangle_a -\\ &- (u_{j_c}v_{j_d}v_{j_a}u_{j_b} + v_{j_c}u_{j_d}u_{j_a}v_{j_b})\,{}_a\langle j_c\bar{j_d}; J|\hat{\vartheta}|j_a\bar{j_b}; J\rangle_a +\\ &+(-1)^{j_a+j_b+J}(v_{j_c}u_{j_d}v_{j_a}u_{j_b} + u_{j_c}v_{j_d}u_{j_a}v_{j_b})\,{}_a\langle j_c\bar{j_d}; J|\hat{\vartheta}|j_b\bar{j_a}; J\rangle_a.\end{aligned} \tag{17}$$

In formulas (16) and (17) ${}_a\langle j_cj_d; J|\hat{\vartheta}|j_aj_b; J\rangle_a$ and ${}_a\langle j_c\bar{j_d}; J|\hat{\vartheta}|j_a\bar{j_b}; J\rangle_a$ are the antisymmetric matrix elements of the effective interaction $\hat{\vartheta}$ in the particle–particle and particle–hole channels with a given spin. They have the form

$$\begin{aligned}{}_a\langle j_cj_d; J|\hat{\vartheta}|j_aj_b; J\rangle_a &= \frac{1}{\sqrt{(1+\delta_{j_cj_d})(1+\delta_{j_aj_b})}}\Big[\langle j_cj_d; J|\hat{\vartheta}|j_aj_b; J\rangle +\\ &+ (-1)^{j_a+j_b+J+1}\langle j_cj_d; J|\hat{\vartheta}|j_bj_a; J\rangle\Big],\end{aligned} \tag{18}$$

$$\begin{aligned}{}_a\langle j_c\bar{j_d}; J|\hat{\vartheta}|j_a\bar{j_b}; J\rangle_a &= -\frac{(-1)^{l_b+l_d}}{\sqrt{(1+\delta_{j_cj_d})(1+\delta_{j_aj_b})}}\sum_{J'}(2J'+1)W[j_bj_aj_cj_d; JJ']\times\\ &\times\Big[\langle j_bj_c; J'|\hat{\vartheta}|j_dj_a; J'\rangle + (-1)^{j_d+j_a+J'+1}\langle j_bj_c; J'|\hat{\vartheta}|j_aj_d; J'\rangle\Big].\end{aligned} \tag{19}$$

If we represent the effective interaction $\hat{\vartheta}$ in the form

$$\hat{\vartheta}(x_1, x_2) = \hat{\vartheta}^{(0)}(\bar{r}_1, \bar{r}_2, \bar{\sigma}_1, \bar{\sigma}_2) + \hat{\vartheta}^{(1)}(\bar{r}_1, \bar{r}_2, \bar{\sigma}_1, \bar{\sigma}_2)\boldsymbol{\tau}_1\boldsymbol{\tau}_2 + + V_c(\mathbf{r}_1, \mathbf{r}_2, t_{z_1}, t_{z_2}), \qquad (20)$$

where $\hat{\vartheta}^{(0)}$ and $\hat{\vartheta}^{(1)}$ are scalars in coordinate–spin space and V_c is a Coulomb interaction, then the pair matrix elements $\langle j_c j_d; J|\hat{\vartheta}|j_a j_b; J\rangle$ entering the right-hand side of relations (18) and (19) are different from zero only in the case of $t_{z_a} + t_{z_b} = t_{z_c} + t_{z_d}$ and have the form

$$\begin{aligned}
\langle j_c(n) j_d(n); J|\hat{\vartheta}|j_a(n) j_b(n); J\rangle &= \langle j_c j_d; J|\hat{\vartheta}^{(0)} + \hat{\vartheta}^{(1)}|j_a j_b; J\rangle, \\
\langle j_c(p) j_d(p); J|\hat{\vartheta}|j_a(p) j_b(p); J\rangle &= \left\langle j_c j_d; J\left|\hat{\vartheta}^{(0)} + \hat{\vartheta}^{(1)} + \frac{e^2}{r_{12}}\right|j_a j_b; J\right\rangle, \qquad (21) \\
\langle j_c(p) j_d(n); J|\hat{\vartheta}|j_a(p) j_b(n); J\rangle &= \langle j_c j_d; J|\hat{\vartheta}^{(0)} - \hat{\vartheta}^{(1)}|j_a j_b; J\rangle, \\
\langle j_c(p,n) j_d(n,p); J|\hat{\vartheta}|j_a(n,p) j_b(p,n); J\rangle &= \langle j_c j_d; J|2\hat{\vartheta}^{(1)}|j_a j_b; J\rangle.
\end{aligned}$$

From the explicit form of the matrix equation (15) one may easily obtain the orthonormalization relation

$$\left|\sum_{a\geq b} X^{n,J}_{j_a j_b} X^{m,J}_{j_a j_b} - \sum_{c\geq d} Y^{n,J}_{j_c j_d} Y^{m,J}_{j_c j_d}\right| = \delta_{mn}, \qquad (22)$$

which in terms of the QPRPA bosons corresponds to condition

$$\langle\tilde{0}|Q_{n,Jm} Q^{+}_{m,Jm}|\tilde{0}\rangle = \delta_{mn}. \qquad (23)$$

The system (15) is applicable for description of systems with any even numbers of protons and neutrons describing the nucleus by the method in which the nucleon numbers Z and N are conserved only on the average. In the case of magic, $Z = Z_0$ and $N = N_0$, or semimagic nuclei, when the pairing vanishes to zero, the system (15) splits into several unlinked subsystems each describing the (Z_0, N_0), ($Z_0 \pm 2$, N_0), (Z_0, $N_0 \pm 2$) nuclei. One should mention that the least two solutions, as conserving the numbers of particles exactly, may be in some sense more adequate than those obtained by using the equation (15) with $Z = Z_0 \pm 2$ or $N = N_0 \pm 2$.

