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Anomalous magnetic moment of ^{9}C and shell quenching in exotic nuclei

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The mechanism is studied which causes an anomalous magnetic moment of ^{9}C, i.e., too large isoscalar spin for the $T_z=\pm 3/2$ mirror pair of $A=9$, from a shell-model viewpoint. Based on the empirically determined shell energies varying from stable to unstable nuclei, the effective $N=8$ shell gap for $Z=3$ is evaluated to be rather narrow in a reasonable way. Under this condition, an additional shell quenching by the Thomas-Ehrman effect, arising for ^{9}C only, allows its ground state to be substantially mixed with the intruder configurations. This mirror asymmetry in the mirror wave functions accounts for the obviously large spin expectation value.

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Isospin is a basic quantum number in nuclei, reflecting the approximate charge independence (or isospin symmetry) of the nucleon-nucleon interaction [1]. Although the Coulomb interaction can violate this symmetry, its effect is predominantly to raise overall energy and then nuclear wave functions still have a sufficiently good isospin yielding various features in the nuclear structure. It is thus of importance to investigate to what extent the isospin symmetry is valid and how its breaking, if it exists, affects the nuclear structure.

For instance, assuming the isospin symmetry, the isoscalar spin expectation value $\langle \sigma_z \rangle$ can be deduced from magnetic moments of mirror nuclei. Experimentally, most of its absolute values in odd nuclei do not exceed the single-particle estimate [2]. However, the magnetic moment of ^{9}C, measured for the first time by Matsuta *et al.* several years ago [3], has proved to lead to anomalously large $\langle \sigma_z \rangle = 1.44$ compared to the single-particle estimate of 1. The magnetic moment of ^{9}C was remeasured later [4], where this anomaly was confirmed. It is noted that the use of the effective nucleon g factors hardly solves this anomaly, partly because their isovector correction does not affect the isoscalar observables, and partly because in general the isoscalar spin correction works unfavorably (for reviews of the effective $M1$ operator, see, e.g., Refs. [5,6]).

Realistic p-shell-model calculations assuming the isospin symmetry generally give $\langle \sigma_z \rangle \sim 1$ due to strong pairing interaction bringing about the spin saturation. A certain improvement of $\langle \sigma_z \rangle$ by 0.09 was obtained by a p-shell model including isospin nonconserving terms [4], but some deviation still remains. Other models, including a microscopic cluster model [7] and the antisymmetrized molecular dynamics (AMD) [8], result in the value comparable to the p-shell model, probably because their wave functions rather resemble one another. These previous results might indicate the possibility of large mirror asymmetry, which was not sufficiently considered so far, between the ground states of ^{9}Li and ^{9}C.

In loosely bound nuclei, it is known that the so-called Thomas-Ehrman effect [9,10] gives rise to the large asymmetry in the energy levels of mirror nuclei associated with a penetrating proton single-particle wave function. In the present case, although the one-proton separation energy (S_p) of ^{9}C is as small as 1.3 MeV [11], such a large asymmetry in the wave function is scarcely expected from the conventional p-shell picture, because the p orbit is less influenced by the loosely binding effect than the s orbit owing to the presence of the centrifugal barrier. In addition, a modification of the radial single-particle wave function by the Thomas-Ehrman effect does not directly change the magnetic moment since its operator has no radial dependence. Accordingly, the origin of the large $\langle \sigma_z \rangle$ has not been cleared yet.

In this study, we investigate, from the viewpoint of a realistic shell-model, the possibility of another mechanism, i.e., the interplay of an evolving shell structure in exotic nuclei. This has not been explored in depth so far in the present context. We now include in this shell-model calculation not only the p orbits but also the full sd ones as a valence shell (i.e., full p-sd valence orbits), and our attention will be focused upon their role.

Since the dimension of the Hamiltonian matrix becomes huge toward the mid-shell, the model space is often truncated. To sufficiently include important correlations, the following calculation is performed now. In the case that only the normal states (i.e., those composed of the configurations having the lowest $\hbar\omega$) are taken into consideration, we take all the $(0+2)\hbar\omega$ and $(1+3)\hbar\omega$ basis states excited from the lowest $\hbar\omega$ for positive- and negative-parity states, respectively. The full calculation is otherwise performed. Now a multi-$\hbar\omega$ valence space is taken, the spurious center-of-mass (c.m.) motion should be removed. In the present study, a prescription by Gloeckner and Lawson [12] is adopted with β taken as $(\beta/A)\hbar\omega=100$ MeV. Unless this c.m. removal is carried out, one sees that unphysical inter-shell mixing occurs even for states regarded as p-shell dominated.

