Probing monopole double giant resonances by dilepton ($E0$) emission

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High lying excitations consisting of particle-hole plus two-particle two-hole configurations coupled to the continuum are analyzed in ^{208}Pb. Partial decay widths are calculated and the effect of the two-particle two-hole excitatons upon the widths are assessed. It is found that double giant resonance excitations are not mixed with other degrees of freedom and that double monopole giant resonances are likely to be detected in ^{208}Pb by measuring dilepton ($E0$) emission. [S0556-2813(98)03602-4]

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Giant resonances provide an invaluable tool to analyze excitations embedded in the continuum part of nuclear spectra. The gross properties of these resonances can be studied suitably by means of bound (e.g., harmonic oscillator) representations [1,2], but when properties closely linked to unbound states, like partial decay widths, are analyzed, the continuum has to be considered from the outset. This is usually achieved in the particle-hole (ph) case by means of the continuum RPA [3,4] extending the energy domain to complex values [5,6]. The coupling of particle-hole to two-particle two-hole excitations (2p2h) would give rise to the spreading width of the giant resonance [7,8]. Besides, the inclusion of the 2p2h excitations would allow us to study two-phonon states at high energies, that is the excitation of the double giant resonances. This is plausible since giant resonances are the most collective vibrational states in nuclei and their double excitations are more likely to occur than the double excitation of low energy vibrational states such as, e.g., the yrast quadrupole states in tin isotopes.

The study of double phonon excitations at high energies has recently been the subject of considerable interest, both theoretically [9–13] and experimentally [13–19]. The double dipole giant resonance (GDR2) was observed in ^{208}Pb by measuring in coincidence the two γ rays that ensue from the decay of the GDR2 to the GDR and from this to the ground state [14]. In this nucleus the cross section in heavy ion inelastic scattering leading to the GDR2 was recently measured [15]. In ^{136}Xe the GDR2 was detected by measuring the emission of neutrons from highly excited states and the subsequent gamma decays [16], while in lighter isotopes traces of the GDR2 were found by analyzing the proton decay pattern from states lying high in the spectrum [17]. Even the deexcitation of the double monopole giant resonance through $E0$ dilepton emission was recently experimentally explored [19].

This intense experimental activity was well matched by theoretical analysis of the origin and mechanisms of decay of double giant resonances. There are a number of calculations where two-particle two-hole excitations are incorporated through bound states [8]. One serious problem in these cases is that the dimensions of the shell-model basis are very large, since very high lying configurations have to be considered. This problem becomes even more acute if the continuum is also included.

The coupling of the bound configurations to the continuum was studied in Ref. [7]. Here the continuum is exactly included, but only in the intermediate particle propagator. To account for the experimental resolution and to reduce the numerical effort an energy-averaging parameter of 400 keV is introduced.

In this paper we will present a formalism to study ph plus 2p2h excitations within the multistep shell model method (MSM) [20] and including the continuum by means of the Berggren representation [6]. After a brief presentation of the formalism (details can be found in the references) we will apply it to calculate excitations in the nucleus ^{208}Pb. In particular, we will analyze the decay of the monopole giant resonance by dilepton decay. This, which has not be done so far, can be an important tool to detect the monopole giant resonance as well as the corresponding double giant resonance.

As a representation to span the shell model space we choose a set of single-particle states which satisfy outgoing boundary conditions corresponding to a Woods-Saxon potential. This set includes bound states and the so-called Gamow resonances. A complete set in the Berggren representation should also include scattering states, as described in Ref. [21]. But the approximation of neglecting the scattering states has been shown to provide a good description of physical properties [6] especially those connected with the resonance behavior of the cross section.

Since wide Gamow states can be considered as part of the proper continuum [22] we included in the basis only bound states and resonances which are not wider than 200 keV. Already this represents a drastic truncation of the basis because most of the high lying states are wide.

