$g_{9/2}$ nuclei and neutron-proton interaction

L. Coraggio,[1] A. Covello,[1,2] A. Gargano,[1,*] and N. Itaco[1,2]

[1]*Istituto Nazionale di Fisica Nucleare, Complesso Universitario di Monte S. Angelo, I-80126 Napoli, Italy*

[2]*Dipartimento di Scienze Fisiche, Università di Napoli Federico II, Complesso Universitario di Monte S. Angelo, I-80126 Napoli, Italy*

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We perform shell-model calculations for nuclei below ^{100}Sn, focusing attention on the two $N = Z$ nuclei, ^{96}Cd and ^{92}Pd, the latter having recently been the subject of great experimental and theoretical interest. We consider nuclei for which the $0g_{9/2}$ orbit plays a dominant role and employ a realistic low-momentum two-body effective interaction derived from the CD-Bonn nucleon-nucleon potential. This implies that no phenomenological input enters our effective Hamiltonian. The calculated results for ^{92}Pd are in very good agreement with the available experimental data, which gives us confidence in our predictions for ^{96}Cd. An analysis of the wave functions of both ^{96}Cd and ^{92}Pd is performed to investigate the role of the isoscalar spin-aligned coupling.

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I. INTRODUCTION

In a recent work [1] three excited states in the $N = Z$ nucleus ^{92}Pd, lying quite far from the stability line, were observed. This is a remarkable achievement, as $N = Z$ nuclei play a special role in our understanding of the nuclear effective interaction, in particular, the interplay and competition between isovector and isoscalar components. In fact, in this case neutrons and protons occupy the same orbitals, which gives rise to a large overlap of their wave functions. In this context, the main question, which is still a matter of debate [2,3], is the existence of strongly correlated $T = 0$ np pairs, similar to the well-known case of neutron and proton pairs coupled to $J = 0$ and $T = 1$.

In Ref. [1] it was pointed out that the main feature of the measured levels of ^{92}Pd is their approximate equidistance. An interpretation was given of this feature in terms of a shell-model calculation with an empirical Hamiltonian. This reproduces very well the experimental levels and yields wave functions built manly from $J = 9$ np pairs, which has been seen as evidence for a spin-aligned np paired phase.

The results in Ref. [1] immediately attracted great attention and the role of isoscalar np pairs in the low-energy structure of $N = Z$ nuclei close to doubly magic ^{100}Sn has been investigated in some very recent theoretical papers [4–6]. These studies confirm substantially the dominance of $J = 9$, $T = 0$ pairs in the low-lying yrast states of $N = Z$ nuclei with four, six, and eight holes below ^{100}Sn. The relevance of the isoscalar component of the np interaction was stressed in a subsequent work [7], where the 16^+ "spin-gap" isomer in ^{96}Cd was identified and its origin explained as being caused by the strong influence of this component.

Another interesting outcome of the above works is that in all nuclei above ^{88}Ru the low-lying yrast states can be essentially described in terms of the single $0g_{9/2}$ shell. This situation is of course reminiscent of the so-called $f_{7/2}$ nuclei, which have been the subject of a large number of theoretical studies since the seminal paper by McCullen, Bayman, and Zamick [8]. It therefore seems appropriate to call these special nuclei below ^{100}Sn "$g_{9/2}$ nuclei."

Some 10 years ago, we performed [9] shell-model calculations for $N = 50$ nuclei immediately below ^{100}Sn employing a realistic effective interaction derived from the Bonn A nucleon-nucleon (NN) potential by means of a G-matrix formalism. In that work we took as model space for the valence proton holes the four levels $f_{5/2}$, $p_{3/2}$, $p_{1/2}$, and $g_{9/2}$ of the 28-50 shell. Our results turned out to be in very good agreement with the available experimental data. Since then, however, a new paradigm for realistic shell-model calculations has been developed which consists in renormalizing the strong short-range repulsion of the bare NN potential through the so-called $V_{\text{low-}k}$ approach [10]. Furthermore, high-precision potentials have been constructed which fit the pp and np scattering data with $\chi^2/\text{datum} \approx 1$.

The exciting new findings mentioned before have stimulated us to perform modern realistic shell-model calculations for nuclei below ^{100}Sn, with particular attention focused on ^{92}Pd and on the heavier $N = Z$ nucleus ^{96}Cd, for which theoretical predictions are likely to be verified in the not too distant future. Based on the dominant role of the $g_{9/2}$ orbit in the low-lying states of the nuclei considered in the present study, we have restricted our calculations to the single $g_{9/2}$ shell. This permits a more transparent analysis of the structure of the wave functions, especially for ^{96}Cd, in terms of either the $[nn] \otimes [pp]$ or the $[np] \otimes [np]$ coupling schemes. Comparison between these two approaches is instrumental to understanding the role of $J = 9$ np pairs. It may also be worth recalling that within the fpg space the choice of single-hole energies is not an easy task [9], as there is no spectroscopic information on the single-hole nuclei ^{99}In and ^{99}Sn.

