Interference between Single-Particle and Core-Polarization Amplitudes in Heavy-Ion Inelastic Scattering

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In the heavy-ion excitation of nuclei near closed shells, an amplitude involving single-particle degrees of freedom interferes with a collective, core-polarization amplitude which has a large Coulomb component. Thus details of the nuclear structure and of the microscopic reaction mechanism may have significant effects on the observed Coulomb-nuclear interference pattern. Calculations are given which demonstrate such effects and resolve the outstanding problem in the heavy-ion excitation of ^{18}O.

A nucleon added to a closed-shell nucleus interacts with the core and distorts it. This "core-polarization" mechanism is a basic aspect of nuclear structure. It underlies the existence of deformed nuclei and has the particular consequence that about half the quadrupole moment of a "single-particle" nucleus is due to the motion induced by the extra nucleon upon the core.[1]

The consequences of core polarization have been studied in light-ion reactions. A number of papers have shown how "valence" amplitudes involving single-particle degrees of freedom interfere with collective core-polarization amplitudes (see particularly Love and Satchler[2]). However, this has not been done for heavy ions. Here a new feature arises since the "core" amplitude has a sizable Coulomb component. This provides a well-known background against which the interplay between collective and single-particle amplitudes can be observed.

In order to discuss such effects, we give a formalism for the heavy-ion excitation of nuclei near closed shells and apply it to the excitation of the 2_1^+ (1.98 MeV) state of ^{18}O. By accounting for the valence and core degrees of freedom for ^{18}O, we overcome the difficulties met by previous studies of this system.[3-9] In our doing so, two particular results emerge: (a) It is important to use a complex interaction to calculate the valence amplitude. (b) The reorientation process —and thus the sign of the 2_1^+ quadrupole moment —plays an important role and depends on the details of the 2_1^+ wave function. In light of these results, it will be of interest to study the heavy-ion excitation of other nuclei near closed shells for which there is relatively little information available.

Consider the excitation of a target A by a structureless projectile a where A consists of a core and valence particles (generalizations are straightforward). We assume that the target states can be written as

$$\Phi_\alpha = \sum_c \varphi_c{}^\alpha \psi_c = \varphi_0{}^\alpha \psi_0 + \Psi_\alpha, \tag{1}$$

where the $\varphi_c{}^\alpha$ describe the valence particles when the core is in the state ψ_c. The core ground state is denoted by ψ_0 and thus Ψ_α includes the effects of core polarization. We also assume that the effective interaction V_{aA} is given by the interaction of a with the valence and core nucleons separately, $V_{aA} = V_{av} + V_{ac}$. Then the matrix element $M_{\beta\alpha} = \langle \Phi_\beta | V_{aA} | \Phi_\alpha \rangle$ splits into a valence and a core part,

$$M_{\beta\alpha}{}^{\text{valence}} = \langle \varphi_0{}^\beta | V_{av} | \varphi_0{}^\alpha \rangle, \tag{2}$$

$$M_{\beta\alpha}{}^{\text{core}} = \langle \varphi_0{}^\beta \psi_0 | V_{ac} | \Psi_\alpha \rangle + \langle \Psi_\beta | V_{ac} | \varphi_0{}^\alpha \psi_0 \rangle + \langle \Psi_\beta | V_{ac} | \Psi_\alpha \rangle. \tag{3}$$

We neglect $\langle \Psi_\beta | V_{av} | \Psi_\alpha \rangle$ since it is second order in the particle-core coupling[2] and should be small compared to Eq. (2). However, $\langle \Psi_\beta | V_{ac} | \Psi_\alpha \rangle$ can be significant if the states have large deformed components—which is the case for ^{18}O.

