Transition rates and microscopic effective charges for ^{16}C exotic nucleus

Raad A. Radhi, Zaheda A. Dakhil, Dheyaa A. A. Al-Ibadi

Department of Physics, College of Science, Baghdad University, Baghdad, Iraq.

E-mail: loi_hil@yahoo.com

Abstract

Quadrupole transitions and effective charges are calculated for ^{16}C exotic nuclei with large basis no core shell model with $(0+2)\hbar\omega$ truncations. Calculations with configuration mixing shell model with these limited model spaces usually underestimate the measured E2 transition strength. Instead of using constant effective charges, excitations out of major shell space are taken into account through a microscopic theory this is called core-polarization effects. The two body Michigan sum of three ranges Yukawa potential (M3Y) is used for the core-polarization matrix elements. The simple harmonic oscillator potential is used to generate the single particle matrix elements of ^{16}C. The present calculations with core polarization (C. P.) effect reproduced the experimental data very well.

Key words

Quadrupole transitions, transition rate, effective charges.

Article info.
Received: Jan. 2014
Accepted: Jan. 2014
Published: Apr. 2014

معدل الأنتقال والشحنة الفعالة للـ ^{16}C

رعد عبد الكريم راضي، زاهده أحمد دخيل، ضياء علوان عبد الحسين العبادي

قسم الفيزياء، كلية العلوم، جامعة بغداد، بغداد، العراق.

الخلاصة

حسبت الانتقالات رباعية القطب والشحنات الفعالة للنواة الغريبة ^{16}C و استخدم أنموذج القشرة بدون قلب الموسع مع قطع $\hbar\omega(2+0)$. حسابات أنموذج القشرة ذو التشكيلات المختلطة مع أنموذج الفضاءات المحدودة عادة تقلل من تقدير قوة الانتقال E2 المقاس. بدلا من استخدام شحنات فعالة ثابتة، فقد أخذت التهيجات خارج فضاءات القشرة الرئيسية بنظر الاعتبار من خلال النظرية المايكروية الذي يسمى تأثير استقطاب القلب. كذلك تم حساب عناصر المصفوفة لاستقطاب القلب بواسطة جهد M3Y. أستخدم جهد المتذبذب التوافقي البسيط لتوليد عناصر المصفوفة للجسيمة المفردة. جهد المتذبذب التوافقي استخدم لتوليد عناصر المصفوفة للجسيمة المفردة لـ ^{16}C. هذه الحسابات مع ألأخذ بنظر الاعتبار تأثير استقطاب القلب اتفقت مع القيم العملية بصوره جيده.

Introduction

The exotic nuclei are phenomena for some light nuclei and it has special conditions – far from stability line or near drip-lines (neutron-rich or proton-rich). In other words, exotic nuclei are nuclei with an extraordinary ratio of protons and neutrons Z/N. Typically; they are very unstable and decay into more stable nuclei, i.e. exotic nuclei so-called rare-isotopes, which have a loosely binding energy, short-lived isotopes and large isospin.

Actually, exotic nuclei are available as secondary beams at many radioactive beam facilities around the world [1]. These unstable nuclei are generally weakly bound with few excited states. These exotic nuclei have a thin cloud of nucleons orbiting at large distances from the others, forming the core [1].

Neutron-rich nuclei have attracted much interest during the past decades [2- 5] and this will continue to be so due to new generation radioactive beam facilities in the world. These nuclei are characterized by a small binding energy, a halo and skin structures with a large spatial extension of the density distribution [6].

One of the most striking features in neutron-rich nuclei is the nuclear deformation. The deformation can be investigated experimentally and theoretically, through their electromagnetic transitions. The general trend of the 2^+ excitation energy $E(2_1^+)$ and the reduced electric quadrupole transition strength between the first excited 2^+ state and the 0^+ ground state, $B(E2;2_1^+ \rightarrow 0_1^+)$ for even-even nuclei are expected to be inversely proportional to one another [7]. However, recent experimental and theoretical studies of the neutron-rich ^{16}C, ^{18}C and ^{20}C isotopes [8-10] show a deviation from the general trends of even-even nuclei. A systematic study of transition rates for these isotopes was recently conducted both theoretically and experimentally [10-13].

