Giant quadrupole excitation in nuclei with neutron skin

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The electric quadrupole excitation of unstable nuclei with the neutron skin is studied in terms of a Hartree-Fock RPA calculation, taking ^{28}O as an example. It is found that there are giant excitation peaks, referred to as giant neutron modes, below the isoscalar giant resonance and these peaks consist of neutron excitations only. Their transition densities are shown to be extended radially in comparison to the isoscalar giant resonance. A strong correlation between $B(E2)$ sum of the giant neutron modes and the number of neutrons in the neutron skin is found.

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The structure of nuclei far from the stability line (i.e., unstable nuclei) has attracted much interest recently. This is largely due to developments in experimental technique for producing radioactive beams, but also due to the discovery of several interesting phenomena such as the neutron halo [1], the neutron skin [2], etc. Especially certain new aspects which have not been seen in stable nuclei open a new era and show us new facets of nuclear structure physics. It has been shown in [3] that, as the neutron number (N) increases from the β stability line, a significant neutron skin emerges in the Skyrme Hartree-Fock (HF) calculation [4–7]. The proton skin arises as the proton number (Z) increases [3]. We shall report in this Brief Report effects of the neutron skin on the basic excitation scheme, taking oxygen isotopes as an example. ^{16}O is on the stability line, and hence there is no neutron skin due to the almost perfect symmetry between protons and neutrons. We shall consider ^{28}O as a typical example of unstable nuclei. The Skyrme SIII interaction is used, while the following results are general, and should be rather insensitive to details of the calculation. ^{28}O, a drip-line nucleus in the SIII interaction, contains 20 neutrons, and exhibits a neutron skin, which, according to [3], is 2.1 fm thick and contains about four neutrons. Since our primary aim is to show the basic features of excitations of unstable nuclei, it is not very relevant whether or not ^{28}O really exists in nature. Certainly, ^{28}O is a tractable case for HF calculations due to its doubly magic shell structure. We shall demonstrate that due to the neutron skin a new excitation mode appears with rather exotic features compared to the excitation scheme of stable nuclei.

The HF potential and related quantities of ^{16}O and ^{28}O show the following characteristic properties [3]. The excess neutrons deepen the proton potential because of a strong proton-neutron interaction. Thus, the proton potential becomes deeper in ^{28}O than in ^{16}O. The neutron potential is raised slightly in ^{28}O compared to ^{16}O. The neutron Fermi level goes up nearly to the threshold, reflecting the fact that ^{28}O is near the neutron-drip line. This happens mainly because we put so many neutrons into this potential, but also because the neutron potential is being neither deepened nor significantly widened by these added neutrons.

Figure 1 presents the proton and neutron density profiles of ^{28}O. In ^{16}O, the proton and neutron densities are quite similar to each other. The ratio between the neutron density (ρ_ν) and the proton density (ρ_π) is about equal to N/Z in the nuclear interior, while the total density is approximately constant [3]. Therefore, in the interior, ρ_ν for ^{28}O is larger than ρ_ν for ^{16}O, producing a deeper HF potential for protons as discussed above. On the contrary, in the interior, ρ_π for ^{28}O is smaller than ρ_π for ^{16}O, and thereby the proton-neutron interaction attracts neutrons to a lesser extent, yielding a shallower potential for neutrons.

The mean potential becomes, in general, less steep in its upper part near the threshold. In low-lying states of stable nuclei, nucleons are deeply bound, and their wave functions are not strongly influenced by this upper part of the potential. The last neutrons of ^{28}O, however, occupy single-particle orbits of which a good fraction is lying in the less steep part of the potential. Thus, their wave functions can be stretched radially to a large extent, producing the neutron skin. Classically, this means that the turning point of single-particle motion is moved farther away. Therefore, contrary to the halo, this can occur for orbits with high orbital angular momenta, for

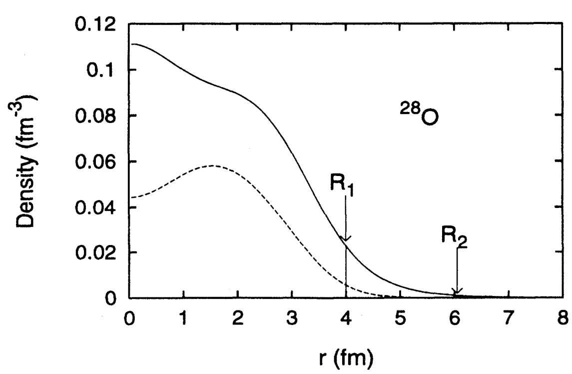

FIG. 1. Density profiles of protons (dotted lines) and neutrons (solid lines) in ^{28}O. The inner (R_1) and outer (R_2) boundaries of the neutron skin [3] are shown by arrows.