We specify the core interaction according to the macroscopic collective model. Thus

$$V_{ac} = \sum_{\lambda\mu} 4\pi \mathfrak{M}_c{}^*(E\lambda, \mu) \left[\frac{Z_a e}{2\lambda + 1} \frac{1}{r_{ac}{}^{\lambda+1}} + \frac{R_c}{3 Z_c e R_{\text{Coul}}{}^\lambda} \frac{d}{dr_{ac}} U_{ac}(r_{ac}) \right] Y_{\lambda\mu}(\hat{r}_{ac}), \tag{4}$$

where $\mathfrak{M}_c(E\lambda, \mu)$ is the electric multipole operator for the core and U_{ac} is the optical potential for the $a + c$ system. Note that Eq. (4) does not necessarily imply a purely vibrational model.

The valence interaction V_{av} is specified microscopically as a sum of nucleon-projectile interactions

which we take to be the optical potentials $U_{ia}(r_{ia})$ for nucleon scattering on system a at the energy $E_{c.m.}/A$. It may be noted that this prescription, when used in folding calculations of the real part of the total optical potential, overestimates the strength by about a factor of 2.[10] On the other hand, this prescription is analogous to the one used conventionally in heavy-ion transfer reactions, since the real part of the nucleon optical potential is similar to the shell-model or average binding potentials. Likewise, it has been used successfully for heavy-ion inelastic scattering by Broglia *et al.*[11] This suggests that the folding procedure for the elastic and inelastic interactions is sensitive to different regions of the nucleon-nucleus potential. (Using the prescription of Rickertsen and Satchler,[12] we have found that the inelastic form factor only decreases by about 20%.) We use nucleon-nucleus optical potentials mainly in order to estimate the imaginary part of the valence interaction which is usually ignored (see Satchler).[13]

Next, we require the form factors $F_{\beta\alpha}{}^{\lambda}(r_{aA})$ which give the reduced radial dependence of the excitation interaction. The valence term will be calculated microscopically, while the core contribution will be determined empirically. Thus using $\vec{r}_{ac} \simeq \vec{r}_{aA}$ in Eq. (4), we obtain

$$F_{\beta\alpha}{}^{\lambda(\mathrm{core})} = \frac{4\pi Z_a e}{2\lambda+1}\,\frac{E_{\beta\alpha}{}^{\lambda}}{r_{aA}{}^{\lambda+1}} - (\beta R)_{\beta\alpha}{}^{\lambda}\,\frac{d}{dr_{aA}}U_{ac}(r_{aA}). \quad (5)$$

Here $E_{\beta\alpha}{}^{\lambda}$ is the reduced electric multipole matrix element corresponding to Eq. (3). The nuclear deformation lengths $(\beta R)_{\beta\alpha}{}^{\lambda}$ are taken to be proportional to the $E_{\beta\alpha}{}^{\lambda}$ in the usual way. For the case of valence neutrons, Coulomb excitation data directly give the values of the $E_{\beta\alpha}{}^{\lambda}$. Thus for ^{18}O, we can fix $F_{\beta\alpha}{}^{\lambda(\mathrm{core})}$ without specifying the details of the core polarization. Equation (5) has essentially been used in previous studies of ^{18}O, but was not identified as being only the core part of the total form factor.

In evaluating the valence interaction we take wave functions for the 0^+ (g.s.) and $2_1{}^+$ (1.98 MeV) states of ^{18}O from the analysis of Lawson, Serduke, and Fortune.[14] With the approximation $\vec{r}_{ia} \simeq |\vec{r}_{ic} - \vec{r}_{aA}|$ in the nucleon-nucleus interaction, the valence form factor reduces to a sum of single-particle integrals (see Ref. 11). The signs of the core form factors are fixed consistently with the valence ones by requiring the matrix elements in Eq. (5) to have the same signs as the total $E2$ matrix elements of Ref. 14.

It can then be shown explicitly that the valence and core terms combine constructively for both the $0^+ \to 2_1{}^+$ and the $2_1{}^+ \to 2_1{}^+$ transitions. This result can be understood in general for the $0^+ \to 2_1{}^+$ transition within the vibrational model.[2] However, for the $2_1{}^+ \to 2_1{}^+$ transition it depends on the details of the $2_1{}^+$ wave function. In this case the sign of the microscopic form factor is determined by the large weight given to the contribution of the $1d_{5/2}2s_{1/2}$ component.