The role of the core and the truncated space can be taken into consideration through a microscopic theory. These effects are essential in describing transitions involving collective modes such as E2 transition between states in the ground-state rotational band, such as in ^{18}O [14]. Umeya [15, 16] calculated effective charges for quadrupole moments and transitions by using first order perturbation. These theoretical results show that the effective charges are smaller than the standard value in light- neutron rich nuclei and imply decoupled quadrupole motions between protons and neutrons in neutron-rich nuclei.

Zheng et al. [17] studied the reaction cross sections for ^{12,16}C had been measured at the energy of 83 MeV by a new experimental method. The larger enhancement of the ^{16}C reaction cross section at the low energy had been used to study the density distribution of ^{16}C. The finite-range Glauber-model calculations for different density distributions had been compared with the experimental data. A large extension of the neutron density distribution to a distance far from the center of the nucleus suggested the formation of neutron halo in the ^{16}C nucleus.

Imai et al. [8] presented a study of the electric quadrupole transition from the first excited 2^+ state to the ground 0^+ state in ^{16}C was studied through measurement of the lifetime by a recoil shadow method applied to inelastically scattered radioactive ^{16}C nuclei. The measured mean lifetime was 77± 14(stat) ±19 (sys.) ps. The central value of mean lifetime corresponded to a $B(E2;2_1^+ \rightarrow 0^+)$ value of 0.63 e^{2}fm^4, or 0.26 Weisskopf units. The transition strength was found to be anomalously small compared to the empirically predicted value.

Hagino and Sagawa [12] applied a three-body model consisting of two valence neutrons and the core nucleus ^{14}C in order to investigate the ground state properties and electric quadrupole transition of the ^{16}C nucleus. The calculated B(E2) value from the first 2^+ state to the ground state showed good agreement with the observed data with the core polarization charge which reproduced the experimental B(E2) value for ^{15}C. It was also showed that their calculations account well for the longitudinal momentum distribution of the ^{15}C fragment from the breakup of the ^{16}C nucleus. They pointed out that the dominant $(d_{5/2})^2$ configurations in the ground state of ^{16}C played a crucial role in these agreements.

Ong et al. [9] presented a studying of

lifetime measurements of first excited states in $^{16,18}C$. In that article, the electric quadrupole transition from the first excited 2^+ state to the ground 0^+ states in ^{18}C was studied through a lifetime measurement by an upgraded recoil shadow method applied to inelastically scattered radioactive ^{18}C nuclei. The measured mean lifetime was 18.9 ± 0.9(stat) ± 4.4(syst) ps, corresponding to a $B(E2;2_1^+ \rightarrow 0_{g.s}^+)$ value of 4.3 ± 0.2 ± 1.0 e^2fm^4, or about 1.5 Weisskopf units. The mean lifetime of the first excited 2^+ state in ^{16}C was re-measured to be 18.3 ± 1.4 ± 4.8 ps, about four times shorter than the value reported previously. The discrepancy between the two results was explained by incorporating the γ-ray angular distribution measured in that work into the previous measurement. The transition strengths were hindered compared to the empirical transition strengths, indicating that the anomalous hindrance observed in ^{16}C persists in ^{18}C.

Wuosmaa et al. [18] studied the ^{15}C (d, p) ^{16}C reaction in inverse kinematics using the Helical Orbit Spectrometer at Argonne National Laboratory. Neutron-adding spectroscopic factors gave a different probe of the wave functions of the relevant states in ^{16}C. Shell-model calculations reproduced both the present transfer data and the previously measured transition rates.

Petri et al. [19] reported the first measurement of the lifetime of the 2_1^+ state was in the near-dripline nucleus ^{20}C. The deduced value of $\tau_{2_1^+}$ = $9.8 \pm 2.8(stat)^{+0.5}_{-1.1}$ (syst) ps gave a reduced transition probability of $B(E2;2_1^+ \rightarrow 0_{g.s}^+) = 7.5^{+3.0}_{-1.7}$ $(stat)^{+1.0}_{-0.4}$ (syst) e^2fm^4 which was in a good agreement with a shell model calculation using isospin-dependent effective charges.

