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

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We have performed shell-model calculations for nuclei below $^{100}$Sn, focusing attention on the two $N = Z$ nuclei $^{96}$Cd and $^{92}$Pd, the latter having been recently the subject of great experimental and theoretical interest. We have considered nuclei for which the $0_{g_{9/2}}$ orbit plays a dominant role and employed 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 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.

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 since the $N = Z$ nuclei play a special role in understanding the nuclear effective interaction, in particular the interplay and competition between the 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 $^2$ $^5$, is the existence of strongly correlated $T = 0$ $np$ pairs, similarly 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 mainly from $J = 9$ $np$ pairs, which has been seen as evidence for a spin-aligned $np$ paired phase.

The results of $^1$ have immediately attracted much attention and the role of the isoscalar $np$ pairs in the low-energy structure of the $N = Z$ 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 the $N = Z$ nuclei with four, six and eight holes below $^{100}$Sn.

The relevance of the isoscalar component of the $np$ interaction has been stressed in a subsequent work $^7$, where the $16^{+}$ “spin-gap” isomer in $^{96}$Cd has been identified and its origin explained as due to 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 $0_{g_{9/2}}$ shell. This situation is of course reminiscent of the so-called $f_{7/2}$ nuclei that have been the subject of a large number of theoretical studies ever since the seminal paper by McCullen, Bayman and Zamick $^8$. It seems therefore appropriate to call these special nuclei below $^{100}$Sn “$g_{9/2}$ nuclei”.

Some ten years ago, we performed $^8$ 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 $^9$. Furthermore, high-precision potentials have been constructed which fit the $pp$ and $np$ scattering data with a $\chi^2$/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 $|np| \otimes |np|$ coupling schemes. The comparison between these two approaches is instrumental to understand the role of the $J = 9$ $np$ pairs. It may also be worth recalling that within the $f_{pg}$ space the choice of the single-hole energies is not an easy task $^8$, since 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. Sec. IV provides a summary and concluding remarks.

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

We assume $^{100}$Sn as 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, that as all modern NN potentials contains high-momentum non-perturbative 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 force acts only through $\hat{Q}$ for $\text{np}$ and $\text{nn}$ matrix elements, relative to the $\text{np}$ one. From 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 of Ref. [4] while this is not the case for interactions defined in the $f_{7/2}$ space, as for instance that developed in [10].

![FIG. 1: Calculated proton hole-neutron hole multiplet in $^{98}$In.](image1)

The $\text{np}$ 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 has unfortunately no experimental counterpart, so cannot be used to test our interaction. We see, however, that the $\text{np}$ 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 the energy values which is due to the greater attractiveness of the $\text{np}$ interaction in lighter systems.

![FIG. 2: (Color on line) Experimental and calculated proton hole-neutron particle multiplet in $^{90}$Nb.](image2)

For a test of our $\text{np}$ 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

| $J$ | $T$ | $pp$ | $mn$ | $np$ |
|-----|-----|------|------|------|
| 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 |

We show 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 $T = 1 \text{ np}$ matrix elements, they differ only by at most 100 keV from the $nn$ ones. This reflects the fact that for $\text{np}$ and $\text{nn}$ valence holes/particles the Coulomb force acts only through $\hat{Q}$-box diagrams starting at second and third order, respectively. From 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 of Ref. [4] while this is not the case for interactions defined in the $f_{7/2}$ space, as for instance that developed in [10].

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matrix elements of the hole-hole \( np \) interaction through the Pandya transformation \[18\]. Note that several \( 1^+ \) states have been observed in \( ^{90}\text{Nb} \) at low excitation energy. Following the suggestion of Ref. \[4\], we report the fifth one at 2.126 MeV while exclude the observed \( 0^+ \) state at 5.008 MeV, which is too high in energy to make trustable our interpretation in terms of the single \( g_{9/2} \) model. Fig. 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.

![FIG. 3](image)

FIG. 3: Experimental and calculated spectra of the \( N = 50 \) isotones \( ^{98}\text{Cd}, \ 97\text{Ag}, \) and \( ^{96}\text{Pd} \).

We now focus attention on the \( pp \) matrix elements by comparing the spectra of the three \( N = 50 \) isotones \( ^{98}\text{Cd}, \ 97\text{Ag}, \) and \( ^{96}\text{Pd} \) with the experimental ones. We include all observed levels for \( ^{98}\text{Cd} \) while positive-parity yrast levels below 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 \( ^{96}\text{Pd} \) 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 of \[1\], where the spectrum of \( ^{96}\text{Pd} \) up to the \( 10^+ \) state was calculated in the \( fpg \) model space.

