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$N = 28$ isotones: Shape coexistence towards proton-deficient side

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Abstract

We have employed RMF+BCS (relativistic mean-field and Bardeen-Cooper-Schrieffer) approach to study the phenomenon of shape coexistence in $N = 28$ isotones towards the proton-deficient side. Our present investigations include single particle energies, deformations, binding energies as well as excitation energies. It is found that towards the proton-deficient side, $N = 28$ shell closure disappears due to reduced gap between neutron $1f_{7/2}$ and $1f_{5/2}$ and the nuclei ^{40}Mg, ^{42}Si, and ^{44}S are found to possess shape coexistence giving further support to weakening of the shell gap. These results are found in excellent match with other theoretical and experimental studies and are fortified with a variety of calculations and parameters.

Keywords: Neutron magic nuclei; Relativistic mean-field plus BCS approach; Shape coexistence; Shell closure.

1. Introduction

The evolution of ground-state shapes in an isotopic or isotonic chain is governed by changes of the shell structure of single-nucleon orbital. In recent past, evolution of the shell structure guiding shape coexistence, has been observed in the $N = 20$ and $N = 28$ isotones around the proton drip line [1–4]. In a more general manner, the major structural features along the isotonic and isotopic chains around the spherical magic numbers 8, 20, 28, 50, 82, and 126 have been reviewed by Sorlin et al. [5] using evolution of the binding energies, trends of first collective states and characterization of single-particle states. A number of experimental investigations have shown [3,6] that in the proton-deficient $N = 28$ isotones below ^{48}Ca the spherical shell gap progressively reduces and the low-energy spectra of ^{46}Ar, ^{44}S, and ^{42}Si display evidence of ground-state deformation and shape coexistence.

Theoretical treatment of evolution of shell closure has been successfully described using mean field theories [7–9] and with their relativistic counterparts [10–15]. The main advantage of the RMF+BCS approach is that it provides the spin–orbit interaction in the entire mass region in a natural way [10–12]. This indeed has proved to be very crucial for the study of nuclei near the drip-line. As nuclei move away from stability and approach the drip-lines, the corresponding Fermi surface gets closer to zero energy at the continuum threshold. A significant number of the available single-particle states then form part of the continuum. Indeed the RMF+BCS scheme [11,14] yields results which are in close agreement with the experimental data and with those of relativistic continuum Hartree–Bogoliubov (RCHB) and other similar mean-field cal-

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culations [15]. Recently, deformed relativistic-Hartree-Bogoliubov (RHB) theory in a continuum has been developed aiming at a proper description of exotic nuclei, particularly for ^{42}Mg [16]. Moreover, the development of the covariant density functional theory in continuum has been introduced for the description of neutron halo phenomena in medium-heavy and heavy nuclei, including the RCHB theory, the relativistic-Hartree-Fock-Bogoliubov (HFB) theory in continuum and the deformed RHB theory in continuum [17].

In recent past, the RMF approach has been successfully used to investigate two-proton radioactivity [18], weakly bound drip-line nuclei [19] and magicity [20]. More recently, relativistic mean field study has been extensively used to describe actinides and superheavy nuclei within covariant density functional theory [21], to calculate decay rates of various proton emitters [22], to study bubble structures [23], to analyze the effects of particle-number fluctuation degree of freedom on symmetric and asymmetric spontaneous fission [24] and to calculate neutron capture cross-sections in nuclei near the $N = 82$ shell closure [25].

In this paper we have investigated the shape coexistence phenomenon for proton deficient $N = 28$ isotones using Relativistic Mean Field plus the Bardeen-Cooper-Schrieffer approach [18–20].

2. Relativistic mean-field model

Our RMF calculations have been carried out using the model of Lagrangian density with nonlinear terms both for the σ and ω mesons along with the TMA parametrization as described in detail in Refs. [12,14,18].