In 2012, Voss et al. [13], presented the lifetime of the first excited 2+ state which was measured with the Köln/NSCL plunger via the recoil distance method to be $\tau\,(2_1^+)$ = $22.4 \pm 0.9(stat)^{+3.3}_{-2.2}$ (syst) ps, which corresponds to a reduced quadrupole transition strength of $B(E2;\ 2_1^+ \rightarrow\ 0_1^+)$ = $3.64^{+0.15}_{-0.14}\ (stat)^{+0.40}_{-0.47}$ e^2fm^4. In addition, an upper limit on the lifetime of a higher-lying state feeding the 2_1^+ state was measured to be τ < 4.6 ps. The results were compared to the large-scale ab initio no-core shell model calculations using two accurate nucleon-nucleon interactions and the importance-truncation scheme. That comparison provided strong evidence that the inclusion of three-body forces was needed to describe the low-lying excited-state properties of that A = 18 system.

The aim of the present work is using the fundamental relations to get the reduced transition strength B(E2) from the first-excited 2^+ state to the ground state for ^{16}C.

Theory

The reduced one –body matrix element for shell-model wave functions of initial spin J_i and final spin J_f for a given multipolarity λ, can be expressed as a linear combination of the single - particle matrix elements:

$$\left\langle J_f \,\|\hat{T}_{JT}\,\| J_i \right\rangle = \sum_{J_i, J_f} OBDM\,(j_i, j_f, J_i, J_f, \lambda) \left\langle J_i \,\|\hat{T_\lambda}\,\| J_f \right\rangle \quad (1)$$

where the one–body density matrix elements (OBDM) are the structure factors. The initial and final single–particle states are denoted by j_i and j_f , respectively.

The reduced single–particle matrix element of the Coulomb (Longitudinal) operator is given by [20]:

$$\left\langle J_f \,\|\hat{T}_{JT}\,\| J_i \right\rangle = \int_0^\infty dr\, r^2\, j_\lambda(qr) \left\langle J_f \,\| Y_\lambda \,\| J_i \right\rangle R_{n_f l_f}(r) R_{n_i l_i}(r) \quad (2)$$

where $j_\lambda(qr)$ the spherical Bessel is function and $R_{nl}(r)$ is the single-particle radial wave function.

Electron scattering Coulomb form factor involving angular momentum l and momentum transfer q, between initial and final nuclear shell model states of spin $J_{i,\ f}$, are [20]:

$$|F(q)|^2 = \frac{4\pi}{Z^2(2J_i+1)} \times \left|\left\langle J_f \left\| \hat{T}_\lambda \right\| J_i \right\rangle\right|^2 \times |F_{c.m}(q)|^2 \times |F_{f.s}(q)|^2 \quad (3)$$

Several corrections must be applied to the nucleus form factor equation to convert them into a representation appropriate for a comparison with the experimental form factor. One of these corrections is the center of mass. The conventional harmonic-oscillator approximation for this correction is given by [21]:

$$F_{c.m} = e^{q^2 b^2/4A} \quad (4)$$

where A is the nuclear mass number.

In the shell-model the nucleon is assumed to be a point, but nucleons are actually of finite size, then the calculated form factors have to be corrected by another correction which takes into account the finite size of the nucleon and is given by [21-23]:

$$F_{f.s}(q) = \left[1 + \left(\frac{q}{4.33} fm^{-1}\right)^2\right]^{-2} \quad (5)$$

The total longitudinal form factor is given by:

$$|F(q)|^2 = \sum_{\lambda \geq 0} |F_\lambda(q)|^2 \quad (6)$$

The electric transition strength is given by:

$$B(C_\lambda, K) = \frac{Z^2}{4\pi}\left[\frac{(2\lambda+1)!!}{K^\lambda}\right]^2 F_\lambda^2(k) \quad (7)$$

where $k = E_x/\hbar c$ with E_x as the excitation energies.

The Relation Between $B(EJ\uparrow)$ and $B(EJ\downarrow)$ is given by:

$$B(EJ\downarrow) = \frac{2J_i+1}{2J_f+1} B(EJ\uparrow) \quad (8)$$

The matter rms radii can be getting by:

$$\left\langle r^2 \right\rangle_m^{\frac{1}{2}} = \sqrt{\frac{4\pi}{A}\int_0^\infty \rho_{0,m}(r) r^4 dr} \quad (9)$$

and $$<r^2>_m = \frac{1}{A}\sum occ\# \left[b^2\left(N+\frac{3}{2}\right)\right] \quad (10)$$

where (A) is matter number A=Z+N.