In Sec. II we focus attention on the $g_{9/2}$ effective interaction employed in our calculations. Results for ^{96}Cd and ^{92}Pd are presented in Sec. III, where we also give a detailed analysis of the structure of wave functions. Section IV provides a summary and concluding remarks.

*gargano@na.infn.it

II. $g_{9/2}$ EFFECTIVE INTERACTION

We assume ^{100}Sn to be a closed core and let the neutron and proton holes move in the single $g_{9/2}$ orbit. Our two-body effective interaction is derived within the framework of the time-dependent degenerate linked-diagram perturbation theory [11] starting from the high-precision CD-Bonn NN potential [12]. This potential, which, as all modern NN potentials, contains high-momentum nonperturbative modes, is renormalized by constructing a low-momentum potential $V_{\text{low-}k}$. This is achieved by integrating out the high-momentum modes of the bare potential down to a cutoff momentum $\Lambda = 2.1$ fm^{-1}. Then the smooth $V_{\text{low-}k}$ potential plus the Coulomb force for protons is used to calculate the two-body matrix elements of the effective interaction by means of the $\hat{Q}$-box folded-diagram expansion [11], with the $\hat{Q}$ box including all diagrams up to third order. These diagrams are computed using the harmonic oscillator basis and considering intermediate states composed of all possible hole states and particle states restricted to the six proton and neutron shells above the Fermi surface. The oscillator parameter $\hbar\omega$ is 8.55 MeV, as obtained from the expression $\hbar\omega = 45A^{-1/3} - 25A^{-2/3}$ for $A = 100$. Shell-model calculations were performed using the NUSHELLX code [13].

We list in Table I the two-body matrix elements of the effective interaction. Owing to the Coulomb force, there is no isospin symmetry, the nn matrix elements being more attractive than the pp ones by about 250–350 keV, which agrees quite well with the results of previous works [14,15] where the effective interaction was determined by a least-squares fit to experimental energies. As regards the $T = 1$ np matrix elements, they differ only by, at most, 100 keV from the nn ones. This reflects the fact that for np and nn valence holes or particles the Coulomb force acts only through $\hat{Q}$-box diagrams starting at second and third order, respectively. In Table I we also see that the most attractive matrix elements correspond to the $J = 0$, $J = 1$, and $J = 9$ states, the $J = 0$ matrix element being the largest one, while the other two have about the same magnitude. This dominance of the $J = 0$ matrix element stems from the fact that our interaction is derived for the single $g_{9/2}$ orbit. The same feature is shown by the three different $g_{9/2}$ interactions in Ref. [4], while this is not the case for interactions defined in the fpg space, as, for instance, that developed in Ref. [16].

TABLE I. Proton-proton, neutron-neutron, and proton-neutron matrix elements of V_{eff} in the $g_{9/2}$ orbit (in MeV).

J	T	pp	nn	np
0	1	−1.836	−2.224	−2.317
1	0			−1.488
2	1	−0.353	−0.662	−0.667
3	0			−0.440
4	1	0.171	−0.088	−0.100
5	0			−0.271
6	1	0.317	0.083	0.066
7	0			−0.404
8	1	0.459	0.221	0.210
9	0			−1.402

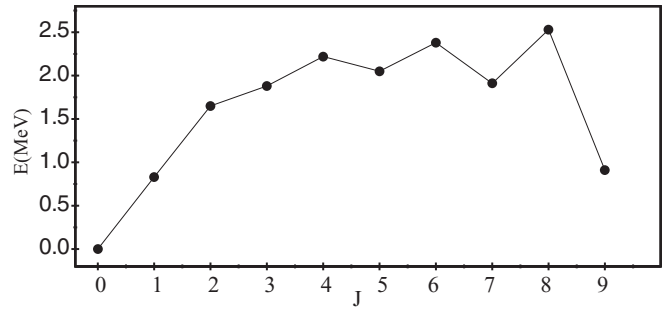

FIG. 1. Calculated proton hole–neutron hole multiplet in ^{98}In.

The np matrix elements, relative to the $J = 0$ energy, correspond to the excitation spectrum of ^{98}In, which we find interesting to show in Fig. 1. This, unfortunately, has no experimental counterpart so cannot be used to test our interaction. We see, however, that the np multiplet exhibits a downward parabolic behavior, which is typical of nuclei with one proton–one neutron valence particles or holes, as, for instance, ^{42}Sc or ^{54}Co, with two particles and holes, respectively, in the $f_{7/2}$ orbit. The observed multiplet in these nuclei shows the same behavior as that we have found for ^{98}In, the only difference being a larger dispersion in energy values, which is caused by the greater attractiveness of the np interaction in lighter systems.