2. PROBABILITIES OF ELECTROMAGNETIC TRANSITIONS

We shall consider below the electromagnetic transitions that are described by the single-particle operator

$$\hat{\mathcal{M}}(\lambda\mu) = \sum_{i,k}\langle i|\hat{m}(\lambda\mu)|k\rangle a^{+}_i a_k. \qquad (24)$$

Here we must distinguish between two different cases, i.e., between the phonon–phonon (between the two excited states) and the phonon–ground state transitions. The latter transition is described by the matrix element

$$\langle\tilde{0}|\hat{\mathcal{M}}(\lambda\mu)|\omega_n, JM\rangle = \left\langle\tilde{0}\left|\left[\hat{\mathcal{M}}(\lambda\mu), Q^{+}_{n,JM}\right]\right|\tilde{0}\right\rangle. \qquad (25)$$

Taking the $\hat{\mathcal{M}}$ in the form (24), using the relations (2), (11) and reducing the results of commutation in (25) to the normal form, one can obtain the final result which looks as follows:

$$\langle\tilde{0}\|\mathcal{M}(\lambda)\|\omega_n,J\rangle=(-1)^{\lambda}\delta(J,\lambda)\delta(\pi_n\pi_\lambda)\times$$
$$\times\left[\sum_{j_a\geq j_b}X^{n,J}_{j_aj_b}(u_{j_a}v_{j_b}\pm v_{j_a}u_{j_b})\frac{(-1)^{l_b}}{\sqrt{1+\delta_{j_aj_b}}}\langle j_a\|\hat{m}(\lambda)\|j_b\rangle-\right.$$
$$\left.-\sum_{j_a\geq j_b}Y^{n,J}_{j_aj_b}(v_{j_a}u_{j_b}\pm u_{j_a}v_{j_b})\frac{(-1)^{l_b}}{\sqrt{1+\delta_{j_aj_b}}}\langle j_a\|\hat{m}(\lambda)\|j_b\rangle\right], \tag{26}$$

where the upper signs refer to T-even ($E\lambda$) while the lower ones to T-odd ($M\lambda$) transitions. We notice here that we define the reduced matrix elements according to relation

$$\langle J'M'|T_{k\kappa}|JM\rangle=(-1)^{J'-M'}\begin{pmatrix}J' & k & J\\ -M' & \kappa & M\end{pmatrix}\langle J'\|T(k)\|J\rangle. \tag{27}$$

One can easily see that the «phonon–phonon» matrix element has the form

$$\langle\omega_n,J'M'|\hat{\mathcal{M}}(\lambda\mu)|\omega_m,JM\rangle=\langle\omega_n,J'M'|[\hat{\mathcal{M}}(\lambda\mu),Q^{+}_{m,JM}]|\tilde{0}\rangle+\delta_{mn}\delta_{\lambda0}\sum_i\hat{m}_{ii}v_i^2. \tag{28}$$

Remaining in (28) only the terms corresponding to the QPRPA space and taking into consideration the case of $\lambda\neq0$, we obtain the result

$$\langle\omega_n,J'\|\hat{\mathcal{M}}(\lambda)\|\omega_m,J\rangle=[(2J+1)(2J'+1)]^{1/2}\sum_{j_a\geq j_b,j_c\geq j_d}\frac{\left[X^{m,J}_{j_aj_b}X^{n,J'}_{j_cj_d}\pm Y^{m,J}_{j_aj_b}Y^{n,J'}_{j_cj_d}\right]}{\sqrt{(1+\delta_{j_aj_b})(1+\delta_{j_cj_d})}}\times$$
$$\times\{\delta_{j_bj_d}W[\lambda j_cJj_b;j_aJ'](u_{j_c}u_{j_a}\mp v_{j_c}v_{j_a})\langle j_c\|\hat{m}(\lambda)\|j_a\rangle-$$
$$-(-1)^{j_c+j_d+J'}\delta_{j_bj_c}W[\lambda j_dJj_b;j_aJ'](u_{j_d}u_{j_a}\mp v_{j_d}v_{j_a})\langle j_d\|\hat{m}(\lambda)\|j_a\rangle-$$
$$-(-1)^{j_a+j_b+J}\delta_{j_aj_d}W[\lambda j_cJj_a;j_bJ'](u_{j_c}u_{j_b}\mp v_{j_c}v_{j_b})\langle j_c\|\hat{m}(\lambda)\|j_b\rangle+$$
$$+(-1)^{j_a+j_b+J+j_c+j_d+J'}\delta_{j_aj_c}W[\lambda j_dJj_a;j_bJ'](u_{j_d}u_{j_b}\mp v_{j_d}v_{j_b})\langle j_d\|\hat{m}(\lambda)\|j_b\rangle\Big\}, \tag{29}$$

where the upper signs refer to $E\lambda$; while lower ones, to $M\lambda$ transitions.

The QPRPA method for excited states takes into account the correlations with the quasi-particle numbers equal to 2, 6, 10 and so on and does not include the correlations with the quasi-particle numbers 4, 8, 12 ... . It means that considering the transitions between the one-phonon states we omit the two-quasi-particle admixtures which are the most responsible for the renormalization of single-particle vertices. One may compensate this shortcoming of the theory by introducing the effective charges (i.e., by the renormalization of single-particle operators $\hat{m}(\lambda\mu)$, entering equation (29)). As to transitions between the excited and the ground states, they proceed with the variation of the quasi-particles number by two and in principle may be described without introducing the effective charge. However, the words «in principle» mean that really we must involve into calculations all the single-particle states, including that of continuum. As we really do not make this in our calculations using the limited basis, the introduction of the effective charge seems also quite reasonable also in this case, especially for positive parity transitions.

3. PARAMETERS OF THE THEORY

As the nuclei under consideration have the magical neutron number $N = 82$ and the values of Z in the interval between the magical (or almost magical at $N = 82$) proton numbers 50 and 64, it is reasonable to carry out the analysis of their properties in the framework of single particle scheme generated by spherical potential. As a such we used the potential of Woods–Saxon type with the parameters defined by us earlier from the best description of single-particle spectrum in the vicinity of ^{132}Sn. This potential has the form

$$U(\mathbf{r},\boldsymbol{\sigma}) = \frac{U}{1+\exp\left[\dfrac{r-R}{a}\right]} + U_{ls} r_0^2 \frac{1}{r}\frac{d}{dr}\left[\frac{1}{1+\exp\left[\dfrac{r-R}{a}\right]}\right] \mathbf{ls}. \tag{30}$$

In formula (30) $U = -U_0\left(1-\beta\dfrac{N-Z}{A}t_z\right)$, $R = r_0 A^{1/3}$, $U_{ls} = \alpha U$, $t_z = 1/2$ for neutrons and $t_z = -1/2$ for protons. In the case of protons we added to (30) the potential of a uniformly charged sphere with $R_c = r_c A^{1/3}$. The parameters of potential (Set1, corresponding to our works [25–27]) were: $U_0 = 51.5$ MeV, $r_0 = 1.27$ fm, $r_c = 1.25$ fm, $\alpha = -0.39$, $\beta = 1.39$, $a = 0.55$ fm for neutrons and $a = 0.67$ fm for protons. For checking the sensitivity of our results to the values of single particle levels we repeated the calculations using some different sets of the mean field parameters: $U_0 = 51.5$ MeV, $r_0 = 1.26$ fm, $r_c = 1.20$ fm, $\alpha = -0.43$, $\beta = 1.39$, $a = 0.60$ fm both for protons and neutrons (Set2). We also used Set3 of parameters which is the same as the Set1, except the α value which in this case is equal to -0.33.