We have considered the case of ^{18}O + ^{58}Ni measured at $E_{\mathrm{lab}} = 63$ MeV.[5] Optical parameters were fixed by fitting the elastic data.[5] $E2$ matrix elements were taken from Flaum.[15] The nuclear deformations were then fixed as in Ref. 5. The single-particle states were bound at half the two-neutron separation energy in a well taken from the analysis of Cooper, Hornyak, and Roos.[16] The n-^{58}Ni optical potential used for the valence interaction was taken from Becchetti and Greenlees.[17] Our main results are summarized in Figs. 1 and 2(a).

One sees in Fig. 1 that the total nuclear form factors are roughly doubled from the core ones alone over the grazing distances of 10–12 fm. The $2_1{}^+ \to 2_1{}^+$ form factors (not drawn) show the

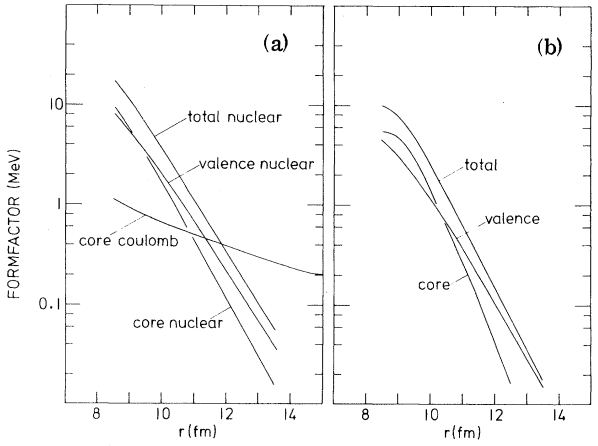

FIG. 1. (a) Real and (b) imaginary parts of the form factors for the reaction ^{58}Ni(^{18}O, ^{18}O($2_1{}^+$)) at $E_{\mathrm{lab}} = 63$ MeV. The core-polarization form factors were obtained macroscopically with deformation parameters derived from electromagnetic properties. The real and imaginary valence form factors were calculated microscopically using a neutron optical potential for the effective interaction.

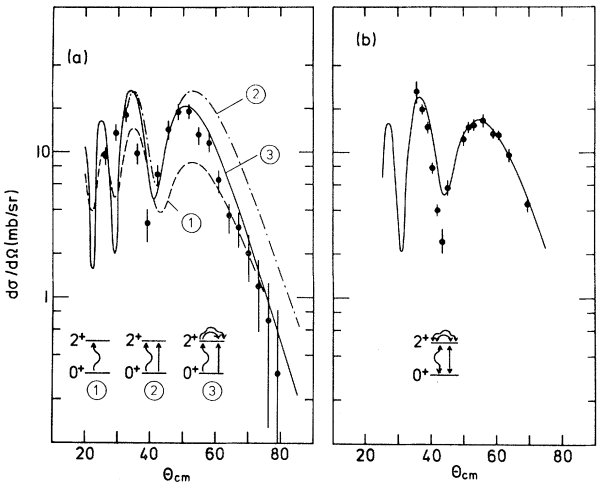

FIG. 2. Angular distributions for the reaction ^{58}Ni(^{18}O, ^{18}O(2_1^+)). (a) E_{lab} = 63 MeV. The data are from Ref. 5. Curve (1) shows the result due to core amplitude, curve (2) includes the valence amplitudes, and curve (3) shows the effect of adding the total reorientation amplitude. The $\lambda = 4$ valence reorientation term was included but it has little effect. (b) Coupled-channels calculation at 60 MeV compared to the data of Ref. 4.

same general features. In Fig. 2(a), the curve which results from the first-order core amplitude shows the well-known discrepancies compared to the data.[5] The last maximum is underpredicted and the interference minimum appears at more backward angles. Including the valence excitation increases the maximum by a factor of 3 and shifts the minimum 2° forward. Without the imaginary part of the valence form factor, there would have been a factor of 2 increase but no shift. Including the second-order reorientation amplitude shifts the minimum an additional 1° forward and shifts and decreases the backward-angle cross section, bringing a reasonable agreement with the data. If the sign of the reorientation amplitude were reversed, the interference minimum would move 1° backwards and there would be a severe disagreement with the backward-angle data.