For a test of our np matrix elements, we compare in Fig. 2 the observed multiplet in ^{90}Nb [17] with the calculated one. The latter is, in fact, directly related to the matrix elements in the hole-hole np interaction through the Pandya transformation [18]. Note that several 1^+ states have been observed in ^{90}Nb at a low excitation energy. Following the suggestion in Ref. [4], we report the fifth one at 2.126 MeV while excluding the observed 0^+ state at 5.008 MeV, which is too high in energy to make our interpretation trustworthy in terms of the single-$g_{9/2}$ model. Figure 2 shows that the agreement between theory and experiment is very good, the calculated energies overestimating the experimental values by at most 150 keV in the case of the 2^+ state. The latter, however, is likely to be admixed with configurations outside our model space.

We now focus attention on the pp matrix elements by comparing the spectra of the three $N = 50$ isotones, ^{98}Cd, ^{97}Ag, and ^{96}Pd, with the experimental ones. We include all observed levels for ^{98}Cd and positive-parity yrast levels below

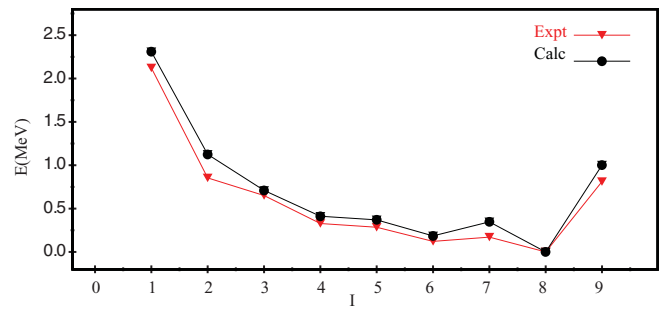

FIG. 2. (Color online) Experimental and calculated proton hole–neutron particle multiplet in ^{90}Nb.

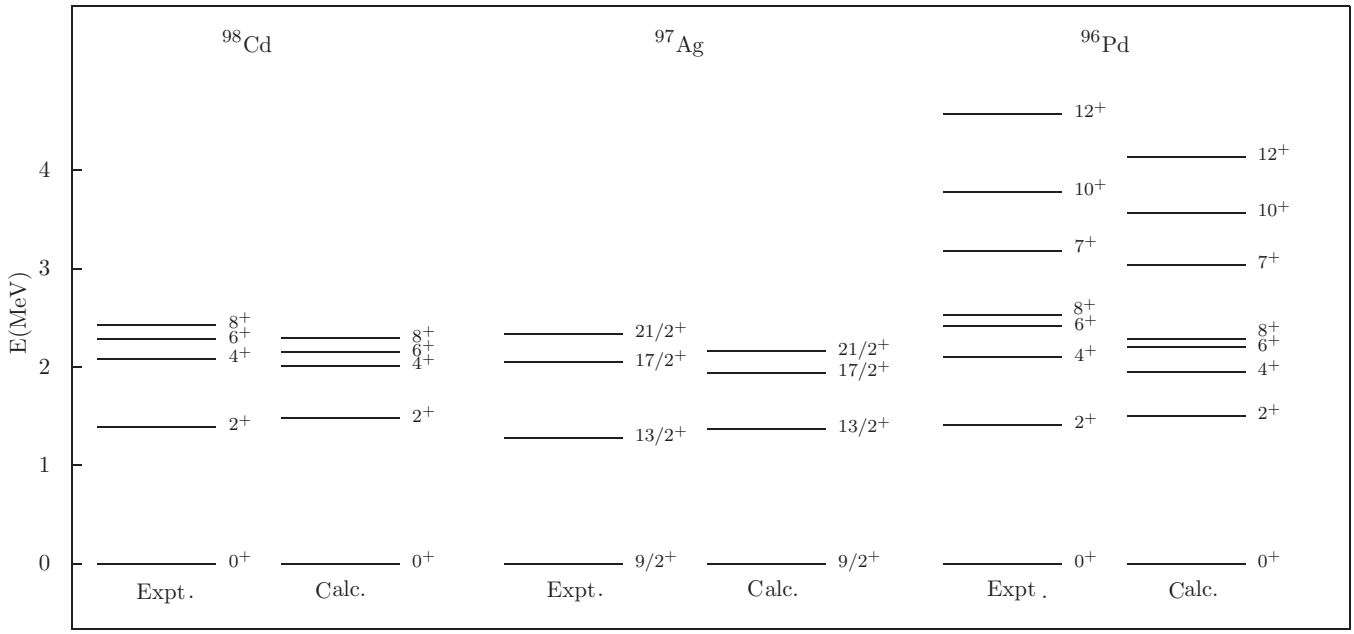

FIG. 3. Experimental and calculated spectra of the $N = 50$ isotones ^{98}Cd, ^{97}Ag, and ^{96}Pd.