The role of the core and the truncated space can be taken into consideration through a microscopic theory, which combines shell model wave functions and configurations with higher energy as first order perturbation to describe EJ excitations: these are called core polarization effects. The reduced matrix elements of the electron scattering operator $\hat{O}_\Lambda$ is expressed as a sum of the model space (MS) contribution and the core polarization (CP) contribution, as follows:

$$\left\langle \Gamma_f \left\|\left| \tilde{\hat{O}}_\Lambda \right\|\right| \Gamma_i \right\rangle = \left\langle \Gamma_f \left\|\left| \hat{O}_\Lambda \right\|\right| \Gamma_i \right\rangle_{MS} + \left\langle \Gamma_f \left\|\left| \Delta\hat{O}_\Lambda \right\|\right| \Gamma_i \right\rangle_{CP} \quad (11)$$

which can be written as:

$$\left\langle \Gamma_f \left\|\left| \tilde{\hat{O}}_\Lambda \right\|\right| \Gamma_i \right\rangle = \sum_{\alpha\beta} X^\Lambda_{\Gamma_f\Gamma_i}(\alpha,\beta)\left[\left\langle \alpha \left\|\left| \hat{O}_\Lambda \right\|\right| \beta \right\rangle + \left\langle \alpha \left\|\left| \Delta\hat{O}_\Lambda \right\|\right| \beta \right\rangle\right] \quad (12)$$

where X is the OBDM elements. The states $|\Gamma_i\rangle$ and $|\Gamma_f\rangle$ are described by the model space wave functions. Greek symbols are used to denote quantum numbers in coordinate space and isospace, i.e. $\Gamma_i \equiv J_i T_i$, $\Gamma_f \equiv J_f T_f$ and $\Lambda \equiv JT$.

According to the first-order perturbation theory, the single particle core-polarization term is given by [24]:

$$\left\langle \alpha \left\|\left| \Delta\hat{O}_\Lambda \right\|\right| \beta \right\rangle = \left\langle \alpha \left\|\left| \hat{O}_\Lambda \frac{Q}{E_i - H_0} V_{res} \right\|\right| \beta \right\rangle + \left\langle \alpha \left\|\left| V_{res} \frac{Q}{E_f - H_0} \hat{O}_\Lambda \right\|\right| \beta \right\rangle \quad (13)$$

where the operator Q is the projection operator onto the space outside the model space. The single particle core-polarization terms given in equation (12) are written as [24]:

$$\left\langle \alpha \left\|\left| \Delta\hat{O}_\Lambda \right\|\right| \beta \right\rangle = \sum_{\alpha_1\alpha_2\Gamma} \frac{(-1)^{\beta+\alpha_2+\Gamma}}{\varepsilon_\beta - \varepsilon_\alpha - \varepsilon_{\alpha_1} + \varepsilon_{\alpha_2}} (2\Gamma+1) \begin{Bmatrix} \alpha & \beta & \Lambda \\ \alpha_2 & \alpha_1 & \Gamma \end{Bmatrix} \sqrt{(1+\delta_{\alpha_1\alpha})(1+\delta_{\alpha_2\beta})}$$
$$\times \left\langle \alpha\alpha_1 | V_{res} | \beta\alpha_2 \right\rangle_\Gamma \left\langle \alpha_2 \left\|\left| \hat{O}_\Lambda \right\|\right| \alpha_1 \right\rangle \quad (14)$$

+terms with α_1 and α_2 exchanged with an overall minus sign, where the index α_1 runs over particle states and α_2 over hole states and ε is the single-particle energy, and is calculated according to [24]:

$$\varepsilon_{n\ell j} = (2n + \ell - 1/2)\hbar\omega + \begin{cases} -1/2(\ell+1)\langle f(r)\rangle_{n\ell} & \text{for } j = \ell - 1/2 \\ 1/2\ell\langle f(r)\rangle_{n\ell} & \text{for } j = \ell + 1/2 \end{cases} \quad (15)$$

with $\langle f(r)\rangle_{n\ell} \approx -20A^{-2/3}$ and $\hbar\omega = 45A^{-1/3} - 25A^{-2/3}$

