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Abstract

Following the reports of candidate chiral doublet bands observed in cesium isotopes, the possible chiral candidates and the evolution of three-dimensional rotation in $^{120-134}textrmCs$ are investigated within the microscopic three-dimensional tilted axis cranking covariant density functional theory (3DTAC-CDFT). By investigating the evolution of the polar angle $ theta$ and azimuth angle $ varphi$ as a function of rotational frequency $ hbar omega$, the transition from the planar rotation to the chiral rotation has been found in $^{121-133}textrmCs$. The corresponding critical rotational frequency $ omega_{ textrm{crit}}$ of the appearance of chiral aplanar rotation decreases as neutron number increases, which can be attributed to the neutrons in (gd) and (sd) shells having smaller angular momentum components along both the short and long axes, and larger components along medium axis, respectively. In comparison, only planar rotation has been obtained in $^{120,134}textrmCs$. With these interpretations, the obtained $Isimhbaromega$ and energy spectra as well as $B(M1)/B(E2)$ values show reasonable agreement with the available experimental data. In addition, the evolution of quadrupole deformation $ beta$ and triaxial deformation $ gamma$ are also discussed.

Full Text

Preamble

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The evolution of the chiral symmetry in cesium isotopes

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Abstract

Following reports of candidate chiral doublet bands observed in cesium isotopes, we investigate possible chiral candidates and the evolution of three-dimensional rotation in $^{120-134}$Cs using microscopic three-dimensional tilted axis cranking covariant density functional theory (3DTAC-CDFT). By examining the evolution of the polar angle and azimuth angle as functions of rotational frequency $\hbar\omega$, we find a transition from planar to chiral rotation in $^{121-133}$Cs. The corresponding critical rotational frequency ω_{crit} for the onset of chiral aplanar rotation decreases with increasing neutron number, which can be attributed to neutrons in the (gd) and (sd) shells having smaller angular momentum components along both the short and long axes, and larger components along the medium axis, respectively. In contrast, only planar rotation is obtained in ^{120},^{134}Cs. With these interpretations, the obtained I $\hbar\omega$ and energy spectra as well as B(M1)/B(E2) values show reasonable agreement with available experimental data. Additionally, the evolution of quadrupole deformation β and triaxial deformation γ are discussed.

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1 Introduction

Chirality has attracted widespread interest across natural sciences, including chemistry, biology, and physics. In nuclear physics, chirality—i.e., spontaneous chiral symmetry breaking—was first proposed by Frauendorf and Meng in 1997 [1]. In a triaxially deformed nucleus with unpaired particle(s) and hole(s) in high-j orbitals, the particle(s) and hole(s) align their angular momentum vectors along the short (s) and long (l) axes, respectively, while the collective core aligns along the medium (m) axis. The total angular momentum vector then lies outside the three principal planes in the intrinsic frame (aplanar rotation) with left- or right-handed orientations, forming a chiral system. Restoration of the spontaneous chiral symmetry breaking in the laboratory frame could give rise to observation of so-called chiral doublet bands—a pair of nearly degenerate ΔI = 1 bands with the same parity [1]. Indeed, a pair of ΔI = 1 bands found in ^{134}Pr with the $\pi h_{11}/_2$ $h_{11}/_2$ configuration, observed earlier in 1996 [2], was reinterpreted in Ref. [1] as a candidate for chiral doubling. Thereafter, in 2001, similar low-lying doublet bands were reported in ^{55}Cs, ^{57}La, and ^{61}Pm—N = 75 isotones of ^{134}Pr—and an island of chiral rotation was suggested in the A 130 mass region [3]. It should be noted that clearly identifying chiral motion is difficult; for example, the chirality in ^{134}Pr mentioned above was later rejected in 2006 [4,5], and in 2011 other chiral candidates in ^{134}Pr were proposed [6].

Since this pioneering work, numerous experimental and theoretical studies related to nuclear chirality have been conducted (see Refs. [7–10] for brief reviews). Nuclear chirality has thus become a hot topic, with related issues such as M D (multiple chiral doublet bands in one nucleus) [11–15], nuclear Chirality-Parity (ChP) violation [16,17], and others being widely discussed.