4.5 MeV for the two latter nuclei. This is done in Fig. 3, where we see that the agreement between theory and experiment is quite satisfactory up to about 3.7 MeV. We overestimate the energy of the first excited state in the three nuclei by less than 100 keV, while the energies of all other levels, except the 12^+ in ^{96}Pd, which is predicted at more than 400 keV below the experimental one, are underestimated by an amount which increases when the mass number decreases, reaching at most 250 keV. In concluding this section, it is worth noting that the accuracy of the present $g_{9/2}$ calculation is similar to that in Ref. [1], where the spectrum of ^{96}Pd up to the 10^+ state was calculated in the fpg model space.

III. RESULTS

The results obtained for ^{90}Nb and $N = 50$ isotones gave us confidence in our effective interaction, at least with regard to its predictive power for low-energy states. We then performed calculations for the two $N = Z$ nuclei ^{96}Cd and ^{92}Pd. The calculated spectra are reported in Figs. 4 and 5(a), respectively, together with a comparison with the experimental data from Ref. [1] for ^{92}Pd. In view of recent experimental findings [7], we have reported yrast states in ^{96}Cd up to $I^\pi = 16^+$, while for ^{92}Pd we have included two more states with respect to those identified in Ref. [1].

The three observed states in ^{92}Pd are very well reproduced by the theory. Of course, one may not expect the same kind of agreement for high-energy states. It is worth mentioning, however, that our results account for the isomeric nature of the 16^+ state identified in ^{96}Cd. More precisely, we find that the location of this state below the $I^\pi = 12^+$ and 14^+ states can be traced to the isoscalar np component of the interaction, in agreement with the results of Ref. [7]. In particular, we have verified that a strong attractive $J = 9$ matrix element is essential to make the 16^+ state isomeric. This is related to the fact that for the 12^+, 14^+, and 16^+ states, the average number of isoscalar $J = 9$ pairs (see Sec. III B for the definition of this quantity) increases with the angular momentum, reaching its maximum value for the 16^+ state. This is, in fact, 2.5, to be compared with 2.2 and 2.0 for the 14^+ and 12^+ states, respectively. Therefore a sufficiently attractive $J = 9$ matrix element pushes the 16^+ state below the two lower spin states. Relevant to this discussion is the fact that the structure of the wave functions for the 12^+, 14^+, and 16^+ states, and consequently the average number of isoscalar $J = 9$ pairs, does not depend on the size of the $J = 9$ matrix element.

Let us now focus on states up to $I^\pi = 10^+$ for ^{96}Cd and ^{92}Pd. For both nuclei, we find an almost regularly spaced level sequence up to $I^\pi = 6^+$, with a slightly reduced 6^+-4^+ spacing. Then the 8^+-6^+ spacing becomes even smaller while the 10^+-8^+ one increases considerably. These features, which, up to the 6^+ state, find a correspondence in the experimental data for ^{92}Pd, are only slightly more pronounced in ^{96}Cd. This similarity may be seen as an indication that the same correlations come into play in their low-energy spectra. We discuss in more detail the results for both nuclei in the two Secs. III A and III B, below.

A. ^{96}Cd

We start with ^{96}Cd, for which a simple and clear analysis of the wave functions can be performed using the orthogonal basis formed by products of nn and pp states. The structure of the wave functions in terms of this basis set is reported in Table II, while their overlaps with the $[(np)9(np)9]_I$ state are listed in Table III, column a.

