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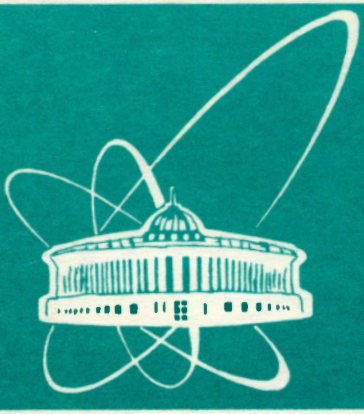

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EFFECT OF GROUND STATE CORRELATIONS ON THE CHARGE TRANSITION DENSITIES OF VIBRATIONAL STATES

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1. Introduction

It is well known that many basic features of the nuclear vibrational states can be described within the Random Phase Approximation (RPA), which enables one to treat some correlations in the ground state.

The ground state correlations (GSC) problem has a long story. This problem was discussed in refs. [1–11]. There is a rather complete list of references on that subject in paper[10]. Higher (renormalized) RPA equations, which include corrections for GSC, have been advanced by several authors [1–11]. To apply these methods one needs the one–particle densities of the correlated ground state. In our calculations we use the approach suggested in [1,2]. The generalization of this approach to take into account the effect of GSC on the pairing has been done in ref.[6]. An alternative approach closely related to the one from refs. [1,2] and differing in the way of evaluating the one-particle densities of the correlated state has been proposed in ref.[11].

Being the spatial overlap between the ground state wave function and the excited state wave function the charge transition density provides a good test for nuclear models. The surface nature of the low-lying collective states has been demonstrated in the experiments on inelastic electron scattering from magic nuclei [12].Such a behavior was predicted by calculations performed within the Hartree-Fock (HF) approach with effective forces [13] and the finite Fermi systems theory [14]. The changes of the nuclear densities due to the zero point fluctuations associated with surface modes in the Ca isotopes were calculated within the nuclear field theory in ref. [15],where the quasiparticle distribution in the ground state was calculated within the RPA. Recent experimental and theoretical (based on the random phase approximation (RPA)) studies of the charge transition densities [16,17] to investigate the interplay between single–particle and collective degrees of freedom in the excitation of the low-lying states in some spherical nuclei are in reasonable agreement, but the theory gives fluctuations of the transition densities in the interior region. In RPA, as in HF, the theoretical fluctuations are too large in the nuclear interior, which indicates a systematic problem of a more fundamental nature (a detailed discussion can be found in refs. [18,19]).

This long standing problem is not solved up to now. In a recent paper [20] we pointed out that an inclusion of the GSC beyond the RPA gives a 30% depletion of the charge transition densities of the first quadrupole state in the interior region of ^{64}Zn. In this case the blocking effect due to the Pauli principle plays an essential role. Nevertheless some discrepancy between the experimental density and the calculated one remains.

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In present paper we extend our approach to take into account the effect of the GSC on the pairing. We consider the charge transition desities of the vibrational states in Zn isotopes.

2. Formalism

We employ a Hamiltonian including an average nuclear field as the Woods-Saxon potential, pairing interactions and isoscalar particle–hole (p–h) residual forces in separable form with the Bohr–Mottelson radial dependence (see for example [21,22]).

$$H = \sum_{\tau}[\sum_{jm}(E_j - \lambda_\tau)a^+_{jm}a_{jm} - \frac{G^{(0)}_\tau}{4} : (P^+_0 \cdot P_0)^\tau : - \sum_{\lambda\mu}\frac{\kappa^{(\lambda)}}{2} : (M^+_{\lambda\mu} \cdot M_{\lambda\mu})^\tau :] \quad (1)$$

The notation $\{\tau = (n,p)\}$ is used. The single-particle states are specified by the quantum numbers (jm); E_j are the single-particle energies; λ_τ is the chemical potential; $G^{(0)}_\tau$ and $\kappa^{(\lambda)}$ are the strengths in the p–p and in the p–h channel, respectively. The pair creation and the multipole operators entering the scalar products in (1) are defined in a standard fashion:

$$P^+_0 = \sum_{jm}(-1)^{j-m}a^+_{jm}a^+_{j-m} \quad (2)$$

$$M^+_{\lambda\mu} = \frac{1}{\sqrt{2\lambda+1}}\sum_{jj'mm'} f^{(\lambda)}_{jj'}(jmj'm' \mid \lambda\mu)a^+_{jm}a_{j'm'} \quad (3)$$

where $f^{(\lambda)}_{jj'}$ stands for the reduced single particle matrix element.