The correlated states that result as the coupling of collective excitations can be studied conveniently within the MSM. In this method one solves the shell model problem in several consecutive steps. In each step one describes a given system in terms of quantities related to the systems evaluated in the previous steps [20]. The resulting violations of the Pauli principle as well as the overcounting of independent shell-model states are corrected by using the overlap matrix among basis states. In our case we have to evaluate the particle-particle (pp), i.e., $|\alpha_2\rangle = P^+(\alpha_2)|0\rangle$, hole-hole (hh), i.e., $|\alpha_{-2}\rangle = P(\alpha_{-2})|0\rangle$, and particle-hole (ph), i.e., $|\alpha_0\rangle$

$=Q^{+}(\alpha_0)|0\rangle$ excitations. The corresponding MSM basis will have, schematically, the form

$$|\alpha_0\rangle = Q^{+}(\alpha_0)|0\rangle = |{}^{208}\mathrm{Pb}\rangle,$$

$$|\beta_0\gamma_0\rangle = Q^{+}(\beta_0)Q^{+}(\gamma_0)|0\rangle = |{}^{208}\mathrm{Pb}\otimes{}^{208}\mathrm{Pb}\rangle, \qquad (1)$$

$$|\alpha_2\alpha_{-2}\rangle = P^{+}(\alpha_2)P(\alpha_{-2})|0\rangle = |{}^{210}\mathrm{Pb}\otimes{}^{206}\mathrm{Pb}\rangle$$

or

$$|{}^{210}\mathrm{Po}\otimes{}^{206}\mathrm{Hg}\rangle.$$

The MSM equations are not complicated for this case. The matrix elements between the 2p2h states (e.g., the overlap $\langle\alpha_0\beta_0|\alpha_2\alpha_{-2}\rangle$) have the same form as the four-particle case of Ref. [23]. One can derive the matrix elements that connect ph with 2p2h states (''scattering vertices'') in a similar fashion.

The nth MSM state carrying angular momentum λ will then have the form

$$|n\lambda\rangle = \sum_{\alpha_0} X(\alpha_0;n\lambda)Q^{+}(\alpha_0)|0\rangle + \sum_{\beta_0\leq\gamma_0} X(\beta_0\gamma_0;n\lambda)(Q^{+}(\beta_0)Q^{+}(\gamma_0))_\lambda|0\rangle + \sum_{\alpha_2\alpha_{-2}} X(\alpha_2\alpha_{-2};n\lambda)(P^{+}(\alpha_2)P(\alpha_{-2}))_\lambda|0\rangle, \qquad (2)$$

where Greek letters label the correlated states as well as the corresponding angular momenta. Note that the angular momentum of the state labeled by α_0 is λ.

The MSM wave function amplitudes X are usually not well defined quantities, since the MSM basis may be overcomplete. But the projections of the MSM states onto the basis vectors, e.g., $\langle n\lambda|(Q^{+}(\beta_0)Q^{+}(\gamma_0))_\lambda|0\rangle$, are well defined. The equivalent to the wave function component in an orthogonal basis is the cosine of the angle θ between the calculated vector and the basis component. For instance, for a basis element of the form $Q^{+}Q^{+}$ that angle θ is defined by

$$\cos\theta = \frac{\langle n\lambda|(Q^{+}(\beta_0)Q^{+}(\gamma_0))_\lambda|0\rangle}{[\langle 0|(Q^{+}(\beta_0)Q^{+}(\gamma_0))_\lambda|(Q^{+}(\beta_0)Q^{+}(\gamma_0))_\lambda|0\rangle]^{1/2}}. \qquad (3)$$

This angle provides information about the structure of the calculated states, as seen below.

The ph states have to be calculated by using the RPA equations. However, at high energies the corresponding backward amplitudes are usually small and therefore one can use the TDA equations instead [24]. This is actually important in the MSM, since the TDA metric allows one to define the angle between two vectors, as given by Eq. (3). For the isoscalar 1^- states, however, we use the RPA in order to isolate the spurious state at zero energy. For the other 1^- states we neglect the backward amplitudes and renormalize the wave functions to obtain the TDA metric.