Other analyses of ^{18}O may be mentioned. Although the complex nature of the microscopic interaction is essential for our results, we do not support the purely imaginary form factors considered in Refs. 5 and 6. Our total form factors go in the direction suggested in Ref. 7 but are not as large as those used empirically in Ref. 7. It has been suggested that a strong two-step transfer contribution ^{18}O-^{17}O-^{18}O* occurs as a result of the many intermediate channels.[8] However, several such transitions have been calculated and were found to be rather small.[9]

The first- and second-order calculations in Fig. 2(a) demonstrate the relative importance of the core and valence amplitudes and the effects of the reorientation, but a coupled-channels (CC) analysis should be compared with the data. This requires readjusting optical parameters to fit the elastic scattering. We will do this later as part of a more systematic study. We have, however, compared to the CC analysis of Ref. 4 for the same reaction at 60 MeV. These authors used a nuclear deformation which was 1.85 times larger than normal for the $0^+ \rightarrow 2_1^+$ transition and a reorientation amplitude which was twice again as large. We see that the factor 1.85 roughly substitutes for the valence amplitudes and that their reorientation is too large.[15] We repeated the CC calculation of Ref. 4 using the same optical potential, normal deformation parameters, and the present valence form factors. The result, shown in Fig. 2(b), is in good agreement with the data. Thus we firmly establish the importance of including single-particle amplitudes based on complex microscopic interactions and support the conclusion that the shift of the interference pattern depends on the sign of the quadrupole moment of the 2_1^+ state.[3,4]

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Shape-Resonance-Enhanced Vibrational Excitation at Intermediate Energies (10–40 eV) in Electron-Molecule Scattering

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Shape-resonant enhancements in the kinetic-energy range 10 to 40 eV are predicted to be a common feature in electron-molecule scattering. These features are usually difficult to detect in elastic and total cross sections, but can emerge clearly in vibrational excitation. A calculation of the e^--CO_2 symmetric-stretch spectrum, just confirmed experimentally by Tronc *et al.*, is given as an example. Such resonances can be predicted by analysis of the corresponding inner-shell photoionization spectra.

The importance of shape resonances in the low-energy portion (< 10 eV) of the electron-molecule scattering spectrum is widely appreciated.[1] The π_g d-wave ($l = 2$) resonance at 2.4 eV in e^--N_2 scattering is perhaps the most thoroughly studied example.[2,3] Such resonances stand out clearly above the nonresonant background, have widths on the order of an eV, and are usually dominated by a single asymptotic angular momentum, which imposes a well-defined angular distribution characteristic of that l. In this Letter we emphasize a new class of shape resonances, falling in the intermediate-energy range (10–40 eV), which have emerged in the course of our current survey studies of electron-molecule scattering, using the multiple-scattering method (MSM).[4,5] Compared to their low-energy counterparts, these resonances are broader, are usually barely detectable in the elastic or total scattering cross sections, and often are composed of significant contributions from several asymptotic angular momenta, whose mixture can vary significantly within the resonance. These properties are consistent with the picture of a quasibound state lying near the top of its centrifugal barrier. These weak resonances are, however, very sensitive to changes in the molecular geometry and can therefore couple significantly to nuclear motion. We have found that these features are thereby enhanced in vibrational excitation and appear prominently, all the more so because the nonresonant background is negligible in the vibrationally inelastic channels.

Experimental evidence is sparse, possibly because shape-resonant structure in the 10–40-eV range was not expected and measurements were concentrated instead in the low-energy range. Several years ago, Pavlovic *et al.*[6] reported a broad enhancement in e^--N_2 vibrational excitation in the 15–35-eV range. We have shown[7] that this