To date, more than 60 candidate chiral nuclei have been reported in the A 80, 100, 130, 160, and 190 mass regions of the nuclear chart (see, e.g., data tables [18]). In particular, the reported candidate chiral nuclei in the A 130 mass region form a large chiral island, where cesium isotopes have the most chiral candidates. So far, chiral doublet bands have been observed in the odd-odd nuclei ^{122,124,126,128,130,132}Cs [3,19–24] with configuration $\pi h_{11/2}$ $h_{11/2}{}^{-1}$, and corresponding chirality has been well explained by different theoretical models [19,20,23–38]. In fact, chiral doublet bands were first identified in the A 130 nuclei including ^{130}Cs [3]. The chiral doublet bands found in ^{128}Cs were proposed as the best example to reveal chiral symmetry breaking [19], and subsequent g-factor measurements can provide important information on the relative orientation of the three angular momentum vectors [39,40].

Meanwhile, lifetimes of excited states belonging to chiral partner bands have been reported in ^{126}Cs, representing the first time a large set of experimental transition probabilities has shown qualitative agreement with all selection rules predicted for the strong chiral symmetry breaking limit [21].

Therefore, it is naturally interesting to search for nuclear chirality in other cesium isotopes. In neighboring odd-odd isotopes ^{120,134}Cs, only one rotational band with the configuration $\pi h_{11/2}$ $h_{11/2}{}^{-1}$ was reported [32,41]. In the heavier ^{134}Cs, experimental observations [32,42] indicate that the band structure is distinctly different from its lighter odd-odd isotopes and was interpreted as a possible magnetic rotation band [42]. Therefore, it is worthwhile to re-examine the rotational modes in ^{120,134}Cs and the possible boundaries of cesium isotopes.

For odd-A cesium isotopes, doublet bands involving the configuration component $\pi h_{11/2}$ with the third neutron quasi-particle or hole in the $g_{7/2}/d_{5/2}$ or $s_{1/2}/d_{3/2}$ orbits have been observed in ^{121,123,125,127,129,131}Cs [43–48]. Theoretically, the constrained triaxial relativistic mean-field (RMF) approach [49,50] has been used to study possible chirality in odd-A cesium isotopes ^{125,127,129,131}Cs, demonstrating and predicting the existence of chirality based on different high-j particle-hole configurations and triaxial deformations. Additionally, a rotational band with the configuration component $\pi h_{11/2}$ $(sd)^{-1}$ has been observed in ^{133}Cs [51]. Therefore, it is also interesting to investigate chirality in other odd-A cesium isotopes, i.e., ^{121,123,133}Cs.

To describe chiral geometry microscopically, we adopt three-dimensional tilted axis cranking based on covariant density functional theory (3DTAC-CDFT). This approach has been developed [52] and applied to investigate nuclear chirality in ^{106}Rh [52], ^{106}Ag [53], ^{135}Nd [54], and $^{102-107}$Rh [55]. These studies support the existence of a critical frequency ω{crit} *corresponding to the transition from planar to aplanar rotation. For example, in Ref. [54], a classical Routhian was extracted by modeling nucleon motion in a rotating mean field as the interplay between single-particle motions of several valence particle(s) and hole(s) and the collective motion of a core-like part. This classical Routhian shows qualitative agreement with the 3DTAC-CDFT result for* ω{crit}. In this paper, following similar procedures outlined in Refs. [52–55], we investigate chirality in $^{120-134}$Cs within 3DTAC-CDFT.

2 Theoretical Framework

The 3DTAC-CDFT based on zero-range point-coupling interaction has been successfully applied to describe chiral rotation [52–55]. The detailed formalism of 3DTAC-CDFT can be found in Ref. [52]; here we present a brief introduction.