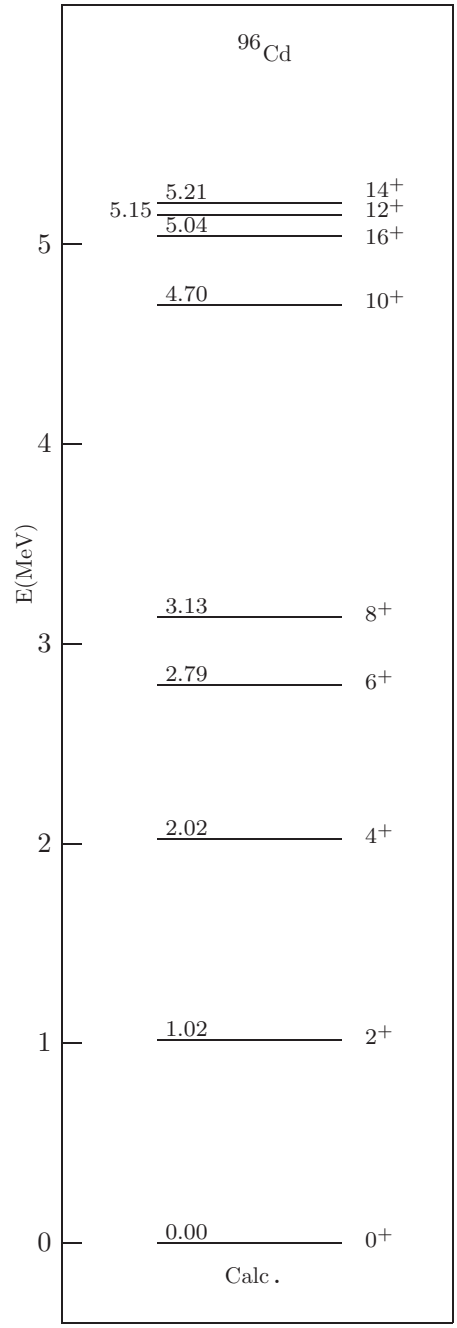

FIG. 4. Calculated spectrum of ^{96}Cd.

Note that for each angular momentum I an orthonormal basis formed by products of two np-pair vectors can been constructed, the $[(np)9(np)9]_I$ state being one of them. These vectors are simply related to the $[nn] \otimes [pp]$ basis through

$$|(np)J_1(np)J_2;I\rangle = \frac{1}{\sqrt{N_{J_1J_2}}} \sum_{J_nJ_p} [\hat{J}_1\hat{J}_2\hat{J}_n\hat{J}_p]^{1/2} \times \begin{Bmatrix} \frac{9}{2} & \frac{9}{2} & J_1 \\ \frac{9}{2} & \frac{9}{2} & J_2 \\ J_n & J_p & I \end{Bmatrix} |(nn)J_n(pp)J_p;I\rangle, \quad (1)$$

where $[\hat{J}] = (2J+1)$ and N denotes the normalization factor. In Table II we see that the wave functions of ^{96}Cd expressed in terms of the $[nn] \otimes [pp]$ basis are strongly fragmented. In fact, large seniority-4 components are present in the ground as well as in the excited $I^\pi = 2^+$, 4^+, 6^+, and 8^+ states. More precisely, their percentage in the ground state is 43%, while in the other states it is not less than 34%. The 10^+ state is characterized by an admixture of different seniority-4 components, each with a weight not exceeding 18%. On the other hand, Table III, column a, shows that, when written in terms of an $[np] \otimes [np]$ basis, the ground and the first two excited states are largely dominated by the $[(np)9(np)9]_I$ component. The weight of this component, however, is significantly smaller for the other three states having the minimum value (6%) for the 8^+ state. These results for ^{96}Cd are in line with those of Ref. [4].

The large fragmentation evidenced in Table II is of course related to the np interaction, and in this connection it is interesting to determine the role of the $J = 9$ matrix element $V_9(np)$. To this end, we redid our calculations with two different values of $V_9(np)$, namely, increasing and reducing $V_9(np)$ by a factor of 2. A similar analysis was done in Refs. [5] and [6].

We find that for the reduced value of $V_9(np)$ the 6^+ and 8^+ states decrease in energy getting close to the 4^+ state. As a matter of fact, in this case the group of the 4^+, 6^+, and 8^+ levels concentrates in a small energy range separated by a large energy gap from both the 2^+ and the 10^+ states. The spectrum of ^{96}Cd (up to $I^\pi = 8^+$) then becomes similar to that of ^{98}Cd, the wave functions still being, however, significantly seniority admixed. The weight of the seniority-4 components is no smaller than 29%. On the contrary, a doubled value of $V_9(np)$ leads to almost-equidistant levels, with an energy separation of about 1 MeV, the only exception being the 8^+-6^+ spacing, which is $\sim$300 keV smaller. This evolution toward an equidistant-level spectrum comes along with a larger fragmentation of the wave functions. The only state which has a less admixed nature is the 10^+ state. We find that the percentage of the $[(nn)0(pp)0]_0$ component in the ground state reduces to 51%, while that of the seniority-2 components, $[(nn)I(pp)0]_I$ and $[(nn)0(pp)I]_I$, in the $I = 2^+$, 4^+, 6^+, and 8^+ states ranges from 33% to 63%. These values, however, remain significantly large, showing that pairing is still in the game.

We now examine the influence of $V_9(np)$ on the dominance of isoscalar $J = 9$ pairs in the wave functions of ^{96}Cd. To this end, we also report in Table III the overlaps with the $[(np)9(np)9]_I$ state obtained using one-half (column b) and twice (column c) the original $V_9(np)$ value. We see that a larger value of $V_9(np)$ leads to an increase in the overlap for all the states. However, the overlap for the 8^+ state does not go beyond 27%. Moreover, it should be noted that for the three lowest-lying states the overlaps do not become significantly smaller even when $V_9(np)$ is reduced. This is related to the structure of the $[(np)9(np)9]_{I=0,2,4}$ states in terms of the $[nn] \otimes [pp]$ basis [see Eq. (1)]. Therefore, as pointed out in Ref. [4], both dynamics and geometry are crucial to the presence of $J = 9$ pairs in the wave functions of ^{96}Cd.