$$\begin{aligned}\mathcal{L} = {} & \bar{\psi}[\imath\gamma^{\mu}\partial_{\mu} - M]\psi \\ & + \frac{1}{2}\,\partial_{\mu}\sigma\partial^{\mu}\sigma - \frac{1}{2}m_{\sigma}^{2}\sigma^{2} - \frac{1}{3}g_{2}\sigma^{3} - \frac{1}{4}g_{3}\sigma^{4} \\ & - g_{\sigma}\bar{\psi}\sigma\psi - \frac{1}{4}H_{\mu\nu}H^{\mu\nu} + \frac{1}{2}m_{\omega}^{2}\omega_{\mu}\omega^{\mu} \\ & + \frac{1}{4}c_{3}(\omega_{\mu}\omega^{\mu})^{2} - g_{\omega}\bar{\psi}\gamma^{\mu}\psi\omega_{\mu} \\ & - \frac{1}{4}G_{\mu\nu}^{a}G^{a\mu\nu} + \frac{1}{2}m_{\rho}^{2}\rho_{\mu}^{a}\rho^{a\mu} - g_{\rho}\bar{\psi}\gamma_{\mu}\tau^{a}\psi\rho^{\mu a} \\ & - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} - e\bar{\psi}\gamma_{\mu}\frac{(1-\tau_{3})}{2}A^{\mu}\psi\,, \end{aligned} \tag{1}$$

where the field tensors H, G and F for the vector fields are defined by

$$\begin{aligned} H_{\mu\nu} &= \partial_{\mu}\omega_{\nu} - \partial_{\nu}\omega_{\mu} \\ G_{\mu\nu}^{a} &= \partial_{\mu}\rho_{\nu}^{a} - \partial_{\nu}\rho_{\mu}^{a} - 2g_{\rho}\,\epsilon^{abc}\rho_{\mu}^{b}\rho_{\nu}^{c} \\ F_{\mu\nu} &= \partial_{\mu}A_{\nu} - \partial_{\nu}A_{\mu} \end{aligned}$$

and other symbols have their usual meaning. Based on the single-particle spectrum calculated by the RMF described above, we perform state-dependent BCS calculations [26,27]. The continuum is replaced by a set of positive energy states generated by enclosing the nucleus in a spherical box. Thus the gap equations have the standard form for all the single particle states, i.e.,

$$\begin{aligned}\Delta_{j_1} = {} & -\frac{1}{2}\frac{1}{\sqrt{2j_1+1}}\sum_{j_2}\frac{\left\langle (j_1{}^2)\,0^+\,|V|\,(j_2{}^2)\,0^+\right\rangle}{\sqrt{\left(\varepsilon_{j_2}-\lambda\right)^2+\Delta_{j_2}^2}} \\ & \times\sqrt{2j_2+1}\,\Delta_{j_2}, \end{aligned} \tag{2}$$

where ε_{j_2} are the single particle energies, and λ is the Fermi energy, whereas the particle number condition is given by $\sum_j (2j+1)v_j^2 = N$. In the calculations we use for the pairing interaction a delta force, i.e., $V = -V_0\delta(r)$ with the same strength $V_0 = 350$ MeV fm^3 for both protons and neutrons [14]. Apart from its simplicity, the applicability and justification of using such a δ-function form of interaction has been discussed in Ref. [28], whereby it has been shown in the context of HFB calculations that the use of a delta force in a finite space simulates the effect of finite range interaction in a phenomenological manner. The pairing matrix element for the δ-function force is given by

$$\begin{aligned}& \left\langle (j_1{}^2)\,0^+\,|V|\,(j_2{}^2)\,0^+\right\rangle \\ & = -\frac{V_0}{8\pi}\sqrt{(2j_1+1)(2j_2+1)}I_R, \end{aligned} \tag{3}$$

where I_R is the radial integral having the form

$$I_R = \int dr\frac{1}{r^2}\left(G_{j_1}^{\star}\,G_{j_2} + F_{j_1}^{\star}\,F_{j_2}\right)^2 \tag{4}$$

Here G_α and F_α denote the radial wave functions for the upper and lower components, respectively, of the nucleon wave function expressed as

$$\psi_\alpha = \frac{1}{r}\begin{pmatrix} i\,G_\alpha\,\mathcal{Y}_{j_\alpha l_\alpha m_\alpha} \\ F_\alpha\,\sigma\cdot\hat{r}\,\mathcal{Y}_{j_\alpha l_\alpha m_\alpha}\end{pmatrix}, \tag{5}$$

and satisfy the normalization condition

$$\int dr\,\{|G_\alpha|^2 + |F_\alpha|^2\} = 1 \tag{6}$$

In Eq. (5) the symbol $\mathcal{Y}_{jlm}$ has been used for the standard spinor spherical harmonics with the phase i^l. The coupled field equations obtained from the Lagrangian density in (1) are finally reduced to a set of simple radial equations which are solved