By performing the canonical Bogolubov transformation

$$a^+_{jm} = u_j\alpha^+_{jm} + (-1)^{j-m}v_j\alpha_{j-m} \quad (4)$$

one can rewrite the Hamiltonian in terms of the bifermion quasiparticle operators and their conjugate ones:

$$B(jj';\lambda\mu) = \sum_{mm'}(-1)^{j'+m'}(jmj'm' \mid \lambda\mu)\alpha^+_{jm}\alpha_{j'-m'}$$

$$A^+(jj';\lambda\mu) = \sum_{mm'}(jmj'm' \mid \lambda\mu)\alpha^+_{jm}\alpha^+_{j'm'} \quad (5)$$

We introduce the phonon creation operators

$$Q^+_{\lambda\mu i} = \frac{1}{2}\sum_{jj'}[\psi^{\lambda i}_{jj'}\,A^+(jj';\lambda\mu) - (-1)^{\lambda-\mu}\,\phi^{\lambda i}_{jj'}\,A(jj';\lambda-\mu)] \quad (6)$$

where the index λ denotes multipolarity and μ denotes its z-projection in the laboratory system.

Using the exact commutation relations between fermion operators one can prove that the following relation is valid:

$$\langle 0 \mid [Q_{\lambda\mu i}, Q^+_{\lambda'\mu' i'}] \mid 0\rangle = \delta_{\lambda\lambda'}\delta_{\mu\mu'}\sum_{jj'}(1 - q_{jj'})\frac{1}{2}[\psi^{\lambda i}_{jj'}\,\psi^{\lambda i'}_{jj'} - \phi^{\lambda i}_{jj'}\,\phi^{\lambda i'}_{jj'}] \quad (7)$$

where $\mid 0\rangle$ is the ground state (we assume that $\langle 0 \mid 0\rangle = 1$), $q_{jj'} = q_j + q_{j'}$ and q_j presents the quasiparticle distribution in the ground state:

$$q_j = \frac{1}{\sqrt{2j+1}}\langle 0 \mid B(jj;00) \mid 0\rangle \quad (8)$$

One assumes that $\mid 0\rangle$ is the phonon vacuum, i.e.

$$Q_{\lambda\mu i} \mid 0\rangle = 0 \quad (9)$$

and the phonon operators satisfy the linearized equation of motion

$$[H, Q^+_{\lambda\mu i}] \mid 0\rangle = \omega_{\lambda i}\,Q^+_{\lambda\mu i} \mid 0\rangle \quad (10)$$

Imposing the condition $q_j = 0$, one derives the well known RPA equations. Ken-ji Hara did the next step beyond the RPA [1] keeping in eqs. (7),(9),(10) all the terms containing the bifermion operators $B(jj';\lambda\mu)$, but replacing the latest by their mean values in the phonon vacuum. In this approach the equations determining the unknowns q_j are derived by using eq.(9). As was shown in refs. [1,20], eqs.(7)–(10) result in the following non-linear system for the phonon energies, amplitudes and quasiparticle distributions in the ground state:

$$(\varepsilon_{jj'} - \omega_{\lambda i})\psi^{\lambda i}_{jj'} - \frac{\kappa^{(\lambda)}f^{(\lambda)}_{jj'}u^{(+)}_{jj'}}{2(2\lambda+1)}\sum_{j_1j'_1}(1 - q_{j_1j'_1})u^{(+)}_{j_1j'_1}f^{(\lambda)}_{j_1j'_1}(\psi^{\lambda i}_{j_1j'_1} + \phi^{\lambda i}_{j_1j'_1}) = 0 \quad (11)$$