We use a separable residual interaction, as in Refs. [6,25,26]. The field in the interaction is the same as the one used in the transition operator and the radial part is the derivative of the Woods-Saxon potential [6].

With this interaction we calculated the TDA ph states in ^{208}Pb, the pp states in ^{210}Pb and ^{210}Po, and the hh states in ^{206}Pb and ^{206}Hg. We found a reasonable agreement between our calculation and the available experimental data. Even the resulting wave functions agree fairly well with those provided by the calculation of Ref. [27], where only one major harmonic oscillator shell was included in the single particle representation. This shows that the effect of high lying configurations does not affect the low lying states significantly.

The output of the calculation is very large. Therefore we will only present in this paper cases for which both the giant resonances and the double giant resonances are of interest at present.

The physical quantities to be compared with the experimental data have to be real quantities. However, our formalism generally provides complex results and one is confronted with the problem of having to interpret complex probabilities. In Ref. [6] this was associated with the fact that the extension of propagators to the complex energy plane [28] is justified only if the corresponding complex poles are close to the real energy axis. That is, only isolated resonances can be associated with physical states [29]. Therefore, a complex number that can be related to a probability, such as e.g., a complex partial decay width, is a manifestation that the corresponding resonance is not isolated. The larger is the imaginary part of that complex number the more is the resonance intermingled in the continuous background.

Berggren found a somewhat similar interpretation of complex probabilities [30]. Within this interpretation the imaginary part of an expectation value of an operator representing a physical quantity (e.g., a cross section) is a measure of the coupling of the resonance with the continuum. Even more akin with this paper is the result that the imaginary part of the matrix element of a Hermitian operator in a resonant state is related to the uncertainty of the real part, which is the expectation value of the operator concerned [31]. Thus, for the Hamiltonian the imaginary part of the complex energy eigenvalue is the decay width.

In order to verify that the main features of the calculation were not altered by neglecting the backward components, we first calculated the EWSR corresponding to the monopole TDA case, i.e., without any mixing with 2p2h excitations. We thus found that the isoscalar monopole giant resonance (GMR) lies at about 13 MeV, but that it is fragmented already at the TDA level, in the same fashion as in the corresponding RPA case [6].

The agreement between the TDA and the RPA calculations was also found for the other resonances. In particular, we confirmed that the isovector dipole giant resonance is at about 13.6 MeV, just in the same region as most of the isoscalar monopole EWSR is concentrated. But the new feature that we can analyze now is the spreading of the giant resonance into the neighboring 2p2h components. In our case, this will be shown by the values of $\cos\theta$, where θ is the angle between the calculated state $n\lambda$ and the vector $|\mathrm{GDR}\rangle$. The giant dipole resonance that we discuss here is very isolated [6] and therefore the imaginary parts of all physical quantities, including $\cos\theta$, are small. We thus found that most of the 2p2h components interfering with

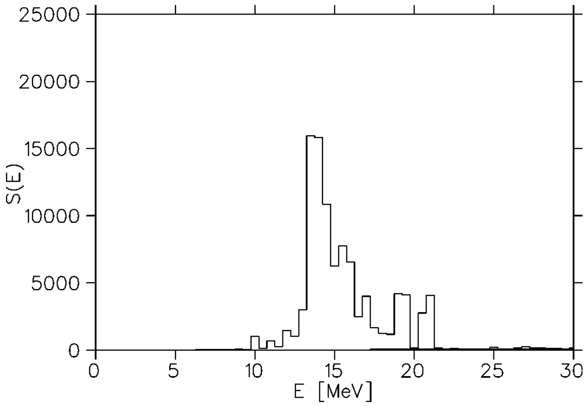

FIG. 1. Strength distribution corresponding to the ^{208}Pb isoscalar monopole mode in arbitrary units.