### III. RESULTS

The results obtained for \( ^{90}\text{Nb} \) and \( N = 50 \) isotones have given us confidence in our effective interaction, at least as regards its predictive power for low energy states. We have then performed calculations for the two \( N = Z \) nuclei \( ^{96}\text{Cd} \) and \( ^{92}\text{Pd} \). The calculated spectra are reported in Figs. 4 and 5a, respectively, together with a comparison with the experimental data of \[4\] for \( ^{92}\text{Pd} \). In view of the recent experimental finding \[17\] we have reported yrast states in \( ^{96}\text{Cd} \) up to \( I^\pi = 16^+ \), while for \( ^{92}\text{Pd} \) we have included two more states with respect to those identified in \[1\].

![FIG. 4](image)

FIG. 4: Calculated spectrum of \( ^{96}\text{Cd} \).

![FIG. 5](image)

FIG. 5: Experimental spectrum of \( ^{92}\text{Pd} \) compared with the results of calculations with: a) all matrix elements, see Table 4 b) \( nn, pp, np \ T = 1 \) matrix elements; c) \( nn, pp, np \ T = 0 \) matrix elements.

The three observed states in \( ^{92}\text{Pd} \) are very well re-
produced 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 [1]. 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 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 shall discuss in more detail the results for both nuclei in the two next subsections.

### 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 shown in Table I while their overlaps with the $[(np)(np)]_I$ state are given in Table III.

**Table I**: Overlap of the calculated $I^\pi = 0^+, 2^+, 4^+, 6^+, 8^+$, and $10^+$ yrast states in $^{96}$Cd with the $[(nn)_{J_1}(pp)_{J_2}]_I$ states, expressed in percentage. Only components with percentage $> 10$ are reported.

| $I^\pi$ | $(J_n, J_p)$ | a | b | c |
|---------|---------------|---|---|---|
| 0$^+$   | (0,0)         | 57 | 30 |   |
| 2$^+$   | (0,0)         | 34 | 32 | 12 |
| 4$^+$   | (2,2)         | 29 | 26 | 28 |
| 6$^+$   | (2,8)         | 33 | 26 | 16 |
| 8$^+$   | (4,6)         | 39 | 25 | 12 |
| 10$^+$  | (4,6)         | 46 | 13 | 94 |

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

$$|(np)_{J_1}(np)_{J_2}; I\rangle = \frac{1}{\sqrt{N_{J_1J_2}}} \sum_J \langle \hat{J}_1 \hat{J}_2 \hat{n} \hat{p}_I \rangle^{1/2} \left\{ \begin{array}{ccc} 0 & 0 & \frac{9}{2} \\ \frac{9}{2} & \frac{9}{2} & J_1 \\ \frac{9}{2} & \frac{9}{2} & J_2 \end{array} \right\} |(nn)_{J_1}(pp)_{J_2}; I\rangle, \tag{1}$$

where $|\hat{J}\rangle = (2J + 1)$ and $N$ denotes the normalization factor. From Table I 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 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 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 $[(np)(np)]_I$ 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 I is of course related to the $np$ interaction and in this connection it is interesting to find out what is the role of the $J = 9$ matrix element $V_{9}(np)$. To this end, we have redone 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 [5, 6].

We find that for the reduced value of $V_{9}(np)$ the $6^+$ and $8^+$ states come down 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 $10^+$ states. The spectrum of $^{96}$Cd (up to $I^\pi = 8^+$) becomes then similar to that of $^{98}$Cd, the wave ef

**Table III**: Overlap of the calculated $I^\pi = 0^+, 2^+, 4^+, 6^+, 8^+$, and $10^+$ yrast states in $^{96}$Cd with the $[(np)(np)]_I$ state, expressed in percentage, obtained using: a) $V_{9}(np)$ of Tab. I; b) one-half the original value of $V_{9}(np)$; c) twice the original value of $V_{9}(np)$.
functions being, however, still 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 $(nn)00(pp)00$ component in the ground state reduces to 51%, while that of the seniority-two components, $(nn)I(pp)0I$ and $(nn)00(pp)I$, 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 have also reported in Table III the overlaps with the $(np)9(np)9 I$ state obtained using one-half (b) and twice (c) the original $V_9(np)$ value. We see that a larger value of $V_9(np)$ leads to an increase of 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 already pointed out in [4], both dynamics and geometry are crucial to the presence of $J = 9$ pairs in the wave functions of $^{96}$Cd.

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 shall first discuss the effects of the $T = 0$ and $T = 1$ components of the interaction on the calculated spectrum. This study was already performed by Cederwall et al. [1], and we have found it worth verifying if a realistic effective interaction confirms their results.