In our previous works [22–27] we defined the finite range effective interaction in nuclei near ^{132}Sn and ^{208}Pb which is suitable for description of particle–particle, particle–hole, even–even and odd–odd nuclei adjacent to the magic core. Here we deal with nuclei whose properties are mainly defined by the interaction of valence protons above the $Z = 50$ shell. The interaction between protons in such system was defined in [13, 14] as having the form

$$\hat{\vartheta} = V_0(P_s + tP_t)\exp(-\alpha r_{12}^2), \tag{31}$$

with $V_0 = -33.2$ MeV, $t = 0.2$, $\alpha = 0.325$ fm^{-2}, P_s and P_t are singlet and triplet spin projectors. As we really include in the consideration not only protons, but also neutron excitations, we must introduce a more general force. Instead of dealing with our previous complicated interaction, we used here a more simple one, which for the particle–particle $T = 1$ channel coincides with that of the works [13, 14], resembling also the interaction of our works [22–27] in other channels. As a such we selected the interaction

$$\hat{\vartheta} = (V + V_{\tau\sigma}\boldsymbol{\sigma}_1\boldsymbol{\sigma}_2\boldsymbol{\tau}_1\boldsymbol{\tau}_2)\exp(-r_{12}^2/r_{00}^2) + \frac{e^2}{r_{12}}\left(\frac{1}{2}-\hat{t}_{z_1}\right)\left(\frac{1}{2}-\hat{t}_{z_2}\right) \tag{32}$$

with $V = -13.3$ MeV, $V_{\tau\sigma} = 6.65$ MeV, $r_{00} = 1.75$ fm, which we used in our calculations. We notice here that the interaction (32) was used by us both in the QPRPA equations (15) and in the pairing equations (8)–(10), giving automatically zero energies of the ghost 0^+ states in the case of developed pairing.

In the work [33] we obtained a good agreement with the experiment in description of electromagnetic properties of the ^{132}Sb nuclei using the values of effective charges for $E2$ transitions equal to $e_{\rm eff}(p) = 1.6|e|$ and $e_{\rm eff}(n) = 0.9|e|$, with effective $M1$ electromagnetic operator in the form

$$m(M1)_\mu = \sqrt{\frac{3}{4\pi}}\mu_N \left\{ g^l_{\rm eff}\hat{\mathbf{l}} + g^s_{\rm eff}\hat{\mathbf{s}} + g_2\hat{t}_z r^2 [Y_2 \otimes \sigma]^1 \right\}_\mu , \tag{33}$$

where $g^l_{\rm eff}(p) = 1.10$, $g^s_{\rm eff}(p) = 3.796$, $g^l_{\rm eff}(n) = -0.005$, $g^s_{\rm eff}(n) = -2.038$, $g_2 = -0.031$ fm^{-2}. The above values of parameters also well describe the $E2$ and $M1$ characteristics of nuclei in the vicinity of ^{208}Pb [23]. In our work [8] we specially defined the values of proton effective charges with respect to $E2$ and $E3$ transitions in the ^{134}Te nuclei, which both were found to be $\sim 1.9|e|$. Here we performed all the calculations by using the values of proton effective charges relatively to $E2$, $E3$ and $E4$ transitions all equal to $1.9|e|$, with the neutron charge equal to $0.9|e|$. As to the $M1$ electromagnetic operator, it was taken in the form (33) with the parameters presented above.

4. RESULTS OF CALCULATIONS

The properties of nuclei considered by us are mainly defined by the valence protons that successively fill the 50–64 (sub)shell. However, for an adequate description of some collective levels, for example, the 3^-_1 state, as well as for a description of nuclei with filled shells, it is necessary to include in the basis more remote proton orbitals as well as the most important neutron single-particle states that form particle–hole neutron configurations with the excitation of neutrons from the filled shell with $N = 82$. Therefore our basis set of single particle states consisted of 11 proton orbitals, namely the 1f7/2, 1f5/2, 2p3/2, 2p1/2, 1g9/2, 1g7/2, 2d5/2, 2d3/2, 1h11/2, 3s1/2 and 2f7/2 levels, and also included 2d3/2, 3s1/2, 1h11/2, 2f7/2, 1h9/2 neutron states that are most close to the neutron $N = 82$ gap. Everywhere we used theoretical values of single-particle energies that corresponded to potential (30).

In Figs. 1–3 one can see the systematics of experimental and calculated energies of 2^+_1, 2^+_2, 0^+_2 and 3^-_1 levels for the variants 1–3 of the mean field parameters. The energies of all levels in $A(Z, N)$ nuclei with a nonzero pairing are defined by the system (15) with the given values of Z and N. The situation becomes more delicate near the closed shells. So, for example for ^{132}Sn nuclei, the energies are defined by the particle–hole branch of solutions of the mentioned system by Z, N = 50, 82. As to the ^{134}Te nuclei, for the description of all its levels except the 3^-_1 state, we used the solution that exactly conserves the number of particles (i.e., proton particle–particle branch of solutions of the system (15) when Z, N = 50, 82). This branch corresponds to the particle–particle RPA method [25] with Z, N = 52, 82. As the mentioned branch does not contain the particle–hole states, for description of the lower 3^-_1 state having large particle–hole admixtures, and only for it, we used the system (15) with Z, $N = 52$, 82 (the next 3^- state has a larger energy and is mainly defined by the particle–particle channel).

One can see from Fig. 1 that all the mentioned variants of the mean field parameterization correctly reproduce the experimental pattern for the 2^+_1 and 2^+_2 levels, which structure essentially alters with the increase of Z. So, in ^{132}Sn the 2^+_1 and 2^+_2 levels have the neutron