Our first task is to obtain in this framework a good description of ^{9}Li (i.e., the mirror nucleus of ^{9}C) where the ordinary shell-model approach is considered to be valid. As for the effective interaction, we start with two different interactions of the WBP and the WBT [13] both of which are often used in this region. Their p-shell and cross-shell matrix elements were determined to reproduce well 216 cross-shell energy levels [13]. Both interactions adopted the USD interaction [14] as the sd-shell part, and in the present study the strength of the USD is fixed for A=10.

Although many cross-shell properties are reproduced by those interactions, there are some reasons to introduce empirical modifications in the present context as follows. First, the original WBP and WBT interactions are designed to be used in a pure $\hbar\omega$ model space, while the present calculation incorporates the mixing with higher excited configurations across the $N=8$ shell gap (often called intruder configurations). Due to the inclusion of the intruder configurations, some changes of the interaction may be needed, and the matrix elements sensitive to the mixing should be examined. Second, energies of the Z(or N)$=3$ and 4 isotopes, which are indeed the target of the present study, are not included in the cross-shell fit of the original interactions, probably due to few experimental data.

With respect to the mixing, the $\langle pp|V|sd\ sd\rangle_{JT}$ matrix elements, particularly the pairing ones [i.e., $(J,T)=(0,1)$], are responsible because they do not conserve $\hbar\omega$. Quite recently, Shimoura *et al.* have found the 0_2^+ state of ^{12}Be [15], whose position is sensitive to the mixing. Those two interactions have almost vanishing cross-shell pairing matrix elements concerning the $1s_{1/2}$, leading to lower excitation energy than the experiment. We thus enlarge them by $\delta G=0.5$ MeV where G is defined by $\langle 0p_j\ 0p_j|V|1s_{1/2}1s_{1/2}\rangle_{J=0,T=1} = G\sqrt{(2j+1)/2}$. The 0_2^+ of ^{12}Be is located at around 2.4 MeV by our final Hamiltonian for Be (see later discussions), compared to the experimental level of 2.24 MeV [15].

We next examine the single-particle structure of $Z=3$ isotopes. It should play an important role in their nuclear structure, but has not been studied in depth by the original interactions, because in Li isotopes there is no direct experimental information about the position of the *sd* orbits such as the abnormal parity states. For $Z\geqslant 4$, however, the abnormal parity states have been observed, and we shall make effective use of this experimental information to fix the location of the *sd* orbits at $Z=3$.

Before proceeding to the concrete method, we introduce the so-called effective single-particle energy (ESPE) [16]. It varies as the proton and neutron numbers move, depending on the strength of the monopole interaction. In general, the neutron ESPE can be rather steeply changed in proportion to the proton number due to the strong proton-neutron interaction. Figure 1 shows, with thin lines, the effective $N=8$ shell gap on the top of the $\nu(0p_{3/2})^4$ core obtained by the WBP and WBT interactions. Their gaps linearly move for Z increasing from 3 to 6 where the proton occupies the $0p_{3/2}$ orbit. The energies of both the $1s_{1/2}$ and the $0d_{5/2}$ are similarly narrowed as Z goes to 3, whose general mechanism has been argued in terms of the spin-isospin dependence of the nucleon-nucleon interaction [17].

Although the $N=8$ shell gap of the WBP and WBT interactions becomes narrow for smaller Z, it turns out that a much narrower shell gap is needed in the present calculation scheme: with the original interactions used in the mixed $\hbar\omega$ configuration space, the parity inversion of ^{11}Be does not occur, and the ground state of ^{12}Be is dominated by the normal configurations bringing about much higher 2_1^+ level than the experimental one. We emphasize that the need for the narrower $N=8$ shell gap at $Z=4$ does not mean a defect of the original interactions but is due to the difference of the adopted model space, since in the pure $0\hbar\omega$ and $1\hbar\omega$ space (to be used in the original interactions) they reproduce the