CDFT starts from a Lagrangian, and the corresponding Kohn-Sham equations take the form of a Dirac equation with effective fields S and V derived from this Lagrangian [56–60]. In the 3DTAC method, these fields are calculated in the intrinsic frame rotating with a constant angular velocity vector ω, pointing in an arbitrary direction in space [52]:

$[\alpha \cdot (\mathrm{p}-\mathrm{V}) + \beta(\mathrm{m}+\mathrm{S}) + \mathrm{V} - \omega \cdot \hat{\mathrm{J}}] \;=\; \;.$ (1)

Here, $\hat{\mathrm{J}}$ is the total angular momentum of the nucleon spinors, and the relativistic fields S, V^0, and V read:

$\mathrm{S}(\mathrm{r}) = \alpha \quad + \beta \quad ^2 + \gamma \quad ^3 + \delta \, \Delta \;,$
$\mathrm{V}^0(\mathrm{r}) = \alpha \quad + \gamma \quad ^3 + \delta \, \Delta \quad + \tau_3 \alpha \qquad + \tau_3 \delta \;\; \Delta \quad + \mathrm{eA}_0,$
$\mathrm{V}(\mathrm{r}) = \alpha \, \mathrm{j} \; + \gamma \, (\mathrm{j} \;)^3 + \delta \, \Delta \mathrm{j} \; + \tau_3 \alpha \;\; \mathrm{j} \quad + \tau_3 \delta \;\; \Delta \mathrm{j} \quad + \mathrm{eA}.$

The S, V^0, and V represent the relativistic scalar field, the time-like component of the vector field, and the space-like components of the vector field, respectively, which are in turn coupled with nucleon densities and current distributions. In the above equations, , , , j , j represent various densities and currents, and α , β , γ , δ , α , γ , δ , α , δ are nine parameters of the Lagrangian. The Dirac equation is solved iteratively in a three-dimensional harmonic oscillator basis, yielding single-nucleon spinors , single-particle Routhians , total energies, expectation values of angular momenta, transition probabilities, and so on. Using the semiclassical cranking condition $\hat{\mathrm{J}} \cdot \hat{\mathrm{J}} = \mathrm{I}(\mathrm{I}+1)$, one can relate the magnitude of the angular velocity ω to the angular momentum quantum number I. Meanwhile, the orientation of ω is determined self-consistently by minimizing the total Routhian.

From the Dirac equation, physical observables including quadrupole moments, magnetic moments, and electromagnetic transition probabilities B(M1) and B(E2) can be calculated.

The quadrupole moments Q_{20} and Q_{22} are:

$\mathrm{Q}_{20} = \; 3\mathrm{z}^2 - \mathrm{r}^2 \; = \mathrm{Qa}_{20},$
$\mathrm{Q}_{22} = \; \mathrm{x}^2 - \mathrm{y}^2 \; = \mathrm{Qa}_{22}.$

The nuclear quadrupole deformation parameters β and triaxial deformation parameters γ can be obtained from:

$\beta = \surd(\mathrm{a}_{20}{}^2 + 2\mathrm{a}_{22}{}^2),$
$\gamma = \arctan(\surd 2 \; \mathrm{a}_{22}/\mathrm{a}_{20}).$

The nuclear magnetic moment in units of the nuclear magneton is given by:

$$= \ d^3r\ [q/(2mc^2)\ \ \dagger(r\times\alpha)\ \ +\ \ \dagger\beta\Sigma\ \ ],\ (5)$$

where the charge q is 1 for protons and 0 for neutrons in units of e, is the nucleon anomalous gyromagnetic factor (= 1.793, = −1.193), and Σ represents the spin operator [61]. In the following calculations, the ratio $mc^2/\hbar c$ in Eq. (5) is taken as 1, as in Refs. [52,62].

The transition probabilities B(M1) and B(E2) are calculated in the semiclassical approximation:

$B(M1) = (3/8\pi)\ [-\ \_z \sin\ + \cos(\ \_x \cos\ +\ \_y \sin\ )]^2 + (\ \_y \cos\ -\ \_x \sin\ )^2$, (6)

$B(E2) = (5/16\pi)\ [(Q_{20} \sin^2\ + \sqrt{2}\ Q_{22} \cos\ \sin 2\ )^2 + (1+\cos^2\ )^2$ cos^{2}2], (7)

where Q_{20} and Q_{22} are the proton quadrupole moments, and and are the orientation angles of the total angular momentum $\hat{J}$ in the intrinsic frame.