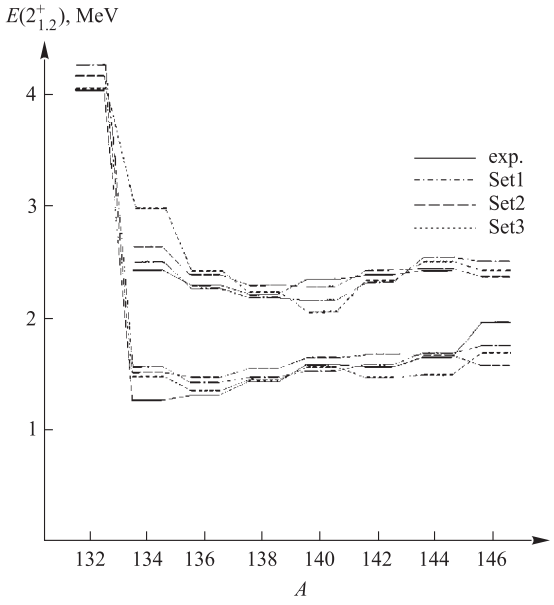

Fig. 1. Systematics of the 2^+_1 and 2^+_2 levels in the $N = 82$ isotones

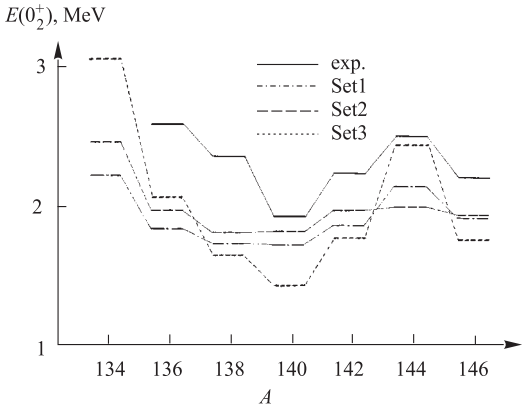

Fig. 2. Systematics of 0^+_2 levels in the $N = 82$ isotones

particle–hole structure with the main configurations correspondingly of the (2f7/2, 1h11/2) and (1h9/2, 1h11/2) types. At $Z \geq 52$, the mentioned levels have mainly the proton two-quasi-particle nature and with the increase of Z from 52 up to 62 the structure of the 2^+_1 level changes from the configuration $(1g7/2)^2$ to $(2d5/2)^2$. At the same time the main component of the 2^+_2 level, presented in ^{134}Te by the proton (1g7/2, 2d5/2) configuration, changes in ^{140}Ce and becomes the mixture of (1g7/2, 2d5/2) and $(2d5/2)^2$ configurations. In ^{146}Gd both the mentioned states have the structure that corresponds to the «crossing» by protons of the $Z = 64$ gap.

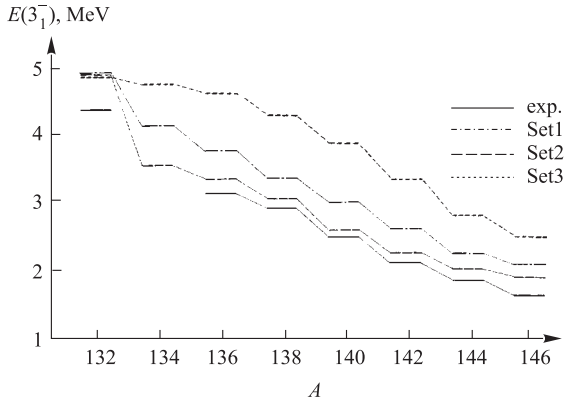

Fig. 3. Systematics of 3_1^- levels in the $N = 82$ isotones

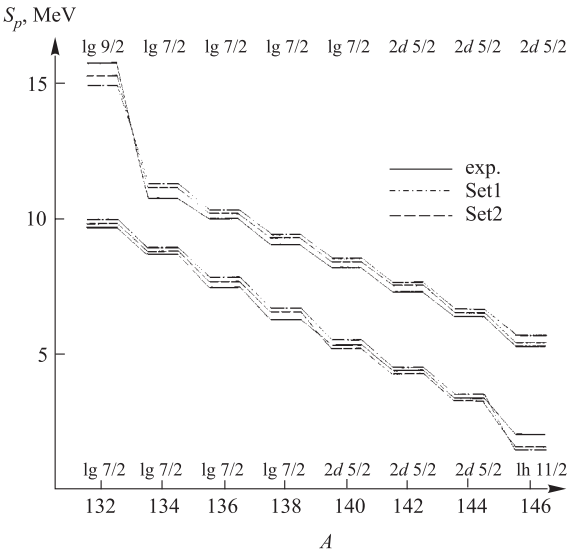

Fig. 4. Proton separate energies S_p from the even–even $Z, N = (A - 82), 82$ (upper curve) and from the proton-odd with $Z, N = (A - 81), 82$ (lower curve) nuclei. Single-particle prescriptions correspond to the remainder proton-odd (upper line) or to the initial proton-odd (lower line) nuclides

As to the pairing–vibrational 0_2^+ states, one can see from Fig. 2 that their energies are systematically underestimated as compared to the experiment, though the general energy systematics is reproduced quite satisfactory in all cases.

One can see another pattern in the case of 3_1^- states. Though the global tendency for the variation of energy levels is reproduced here for all variants, «Set1» and «Set2» are found to be more preferable. In all cases the wave function of the 3_1^- level in ^{132}Sn is a collective particle–hole mixture. As the number of protons increases, the main configuration of this level successively occurs to be of the (1g7/2, 1h11/2) and (2d5/2, 1h11/2) types.

Tables 1–8 present experimental and theoretical energy levels for all nuclei considered by us. For each given J^π one can find there all the calculated levels with energies less than the maximal, presented in the Tables for given J^π. For large excitation energies the correspondence of experimental and theoretical energies may be rather arbitrary as not all levels may be known from the experiment by the present time. On the whole, the coincidence with the experiment in the energies of positive parity levels proves to be the best in case of the «Set2» variant remaining quite acceptable for the «Set1» variant of the mean field parameters, that optimizes single-particle spectra in the vicinity of ^{132}Sn. The negative parity levels are better described by the mean field parameters of the variant «Set1» and a little worse — by those of the «Set2» variant. From the analysis of Tables 1–8 one can also see that there is a nonsatisfactory agreement with the experiment in the case of «Set3» variant.

The global properties of both the mean field parameters and the residual interaction used in calculations may be checked in description of the one-nucleon separation energies. If Z, N are even numbers, then in the framework of theory that takes into account pairing correlations of superconducting type we have the formula describing the (one-proton for example) separation energy from the $A(Z, N)$ nuclei:

$$S_p^j(Z, N) = -\lambda_p(Z, N) + E_j^p(Z, N), \tag{34}$$

where the index «j» refers to the final odd nuclei. At the same time the (proton) separation energy from the proton-odd $(Z + 1, N)$ nuclei that stays in the state «j» is given by the expression

$$S_p^{j'}(Z + 1, N) = -\lambda_p(Z, N) - E_{j'}^p(Z, N). \tag{35}$$

One should note that in the absence of pairing correlations $E_j \to (\lambda - \varepsilon_j)$ and $E_{j'} \to (\varepsilon_{j'} - \lambda)$ (here j refers to holes; while j', to particles), and thus the relations (34) and (35) turn into those presenting the Coopman theorem in the case of the Hartree–Fock procedure.

Comparison of experimental and calculated proton separation energies is presented in Fig. 4. One can see that with the differences in the calculated and experimental separation energies at most of about one MeV (usually less) the calculations correctly reproduce both the values of S_p and the global tendencies of their systematics, including the peculiarities near the closed shells.

In Table 9 one can see the results of our calculations for the reduced transition probabilities and nuclear moments as well as comparison with the experiment.