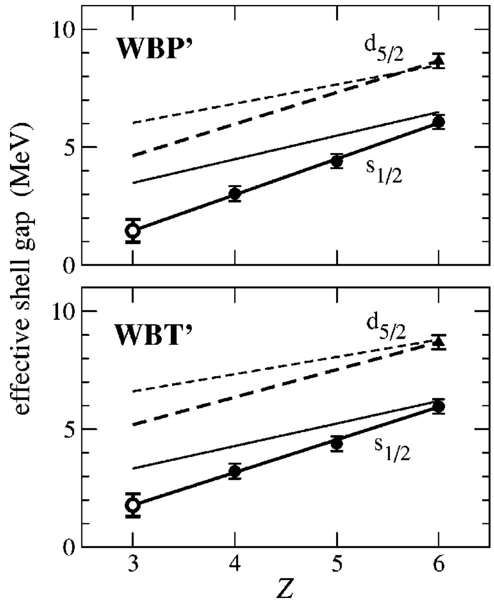

FIG. 1. Effective neutron single-particle energies of the $1s_{1/2}$ (solid line) and the $0d_{5/2}$ (dashed line) on the top of the $N=6$ $(\nu(0p_{3/2})^4)$ core, measured from the $0p_{1/2}$. The thin lines are those of the original WBP and WBT interactions, while the thick ones are drawn to fit experimental levels as presented by filled symbols. See the text in detail.

parity inversion of ^{11}Be. We now assume the ^{4}He core to be inert for both parity states equally, and include the mixing with configurations excited across the $N=8$ gap.

We then empirically search for the ESPE suitable to the present model space on an isotope by isotope basis as follows. First, the single-particle energies of the $1s_{1/2}$ and the $0d_{5/2}$ are determined to reproduce the $1/2_1^+$ and $5/2_1^+$ levels of ^{13}C, and the $0d_{3/2}$ is fixed to obtain experimental $5/2_1^+$-$3/2_1^+$ splitting of ^{17}O. Next, the ESPE of the *sd* orbits for $Z=4$ and 5 is adjusted to fit the experimental abnormal (negative) parity levels of $N=7$ isotones. There are two similar methods for this purpose, i.e., the adjustment of the monopole interactions or the single-particle energies, and we now take the latter for simplicity. This adjustment of the *sd* orbits is controlled by one parameter: all the ESPE's of the *sd* orbits are equally shifted from those of the C isotopes. For $Z=4$, the energy difference between the $1/2_1^+$ and the $1/2_1^-$ is used in the adjustment. As for $Z=5$, in order to weaken the effect by a specific proton-neutron coupling matrix element, we consider the angular-momentum averaged energy defined by $\Sigma_J(2J+1)E_J/\Sigma_J(2J+1)$. As positive- and negative-parity energies, we, respectively, take $J=1_1^+, 2_1^+$ and $J=1_1^-, 2_1^-$.

The empirically adjusted ESPE of the $1s_{1/2}$ is presented by the filled circles in Fig. 1. The error bar in Fig. 1 is determined to obtain the experimental energy within accuracy of 0.3 MeV which is a typical deviation of the realistic shell model [13]. The agreement between the original and the empirical Hamiltonians is good at $Z=6$ as expected, whereas the difference between them is enhanced toward smaller Z: the shell gap should vary more transparently in the sufficiently extended model space. Note that this situation is similar to the quenching of the $N=20$ shell gap for neutronrich nuclei discussed with the Monte Carlo shell model calculation [16]. Those three empirical ESPE's are almost linearly placed for both interactions as shown in Fig. 1, being consistent with the general trend of the ESPE given by the shell model. Therefore, it is most likely that the ESPE of Li isotopes with $Z=3$ is located at values where those ESPE's are extrapolated. Figure 1 also shows, with the unfilled circles, the extrapolated ESPE at $Z=3$ with the χ-square fitting. The estimated error bar at $Z=3$ is $\sim$0.5 MeV. Hereafter, the Hamiltonians with the extrapolated ESPE are referred to as WBP' and WBT'. The WBP' and WBT' Hamiltonians are hereafter used for $N \geqslant Z$ nuclei, while for $N < Z$ nuclei the Hamiltonians of their mirror partner are used to keep the mirror symmetry at this stage.

In the calculation of the magnetic moment, the free-nucleon g factors are adopted. When a complete major shell is taken as valence orbits, the magnetic moment is usually fairly well described by them (see, e.g., [14] for the sd-shell calculation). Under the assumption that the mirror symmetry is valid between two mirror states, their spin expectation value is deduced [2] as

$$\langle \sigma_z \rangle = \frac{\mu(T_z = +T) + \mu(T_z = -T) - J}{(g_s^{\mathrm{p}} + g_s^{\mathrm{n}} - 1)/2}, \quad (1)$$

where μ and J are their magnetic moment and the total angular momentum, respectively. The $g_s^{\mathrm{p,n}}$ are the spin g factors of the proton and the neutron, respectively, producing the denominator of Eq. (1) to be 0.3796 with the free-nucleon ones. Regardless of the validity of the mirror symmetry, this expression for $\langle \sigma_z \rangle$ is hereafter used to easily compare the result with the experiment.