In the present work, we adopt the point-coupling Lagrangian PC-PK1 [61] and the calculations are free of additional parameters. The Dirac equation (1) is solved in a three-dimensional Cartesian harmonic oscillator basis with 10 major shells. Pairing correlations are neglected in the calculations, though one should bear in mind that pairing might influence the description of critical frequency [37] as well as total angular momentum and B(M1) values [63,64].

3 Results and Discussion

In our study of $^{120-134}$Cs, we assign two-quasi-particle configurations $\pi h_{11}/_2\ \ h_{11}/_2{}^{-1}$ for odd-odd nuclei and three-quasi-particle configurations $\pi h_{11}/_2\ \ h_{11}/_2{}^{-1}(gd)^1$ and $\pi h_{11}/_2\ \ h_{11}/_2{}^{-1}(sd)^1$ for odd-A nuclei, similar to earlier studies [3,19–24,32,41,43–47,49–51,65]. In the calculations, an unpaired proton always occupies the bottom of the $h_{11}/_2$ orbital, and the occupation of valence neutrons in the $h_{11}/_2$ orbital is fixed by tracing single-particle levels with increasing frequency as shown in Fig. 1, while other nucleons are treated self-consistently by filling orbits according to their energies from the bottom of the potential well. The (gd) notation represents low-j neutron orbits $1g_7/_2$ and $2d_5/_2$, while (sd) represents $2d_3/_2$ and $3s_1/_2$, as it is difficult to distinguish them due to strong mixing. Note that six or eight neutrons in the $h_{11}/_2$ orbital are antialigned, giving the neutron configuration $h_{11}/_2{}^{-1}$. Thus, the valence nucleon configurations in our 3DTAC-CDFT calculations are $\pi h_{11}/_2\ \ h_{11}/_2{}^{-1}(gd)^n$ (n = 8–14) for $^{120-126}$Cs, $\pi h_{11}/_2\ \ h_{11}/_2{}^{-1}(sd)^n$ (n = 1–2) for $^{127-128}$Cs, and $\pi h_{11}/_2\ \ h_{11}/_2{}^{-1}(gd)^4\ \ h_{11}/_2{}^{-1}(sd)^n$ (n = 1–6) for $^{129-134}$Cs.

In 3DTAC-CDFT calculations, the orientation of the angular velocity ω with respect to the principal axes can be determined self-consistently either by minimizing the total Routhian or by requiring that ω is parallel to the total angular momentum J at a fixed ω value. We use the polar angle and azimuth angle to denote the direction of $\omega = \omega(\sin\ \cos\ , \sin\ \sin\ , \cos\ )$. The angle is between ω and the long (l) axis, while is between the projection of ω onto the short-medium (sm) plane and the short (s) axis. It has been noted that can characterize the chirality of a rotational system [64].

To examine possible chiral geometry, Fig. 2 shows the self-consistently obtained orientation angles and of the total angular momentum J in the intrinsic frame as functions of rotational frequency $\hbar\omega$. The polar angle in $^{120-134}$Cs shows similar behavior, increasing with rotational frequency. Nevertheless, remains larger than 50° in all cesium isotopes, attributed to the angular momentum alignment along the s axis from proton particles in the $h_{11}/_{2}$ orbital being much larger than that along the l axis from neutron holes in the $h_{11}/_{2}$ orbital (cf. Fig. 4).

For comparison, the azimuth angle for $^{121-133}$Cs is zero at low rotational frequencies, corresponding to planar rotation in the sl plane. Above a limiting rotational frequency corresponding to the critical frequency ω_{crit} for chiral rotation [37,66], becomes nonzero, resulting in a transition from planar to aplanar rotation. Note that kinks appear in the curves for $^{121-133}$Cs at the critical frequency.