$$(\varepsilon_{jj'} + \omega_{\lambda i})\phi^{\lambda i}_{jj'} - \frac{\kappa^{(\lambda)}f^{(\lambda)}_{jj'}u^{(+)}_{jj'}}{2(2\lambda+1)}\sum_{j_1j'_1}(1 - q_{j_1j'_1})u^{(+)}_{j_1j'_1}f^{(\lambda)}_{j_1j'_1}(\psi^{\lambda i}_{j_1j'_1} + \phi^{\lambda i}_{j_1j'_1}) = 0 \quad (12)$$

$$q_j = \sum_{\lambda ij'}\frac{2\lambda+1}{2j+1}(1 - q_{jj'})(\phi^{\lambda i}_{jj'})^2 \quad (13)$$

$$\sum_{jj'}(1 - q_{jj'})[(\psi^{\lambda i}_{jj'})^2 - (\phi^{\lambda i}_{jj'})^2] - 2 = 0 \quad (14)$$

Here $\varepsilon_{jj'}$ are the two-quasiparticle energies and $u^{(+)}_{jj'} = u_jv_{j'} + v_ju_{j'}$. To take into account the effect of GSC on the pairing one needs to modify the standardt BCS equations [1,6]. For the gap

$$\Delta = G^0_\tau\sum_j(2j+1)\,u_j\,v_j(1 - 2q_j) \quad (15)$$

one gets the equation

$$\frac{2}{G^0} = \sum_j \frac{(j+1/2)(1-2q_j)}{\sqrt{\Delta^2+(E_j-\lambda)^2}} \tag{16}$$

and from the particle-number condition:

$$\langle 0 \mid \hat{N} \mid 0 \rangle = N = \sum_j (2j+1)[v_j^2 + (u_j^2 - v_j^2)q_j] \tag{17}$$

it is possible to determine the chemical potential λ:

$$(j+1/2) - N = \sum_{j'} \frac{(j'+1/2)(E_{j'}-\lambda)(1-2q_{j'})}{\sqrt{\Delta^2+(E_{j'}-\lambda)^2}} \tag{18}$$

The quasiparticle energies ε_j and coefficients u_j, v_j can be calculated in the same way as within BCS. The factors $(1 - q_{jj'})$ take into account the blocking effect due to the Pauli principle.

As a result we get the self-consistent system of equations to treat the GSC. In the case of $q_j = 0$ these equations reduce to the usual RPA and BCS equations.

The system of non-linear equations (11)-(18) is solved numerically by means of the auto-regularized Gauss-Newton iteration process [23,24], executed by the program-package REGN [25]. In this way we construct stable solutions of the system with high accuracy.

In the present paper we study the effect of the GSC on the charge transition density for the vibrational states in spherical nuclei. As an example we consider the chain of the even Zn isotopes with $A = 64 - -70$. The charge transition density is calculated by the formula

$$\rho_i^\lambda(r) = \frac{1}{2}\sum_{jj'}(1-q_{jj'})\,u_{jj'}^{(+)}\,(\psi_{jj'}^{\lambda i} + \phi_{jj'}^{\lambda i})\,\rho_{jj'}^\lambda(r) \tag{19}$$

The expression for the two-quasiparticle charge transition density $\rho_{jj'}^\lambda(r)$ can be found for example in [16]. The calculated charge transition densities $\rho_i^\lambda(r)$ are folded with the formfactor of the proton charge distribution [26]. Knowing the charge transition density it is possible to calculate the reduced transition probabilities [27].

3. Details of calculations and results

We performed numerical calculations for the Zn isotopes for which the experimental data [28] and the RPA results [29] are known. The Woods-Saxon potential parameters in use are basically the same as in our previous paper [20]. They were chosen to reach a reasonable theoretical description of the experimental ground state density and the r.m.s. radius [30]. The pairing constants G_τ^0 are fixed so as to reproduce the odd-even mass difference of neighboring nuclei. We always adjust the strength parameters $\kappa^{(\lambda)}$ so that the $B(E\lambda)$ values were close to the experimental ones.

The system of eqs. (11)-(18), which treats the GSC self-consistently, describes (via eq.(13)) the coupling between different vibrations and between all the phonon roots of a certain multipolarity. Our studies [20] show that the interplay of different roots of the system of eqs.(11)-(14) is not essential: the contribution from the second root for example affects the q_j by no more than 2 %. Thus, one can restrict the sum in eq. (13) to only the first (collective) root without substantial loss of accuracy.