In Fig. 5 the results of the full-interaction calculation (a) are compared with those obtained by removing separately the $T = 0$ (b) and $T = 1$ (c) $np$ matrix elements. We see that in case (b) the excited states up to $I^\pi = 8^+$ are compressed in a smaller energy interval, about 1 MeV to be compared with 2 MeV of the full calculation. Actually, the spectrum of $^{92}$Pd evolves toward that of the neutron closed shell nucleus, as it 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 of Fig. 5. 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 of Fig. 5a. To this end, we have used the relation

$$E_I = \sum_j \left[ C_j^I(np)V_J(np)+C_j^I(pp)V_J(pp)+C_j^I(nn)V_J(nn) \right],$$

(2)

where the energy of a given state is written in terms of the average numbers of $nn$, $pp$, and $np$ pairs, $C_j^I$’s, defined as $C_j^I(\delta) = \langle \psi_I^{(92)Cd} | (|a_i a_j^\dagger\rangle_{J=0} \times (a_i a_j^\dagger)_{J=1} \psi_I^{(92)Cd} >$. In Eq. (2) the matrix elements in the three different channels appear explicitly, since, as mentioned in Sec. II, our effective interaction includes the Coulomb force. For the sake of simplicity, in the following we shall not distinguish between $nn$, $pp$, and $np$ isovector pairs and take as $V_I$, with $J$ even, the mean value of the three corresponding matrix elements.

![FIG. 6: (Color on line) Average number of the isoscalar (a) and isovector (b) $(g_{9/2})^2$ $J$ pairs, $C_j^I$, as a function of the angular momentum $I$ of the yrast states in $^{92}$Pd.](image)

In Figs. 6a and 6b, we show the average number of the 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 was already observed in [3]. We draw attention here on the fact that the curves of Fig. 6b are almost flat when compared to those of Fig. 6a. 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 other ones.
TABLE IV: Contributions (in MeV) of isoscalar and isovector pairs to the energy level spacings of 92\(^{\text{Pd}}\). Values < 0.13 are omitted. Column \( S \) gives the sum of contributions while column \( \Delta E \) the energy spacing obtained from the full calculations. See text for details.

| \( J^T - I^T \) | \( J = 0 \) | \( J = 2 \) | \( J = 8 \) | \( J = 1 \) | \( J = 9 \) | \( S \) | \( \Delta E \) |
|----------------|---------|---------|---------|---------|---------|------|--------|
| 2\(^T\) - 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 |        |

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 two last columns allows to compare the calculated spacings of Fig. 5 with those obtained by summing the contributions reported in the table. We see that they do not differ significantly. From 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 structure of the wave functions, and consequently the specific action of the interaction in the \( J = 0 \) channels, is strongly influenced by the size of the \( J = 9 \) matrix element.

In concluding this discussion, we should mention that the content of isoscalar spin-aligned pairs in the wave functions of 92\(^{\text{Pd}}\) is still significant even when \( V_0(np) \) is suppressed. This, as already emphasized for 96\(^{\text{Cd}}\), is related to geometrical features.

IV. SUMMARY AND CONCLUDING REMARKS

In this work, we have performed a shell-model study of \( N = Z \) nuclei below \( ^{160}\text{Sn} \) 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\(^{\text{Pd}}\) observed in a recent experiment, and therefore provides confidence in our predictions for 96\(^{\text{Cd}}\) and 98\(^{\text{In}}\).

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 the almost equidistant levels in the 92\(^{\text{Pd}}\) 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 like valence particle nuclei. We have thus decided to investigate the role played by the \( np \) interaction in the \( J = 9, T = 0 \) channel.

As regards 96\(^{\text{Cd}}\), we have found that the matrix element \( V_0(np) \) has a direct influence on energies of the high-spin states \( 12^+, 14^+ \) and \( 16^+ \), making the latter isomeric, while for the states with \( I^T \) 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 \( I^T = 8^+ \), 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 \( I^T = 0^+ \), \( 2^+ \) and \( 4^+ \) arises also from the significant overlap of the \( [(np)9(np)]_I \) component with \( [(nn)0(pp)]_I \) and \( [(nn)(pp)]_I \).

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\(^{\text{Cd}}\) and 92\(^{\text{Pd}}\). However, as shown in some detail for 92\(^{\text{Pd}}\), 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 say, the pairing force pushed out the door comes back through the window.

In summary, we confirm the relevant role of the isoscalar spin-aligned coupling in the low-energy states of 92\(^{\text{Pd}}\) and 92\(^{\text{Cd}}\). Based on our study and the presently available data for 92\(^{\text{Pd}}\), 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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