Among the transitions entering Table 9 there exist those proceeding between the excited and ground states. The latter $E\lambda$ transitions are described by the formula (26) and are rather weakly dependent on the single-particle spectra. Such a dependence is mainly due only to the X, Y amplitudes entering Eq. (26), while the factors $(uv' + vu')$ are only slightly dependent on single particle energies. One can see that the mentioned transitions are rather well described by our calculations, what really confirms the approximate equality of the $E2$, $E3$ (and $E4$) effective charges in the region of ^{132}Sn. It is useful to mention that the energies of 2_1^+, 3_1^- and 4_1^+ states in ^{132}Sn are also close to each other (see Table 1). Such a «degeneracies» of proton effective charges and phonon energies for different $E\lambda$ vibrations in ^{132}Sn manifest about some additional symmetry of collective Hamiltonian of this nucleus.

Turning to description of transitions between the excited states (transitions of the type «phonon–phonon») we see that the Eλ transition rates are extremely dependent on the variations of single particle spectrum. Such sensitivity is due to the factor $(uu' - vv')$ in formula

Table 1. **The level energies of ^{132}Sn, MeV**

J^π	Exp.	Set1	Set2	Set3	J^π	Exp.	Set1	Set2	Set3
1+		6.184	6.702	5.418	1–		5.475	5.440	5.534
2+	4.04	4.274	4.163	4.045	1–		7.466	7.539	7.549
2+		5.310	5.816	4.579	2–		6.298	5.985	6.630
2+		5.712	6.160	5.203	2–		6.828	6.796	7.073
3+		5.071	4.741	4.377	3–	4.35	4.877	4.870	4.810
3+		5.127	5.632	4.830	3–		5.654	5.468	5.743
3+		5.492	6.157	5.062	3–		5.880	5.705	5.958
4+	4.42	4.746	4.482	4.194	4–	4.83	5.694	5.428	5.706
4+		4.990	5.463	4.546	4–	5.39	5.963	5.889	5.881
4+		5.335	5.941	4.841	4–		6.076	5.941	5.950
5+		4.919	4.692	4.168	5–		5.647	5.391	5.515
5+		5.038	5.421	4.758	5–		6.022	6.518	5.806
5+		5.467	6.132	4.846	5–		6.095	6.836	6.103
6+	4.72	4.792	4.551	4.076	6–	6.60	6.434	7.001	6.183
6+	5.40	4.930	5.334	4.627	6–	6.71	7.599	7.540	7.753
6+	6.17	5.314	5.943	4.733	6–	7.21	9.145	9.273	9.014
6+	6.63	5.564	6.129	5.146	7–	7.24	7.461	7.399	7.616
6+		7.803	7.858	7.235	7–		11.955	12.354	11.334
7+	4.92	4.826	4.709	4.074					
7+	5.63	5.049	5.328	4.686					
7+	6.23	5.517	6.052	4.826					
7+	6.90	5.695	6.352	5.482					
7+		10.19	9.544	10.639					
8+	4.85	4.729	4.594	3.974					
8+	5.48	4.958	5.250	4.706					
8+		5.527	6.059	4.724					
9+		5.272	4.933	4.641					
9+		5.475	6.062	5.050					
10+		5.531	6.072	4.709					

(29), the magnitude of which is strongly dependent on the structure of single particle spectrum near the Fermi level. It follows from Table 9 that with the increase of Z the experimental ratio of the $B(E2; 4_1^+ \to 2_1^+)$ to $B(E2; 6_1^+ \to 4_1^+)$ values changes in a rather irregular manner. One should also pay attention to the irregularity of the relative distance between the 6_1^+ and 4_1^+ levels with the increase of Z. The mentioned facts are due to variation of the structure of the lower 2_1^+, 4_1^+ and 6_1^+ levels. So, if in ^{134}Te and ^{136}Xe all the mentioned states have the structure of the $(1g7/2)^2$ type, in ^{138}Ba the 6_1^+ level presents the mixture of the $(1g7/2)^2$ and (1g7/2, 2d5/2) configurations, while in ^{144}Sm the mentioned 2_1^+, 4_1^+ and 6_1^+ levels have the structure of the $(2d5/2)^2$, $(2d5/2)^2$ and (1g7/2, 2d5/2) types. The amplitudes «X» and «Y» in formula (29) and especially the factor $(uu' - vv')$ are extremely dependent on the relative position of single-particle proton states in the 50–82 proton shell. Potential (30) with its global parameterizations gives only some averaged properties of single particle spectra not reproducing in detail the structure of single particle spectrum in each given nucleus. This

Table 2. **The level energies of ^{134}Te, MeV**

J^π	E	Set1	Set2	Set3	J^π	E	Set1	Set2	Set3
0+		2.216	2.460	3.052	5+		5.029	5.439	5.095
1+	2.600	2.557	2.678	2.996	6+	1.665	1.921	1.876	1.838
2+	1.265	1.562	1.515	1.480	6+	2.350	2.250	2.358	2.669
2+	2.435	2.516	2.641	2.996	3–		4.107	3.484	4.733
2+	2.900	2.825	3.117	4.643	4–		4.594	4.061	5.491
2+		4.517	4.895	5.894	5–		4.471	3.943	5.377
3+		2.591	2.712	3.030	6–		5.006	4.072	5.502
4+	1.550	1.838	1.787	1.744	7–	4.440	4.363	3.835	5.268
4+	2.520	2.484	2.605	2.918	8–		4.619	4.085	5.515
4+		3.005	3.305	3.970	9–	4.170	3.905	3.371	4.801
5+	2.700	2.594	2.688	3.033					

Table 3. **The level energies of ^{136}Xe, MeV**

Experiment		Theory				Experiment		Theory			
J^π	E	J^π	Set1	Set2	Set3	J^π	E	J^π	Set1	Set2	Set3
0+	2.582	0+	1.820	1.955	2.055			5+	4.508	4.900	4.408
1, 2(+)	3.211	1+	2.442	2.556	2.634	6+	1.892	6+	1.875	1.912	1.761
2+	1.313	2+	1.430	1.464	1.334	6+	2.262	6+	2.073	2.152	2.182
2+	2.289	2+	2.270	2.386	2.423			7+	5.284	6.033	4.985
2+	2.414	2+	2.379	2.555	2.776			8+	5.101	4.854	4.334
2+	2.634	2+	3.681	3.882	3.550			8+	5.207	5.018	4.778
3+,4+	2.125	3+	2.334	2.438	2.431	3–	3.275	3–	3.727	3.318	4.472
2, 3, 4(+)	3.873	3+	4.408	4.790	4.295			4–	4.154	3.598	4.761
4+	1.694	4+	1.796	1.816	1.655			5–	3.967	3.432	4.593
(4+)	2.465	4+	2.223	2.324	2.361			6–	4.115	3.559	4.709
4+	2.560	4+	2.511	2.715	2.972			7–	3.907	3.356	4.530
4+	2.608	4+	4.085	4.398	4.080			8–	4.099	3.544	4.687
5	2.444	5+	2.306	2.409	2.447			9–	3.620	3.062	4.268

structure was in particular defined from the experiment in [15] using the procedure of inverse gap equations and it was shown the isolation of a $Z = 64$ (sub)shell in the ^{146}Gd nuclide, with the energy gap between the groups of 1g7/2, 2d5/2 and 1h11/2, 3s1/2, 2d3/2 levels of about 3 MeV.