To study the coupling between different vibrations we took into account $\lambda = 2,3,4$ terms in eqs. (11)-(13), using the same values for the quadrupole, octupole and hexadecapole constants. According to our calculations for ^{64}Zn the most essential role for the GSC is played by the quadrupole and octupole vibrations. The values of q_j calculated for the pure quadrupole vibrations(i.e. $\lambda = 2$, only) are as a rule a few times higher than those for the pure octupole vibrations ($\lambda = 3$, only). However, in the case of coupled ($\lambda = 2,3$) vibrations the resulting q_j are larger than the sum of q_j ($\lambda = 2$) and q_j ($\lambda = 3$) contrary to the RPA case because of the non- linear character of the λ-mixing. The admixture of the hexadecapole vibrations changes q_j by no more than 5 %. So in what follows we take into account the mixing of the quadrupole and octupole modes only.

The results of our calculations for the quasiparticle distribution in the ground state of ^{66}Zn are shown in the table 1. The table 1 contains the results when we take into account the effect of GSC on q_j but neglect the blocking in the pairing due to GSC ($q_j(Hara)$) and when the blocking is included ($q_j(B-Hara)$). The q_j calculated within the RPA [3] as

$$q_j = \frac{1}{2}\sum_{\lambda i j'} \frac{2\lambda+1}{2j+1}(\phi_{jj'}^{\lambda i})^2 \tag{20}$$

are given in the same table too. As one can see from the table 1 the q_j have large values for the subshells near the Fermi surface only and the Hara approach gives stronger correlations in comparison with the RPA. The inclusion of the blocking in the pairing due to GSC results in an additional enhancement of correlations. A similar behaviour of q_j's was found for all Zn isotopes.

The theoretical and experimental values [28-29] for energies and the reduced transition probabilities of the first quadrupole and octupole states in the Zn isotopes are shown in tables 2,3. As one can see from the table 2 the energy of the 2_1^+ state calculated within the RPA with the above choice of parameters is usally higher than the experimental value. In the same time the solution of the GSC problem beyond the RPA exhibits a shift up of the energy $\omega_{2_1^+}$ and an essential reduction of the B(E2)-value. To restore the B(E2) within the Hara approach we had to increase, for example, the value of $\kappa^{(2)}$ by 38% and 23% in ^{64}Zn and ^{68}Zn, respectively. In the last case the resulting energies are shifted down a little. The selfconsistent calculations taking into account the blocking in the pairing give

an additional shift down in energies, which are in a better agreement with experimental data. The increase of $\kappa^{(2)}$ is about 12% for all isotopes in comparison with the RPA values.

It is worth mentioning that the value of the $\kappa^{(2)}$ for the non–linear problem is quite larger than the critical RPA constant where the RPA solution becomes complex. For such values of $\kappa^{(2)}$ the use of the procedure [7] where one solves first eqs. (11)–(12) with $q_j = 0$ and then eq.(13) becomes non–applicable.

It is seen from the table 3 that the GSC with the coupled vibrations change drastically the 3^-_1 energies. The renormalized values of $\kappa^{(3)}$ exceed the RPA ones by 11–19 % and they are smaller than the critical values, in contrast to the quadrupole constants. In the cases of ^{66}Zn and ^{68}Zn, where the experimental B(E3)-values are known, our choice of parameters enables us to reproduce energies and transition probabilities simultaneously. We overesimate apparently the collectivity of the 3^-_1 state in ^{70}Zn.

It should be noted also that according to our calculations the selfconsistent inclusion of the GSC in the pairing problem (see eqs. (15)–(18)) garantees the number conservation with a high accuracy. Calculations without taking into account the GSC effect on the pairing lead sometimes up to a 3% difference from the exact particle number.

The figures 1–4 show the transition charge densities from the ground to the first 2^+ states in the Zn isotopes. Our calculations for the $\rho^\lambda_i(r)$ give results which are similar to the ones of [29], obtained in the RPA with the Skyrme forces, but in contrast with [29] we did not assign occupation probability to each single-particle orbital empirically. As it was pointed in [29] the authors were enforced to destroy the self-consistency shifting the single-particle spectrum for the unoccupied orbitals with respect to the occupied ones to reproduce the experimental value for the first 2^+ state.