the GDR are of the forms $(Q^+(2^+_n)Q^+(3^-_m))_{1^-}|0\rangle$, $(Q^+(3^-_n)Q^+(4^+_m))_{1^-}|0\rangle$, and $(Q^+(4^+_n)Q^+(5^-_m))_{1^-}|0\rangle$, where n and m label the correlated ph states. For the GMR, instead, the majority of states contributing to the splitting of the wave function are of the pairing form, i.e., $|\pi\rangle = (P^+(0^+_n)P(0^+_m))|0\rangle$. This is due to the fact that there is a large number of such states in the region of the isoscalar monopole giant resonance and that the angular momentum recoupling coefficients are relatively large since all angular momenta involved are zero. Moreover, there is not any Pauli blocking in the MSM basis states of the form $|\pi\rangle$, since the corresponding overlap matrix is the unit matrix.

To analysis the effect of 2p2h excitations on the width of the GMR we evaluated the strength function $S(E)$ [26], which is essentially the imaginary part of the response function.

As seen in Fig. 1, the GMR strength function does not show the pronounced peaks that one obtains when only ph excitations are allowed, as provided by the continuum RPA [6,26]. This is due to the fragmentation of the GMR into 2p2h components. These many pieces contribute to the spreading width, as seen in Fig. 1. The experimental distribution corresponding to the inelastic excitation of the GMR [32] is well reproduced by Fig. 1.

The EWSR follows a similar pattern as the one seen in Fig. 1. It is worthwhile to point out that the EWSR corresponding to ph excitations is concentrated in states lying at energies which are about the same as the peaks in Fig. 1. As can be seen by comparing with the ph calculation shown in Table I of Ref. [6], the large bump at 13.6 MeV corresponds to the main piece of the ph EWSR, while the bumps at about 20 MeV are due to the pieces of the ph EWSR at about 21 MeV. This rather high energy value of the EWSR (or the strength in Fig. 1) is a result of the coupling between the isoscalar and isovector modes in this heavy nucleus.

It is important to know whether there are states for which the dominant MSM component is of the form $|\mathrm{GR2}\rangle = (Q^+(\mathrm{GR}_1)Q^+(\mathrm{GR}_2))_\lambda|0\rangle$, where GR_i labels a giant resonance. We have found a few states where this is indeed the case. Thus for the monopole case there are the three states

$$|0^+(E=(27.332-0.543i)\ \ \mathrm{MeV})\rangle = (0.975-0.068i)|\mathrm{GMR}\otimes\mathrm{GMR}\rangle,$$

$$|0^+(E=(27.292-0.209i)\ \ \mathrm{MeV})\rangle = (1.000+0.000i)|\mathrm{GDR}\otimes\mathrm{GDR}\rangle,$$

$$|0^+(E=(20.741-0.144i)\ \ \mathrm{MeV})\rangle = (0.998-0.002i)|\mathrm{GQR}\otimes\mathrm{GQR}\rangle,$$

while for the quadrupole case they are the three following states:

$$|2^+(E=(27.257-0.199i)\ \ \mathrm{MeV})\rangle = (0.991-0.001i)|\mathrm{GDR}\otimes\mathrm{GDR}\rangle,$$

$$|2^+(E=(20.737-0.145i)\ \ \mathrm{MeV})\rangle = (0.999+0.000i)|\mathrm{GQR}\otimes\mathrm{GQR}\rangle$$

$$|2^+(E=(23.970-0.286i)\ \ \mathrm{MeV})\rangle = (1.000+0.000i)|\mathrm{GMR}\otimes\mathrm{GQR}\rangle,$$

where GQR labels the isoscalar quadrupole giant resonance.

The widths of these double giant resonances are about $\sqrt{2}\times(\Gamma_1\Gamma_2)^{1/2}$, where Γ_i is the width of the giant resonance GR_i, as expected for such pure states [11]. This feature was confirmed experimentally [14] for the case $\mathrm{GR}_1=\mathrm{GR}_2=\mathrm{GDR}$.