With the reservations presented above we can say that the coincidence with the experiment for $E2$ transitions between the one-phonon levels for the Set1 of single-particle spectrum, which parameters were defined by us from the description of «one quasi-particle above the core» nuclei near ^{132}Sn, seems to be quite acceptable. It remains the same for «Set2» variant, where we used the same diffuseness parameter $a = 0.60$ fm (see Eq. (30)) both for protons and neutrons. The largest discrepancies with the experiment are observed for the $4_1^+ \to 2_1^+$ and $6_1^+ \to 4_1^+$ transitions in ^{138}Ba, ^{140}Ce and ^{142}Nd where there happens the change of a valence 1g7/2 to 2d5/2 orbitals. It is worth of mentioning that though the «Set3» variant gives in

Table 4. **The level energies of ^{138}Ba, MeV**

Experiment		Theory				Experiment		Theory			
J^π	E	J^π	Set1	Set2	Set3	J^π	E	J^π	Set1	Set2	Set3
0+	2.340	0+	1.711	1.794	1.627	4+	3.155	4+	4.097	4.453	4.017
(0+)	3.612	0+	4.436	3.569	5.224	5+	2.416	5+	2.258	2.340	2.250
(1, 2, 3)	2.190	1+	2.444	2.544	2.520			5+	4.302	4.716	4.141
(1+, 2+)	2.583	1+	4.494	4.951	4.488	6+	2.091	6+	1.919	1.996	1.878
2+	1.436	2+	1.473	1.541	1.440	6+	2.203	6+	2.148	2.192	2.064
2+	2.218	2+	2.184	2.296	2.241			6+	5.225	4.415	4.639
2+	2.640	2+	2.297	2.392	2.365	(7+)	3.360	7+	6.025	6.019	5.170
2+	3.049					8+	3.184	8+	5.229	4.449	4.541
2+	3.339	2+	3.460	3.652	3.297	3–	2.881	3–	3.339	2.921	4.046
3+	2.446	3+	2.295	2.379	2.298	(4)–	3.561	4–	3.913	3.394	4.448
3+	2.991	3+	4.171	4.572	3.995	(5)–	3.857	5–	3.671	3.183	4.259
4+	1.899	4+	1.897	1.950	1.793			6–	3.853	3.333	4.365
4+	2.308	4+	2.204	2.280	2.188			7–	3.657	3.141	4.214
4+	2.583	4+	2.319	2.466	2.510	(8–)	3.678	8–	3.825	3.305	4.327
4+	2.779	4+	3.815	4.105	3.763	(9–)	3.633	9–	3.449	2.922	4.044

Table 5. **The level energies of ^{140}Ce, MeV**

Experiment		Theory				Experiment		Theory			
J^π	E	J^π	Set1	Set2	Set3	J^π	E	J^π	Set1	Set2	Set3
0+	1.903	0+	1.706	1.798	1.412			5+	4.116	4.563	3.928
0+	3.017	0+	3.772	3.055	4.371	6+	2.108	6+	1.981	2.095	1.939
1+	2.547	1+	2.493	2.619	2.495	6+	2.629	6+	2.350	2.443	2.318
2+	1.596	2+	1.531	1.644	1.576	6+	3.484	6+	4.872	3.885	5.117
2+	2.348	2+	2.172	2.281	2.073	7+	3.433	7+	6.419	5.175	5.262
2+	2.521	2+	2.342	2.450	2.269	8+	3.513	8+	4.906	3.918	4.549
2+	2.900	2+	3.256	3.373	3.086	8+	3.621	8+	5.319	5.063	5.152
2+	3.001					3–	2.464	3–	2.962	2.563	3.468
2+	3.118					3–	3.040	3–	3.724	3.277	4.201
3+	2.412	3+	2.329	2.437	2.244	3–	3.473	3–	5.266	5.271	5.245
		3+	3.957	4.395	3.754	(4–)	3.395	4–	3.597	3.177	4.076
4+	2.083	4+	1.988	2.111	1.958	5–	3.256	5–	3.345	2.925	3.837
4+	2.481	4+	2.229	2.353	2.150			6–	3.616	3.183	4.071
3+, 4+	2.515	4+	2.313	2.410	2.258	7–	3.425	7–	3.415	2.973	3.911
4+	3.331	4+	3.533	3.707	3.365	8–	3.477	8–	3.590	3.149	4.022
(4+)	3.395					9–	3.493	9–	3.305	2.837	3.856
5+	2.350	5+	2.296	2.394	2.191						

some cases better description with the experiment for $E2$ transitions, the overall description of energy levels given by this variant seems to be not acceptable.

Table 6. **The level energies of ^{142}Nd, MeV**

Experiment		Theory				Experiment		Theory			
J^π	E	J^π	Set1	Set2	Set3	J^π	E	J^π	Set1	Set2	Set3
0+	2.217	0+	1.852	1.939	1.739			5+	3.964	4.465	3.832
0+	2.978	0+	3.057	2.612	3.433	6+	2.210	6+	2.126	2.268	2.124
1(+)	2.583	1+	2.623	2.782	2.650	6+	2.887	6+	2.695	2.854	2.922
2+	1.576	2+	1.584	1.680	1.457			6+	4.198	3.351	4.601
2+	2.385	2+	2.320	2.427	2.334	(7+)	3.520	7+	5.395	5.975	6.190
2+	2.846	2+	2.522	2.658	2.608	8+	3.454	8+	4.235	3.358	5.015
(2)+	3.046	2+	3.062	3.030	2.938			8+	5.330	5.153	5.232
(1,2+)	3.128					3–	2.085	3–	2.579	2.220	2.882
2+	3.487					4–	2.960	4–	3.210	2.848	3.481
3+	2.548	3+	2.475	2.611	2.437	(4)–	3.244	4–	3.541	3.189	4.027
2+, 3+	3.300	3+	3.690	4.185	3.347	(4)–	3.296	4–	4.754	4.741	4.905
4+	2.101	4+	2.011	2.142	1.783	(5–)	2.976	5–	2.973	2.617	3.261
4+	2.438	4+	2.329	2.459	2.298			6–	3.237	2.874	3.513
4	2.738	4+	2.616	2.758	2.806	7–	3.243	7–	3.091	2.732	3.380
		4+	3.247	3.222	2.940	8–	3.457	8–	3.302	2.925	3.575
5+	2.514	5+	2.440	2.571	2.393	9–	3.486	9–	3.177	2.786	3.839