As one can see from the figures 1–4 the RPA reproduces the behaviour of the charge transition densities qualitatively but it overestimates the interior part of the $\rho^\lambda_i(r)$. The inclusion of the GSC beyond the RPA (the Hara approach) gives a 30% –35% depletion of the maximum of the $\rho^\lambda_i(r)$ in the interior region of the nuclei. The selfconsistent treatment of the GSC results in a suppression of interior oscilations by factors 2.3 (for ^{68}Zn) – 3.4 (for ^{70}Zn). The calculated $\rho^\lambda_i(r)$ for ^{64}Zn and ^{66}Zn are very close to the experimental data. Such a depletion is related with the Pauli blocking effect for the proton two-quasiparticle configuration $\{2p_{3/2}, 2p_{3/2}\}$, which is mainly responsible for the interior bump in the charge transition densities in the Zn isotopes. According to our RPA calculations the proton two-quasiparticle configuration $\{2p_{3/2}, 2p_{3/2}\}$ gives a contribution about 38% and about 46% into the norm of the first quadrupole phonon in ^{66}Zn and in ^{68}Zn, respectively. The

inclusion of the GSC redistributes the strength of this configuration over many phonon roots and as result the contribution into the first root becomes 8.8% in ^{66}Zn and 15.6% in ^{68}Zn. As it follows from the eq. (19) the GSC supresses the contribution of the partial two-quasiparticle transition densities having big $q_{jj'}$. It is seen from the table 1 that the $q_{2p_{3/2}}$ is the biggest for protons and as it was mentioned above plays an essential role in the structure of the interior part of the transition density for the 2^+_1 states. The configuration $\{1f_{5/2}, 1f_{5/2}\}$ gives some contribution in the interior part too and the same mechanism of a supression takes place for it. The amplitudes of the oscilations for the configurations with low orbital momenta are bigger than for the ones with high orbital momenta.

We would like to emphasize that this effect can not be reproduced by any renormalization of the $\kappa^{(2)}$ in the RPA.

The behaviour of the neutron transition densities differs from the proton ones (see figs.5,6). In the case of ^{66}Zn, for example, the influence of the GSC on the interior part of the transition density is very weak because the neutron configuration $\{2p_{3/2}, 2p_{3/2}\}$ contributes not more than 1.5% in the norm of the 2^+_1 state, and the contribution of the configuration $\{1f_{5/2}, 1f_{5/2}\}$ remains practically the same (about 8%) in all cases. An extra four neutrons in ^{70}Zn changes the structure of the neutron part of the 2^+_1 wave function and the GSC affects mainly the surface part of the neutron transition density. In the this case the configurations which are responsible of the interior oscillations in ^{66}Zn have very small contribution to the structure of the first quadrupole state.

We calculated the charge transition density for the one-phonon 4^+ state and did not find any essential oscillations in the interior region. That is due to the lack of the configuration $\{2p_{3/2}, 2p_{3/2}\}$ in such states because of the angular momentum coupling rule.

The fig.7 presents the charge transition density for the 3^-_1 state in ^{66}Zn. This density has a clear surface nature and there are no strong oscillations in the interior region of the nucleus because of a destructive intenference of the two-quasiparticle partial transition densities constructed from the single-particle wave functions with different parity. The same picture takes place in the other Zn isotopes and it is typical for the transition densities of the octupole vibration states (see [16–19]).

It is interesting to note that the GSC increase slightly the ground state r.m.s. charge radius up to 4.05 fm in ^{66}Zn. Its RPA value is 4.02 fm, compared to the experimental one 3.99 fm [30].