When the main component of the GR2 is $|\mathrm{GMR}\otimes\mathrm{GMR}\rangle$ we considered the possibility of observing its decay through dilepton $(E0)$ emission.

Monopole excitations can be induced by heavy-ion compound fusion reactions. But this populates also other giant resonances, especially the isoscalar giant quadrupole resonance and even the GDR. The deexcitation of these resonances proceeds mainly by photon emission but also by emission of dileptons with energy spectra that is rather similar for all excitations. In principle one can extract the dilepton yield corresponding to giant resonances other than monopole, i.e., with multipolarities $\lambda \geqslant 1$, by measuring simultaneously the dilepton and photon spectra. Such a method has been used recently in analyzing the dilepton decays in ^{28}Si [19]. However, for this light isotope the excess of dileptons corresponding to $E0$ transitions is very small. This implies that, with present experimental resolution techniques, the separation of the $E0$ dilepton yield is a very difficult undertaking. But the probability of $E0$ dilepton

TABLE I. Partial decay widths $\Gamma_{n=\mathrm{GDR},ch}$ (in MeV) corresponding to the decay of the GDR calculated within the TDA and within the MSM, i.e., including ph as well as 2p2h excitations. The first column, labeled ch, shows the spin and parity of the neutron hole channels (states in ^{207}Pb) to which the GDR decays.

ch	TDA	MSM
$9/2^-$	$(0.002,-0.001)$	$(0.004,-0.001)$
$7/2^-$	$(0.107,-0.007)$	$(0.118,-0.084)$
$13/2^+$	$(0.000,-0.000)$	$(0.000,-0.000)$
$3/2^-$	$(0.051,-0.008)$	$(0.063,0.039)$
$5/2^-$	$(0.091,-0.011)$	$(0.109,0.070)$
$1/2^-$	$(0.053,-0.008)$	$(0.066,0.042)$

emission increases strongly with Z and with the energy of the transition. We have found that for the case of ^{208}Pb at giant resonance energies, the ratio between the probability of emitting $E1$ photons from the GDR and the probability of $E0$ dileptons emission from the GMR is about 40. Taking for the $E1$ transition a pair production branching of about 10^{-3} [33], one obtains that the $E0$ dilepton yield is larger than the one coming from $E1$ transitions. This is an important result since it shows that for ^{208}Pb dilepton detection may be a proper tool to detect the decay of giant monopole excitations, very particularly the double GMR.

We have also calculated the partial decay widths $\Gamma_{n,ch}$ of the resonances. For the GDR the quantities $\Gamma_{n,ch}$ are sound physical quantities, since they are practically real numbers within the continuum RPA [6] as well as in our TDA case, as seen in Table I. It is interesting to investigate whether such a feature is still present when the mixing with 2p2h excitations is included. We present in Table I the calculated partial decay widths including only ph excitations (TDA) and ph plus 2p2h excitations (MSM). One sees in this table that indeed the 2p2h mixing tends to increase the imaginary parts of the partial decay widths. This indicates that even when ph calculations predict detectable values for physical quantities, more complicated excitations tend to spoil that result. In our case, however, the imaginary parts are still rather small and the difficult undertaking of measuring partial decay widths [34] may still be a sensible task for the GDR in ^{208}Pb. This has not been done yet, although some exploratory attempts have already been considered [35].

In conclusion we have studied in this paper the coupling of the continuum to particle-hole plus two-particle two-hole excitations. We have found that that coupling contributes significantly to the width of the giant resonances as well as to the corresponding particle decay. However, we have found that double giant resonances are not mixed with other degrees of freedom and, as a result, their width is about $\sqrt{2}$ times the width of the corresponding giant resonance. We have studied the decay pattern of the double giant resonance excitations and found that in ^{208}Pb the monopole giant resonance and even the double monopole giant resonance can be detected by measuring dilepton (e^+,e^-) emission.

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