The ground state of ^{9}Li calculated with the WBP' (and WBT') is dominated by the normal state even with the narrowing shell gap: the probability of the $0\hbar\omega$ configurations is 88%(88%) for the WBP' (WBT'). Assuming the mirror symmetry, the $\langle \sigma_z \rangle$ is calculated to be 1.02 (1.01) comparable to the $0\hbar\omega$ value of 1.03 (1.02) with the same interaction. Namely, as far as the wave function with the mirror symmetry is adopted, the mixing with the sd-shell configurations does not account for the experimental $\langle \sigma_z \rangle$. Note that a similar conclusion was drawn in a recent shell-model study [18] with the same valence shell. In order to confirm that the present interactions do not give a too narrow $0p_{3/2}$-$1s_{1/2}$ gap, it would be helpful to calculate the $1/2_1^+$ level of ^{9}Li. It is located as high as at 4.8(5.0) MeV with the WBP' (WBT') interaction, being consistent with the experimental fact that there has not found a low-lying positive-parity state in ^{9}Li.

While ^{9}Li is a moderately bound nucleus with $S_n = 4.06$ MeV, the S_p of ^{9}C is only 1.30 MeV [11]. Thus, the proton single-particle levels for ^{9}C, in particular the $1s_{1/2}$, can be shifted from those of neutrons for ^{9}Li by the Thomas-Ehrman effect [10]. Indeed, such a shift has been experimentally studied for the ^{11}Be-^{11}N mirror pair in this neighborhood [19]: the $1/2^+$ and the $5/2^+$ levels of ^{11}N are, respectively, lowered by 0.4 and 0.1 MeV from the $1/2^-$ state, and the shift was studied theoretically by the complex scaling method [20]. In the shell-model, these energy shifts are reproduced by lowering the $1s_{1/2}$ energy by 0.6 MeV.

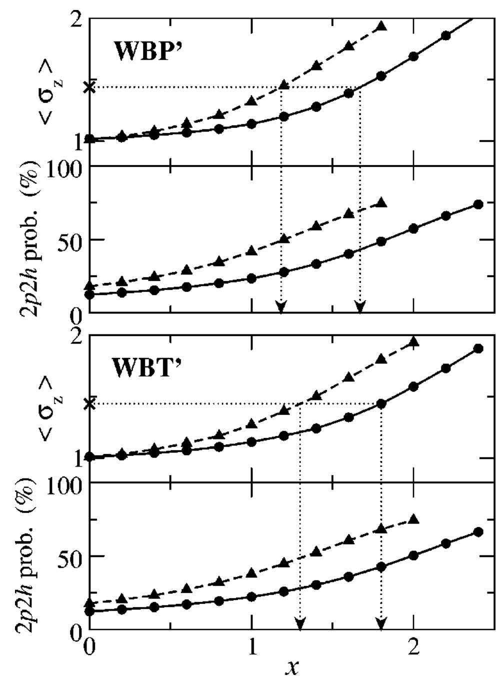

FIG. 2. Spin expectation value of the ^{9}Li-^{9}C mirror pair and the $2\hbar\omega$ probability in the ground state of ^{9}C as a function of x by the WBP' and WBT' interactions. As the ESPE of ^{9}Li, we adopt and compare two values: one is the extrapolated value in Fig. 1, i.e., the unfilled circle in Fig. 1, and the other is the lower limit of its error bar. The circles and the triangles correspond to the former and the latter ESPE's, respectively. The crosses denote the experiment (Refs. [3,4]).

The $0p_{3/2}$-$0p_{1/2}$ energy difference can be varied also since only the $0p_{1/2}$ of ^{9}C will not be bound, and its shift is evaluated to be 0.5 MeV by the Woods-Saxon potential with the standard potential parameters [1] for the ^{8}Li and ^{8}B cores. We can make use of these values to estimate the single-particle levels of ^{9}C: the ESPE of ^{9}C is shifted from that of ^{9}Li with a parameter x as $\delta(\epsilon(1s_{1/2}) - \epsilon(0p_{3/2})) = 0.6x$ MeV and $\delta(\epsilon(0p_{3/2}) - \epsilon(0p_{1/2})) = 0.5x$ MeV, where δ denotes the difference between ^{9}Li and ^{9}C. Note that the precise evaluation of the shift would need to incorporate the many-body effect, which is rather difficult to carry out though, but that this kind of the simple single-particle evaluation does not give so large deviation when adopted around the ^{16}O core.