Table 1. The quasiparticle distribution in the ground state of ^{66}Zn

nlj	neutrons			protons		
	q_j(RPA)	q_j(Hara)	q_j(B-Hara)	q_j(RPA)	q_j(Hara)	q_j(B-Hara)
$2s_{1/2}$	0.0052	0.0170	0.0249	0.0132	0.0444	0.0674
$1f_{7/2}$	0.0155	0.0492	0.0723	0.0307	0.0907	0.1394
$2p_{3/2}$	0.0561	0.1640	0.2286	0.0894	0.2195	0.2425
$1f_{5/2}$	0.0438	0.1180	0.1688	0.0333	0.0959	0.1039
$2p_{1/2}$	0.0722	0.1791	0.2476	0.0367	0.0960	0.0930
$1g_{9/2}$	0.0204	0.0685	0.0997	0.0164	0.0564	0.0701

Table 2. Energies and transition probabilities for the 2_1^+ states in the Zn isotopes

A	Experiment		Calculation			
	ω, MeV	$B(E2\uparrow), e^2fm^4$	ω_{RPA}, MeV	ω_H, MeV	ω_{B-H}, MeV	$B(E2\uparrow), e^2fm^4$
64	0.992	1597	1.16	1.09	0.70	1597
66	1.039	1426	1.33	1.27	1.02	1426
68	1.077	1360	1.39	1.33	1.25	1360
70	0.885	2050	1.32	1.25	1.05	2050

Table 3. Energies and transition probabilities for the 3_1^- states in the Zn isotopes

A	Experiment		Calculation			
	ω, MeV	$B(E3\uparrow), e^2fm^6$	ω_{RPA}, MeV	ω_H, MeV	ω_{B-H}, MeV	$B(E3\uparrow), e^2fm^6$
64	2.999	-	3.90	3.03	2.69	44000
66	2.830	41660	3.89	3.19	2.83	41660
68	2.751	38460	3.62	3.06	2.75	38460
70	2.839	-	3.55	2.77	2.32	36000

4.Conclusion

In conclusion, we point out that the system of non–linear equations, describing the ground state correlations beyond RPA selfconsistently was solved in a realistic case. It was found that such a proper treatment of the GSC leads to a suppression by factors 2.3–3.4 in comparison to the RPA of the charge transition density in the interior region of the Zn isotopes. It follows from our study that the inclusion of the blocking effect due the Pauli prinsiple enables us to solve the long standing problem of the theoretical overstimation of the charge transition densities in the interior region of nuclei for the positive parity vibration states compared to experimental data.

Our preliminary investigation shows that taking into account the blocking effect for the proton configuration $\{2d_{5/2}, 2d_{5/2}\}$ it is possible to get a good description of the interior patr of the charge transition densities of $2^+, 4^+$ states in the Nd isotopes. To describe the surface part one has to take into account in addition the interplay between the one-phonon and more complex configurations, so the system of non–linear equations should be extended. Such calculations are in progress now.

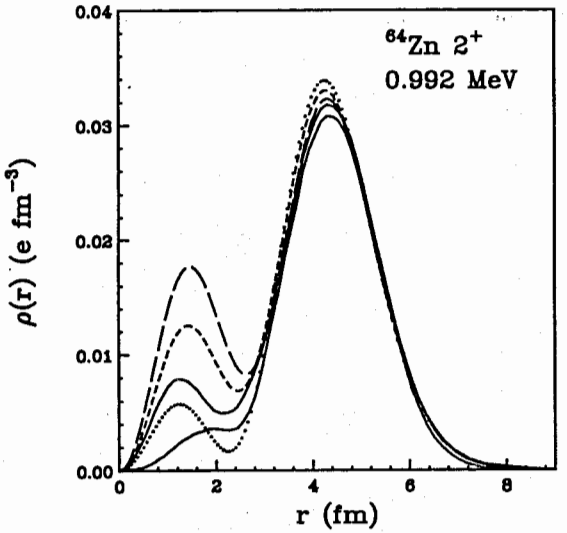

Fig.1: Transition charge density from the ground state to the first 2_1^+ state in ^{64}Zn.
solid curves – experimental data [28]; long dashed curve – RPA calculation;
dashed curve – results with GSC beyond the RPA without the blocking in the pairing;
dotted curve – results with the selfconsistent inclusion of GSC beyond the RPA.