Higher energy configurations are taken into consideration through 1*p*-1*h* 2 $\hbar\omega$ excitations. For the residual two-body interaction V_{res}, the M3Y interaction of Bertsch et al. [25] is adopted. The form of the potential is defined in equations (1-3) in Ref. [25]. The parameters of 'Elliot' are used which are given in Table 1 of the mentioned reference. A transformation between LS and jj is used to get the relation between the two-body shell model matrix elements and the relative and center of mass coordinates, using the harmonic oscillator radial wave functions with Talmi-Moshinsky transformation. The single particle matrix elements reduced in both spin and isospin is given by:

$$\langle \alpha \||| \hat{T}_{JT} ||| \beta \rangle = \sqrt{\frac{2T+1}{2}} \sum_{t_z} I_T(t_z) \langle \alpha \| \hat{T}_{Jt_z} \| \beta \rangle \quad (16)$$

where,

$$I_T(t_z) = \begin{cases} 1 & for\ T = 0 \\ (-1)^{\frac{1}{2}-t_z} & for\ T = 1 \end{cases} \quad (17)$$

The reduced electric transition strength is given by:

$$B(EJ) = \frac{1}{(2J_i+1)} \left| \sum_{T=0,1} \begin{pmatrix} T_f & T & T_i \\ -T_z & 0 & T_z \end{pmatrix} \langle J_f T_f ||| \tilde{\hat{O}}_{JT} ||| J_i T_i \rangle \right|^2 \quad (18)$$

which can be written as:

$$B(EJ) = \frac{1}{(2J_i+1)} \left| \sum_{T=0,1} e_T \begin{pmatrix} T_f & T & T_i \\ -T_z & 0 & T_z \end{pmatrix} \tilde{M}_{JT} \right|^2 \quad (19)$$

where $\tilde{M}_{JT} = \langle J_f T_f ||| \tilde{\tilde{M}}_{JT} ||| J_i T_i \rangle$ The isoscalar (T=0) and isovector (T=1) charges are given by $e_0 = e_{IS} = \frac{1}{2}e$, $e_1 = e_{IV} = \frac{1}{2}e$.

The B(E2) value can be represented in terms of only the model space matrix elements as:

$$B(EJ) = \frac{1}{(2J_i+1)} \left| \sum_{T=0,1} e_T^{\text{eff}} \begin{pmatrix} T_f & T & T_i \\ -T_z & 0 & T_z \end{pmatrix} M_{JT} \right|^2 \quad (20)$$

Then the isoscalar and isovector effective charges are given by:

$$e_T^{\text{eff}} = \frac{M_{JT} + \Delta M_{JT}}{2M_{JT}} e = \frac{e_p + (-1)^T e_n}{2} \quad (21)$$

The proton and neutron effective charges can be obtained as follows:

$e_p = e_0^{eff} + e_1^{eff}$ and $e_n = e_0^{eff} - e_1^{eff}$.

The above effective charges work for mixed isoscalar and isovector transitions. For pure isoscalar transition (for ^{12}C), the polarization charge is given by:

$$\delta e = \frac{\Delta M_J}{2M_J} e \quad (22)$$

and the effective charges for the proton and neutron becomes:

$$e_p^{\text{eff}} = e + \delta e, \quad e_n^{\text{eff}} = \delta e \quad (23)$$

Results and Discussion

Large-basis no core model space is used in this study. This space covered the four shells 1*s*, 1*p*, 2*s*–1*d* and 2*p*–1*f* with $(0+2)\hbar\omega$ truncations. Shell model interactions encompassing the four oscillator shells have been constructed by Warburton and Brown [26]. These interactions are based interactions for the 1*p*2*s*1*d* shells determined by a least square fit to 216 energy levels in the A = 10–22 region assuming no mixing of $n\hbar\omega$ and $(0+2)\hbar\omega$ configurations. The 1*p*2*s*1*d* part of the interaction (cited in Ref. [26] as WBP) results from a fit to two-body matrix elements and single-particle energies for the *p* shell and a potential representation of the 1*p*–2*s*1*d* cross shell interaction. The WBP model space was expanded to include the 1*s* and 2*p*1*f* major shells by adding the appropriate 2*p*1*f* and cross-shell 2*s*1*d*–2*p*1*f* two-body matrix element of the Warburton–Becker–Milliner–Brown (WBMB) interaction [27] and all the other necessary matrix elements from the bare G-matrix potential of Hosaka, Kubo and Toki [28]. The $2s1d$ shell interaction of Wildenthal [29] used in WBP interaction is replaced in this study by a new interaction referred as USDB (Universal *sd*-shell B) [30], where the derivation of the USD Hamiltonian [29] has been refined with an up dated and complete set of energy data. The new Hamiltonian USDB leads to a new level of precision for realistic shell-model wave functions.