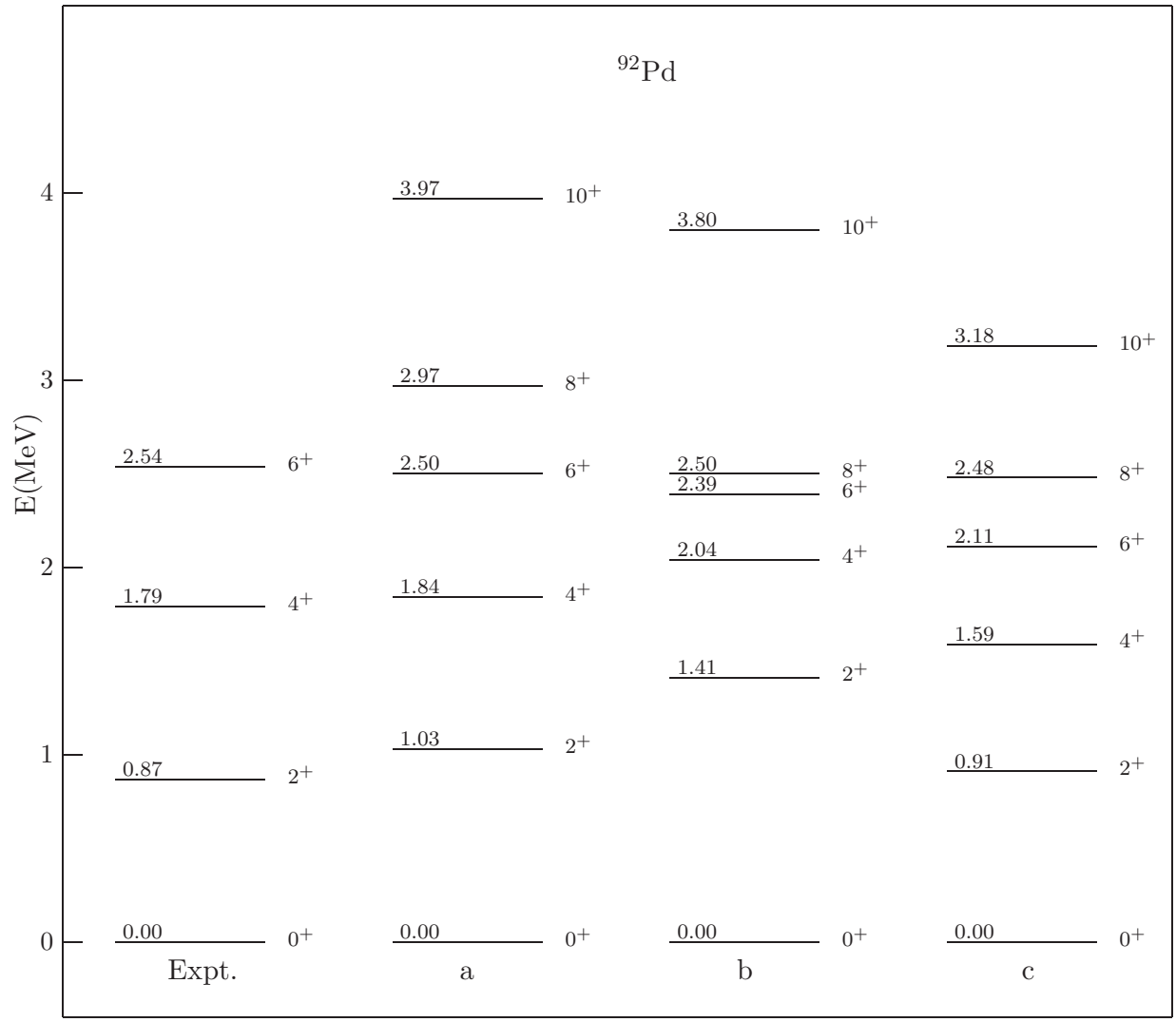

FIG. 5. Experimental spectrum of ^{92}Pd compared with the results of calculations with (a) all matrix elements (see Table I); (b) nn, pp, and np $T = 1$ matrix elements; and (c) nn, pp, and np $T = 0$ matrix elements.

B. ^{92}Pd

The simple analysis done for ^{96}Cd cannot be performed for ^{92}Pd with four neutron and four proton holes. In this case, we first discuss the effects of the $T = 0$ and $T = 1$ components of the interaction on the calculated spectrum. These effects were studied by Cederwall *et al.* using their empirical Hamiltonian [1], and we found it worth verifying whether a realistic effective interaction would confirm their results.