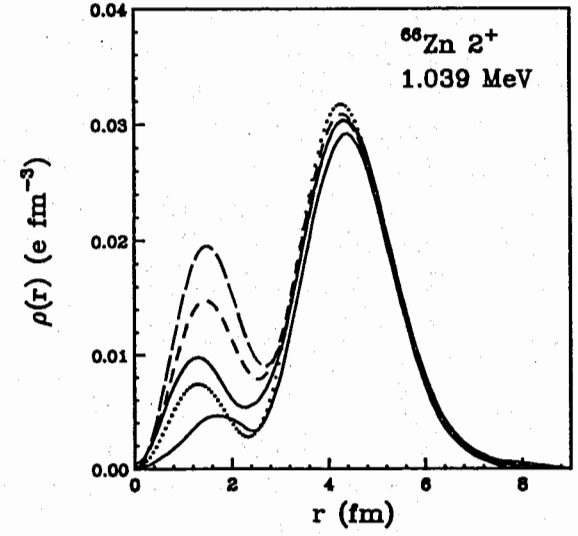

Fig.2: Transition charge density from the ground state to the first 2_1^+ state in ^{66}Zn.
Notations are the same as in fig.1.

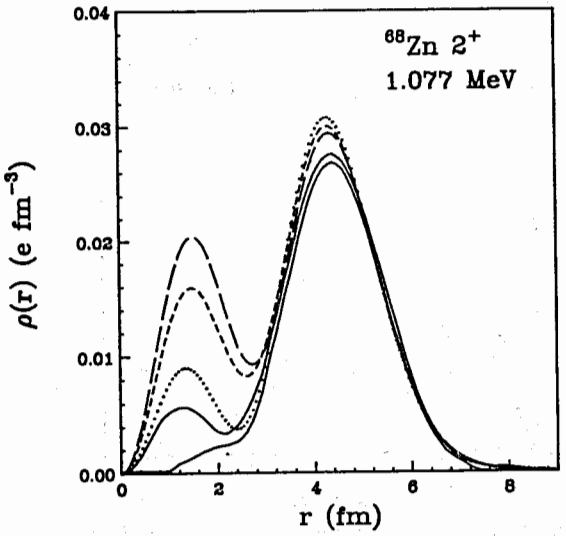

Fig.3: Transition charge density from the ground state to the first 2_1^+ state in ^{68}Zn.
Notations are the same as in fig.1.

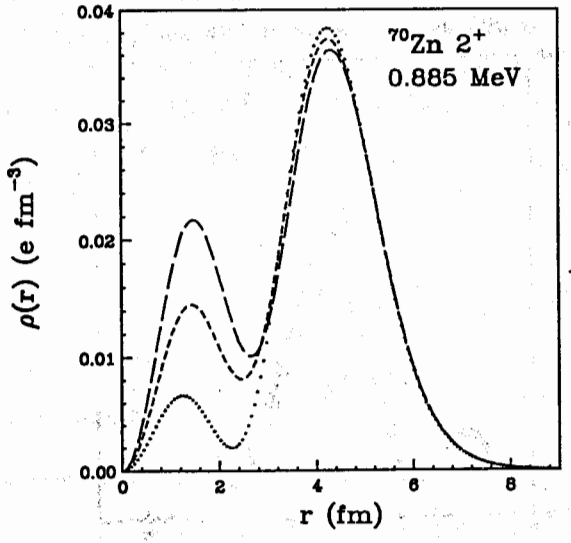

Fig.4: Transition charge density from the ground state to the first 2_1^+ state in ^{70}Zn.
Notations are the same as in fig.1.