The b value of this isotope is adjusted to reproduce the experimental matter radius. The b average value for ^{16}C is 1.78 fm. The experimental and theoretical Rm radii are tabulated in Table 1. The experimental value of Rm is taken from Ref. [31].

Table 1: The calculated root mean square matter radii of ^{16}C compared with the experimental data.

A	***b*** *(fm)*	*Root mean Square matter radii (fm)*	
		spsdpf $(0+2)\hbar\omega$	*Exp. [31]*
16	*1.78*	*2.74*	*2.7±0.03*

Shell model calculations were performed with the shell-model code OXBASH [32], where the OBDM elements given in equation (12) were obtained. The first term in this equation is the zero-order contribution, which gives the single-particle matrix element for the model space (MS) contribution. The second and third terms are the first-order contributions which account for the higher energy configurations (core-polarization effects). These configurations are taken through $1p$–$1h$ excitations from the core and MS orbits into all higher orbits with $2\hbar\omega$ excitations. For the residual interaction V_{res}, the M3Y interaction of Bertsch et al. [25] is adopted.

The ground state of ^{16}C is with $J^{\pi}T = 0^{+}2$, with half life= 0.747s. The first excited 2^+ state of this exotic nucleus is at 1766 (10) keV [9].

The recent measurements of the $B(E2; 2_1^+ \rightarrow 0_1^+)$ values are 4.15±0.73 e^{2}fm^4 [30] and 2.6±0.9 e^{2}fm^4 [9]. The calculations are performed with *spsdpf* model space with $(0+2)\hbar\omega$. For the ground state configurations, two neutrons are distributed over the *sd* shell orbits with 62% over $1d_{5/2}$ orbit, 6% over $1d_{3/2}$ orbit and 32% over the $2s$ orbit. The calculated $B(E2; 2_1^+ \rightarrow 0_1^+)$ is 0.51 e^{2}fm^4 without CP effects and that with CP effects is 3.3 e^{2}fm^4, which agrees very well with both experimental values. The effective charges are calculated to be equal to 1.1 e and 0.24 e, for the proton and neutron, respectively. The results of B(E2) and effective charges are tabulated in Table 2. The analysis of the above calculations shows a major contribution of $1d_{5/2}$ for valence two neutrons. The configurations resulting from the *spsdpf* model space with $(0+2)\hbar\omega$ gives a simple structure of ^{16}C= ^{14}C+n+n. A large extension of the neutron density distribution to a distance far from the center of the nucleus suggests the formation of neutron halo in the ^{16}C nucleus [17]. Due to the relatively small value of B(E2), and the distribution of the two neutrons over the *sd* shell out side ^{14}C may support the decoupling of neutrons from protons to form neutron halo. The result of Forss´en et al. is 2.2(6) e^2fm^4 [33] using a large-scale ab initio no-core shell model (NCSM).

Table 2: The calculated effective charges and B(E2) values of ^{16}C compared with the experimental data.

Model Space	Effective Charges (e) e_p^{eff}, e_n^{eff}	$B(E2; 2_1^+ \rightarrow 0_1^+)\ e^2fm^4$ Theo. (No CP)	Theo. (With CP)	Exp. [Ref]
spsdpf (0+2) $\hbar\omega$	1.1, 0.24	0.51	3.33	4.15±0.73 [30] 2.6±0.9 [25]

Conclusions

Shell model calculations are performed for ^{16}C including core-polarization effects through first-order perturbation theory, where 1*p*–1*h* with $2\hbar\omega$ excitation are taken into considerations. In general, there are some notes have been indicated from the present work which can be explained as:

The $0\hbar\omega$ and (0 + 2) $\hbar\omega$ calculations which succeed in describing energy levels and other static properties, are less successful for describing dynamical properties such as transition strengths B(E2). The core contributions cannot be ignored in such transitions and the core polarization effects play a major role for describing such dynamical property.

The small value obtained for the transition strength in ^{16}C which agrees very well with the measured values suggests a possible proton-shell closure with a simple ^{14}C+n+n structure. These two neutrons may form a halo around the ^{14}C nucleus, but cannot be considered as a Borromean since ^{14}C exists as a long lived isotope. The present calculations of the neutron-rich ^{16}C isotope show a deviation from the general trends of even-even nuclei in accordance with experimental and other theoretical studies.