Since the ESPE of ^{9}Li is somewhat ambiguous due to the extrapolation and model errors (see Fig. 1 and related discussions), we take two sets of the ESPE for ^{9}Li, i.e., the central value and the lower limit in Fig. 1. Once a fixed ESPE for ^{9}Li is taken, the ESPE of ^{9}C is shifted from it with the varying parameter x. In Fig. 2, we show the effect of the narrowing shell gap by the Thomas-Ehrman shift on the magnetic moment. The $\langle \sigma_z \rangle$ increases as x is larger, and the experimental value is reproduced between $x=1$ and 2 by both interactions. In the case of the lower ESPE, the fitted x is rather close to unity where the Thomas-Ehrman shift is estimated by the simple single-particle picture as discussed above. As Fig. 2 shows, this agreement with the experiment coincides, in the present model, with a large mixing ($\sim$40%) with the intruder configurations in ^{9}C caused by the

TABLE I. Expectation value of the angular-momentum operators for the ground state of ^{9}C, compared between the single-particle $\nu(0p_{3/2})^1$ state (SP), the $0\hbar\omega$ and $2\hbar\omega$ shell-model calculations without the mixing, the $(0+2)\hbar\omega$ mixed ones, and the experiment (Exp.). For the $(0+2)\hbar\omega$ calculations, the values at $x=0$ (i.e., without the Thomas-Ehrman effect) and the optimum x for $\langle\sigma_z\rangle$ are presented. The WBP' is used.

	SP	$0\hbar\omega$	$2\hbar\omega$	$(0+2)\hbar\omega$		Exp.
				$x=0$	$x=$optimum	
$\langle l_z^{\mathrm{p}}\rangle$	0	0.14	0.83	0.24	0.44	
$\langle l_z^{\mathrm{n}}\rangle$	1	0.84	0.15	0.75	0.56	
$\langle s_z^{\mathrm{p}}\rangle$	0	0.02	0.02	0.02	0.01	
$\langle s_z^{\mathrm{n}}\rangle$	0.5	0.50	0.50	0.49	0.49	
$\langle\mu\rangle$	−1.91	−1.66	−0.95	−1.54	−1.38	(−)1.3914(5)[a]
						(−)1.396(3)[b]

[a]Taken from Ref. [3].
[b]Taken from Ref. [4].

proton shell quenching. This narrowing shell gap against stable nuclei is produced predominantly by the character of the nucleon-nucleon interaction affecting both mirror nuclei (see Fig. 1), and the additional Thomas-Ehrman effect working only on ^{9}C triggers the mixing.

We next discuss how the large mixing of the intruder configurations in ^{9}C can account for the anomalous spin expectation value. Table I shows the calculated magnetic moment of ^{9}C decomposing into its orbital and spin components with the WBP' interaction. To clarify the effect of the intruder state, the pure $0\hbar\omega$ and $2\hbar\omega$ calculations are also presented as well as the $(0+2)\hbar\omega$ ones with and without the Thomas-Ehrman effect. In the single-particle model, the total angular momentum $J=3/2$ is decomposed into the neutron orbital angular momentum 1 and spin 1/2 leading to $\langle\sigma_z\rangle=1$. Due to the strong pairing correlation all the calculations give the spin part quite close to one another, from which the conventional shell-model calculation results in an almost equal spin expectation value. Thus, the orbital part seems to play an essential role, which we shall examine in detail. We first consider two limit cases, i.e., the $0\hbar\omega$ and $2\hbar\omega$ lowest states without the mixing. The $0\hbar\omega$ calculation gives the distribution of the orbital angular momentum basically similar to the single-particle model, although little fraction of the orbital angular momentum is taken by the protons. On the other hand, the $2\hbar\omega$ lowest state carries a large proton orbital angular momentum. This happens because the extensive valence space allows the nucleus to be strongly deformed, yielding the orbital angular-momentum distribution nearly proportional to the number of the valence nucleons (4:1 in this case). Since g_l^{p} is positive, the magnetic moment moves positively beyond the experimental value.