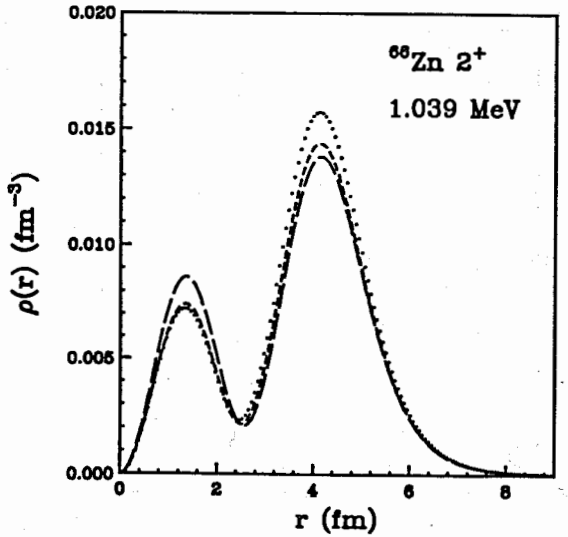

Fig.5: Neutron transition density from the ground state to the first 2_1^+ state in ^{66}Zn.
long dashed curve – RPA calculation;
dashed curve – results with GSC beyond the RPA without the blocking in the pairing;
dotted curve – results with the selfconsistent inclusion of GSC beyond the RPA.

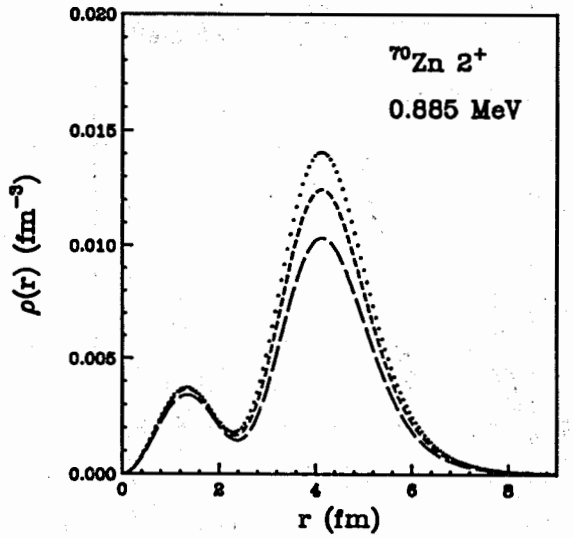

Fig.6: Neutron transition density from the ground state to the first 2_1^+ state in ^{70}Zn.
Notations are the same as in fig.5.

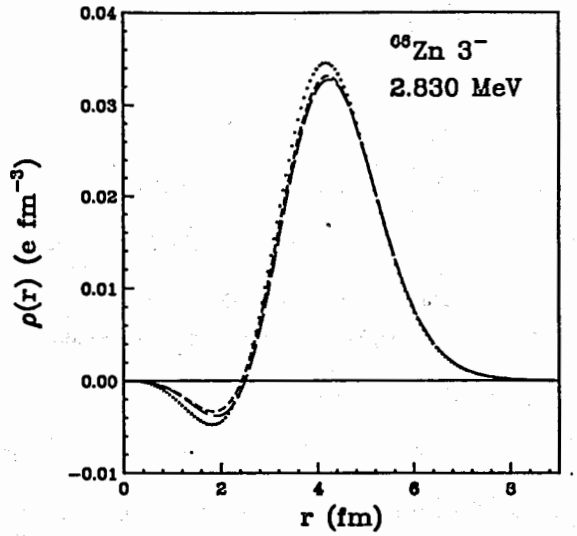

Fig.7: Transition charge density from the ground state to the first 3_1^- state in ^{66}Zn.
Notations are the same as in fig.5.

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Влияние корреляций в основных состояниях
на зарядовые переходные плотности вибрационных состояний

Изучено влияние корреляций в основных состояниях на зарядовые переходные плотности вибрационных состояний в сферических ядрах. Выход за приближение случайных фаз приводит к необходимости решать нелинейную систему уравнений. При этом также учитывается влияние корреляций на спаривание. Показано, что учет корреляций в основных состояниях приводит к существенному подавлению зарядовых переходных плотностей во внутренней области ядер по сравнению с расчетами в приближении случайных фаз и позволяет описать экспериментальные данные.

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Effect of Ground State Correlations
on the Charge Transition Densities of Vibrational States

The effect of ground state correlations on the charge transition densities of vibrational states in spherical nuclei is studied. The problem for the ground state correlations beyond RPA leads to a non-linear system of equations, which is solved numerically. The influence of the correlations on the pairing is taken into account too. The inclusion of ground state correlations beyond RPA results in an essential suppression of the charge transition density in the nuclear interior in comparison with the RPA calculations and enables one to reproduce the experimental data.

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

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