The azimuth angle remains zero for ^{120},^{134}Cs, corresponding to planar rotation in the sl plane, even when the rotational frequency reaches $\hbar\omega = 0.55$ MeV for ^{120}Cs and $\hbar\omega = 0.45$ MeV for ^{134}Cs. With increasing rotational frequency, the calculations fail to converge for the configuration $\pi h_{11}/_{2}\quad h_{11}/_{2}{}^{-1}(gd)^{8}$ in ^{120}Cs and $\pi h_{11}/_{2}\quad h_{11}/_{2}{}^{-1}(sd)^{6}$ in ^{134}Cs. This suggests that $\hbar\omega$_{crit} in the present 3DTAC-CDFT calculations might be larger than $\hbar\omega = 0.55$ MeV for ^{120}Cs and $\hbar\omega = 0.45$ MeV for ^{134}Cs, even if it exists, corresponding to I 24ℏ—beyond the currently observed spin range [32,42,47]. Thus, our theoretical analysis does not support the existence of chirality in ^{120},^{134}Cs.

To understand this behavior, we must analyze the angular momentum geometry. In 3DTAC-CDFT calculations, the total angular momentum arises from individual nucleons in a coherent superposition. For all cesium isotopes studied, a single proton occupying the bottom of the $h_{11}/_{2}$ orbital contributes roughly 5.5ℏ of angular momentum along the s axis. It should be emphasized that the proton angular momentum contribution is approximately along the s axis at low rotational frequency. Therefore, the angular momentum components along the l and m axes come mainly from the neutron part.

As a demonstration, we study the angular momentum contributions of valence neutrons in the $h_{11}/_{2}$, (gd), and (sd) orbits along the s, m, and l axes using ^{124}Cs, ^{125}Cs, and ^{126}Cs as examples, shown in Fig. 4. In contrast, five neutron holes at the top of the $h_{11}/_{2}$ orbital contribute about 3.5ℏ to the angular momentum along the l axis since four of them are antialigned. The only difference is the distribution of neutrons over the (gd) orbits, resulting in different configurations $\pi h_{11}/_{2}\quad h_{11}/_{2}{}^{-1}(gd)^{n}$ (n = 12–14) for $^{124-126}$Cs.

Previous studies indicate a possible magnetic rotation nature for the configuration $\pi h_{11}/_2$ $h_{11}/_2{}^{-1}$ in ^{134}Cs based on the decreasing trend of B(M1) [42]. Figure 3 shows the composition of the total angular momentum at different rotational frequencies in 3DTAC-CDFT calculations for ^{120}Cs and ^{134}Cs with this configuration. In our 3DTAC-CDFT calculations for ^{120},^{134}Cs, as the four protons in the $g_7/_2$ orbital are paired, the proton angular momentum comes mainly from particles in the $h_{11}/_2$ orbital, which aligns along the s axis. The neutron hole in the $h_{11}/_2$ orbital and neutrons in low-j (gd)/(sd) orbits give substantial contributions to the neutron angular momentum, leading to large components in both the l and s axes. Higher spin states in the band are created by aligning the neutron angular momentum toward the s axis, while the proton angular momentum remains unchanged along the s axis. Thus, ^{120},^{134}Cs likely exhibit planar rotation.

Chirality in nuclei with stable triaxial deformation arises from aplanar rotation formed by valence particle(s), valence hole(s), and collective core angular momentum vectors [1]. Therefore, an important consideration is the transition point from planar to chiral rotation—the critical frequency *ω{crit}. As shown in Fig. 2, the calculated critical frequencies* ω{crit} all decrease with increasing neutron number, which can also be observed in the behaviors of I and B(M1)/B(E2) versus $\hbar\omega$.

Taking ^{124}Cs, ^{125}Cs, and ^{126}Cs as examples to study the angular momentum contributions of valence neutrons in the $h_{11}/_2$, (gd), and (sd) orbits along the s, m, and l axes, Fig. 4 shows that the angular momentum increment of the (gd) or (sd) orbits along the s and l axes becomes smaller with increasing neutron number, while the increment along the m axis becomes larger. This makes it easier for the nucleus to form chiral rotation, and the corresponding critical frequency ω_{crit} decreases. Therefore, it is clear that the critical frequency for the same configuration in Fig. 2 always tends to shift leftward, similar to earlier studies [55]. It should be noted that in ^{133}Cs, the neutron in the (sd) orbits transforms from particle-state character to hole-state behavior, showing a different change compared to other nuclei.