In Fig. 5, the results of the full-interaction calculation [Fig. 5(a)] are compared with those obtained by removing separately the $T = 0$ [Fig. 5(b)] and $T = 1$ [Fig. 5(c)] np matrix elements. We see that in Fig. 5(b) the excited states up to $I^\pi = 8^+$ are compressed in a smaller energy interval, about 1 MeV, compared with 2 MeV for the full calculation. Actually, the spectrum of ^{92}Pd evolves toward that of the neutron closed-shell nucleus, as was the case for ^{96}Cd when using a reduced value of $V_9(np)$. On the other hand, when we exclude the $T = 1$ np matrix elements, all the excited levels move down, but the spectrum keeps the same structure as that obtained from the full calculation.

Our findings are in line with those of Ref. [1], confirming the more relevant role of the $T = 0$ versus the $T = 1$ np component. More specifically, we have verified that the addition of the sole $J = 9$ np matrix element to the interaction between identical particles is sufficient to produce a spectrum very similar to that in Fig. 5(c). We may therefore conclude that the structural makeup of the ^{92}Pd spectrum is mainly determined by the combined action of $J = 9$ np, nn, and pp matrix elements. Needless to say, a quite distorted, highly compressed spectrum would result from ignoring the interaction between identical particles.

In this context we have tried to better understand how different isoscalar and isovector pairs contribute to produce the spectrum in Fig. 5(a). To this end, we have used the relation

$$E_I = \sum_J \left[C_J^I(np)V_J(np) + C_J^I(pp)V_J(pp) + C_J^I(nn)V_J(nn)\right], \quad (2)$$

TABLE II. Overlap of the calculated $I^\pi = 0^+, 2^+, 4^+, 6^+, 8^+$, and 10^+ yrast states in ^{96}Cd with the $[(nn)J_n(pp)J_p]_I$ states, expressed as a percentage. Only components with a percentage >10 are reported.

	(J_n, J_p)									
I^π	(0,0)	(0, J)	(J, 0)	(2,2)	(2,4)	(4,2)	(2,8)	(8,2)	(4,6)	(6,4)
0^+	57			30						
2^+		34	32	12						
4^+		29	26	28						
6^+		33	26		16	15				
8^+		39	25				12			
10^+							17	15	18	17

where the energy of a given state is written in terms of the average numbers of nn, pp, and np pairs, $C^I_J(ij)$'s, defined as $C^I_J(ij) = \langle\psi_I\,(^{92}\text{Pd})\,|\,[(a^\dagger_i\,a^\dagger_j)_J \times (a_i\,a_j)_J]_0\,|\,\psi_I\,(^{92}\text{Pd})\rangle$. In Eq. (2) the matrix elements in the three different channels appear explicitly, because, as mentioned in Sec. II, our effective interaction includes the Coulomb force. For the sake of simplicity, in the following we do not distinguish among nn, pp, and np isovector pairs and take as V_J, with J even, the mean value of the three corresponding matrix elements.

In Figs. 6(a) and 6(b), we show the average number of isoscalar and isovector pairs for the six considered yrast states. We see that the curves corresponding to $J = 8$ and 9 lie significantly higher than the others, as observed in Ref. [5]. We draw attention here to the fact that the curves in Fig. 6(a) are almost flat compared to those in Fig. 6(b). This means [see Eq. (2)] that the isovector pairs contribute to the level spacings more significantly than the isoscalar ones. In particular, a main role is played by the $J = 0$ and 2 pairs, whose corresponding matrix elements are much larger than the others.

The contributions to the five spacings arising from the different pairs are reported in Table IV, where we only include values larger than 130 keV. The last two columns allow us to compare the calculated spacings in Fig. 5(a) with those obtained by summing the contributions reported in the table. We see that they do not differ significantly. In Table IV it appears that the most important contributions to the energy spacings arise from the $J = 0$ pairs, although those from $J = 2$ and 9 cannot be ignored at all, being particularly relevant for the 6-4 and 8-6 spacings. In this regard, it should be kept in mind that the size of the $J = 9$ matrix element strongly influences the structure of the wave functions and, consequently, the specific action of the interaction in the $J = 0$ and 2 channels.

In concluding this discussion, we should mention that the content of isoscalar spin-aligned pairs in the wave functions of ^{92}Pd is still significant even when $V_9(np)$ is suppressed. This, as already emphasized for ^{96}Cd, is related to geometrical features.