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instance, the $d_{3/2}$ orbit in the present case. The tunneling can expand the wave functions additionally, particularly for an s orbit. Thus, the neutron skin is related to the shape of the potential. The potential becomes less steep in its upper part in the presence of the neutron skin in a self-consistent manner.

We now discuss giant resonance(-type) excitations in unstable nuclei. Because of the limited length, we shall restrict ourselves, in this Brief Report, to electric quadrupole ($E2$) type excitations. The giant quadrupole resonance (GQR) has played a very important role in clarifying the basic structure of the nucleus [8,9]. It has been established that the isoscalar (IS) GQR appears at $E_{2^+} \simeq (60-65)A^{-1/3}$ MeV [8,9]. We shall demonstrate that there should be huge excitation peaks below the IS-GQR in unstable nuclei.

The excitations from the HF ground state are studied in terms of the random phase approximation (RPA). As is well known, such a RPA calculation is quite suitable for studying the GQR and related phenomena [10,11]. The single-particle states with positive energies are obtained in the bound-state approximation [12], and we solve the RPA equations in terms of a sufficiently large number of harmonic oscillator bases.

We consider a quadrupole operator

$$Q_{2\mu} = \sum_{i=1}^{A} e_i {r_i}^2 Y_{2\mu}(\hat{r}_i), \tag{1}$$

where e_i is the nucleon charge, r denotes the radius from the center of nucleus, and Y_2 is the spherical harmonics. We use mainly the isoscalar Q operator, referred to as $Q_{\rm IS}$, where $e_i = e/2$ for all nucleons. This operator seems to be reasonable for describing excitations due to the nuclear processes, for instance, reactions such as (p,p'), (α,α'), etc., which can be carried out by bombarding these targets with unstable nuclei as secondary radioactive beams. In addition, in some illustrative calculations, we employ the proton (neutron) Q operator which is denoted as Q_π (Q_ν) and refers to $e_i = e/2$ for protons (neutrons) and $e_i = 0$ for neutrons (protons).

The reduced transition probability for the 2_i^+ state,

$$B(E2; 0^+ \to 2_i^+) = \sum_\mu |\langle 2_i\mu|Q_{2\mu}|00\rangle|^2, \tag{2}$$

is then calculated with $Q_{\rm IS}$, Q_π, or Q_ν. Figure 2(a) shows $B(E2)$'s for $Q_{\rm IS}$ in ^{16}O. One sees several peaks around $E_x \sim 21$ MeV. In order to see global features, each discrete peak obtained by the bound-state approximation is smeared out, hereafter, by the Gaussian function with a fixed width (1.0 MeV in this case) and the height determined so that the covered area of $dB(E2)/dE_x$ is equal to $B(E2)$ of the corresponding discrete peak. The Gaussian functions from different peaks are superimposed. One then obtains, for ^{16}O, a single large peak representing the ISGQR in good agreement with experimental data, $E_x \sim 22$ MeV [8].

We next calculate $B(E2)$'s for $Q_{\rm IS}$ in ^{28}O. Figure 2(b) displays such $B(E2)$'s after the smearing mentioned above. One then sees a striking feature: Instead of a single peak there are several huge ones with strengths comparable to each other. In order to see their structure in more detail, we calculate $B(E2)$'s for Q_π and Q_ν, denoted, respectively, as $B_\pi(E2)$ and $B_\nu(E2)$. Figure 2(c) presents such results. For the peak at $E_x \sim 21$ MeV, both $B_\pi(E2)$ and $B_\nu(E2)$ are of sizable magnitude. Indeed, $B_\pi(E2)/B_\nu(E2) \sim Z/N$ holds for this peak, suggesting its isoscalar character. Thus, this peak belongs to the same category as the ISGQR in ^{16}O.

We now survey the lower energy region. One then finds that only $B_\nu(E2)$'s take large values. In other words, the lower-lying peaks, which have heights comparable to the isoscalar peak, consist practically of neutron transitions only. This is quite remarkable, because their $B(E2)$'s are even larger than the isoscalar one and they are located so low in energy. We state that ^{28}O is a doubly closed nucleus, and hence there is no impurity of the valence-shell transitions. In this sense, although the excitation energies are quite low, all peaks in Figs. 2(b), 2(c) resemble $2\hbar\omega$ (i.e., cross shell) excitations in the harmonic oscillator potential.