The experimental values are very well reproduced confirming the anomalous suppression in ^{16}C. The configurations arises from the shell model calculations with core-polarization effects which reproduce the experimental B(E2) values, and give small effective charges confirm the formation of proton-shell closure for ^{16}c.

References

[1] E. Kwan, C. Y. Wu, N. C. Summers, G. Hackman, T. E. Drake, C. Andreoiu, R. Ashley, G. C. Ball, P. C. Bender, A. J. Boston, H. C. Boston, A. Chester, A. Close, D. Cline, D. S. Cross, R. Dunlop, A. Finley, A. Garnsworthy, A. B. Hayes, A. Laffoley, T. Nano, P. Navr´atil, C. J. Pearson, J. Pore, K. Starosta4, I. J. Thompson1, P. Voss, S. J. Williams, Z. M. Wang, Phys. Lett., B 650 (2007) 124.

[2] M. V. Zhukov, B. V. Danilin, D. V. Fedorov, J. M. Bang, I. J. Thompson and J. S. Vaagen, Phys. Rep., 231 (1993) 151.

[3] G. Raimann, A. Ozawa, R. N. Boyd, F. R. Chloupek, M. Fujimaki, K. Kimura, T. Kobayashi, J. J. Kolata, S. Kubono, I. Tanihata, Y. Watanabe, K. Yoshida, Prog. Part. Nucl. Phys., A 35 (1995) 505.
[4] B. Jonson, Phys. Rep., 389 (2004) 1.
[5] A. S Jensen, K. Riisageer, D. V. Fedorov and E. Garrido, Rev. Mod. Phys., 76 (2004) 215.
[6] I. Tanihata, H. Hamagaki, O. Hashimoto, Y. Shida, N. Yoshikawa, K. Sugimoto, O. Yamakawa, T. Kobayashi and N. Takahashi, Phys. Rev. Lett., 55 (1985) 2676.
[7] S.Raman, C. W. Nestor, J.R. and P. Tikkanen, Atomic Data and Nuclear Data Table, 78 (2001) 1.
[8] N. Imai, H.J. Ong, N. Aoi, H. Sakurai, K. Demichi, H. Kawasaki, H. Baba, Zs. Dombrádi, Z. Elekes, N. Fukuda, Zs. Fülöp, A. Gelberg, T. Gomi, H. Hasegawa, K. Ishikawa, H. Iwasaki, E. Kaneko, S. Kanno, T. Kishida, Y. Kondo, T. Kubo, K. Kurita, S. Michimasa, T. Minemura, M. Miura, T. Motobayashi, T. Nakamura, M. Notani, T. K. Onishi, A. Saito, S. Shimoura, T. Sugimoto, M. K.Suzuki, E. Takeshita, S. Takeuchi, M. Tamaki, K. Yamada, K. Yoneda, H. Watanabe and M. Ishihara, Phys. Rev. Lett., 92 (2004) 6.
[9] H.J. Ong, N. Imai, D. Suzuki, H. Iwasaki, H. Sakurai, T.K. Onishi, M.K. Suzuki, S. Ota, S. Takeuchi, T. Nakao, Y. Togano, Y. Kondo, N. Aoi, H. Baba, S. Bishop, Y. Ichikawa, M. Ishihara, T. Kubo, K. Kurita, T. Motobayashi, T. Nakamura, T. Okumura and Y. Yanagisawa, Phys. Rev., C 78 (2008) 014308.
[10] Z. Elekes, Zs. Dombrádi, T. Aiba, N. Aoi, H. Baba, D. Bemmerer, B.A. Brown, T. Furumoto, Zs. Fülöp, N. Iwasa, Á. Kiss, T. Kobayashi, Y. Kondo, T. Motobayashi, T. Nakabayashi, T. Nannichi, Y. Sakuragi, H. Sakurai, D. Sohler, M. Takashina, S. Takeuchi, K. Tanaka, Y. Togano, K. Yamada, M. Yamaguchi and K. Yoneda, Phys. Rev., C 79 (2009) 011302(R).
[11] H. Sagawa, X. R. Zhou, X. Z. Zhang and T. Suzuki, Phys. Rev., C 70 (2004) 054316.
[12] K. Hagino and H. Sagawa, Phys. Rev., C 75 (2007) 021301(R).
[13] P. Voss, T. Baugher, D. Bazin, R. M. Clark, H. L. Crawford, A. Dewald, P. Fallon, A. Gade, G. F. Grinyer, H. Iwasaki, A. O. Macchiavelli, S. McDaniel, D. Miller, M. Petri, A. Ratkiewicz, W. Rother, K. Starosta, K. A. Walsh, D. Weisshaar, C. Forss´en, R. Roth and P. Navr´atil, Phys. Rev., C 86 (2012) 011303(R).
[14] R.A. Radhi and E.A. Salman, Nucl. Phys., A 806 (2008) 179.
[15] A. Umeya, G. Kaneko and K. Muto, Nucl. Phys., A 829 (2009) 13.
[16] A. Umeya and K. Muto, Nucl. Phys., A 722 (2003) 558 C.
[17] T. Zheng , T. Yamaguchi , A. Ozawaa, M. Chibaa, R. Kanungo,T. Kato , K. Katori , K. Morimoto , T. Ohnishi , T. Sudaa, I. Tanihata , Y. Yamaguchi , A. Yoshida , K. Yoshida , H. Toki and N. Nakajima , Nucl. Phys., A 709 (2002) 103.
[18] A. H. Wuosmaa, B. B. Back, S. Baker, B. A. Brown, C. M. Deibel, P. Fallon, C. R. Hoffman, B. P. Kay, H.Y. Lee, J. C. Lighthall, A. O. Macchiavelli, S. T. Marley, R. C. Pardo, K. E. Rehm, J. P. Schiffer, D.V. Shetty and M. Wiedeking, Phys. Rev. Lett., 105 (2010) 132501.
[19] M. Petri, P. Fallon, A. O. Macchiavelli, S. Paschalis, K. Starosta, T. Baugher, D. Bazin, L. Cartegni, R. M. Clark, H. L. Crawford, M. Cromaz, A. Dewald, A. Gade, G. F. Grinyer, S. Gros, M. Hackstein, H. B. Jeppesen, I. Y. Lee, S. McDanie D. Miller, M. M. Rajabali, A. Ratkiewicz, W. Rother, P. Voss, K. A. Walsh, D. Weisshaar, M. Wiedeking, and B. A. Brown, Phys. Rev. Lett., 107 (2011) 102501.
[20] T.de Forest Jr., and J.D.Walecka, Adv. Phys., 15 (1966) 1.