Table 7. **The level energies of ^{144}Sm, MeV**

Experiment		Theory				Experiment		Theory			
J^π	E	J^π	Set1	Set2	Set3	J^π	E	J^π	Set1	Set2	Set3
0+	2.478	0+	2.116	1.962	2.409	6+	2.323	6+	2.430	2.521	2.532
0+	2.827	0+	2.436	2.633	2.493	(6+)	2.729	6+	3.220	2.864	3.710
		1+	2.902	3.044	3.009	6+	3.308	6+	3.457	3.367	4.018
2+	1.660	2+	1.686	1.680	1.487			7+	5.366	5.342	6.606
2+	2.423	2+	2.542	2.440	2.509			8+	3.489	2.874	4.057
2+	2.799	2+	2.846	2.882	2.883	3–	1.810	3–	2.201	1.977	2.394
2+	3.318	2+	3.029	3.158	2.936	3–	3.228	3–	3.472	3.212	3.964
		3+	2.796	2.903	2.879	(4–)	3.205	4–	2.848	2.627	2.991
		3+	3.347	3.902	2.912			4–	3.455	3.199	3.936
4+	2.191	4+	2.108	2.199	1.815	5–	2.825	5–	2.609	2.393	2.762
4+	2.588	4+	2.621	2.619	2.523			6–	2.870	2.647	3.016
4+	2.884	4+	2.954	2.849	2.786	7–	3.124	7–	2.734	2.521	2.880
		4+	3.160	3.288	3.465	8–	3.377	8–	2.957	2.709	3.110
		5+	2.770	2.871	2.854			9–	3.079	2.797	3.666
		5+	3.864	4.419	3.808						

Table 9 presents also some data about the $M1$ properties of nuclei in question. One can see that the theory correctly reproduces the global characteristics of this type. The situation is different for the quadrupole moments another time manifesting on the sensitivity of this characteristic to the peculiarities of single particle spectrum.

Table 8. **The level energies of ^{146}Gd, MeV**

Experiment		Theory				Experiment		Theory			
J^π	E	J^π	Set1	Set2	Set3	J^π	E	J^π	Set1	Set2	Set3
0+	2.165	0+	1.885	1.902	1.728	6+	3.457	6+	2.714	2.436	2.972
0+	3.020	0+	2.843	3.054	2.568	6+	3.486	6+	3.012	2.925	3.301
0+	3.485	0+	3.149	4.090	2.882	6+	3.660	6+	3.973	3.891	4.549
		1+	3.431	3.366	3.029			7+	5.328	5.411	7.353
2+	1.972	2+	1.759	1.592	1.715	8+	3.779	8+	2.782	2.513	3.024
		2+	2.513	2.385	2.459	8+	4.107	8+	5.281	5.298	5.365
		2+	2.886	3.144	2.717	3–	1.579	3–	2.038	1.853	2.095
		2+	3.174	3.377	2.786	4–	2.997	4–	2.681	2.524	2.696
3+	3.031	3+	3.150	3.257	2.682	5–	2.658	5–	2.416	2.281	2.432
(3)	3.389	3+	3.367	3.698	3.108	6–	3.098	6–	2.699	2.541	2.719
(3)	3.436	3+	3.408	3.935	3.653	7–	2.982	7–	2.551	2.418	2.539
4+	2.612	4+	2.346	2.206	2.126	7–	3.290	7–	3.030	3.053	2.812
4+ (2+)	2.967	4+	2.665	2.488	2.582	8–	3.183	8–	2.797	2.607	2.837
4(+)	3.412	4+	2.942	3.117	2.905	8–	3.293	8–	3.288	3.168	3.748
4(+)	3.411	4+	3.244	3.380	3.461	9–	3.428	9–	3.097	2.863	3.660
(3, 5)+	3.287	5+	3.354	3.231	3.662						

CONCLUSION

The calculations presented above manifest that the bulk of experimental data on the properties of $N = 82$ even–even isotones may be rather successfully explained in the framework of the two-quasi-particle RPA (QPRPA) method that treats the pairing force and residual interaction in other channels on equal grounds. In our calculations we used the minimal amount of free parameters. The «Set1» variant of them that practically borrows their values from our previous calculations and the «Set2» variant, employing some fitting of single-particle potential with the same diffusenesses for protons and neutrons, give practically the same quality of description of experimental data. The sensitivity of the calculation scheme to the input parameters is demonstrated on the example of the «Set3» parameters, where the 15 % decrease of the spin-orbit splitting leads to a better description of transition rates but to unsatisfactory agreement in the energies of levels. The divergence between the predictions of theory and the experimental data in each specific nucleus may be often diminished by the variations of single particle energies close to the Fermi levels, i.e., by considering the mentioned energies as free parameters in each nuclide. In our calculations we used a more general approach in which the mentioned energies were generated by a mean field potential which parameters are the same (except the Z, N values) for all nuclides considered by us. The other way to improve the agreement with the experiment may consist in using the method in which the particle numbers conserve not on the average, but exactly. This will be made by us in the future.

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Table 9. **Probabilities of the** $E2$**,** $E3$ **and** $M1$ **transitions and the values of electrical quadrupole and magnetic dipole moments for the even–even** $N = 82$ **nuclei. The** $B(E, M\lambda)$ **values are given in Weisskopf units:** $B(E\lambda)_W = \dfrac{9e^2R^{2\lambda}}{4\pi(\lambda+3)^2}$ **and** $B(M\lambda)_W = \dfrac{10}{\pi}\left(\dfrac{3}{\lambda+3}\right)^2 R^{2\lambda-2}\mu_N^2$ **with** $R = 1.2A^{1/3}$ **fm. The values of the quadrupole moments are presented in the units of e.fm^2, whereas the magnetic moments are in the units of** μ_N