The critical frequency also affects the experimental spin-rotational frequency I($\hbar\omega$) relationship. As shown in Fig. 5, the calculated spin values I($\hbar\omega$) agree well with experimental data for $^{120-123}$Cs but overestimate the data for $^{124-127}$,$^{129-131}$,^{133}Cs. Additionally, the calculated I($\hbar\omega$) values are much higher than experimental data for ^{128},^{132},^{134}Cs. For $^{124-134}$Cs, effective neutron and proton pairing correlations and beyond-mean-field effects not considered here may play important roles, as confirmed in Refs. [67,68]. Furthermore, the backbending phenomenon in ^{128},^{132},^{134}Cs, also not considered here, may significantly affect the present calculated results.

Obviously, the calculated and experimental rotational frequencies $\hbar\omega$ for the yrast bands in $^{121-133}$Cs behave similarly with respect to angular momentum. The bands show nearly straight lines at low rotational frequency. A kink appears when the spin reaches the critical frequency *ω{crit} in $^{121-133}$Cs. This kink occurs due to the transition from planar to chiral solutions, with the angular* *momentum component increasing along the m axis, causing substantial changes in angular momentum. The corresponding kink in experimental I(ħω) appears almost at the same position as the calculated ħω*{crit} results. Further experimental efforts for ^{133}Cs are encouraged to verify our conclusion on ħω_{crit}. For ^{120},^{134}Cs, previous experimental investigations do not support a static chirality interpretation in the observed spin range [32,42,47], which agrees with our theoretical result showing no kink.

As mentioned, positive-parity rotational bands have been observed in $^{120-134}$Cs in previous works [3,19–24,43–46,49,50]. Comparisons between 3DTAC-CDFT excitation energies and available experimental data are shown in Fig. 6. At the present mean-field level, neither chiral vibrations nor tunneling between left- and right-handed sectors are taken into account; therefore, only the band with lower excitation energies can be reproduced. However, Fig. 6 shows that experimental excitation energies of the lower band are well reproduced. To describe the partner band, one needs to go beyond mean-field calculations by combining, for example, random phase approximation methods [35] or collective Hamiltonian with CDFT [64,69].

Figure 7 displays experimental B(M1)/B(E2) data for $^{120-134}$Cs [3,19–24,43–46,49,50] compared with 3DTAC-CDFT results derived in the semiclassical approximation from magnetic and quadrupole moments. Note that deformation parameters (β, γ) change only slightly along the band, so corresponding B(E2) values are almost constant. However, B(M1) values decrease smoothly due to continuous variation of rotational frequency and closing of neutron and proton angular momentum vectors, which mainly align along the short and long axes, respectively. This leads to a smoothly decreasing trend for B(M1)/B(E2) ratios. Therefore, the 3DTAC-CDFT ratios differ from experimental data, though the smooth-decreasing tendency is reproduced. In particular, magnetic moments are derived from the relativistic expression of the effective current operator as in Ref. [57]. As shown in Fig. 7, 3DTAC-CDFT results for ^{121},^{123},$^{125-134}$Cs show good agreement with data, while calculations overestimate data for ^{120},^{122},^{124}Cs. The theoretical results show that B(M1)/B(E2) for $^{121-133}$Cs has a smooth falling behavior in planar rotation, with the falling tendency slowing down above critical frequency *ω{crit}. This again proves that nuclei transition from planar to chiral rotation above ω*{crit}. The falling tendency becomes steeper with increasing neutron number. Further efforts to include neutron and proton pairing correlations will help justify these results.

Calculated quadrupole deformation parameters β and triaxial deformation parameters γ are given in Fig. 8. We find that β and γ are stable with only slight changes as rotational frequency increases in $^{120-134}$Cs. With increasing neutron number, quadrupole deformation β in cesium isotopes gradually decreases, consistent with results in Refs. [70,71]. This can be understood because the neutron number in ^{120}Cs is midway between magic numbers, so quadrupole deformation β gradually decreases. For triaxial deformation γ, it increases with neutron number in $^{120-126}$Cs, then decreases with neutron number in $^{127-134}$Cs.