[21] L. J. Tassie and F. C. Barker; Phys. Rev., 111 (1958) 940.
[22] H. Chandra and G. Sauer; Phys. Rev., C 13 (1976) 245.
[23] J. P. Glickman, W. Bertozzi, T. N. Buti, S. Dixit, F. W. Hersman, C. E. Hyde-Wright, M. V. Hynes, R. W. Lourie, B. E. Norum, J. J. Kelly, B. L. Herman and D. J. Millener. Phys. Rev., C 43 (1990) 1740.
[24] P. J. Brussaard and P. W. M. Glaudemans, "Shell-model Applications in Nuclear Spectroscopy", North-Holland Publishing Company, Amsterdam (1977).
[25] G. Berstch, J. Borysowicz, H. McManus, W. G. Love, Nucl. Phys., A 284 (1977) 399.
[26] E. K. Warburton and B.A. Brown, Phys. Rev. C 46 (1992) 923.
[27] E.K. Warburton, J.A. Becker and B.A. Brown, Phys. Rev., C 41 (1990) 1147.
[28] A. Hosaka, K.-I. Kubo and H. Toki, Nucl. Phys., A 244 (1985) 76.
[29] B.A. Brown, B.H. Wildental, Annu. Rev. Nucl. Part. Sci., 38 (1988) 29.
[30] B.A. Brown, W.A. Richter, Phys. Rev., C 74 (2006) 034315.
[31] A. Ozawa,T.Suzuli and I. Tanihata, Nucl. Phys. A 693 (2001) 32.
[32] B. A. Brown, A. Etchegoyen, N. S. Godwin, W. D. M. Rae, W. A. Richter, W. E. Ormand, E. K. Warburton, J. S. Winfield, L. Zhao, C. H. Zimmerman, MSU-NSCL report number 1289 (2005) version.
[33] C. Forss´en, R. Roth,P. Navr´atil, J. Phys.G:Nucl. Part. Phys., 40 (2012) 055105.