When the large mixing occurs due to the narrowing shell gap, one can expect that the expectation values of the angular-momentum operator are located in between. In fact, in the case that the Thomas-Ehrman effect is switched on (see $x=$optimum in Table I), the orbital angular momentum is approximately halved by protons and neutrons. This behavior reflects that the $2\hbar\omega$ component of the mixed ground state predominantly consists of the lowest state of the $2\hbar\omega$ calculation: the overlap probability of the mixed ground state with the lowest $2\hbar\omega$ eigenstate reaches 36%, accounting for 86% of the total $2\hbar\omega$ component. Since the ground state of ^{9}Li corresponds to $x=0$ by exchanging between the proton and neutron, there is a large mirror asymmetry in the angular-momentum distribution. The resulting $\langle\sigma_z\rangle$ by Eq. (1) then becomes seemingly large due to the mirror asymmetry, while the actual $\langle\sigma_z\rangle$ calculated directly from the wave function is still normal. The property about the angular-momentum distribution in Table I is considered to be basically independent of the detail of the shell-model interaction.

Finally, it may be interesting to compare the ^{9}Li-^{9}C pair with another Z (or N)=3 mirror pair of ^{7}Li-^{7}Be. The magnetic moment of ^{7}Be has been recently measured to be $-1.398(15)\mu_N$ [21] quite close to that of ^{9}C, while the experimental magnetic moment of ^{7}Li ($\mu=+3.256\mu_N$) is somewhat different from that of ^{9}Li ($\mu=+3.439\mu_N$) [11]. As a result, the experimental $\langle\sigma_z\rangle$ value of the ^{7}Li-^{7}Be pair is a normal one, 0.94. A large difference between those two pairs can be seen in the binding property: the experimental S_n of ^{7}Li and S_p of ^{7}Be are, respectively, 7.25 and 5.61 MeV [11], which are rather close to each other and larger than the S_n of ^{9}Li (4.06 MeV). In this relatively well bound case, effects of the *sd* orbits should be quite small when these orbits lie as loosely-bound or resonance states. Furthermore, for the ^{7}Li-^{7}Be pair all nonspurious c.m. states with the $2\hbar\omega$ excitation involve the excitation to the *pf* shell or the breaking of the ^{4}He core (out of the present model space) whose excitation energy should be rather high. On the other hand, for the ^{9}Li-^{9}C pair, many nonspurious c.m. states with the $2\hbar\omega$ excitation can be constructed within the present *p*-*sd*-shell space: additional two nucleons in the *p* shell can block the *sd*- to *p*-shell step-down operation by B_μ (see, e.g., [22] in detail) due to the Pauli principle. Thus, for the $A=7$ pair, it is reasonable to calculate the ground state within the *p* shell. Even though the difference of the single-particle energies between the $0p_{3/2}$ and the $0p_{1/2}$ is narrowed by 0.5 MeV for ^{7}Be (a typical value, but it may be overestimated because of $\Delta Z=1$), the magnetic moment is changed by only $0.02\mu_N$ with the WBP' interaction. The calculated $\langle\sigma_z\rangle$ value is thus still a nearly normal one, 1.05. On the other hand, the isovector moment of the ^{7}Li-^{7}Be pair defined as $(\mu(T_z=+T)-\mu(T_z=-T))/2$ is calculated to be smaller by $0.11\mu_N$ than the experimental one, which may indicate the need for the refinement of the isovector g factors.

In conclusion, we investigated with a realistic shell-model calculation the mechanism which causes the anomalous isoscalar spin expectation value deduced from magnetic moments of the ^{9}Li-^{9}C mirror pair. The shell structure evolving from stable to unstable nuclei was carefully examined. It is pointed out that the $N=8$ shell gap for $Z=3$ can be empirically well determined from its trend of $Z=4-6$ isotopes and should be rather narrow. With this shell structure, the ground state of ^{9}Li is still dominated by the *p*-shell configurations, whereas its mirror nucleus ^{9}C would have the ground state considerably mixed with the intruder configurations due to the additional shell quenching by the Thomas-Ehrman effect. This mirror asymmetry can account for the anomalous spin expectation value: in the present calculation, about 40% mixing in ^{9}C accounts for the experiment. The mirror asymmetry which has been studied so far was mainly just a modification of the single-particle wave function in a many-body state. In the present case, on the other hand, the component of many-body configurations is changed too, which enables the isoscalar magnetic moments to be shifted substantially. This kind of the Thomas-Ehrman effect, referred to now as "Thomas-Ehrman mixing," may inspire theoreticians with exploring a more sophisticated approach to reconcile the single-particle property and the many-body correlation.

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