This can be understood as follows: in $^{120-126}$Cs, triaxial deformation γ gradually increases with neutron number, the neutron hole contribution becomes stronger, and the deformation changes from prolate to oblate ellipsoid, leading to increasing γ. Similarly, the neutron particle contribution weakens with increasing neutron number in $^{127-134}$Cs, and the deformation evolves from oblate to prolate ellipsoid, causing γ to decrease. However, the sudden increase in triaxial deformation γ between ^{128}Cs and ^{129}Cs is due to increasing neutron number in the $\mathrm{h}_{11}/_{2}$ orbital.

Figure 8 also shows that the triaxial deformation in ^{134}Cs is approximately zero, further indicating no chirality in ^{134}Cs, in agreement with previous results [32]. The triaxial deformation parameter γ in ^{120}Cs shows a relatively large value (15°), although planar rotation with zero azimuth angle is obtained. Therefore, the existence of chirality cannot be definitively established based solely on high-j particle-hole configurations and stable triaxial deformation [49].

4 Summary

In summary, we have performed a fully self-consistent and microscopic investigation of chirality evolution in cesium isotopes $^{120-134}$Cs using 3DTAC-CDFT. By examining the evolution of polar angle and azimuth angle for the total angular momentum J as functions of increasing rotational frequency $\hbar\omega$, we find a transition from planar to chiral rotation in $^{121-133}$Cs. In contrast, only planar rotation is obtained in ^{120},^{134}Cs, which is probably not magnetic rotation.

Moreover, the critical frequency *ω{crit} in $^{121-133}$Cs decreases with increasing neutron number, which can also be observed in the behaviors of I and B(M1)/B(E2) versus $\hbar\omega$. Taking ^{124}Cs, ^{125}Cs, and ^{126}Cs as examples to study angular momentum contributions of valence neutrons in the $h_{11}/_{2}$, (gd), and (sd) orbits along the s, m, and l axes, we find that the angular momentum increment of the (gd) or (sd) orbits along the s and l axes becomes smaller with increasing neutron number, while the increment along the m axis becomes larger. This makes it easier for the nucleus to form chiral rotation, and the corresponding critical frequency* ω{crit} decreases.

Additionally, the calculated I $\hbar\omega$ show good agreement with available experimental data for $^{120-123}$Cs but overestimate data for $^{124-134}$Cs. The calculated energy spectra for yrast bands of $^{120-134}$Cs are reasonable compared with corresponding experimental data. The 3DTAC-CDFT results show good agreement with B(M1)/B(E2) data for ^{121},^{123},$^{125-134}$Cs, while overestimating data for ^{120},^{122},^{124}Cs. The quadrupole deformation parameters β and triaxial deformation parameters γ in $^{120-134}$Cs show clear systematic trends: β decreases with increasing neutron number in $^{120-134}$Cs; γ increases in $^{120-126}$Cs and decreases in $^{127-134}$Cs, related to neutron occupation. Triaxial deformation in ^{120}Cs and ^{134}Cs indicates that chirality cannot be definitively established based solely on high-j particle-hole configurations and stable triaxial deformation.

Recently, chiral bands in ^{119}Cs [72] based on configurations with only protons were observed for the first time in the A 130 mass region—a configuration with three protons, one in the strongly coupled [404]9/2$^+$ orbital which does not change orientation with increasing rotational frequency, and two in the $h_{11}/_2$ orbital which reorient to the rotation axis. This differs from the common chiral geometry picture involving both protons and neutrons. This is excellent and significant work. Preliminary 3DTAC-CDFT calculations do not support the existence of aplanar rotation with such a configuration. Meanwhile, not only TAC-CDFT but also other theoretical models such as the particle-rotor model would be helpful to justify these results and understand the underlying rotational structure in ^{119}Cs. Investigations in these directions are in progress. Additionally, more efforts on investigating nuclear chirality in the A 130 mass region are still needed, including different theoretical methods, evolution of isotones, possible configurations, and so on.

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