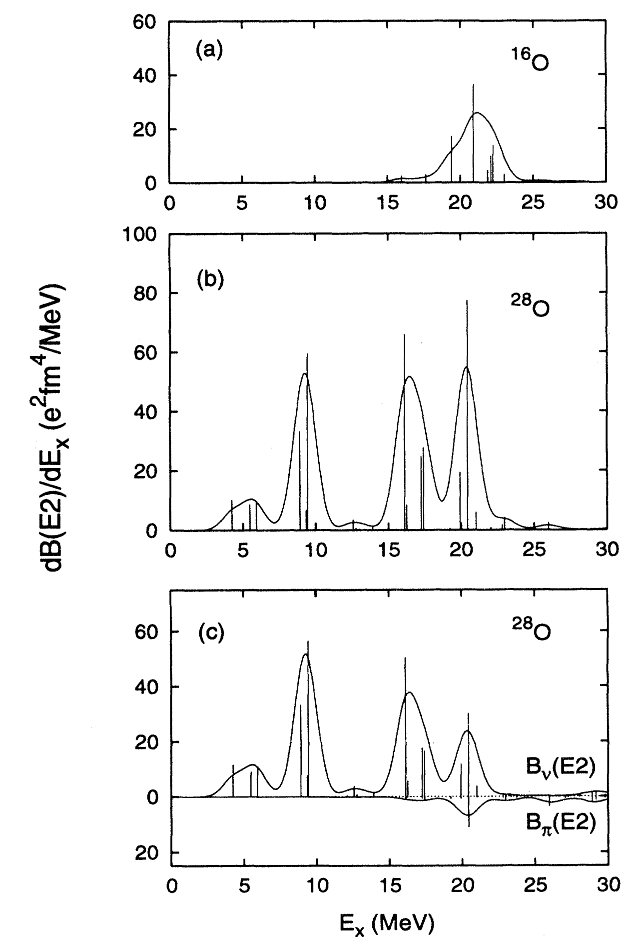

FIG. 2. $dB(E2)/dE_x$ as a function of excitation energy in the HF+RPA calculation for (a) ^{16}O and (b), (c) ^{28}O. See the text for the discrete peaks and their smearing. In (a) and (b), the isoscalar $E2$ operator is taken (see the text). In (c), the upper (lower) part is for the neutron (proton) $E2$ operator. Notice that the scale is downward in the lower part of (c).

A closer inspection of the transition matrix elements clearly shows that these lower-lying peaks are constituted almost only by the neutron particle-hole excitations from the highest single-particle orbits which actually form the neutron skin. Thus, these lower-lying peaks are related essentially to the neutron skin, and will be referred to as giant neutron modes. They are indeed giant excitations in the sense that their excitation strengths are comparable to that of the ISGQR. Figure 2(b) indicates that there can be several giant neutron modes in a nucleus. As discussed later, at present, whether they are "giant resonances" or not needs to be confirmed. Thus, we refer to them as giant neutron modes. The first report of the giant neutron mode was made in [13]. The next report was presented by Tohyama [14] in terms of the time-dependent (TD) HF and density matrix calculations for ^{22}O.

Similar calculations are carried out for ^{22}O and ^{24}O assuming the $1d_{5/2}$ and $2s_{1/2}$ subshell closure, respectively. We then obtain the $B(E2)$ sum of the $Q_{\rm IS}$ operator for the giant neutron modes, which will be called S_ν. Note that the isoscalar peak and the giant neutron mode peaks are rather well separated in all cases. On the other hand, valence-shell transitions (i.e., transitions within the valence shell) are mixed strongly in the lowest peaks of ^{22}O and ^{24}O. Note that the valence shell here means the *sd*-shell. We remove valence-shell transitions by eliminating the *sd*-shell single-particle matrix elements from the $E2$ operators. This is needed because such transitions can be seen also in stable nuclei. In fact, the two lowest peaks of ^{22}O and the lowest peak of ^{24}O are nearly pure valence-shell excitations. The TDHF analysis in [14] was made for ^{22}O and seems to contain the valence-shell contributions, although the formalism is slightly different. Figure 3 presents the S_ν of ^{16,22,24,28}O as a function of the number of neutrons in the neutron skin. We find a remarkable linear dependence of the S_ν on this number. A deeper understanding of this salient correlation is being developed. Figure 3 includes $B(E2)$'s of the ISGQR peaks, which stay rather constant. Note that S_ν exceeds it in ^{28}O.

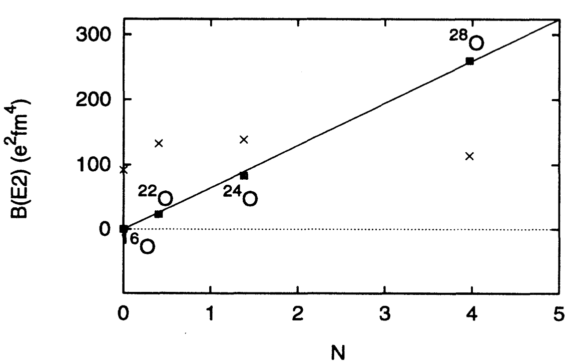

FIG. 3. $B(E2)$ sum (squares) of the giant neutron modes (S_ν in the text) as a function of the number of neutrons in the neutron skin. The solid line indicates a linear fit. The crosses indicate $B(E2)$'s of the ISGQR peaks. The results for ^{16,22,24,28}O are shown.