Nucleus	Value	Exp.	Set1	Set2	Set3
^{132}Sn	$B(E2, 4_1^+ \to 2_1^+)$	0.426	0.413	0.347	0.075
	$B(E2, 6_1^+ \to 4_1^+)$	0.263	0.283	0.389	0.118
	$B(E2, 8_1^+ \to 6_1^+)$	0.110	0.076	0.180	0.032
	$B(E2, 5_1^- \to 3_1^-)$	0.275	0.749	0.654	0.660
	$B(M1, 4_1^- \to 3_1^-)$	0.0078	0.0058	10^{-6}	7.10^{-5}
	$B(M1, 7_1^+ \to 8_1^+)$	0.037	0.101	0.0074	0.106
	$B(M1, 7_1^+ \to 6_1^+)$	0.059	0.123	0.0091	0.189
	$B(M1, 5_1^- \to 4_1^-)$	0.0307	0.038	0.068	0.183
^{134}Te	$B(E2, 2_1^+ \to 0_1^+)$	–	4.52	4.39	4.62
	$B(E4, 4_1^+ \to 0_1^+)$	–	2.81	2.65	2.78
	$B(E2, 4_1^+ \to 2_1^+)$	4.3	4.38	4.32	4.33
	$B(E2, 6_1^+ \to 4_1^+)$	2.05	1.95	1.94	1.96
	$B(E2, 2_1^+ \to 0_2^+)$	–	0.24	0.22	0.13
	$B(E2, 6_2^+ \to 4_2^+)$	–	1.66	1.54	1.54
	$B(E2, 6_2^+ \to 6_1^+)$	–	0.53	0.415	0.31
	$B(E3, 3_1^- \to 0_1^+)$	–	13.6	6.8	45.2
^{136}Xe	$B(E2, 2_1^+ \to 0_1^+)$	9 ± 4	14.88	13.09	13.08
	$B(E4, 4_1^+ \to 0_1^+)$	–	8.79	6.92	7.13
	$B(E2, 4_1^+ \to 2_1^+)$	1.25 ± 0.06	0.473	0.340	0.101
	$B(E2, 6_1^+ \to 4_1^+)$	0.0132(5)	0.228	0.147	0.037
	$B(E2, 2_1^+ \to 0_2^+)$	–	0.054	0.045	0.069
	$B(E2, 6_2^+ \to 4_2^+)$	>0.28	0.444	0.508	0.594
	$B(E2, 6_2^+ \to 6_1^+)$	>0.0003	0.320	0.324	0.238
	$B(E3, 3_1^- \to 0_1^+)$	16.91	16.99	12.34	22.9
^{138}Ba	$B(E2, 2_1^+ \to 0_1^+)$	11.4 ± 0.3	19.9	17.7	15.6
	$B(E4, 4_1^+ \to 0_1^+)$	–	13.2	10.3	9.1
	$B(E2, 4_1^+ \to 2_1^+)$	0.286(11)	0.077	0.017	0.106
	$B(E2, 6_1^+ \to 4_1^+)$	0.053(7)	0.076	0.029	0.028
	$B(E2, 2_1^+ \to 0_2^+)$	–	0.331	0.337	0.54
	$B(E2, 4_2^+ \to 2_1^+)$	>0.2	0.217	0.215	0.286
	$B(E2, 2_2^+ \to 0_1^+)$	1.98(24)	0.46	0.92	2.05
	$B(E3, 3_1^- \to 0_1^+)$	16.8	19.1	15.1	18.0
	$B(M1, 3_1^+ \to 2_3^+)$	0.135	0.354	0.325	0.19
	$Q(2_1^+)$	–14(7)	–5.27	–4.24	–8.19
	$\mu(4_1^+)$	+3.2(6)	3.71	3.52	3.37
	$\mu(6_1^+)$	+5.86	6.40	6.13	5.46

Ending of Table 9

Nucleus	Value	Exp.	Set1	Set2	Set3
^{140}Ce	$B(E2, 2_1^+ \to 0_1^+)$	16.6 ± 2.4	22.7	21.3	17.7
	$B(E4, 4_1^+ \to 0_1^+)$	(12.1)	17.0	14.8	14.3
	$B(E2, 4_1^+ \to 2_1^+)$	0.137(1)	0.062	0.075	0.159
	$B(E2, 6_1^+ \to 4_1^+)$	0.28(6)	0.025	0.033	0.025
	$B(E2, 2_1^+ \to 0_2^+)$	2.3(2)	0.54	0.64	1.51
	$B(E2, 6_2^+ \to 4_2^+)$	–	0.020	0.022	0.048
	$B(E2, 6_2^+ \to 6_1^+)$	–	0.024	0.036	0.056
	$B(E2, 2_4^+ \to 0_1^+)$	$\sum_4^6 = 5.2$	7.75	6.36	8.40
	$B(E3, 3_1^- \to 0_1^+)$	25.7	21.8	18.3	16.2
	$\mu(4_1^+)$	+4.35	4.72	4.59	4.89
^{142}Nd	$B(E2, 2_1^+ \to 0_1^+)$	12.04	22.2	23.3	15.5
	$B(E4, 4_1^+ \to 0_1^+)$	13 ± 4	13.8	12.5	9.7
	$B(E2, 4_1^+ \to 2_1^+)$	0.50	0.039	0.174	0.26
	$B(E2, 6_1^+ \to 4_1^+)$	0.018	0.0022	0.0015	0.0084
	$B(E2, 2_1^+ \to 0_2^+)$	–	0.26	0.20	0.37
	$B(E2, 6_2^+ \to 4_2^+)$	–	0.010	0.013	0.0017
	$B(E2, 6_2^+ \to 6_1^+)$	–	0.018	0.017	0.036
	$B(E2, 2_2^+ \to 0_1^+)$	1.0	2.23	2.53	4.3
	$B(E3, 3_1^- \to 0_1^+)$	28.4	25.2	22.5	19.7
^{144}Sm	$B(E2, 2_1^+ \to 0_1^+)$	11.9	22.7	27.5	16.3
	$B(E4, 4_1^+ \to 0_1^+)$	–	13.4	14.1	11.1
	$B(E2, 4_1^+ \to 2_1^+)$	–	0.003	0.12	0.0024
	$B(E2, 6_1^+ \to 4_1^+)$	0.19	0.042	0.0019	0.088
	$B(E2, 2_1^+ \to 0_2^+)$	–	0.17	0.11	0.15
	$B(E2, 8_1^- \to 7_1^-)$	≤ 3.6	1.73	0.67	1.80
	$B(E3, 3_1^- \to 0_1^+)$	31.3	29.2	26.5	25.7
^{146}Gd	$B(E2, 2_1^+ \to 0_1^+)$	>0.4	29.8	33.0	23.8
	$B(E4, 4_1^+ \to 0_1^+)$	–	25.9	9.9	22.0
	$B(E2, 4_1^+ \to 2_1^+)$	–	0.49	0.76	0.21
	$B(E2, 6_1^+ \to 4_1^+)$	–	0.60	0.94	0.076
	$B(E2, 2_1^+ \to 0_2^+)$	–	1.58	0.53	1.66
	$B(E2, 7_1^- \to 5_1^-)$	0.46	0.024	0.0026	2.10^{-5}
	$B(E3, 3_1^- \to 0_1^+)$	37(4)	32.4	29.0	31.1
	$\mu(3_1^-)$	+2.1(9)	3.34	3.40	3.32
		1.86(36)			
	$\mu(7_1^-)$	+ 8.98(19)	9.56	9.78	9.09
		+ 7.56(56)			
		+ 8.26(35)			

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