TABLE III. Overlap of the calculated $I^\pi = 0^+, 2^+, 4^+, 6^+, 8^+$, and 10^+ yrast states in ^{96}Cd with the $[(np)9(np)9]_I$ state, expressed as a percentage, obtained using (a) $V_9(np)$ in Table I, (b) one-half the original value of $V_9(np)$, and (c) twice the original value of $V_9(np)$.

I^π	a	b	c
0^+	90	82	96
2^+	97	94	99
4^+	85	73	94
6^+	48	27	84
8^+	6	3	27
10^+	46	13	94

IV. SUMMARY AND CONCLUDING REMARKS

In this work, we have performed a shell-model study of $N = Z$ nuclei below ^{100}Sn that can be described in terms of the single $g_{9/2}$ orbit. The effective interaction for this orbit has been derived from the CD-Bonn NN potential without using any adjustable parameter. This approach reproduces very well the excited states of ^{92}Pd observed in a recent experiment and, therefore, provides support for our predictions for ^{96}Cd and ^{98}In.

Aside from the intrinsic interest in employing a realistic effective interaction to describe these neutron-deficient nuclei, the present work was motivated by the suggestive interpretation given in Ref. [1] of almost-equidistant levels in the ^{92}Pd spectrum. It was argued, in fact, that this feature may be traced to an isoscalar spin-aligned np coupling scheme replacing the normal isovector $J = 0$ pairing which is dominant for

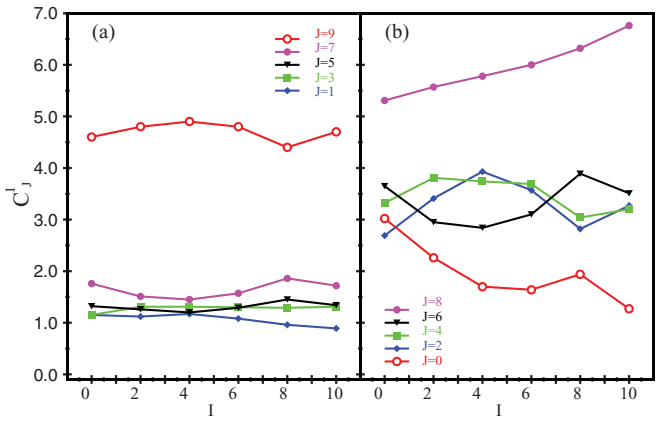

FIG. 6. (Color online) Average number of (a) isoscalar and (b) isovector $(g_{9/2})^2 J$ pairs, C^I_J, as a function of the angular momentum I of the yrast states in ^{92}Pd.

TABLE IV. Contributions (in MeV) of isoscalar and isovector pairs to the energy level spacings of ^{92}Pd. Values <0.13 are omitted. Column S gives the sum of contributions; column ΔE, the energy spacing obtained from the full calculations. See text for details.

I_i^π-I_j^π	$J=0$	$J=2$	$J=8$	$J=1$	$J=9$	S	ΔE
2^+-0^+	1.62	−0.40			−0.28	0.94	1.03
4^+-2^+	1.19	−0.29			−0.14	0.76	0.81
6^+-4^+	0.13	0.20		0.13	0.14	0.60	0.66
8^+-6^+	−0.64	0.42		0.18	0.56	0.52	0.47
10^+-8^+	1.42	−0.26	0.13		−0.42	0.88	1.00

like-valence particle nuclei. We thus decided to investigate the role played by the np interaction in the $J=9$, $T=0$ channel.

As regards ^{96}Cd, we have found that the matrix element $V_9(np)$ has a direct influence on energies of the high-spin states 12^+, 14^+ and 16^+, making the latter isomeric, while for the states with I^π from 0^+ to 10^+ it also substantially affects the structure of the wave functions. More precisely, a more attractive matrix element leads to a larger content of $J=9$ pairs and, for states up to $I^\pi=8^+$, to a larger fragmentation in terms of the $[nn]\otimes[pp]$ basis. It should be mentioned, however, that the large content of $J=9$ pairs for $I^\pi=0^+$, 2^+, and 4^+ arises also from the significant overlap of the $[(np)9(np)9]_I$ component with $[(nn)0(pp)I]_I$ and $[(nn)I(pp)0]_I$.

In agreement with previous papers, we have found that the $J=9$ matrix element plays an important role in determining the low-energy spectrum of both ^{96}Cd and ^{92}Pd. However, as shown in some detail for ^{92}Pd, it does not contribute to the energy spacings directly but, rather, by changing the weights of the isovector contributions through the induced fragmentation. So to speak, the pairing force pushed out the door comes back through the window.

In summary, we have confirmed the relevant role of isoscalar spin-aligned coupling in the low-energy states of ^{92}Pd and ^{92}Cd. Based on our study and the data presently available for ^{92}Pd, we feel, however, that one can hardly speak of a new phase of nuclear matter similar to the well-known one induced by the strong pairing correlations between identical particles.

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