The energy-weighted sum rule becomes $6150e^2$ fm^4 MeV for ^{28}O, of which the result in Fig. 2(b) exhausts 91% with $E_x = 0-30$ MeV. The ISGQR accounts for 34% of this value.

Figure 4 shows the transition densities in ^{28}O for the two (discrete) peaks at $E_x = 20.4$ MeV and $E_x = 9.4$ MeV. The former belongs to the isoscalar mode, while the latter to the giant neutron modes. It should be noticed that the giant neutron mode has a quite extended transition density, reflecting the stretched radial density distribution of the neutron skin.

We note that the excitation energy of the giant neutron mode is determined almost only by the particle-hole single-particle energy difference, which may suggest that this mode is rather isolated. This is in contrast to the excitation energy of the isoscalar peak which is lowered by about 3 MeV as an effect of the residual interaction on a coherent combination of various particle-hole components. The hole states contributing to the giant neutron modes are loosely bound, and their wave functions are sensitive to the upper part of the mean potential. This part is less steep than the lower one. The excitation energy determined primarily by this upper part should be smaller than that determined by the lower (steeper) part. The normal $E2$ giant resonance is characterized basically as a $2\hbar\omega$ excitation in the harmonic oscillator potential. The less steep part of the HF potential can then be interpreted in terms of a smaller effective $\hbar\omega$, which results in lower excitation energies.

The escaping width of the giant neutron mode, in comparison to that of the ISGQR, is influenced by competing mechanisms: lower excitation energy makes it smaller, whereas more extended wave functions enlarge it. Therefore, this width is expected to be of the same order of magnitude as that of ISGQR, of which a more accurate evaluation is under way in terms of the continuum RPA. The spreading width has been reported in [14] to be smaller, in the case of ^{22}O, than that for the ISGQR of stable nuclei. We then expect that the giant neutron modes have a width structure similar to the giant resonances. As regards dynamical aspects, the giant neutron modes are probably quadrupole oscillations of the neutron skin, which are rather mutually decoupled, and are

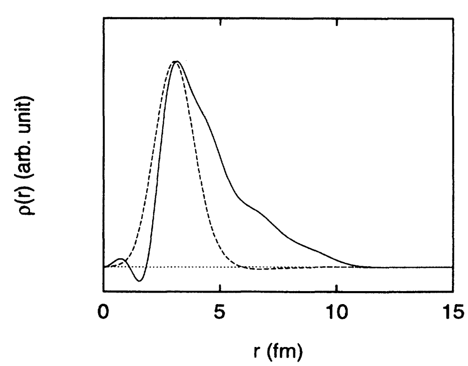

FIG. 4. Transition densities of the isoscalar giant resonance (dashed line) and of a giant neutron mode (solid line) in ^{28}O. The unit is arbitrary so that the heights are adjusted.

subject to a weaker restoring force. This point is to be studied theoretically and experimentally.

We note that the excitation from the neutron halo occurs mainly in the region where the mean potential vanishes, because the halo is nothing but tunneling. Therefore, a neutron excited from the halo runs away quickly. On the other hand, to a sizable extent, a neutron in the neutron skin is excited to the region where the potential is still attractive. This neutron is less movable than the one excited from the halo. The giant neuron mode can thus be considered to be similar to the giant resonance, in contrast to the excitation from the halo.

The giant proton mode can be seen in nuclei with a proton-skin(-like) structure, for instance, in ^{34}Ca. The giant neutron mode can thus be found in other multipolarities. In the case of $E1$, it is usually referred to as the soft dipole mode. However, because of the $E1$ effective charge $-\frac{Z}{A}e$ for neutrons and $\frac{Z}{A} \ll 1$ for nuclei with a neutron skin, the excitation is not a giant one for $E1$. This is a distinct difference between $E1$ and other multipolarities. Results on other multipolarities will be reported elsewhere, together with more detailed description on the $E2$ excitations.

In summary, we have presented a brief sketch of the $E2$ excitation scheme in the oxygen neutron-rich unstable isotopes. There are exotic features such as appearance of the giant neutron modes, their extraordinarily large strengths, the low excitation energies, and the quite extended transition densities. These are all salient features characterizing the excitation from nuclei with a neutron skin. Because these features possess rather general aspects, we expect that they can be seen in general for nuclei having a neutron skin with sizable thickness. Further studies are in progress on this intriguing issue.

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