Exit Doorway Model for Nuclear Breakup of Weakly Bound Projectiles

M. S. Hussein\textsuperscript{1, 2}, R. Lichtenthaler\textsuperscript{2}

\textsuperscript{1}Max-Planck-Institut für Physik komplexer Systeme, Nöthnitzer Str. 38, D-01187, Dresden, Germany
\textsuperscript{2}Instituto de Fisica, Universidade de São Paulo, C.P. 55318, 05315-970, São Paulo, SP, Brazil

We derive closed expressions for the nuclear breakup cross sections in the adiabatic limit using the Austern-Blair theory. These expressions are appropriate for the breakup of weakly bound nuclei. The concept of an exit doorway that mediates the coupling between the entrance channel and the breakup continuum is used. We prove the validity of the scaling law that dictates that the nuclear breakup cross section scales linearly with the radius of the target. We also compare our results for the nuclear breakup cross section of $^{11}$Be, $^{8}$B on several targets with recent CDCC calculation.

PACS numbers: 25.60.Dz, 25.70.De, 24.10.Eq

The breakup of nuclei is a common occurrence when the bombarding energy is high enough and/or the binding energies are sufficiently low. In the case of weakly bound nuclei the threshold for breakup is small and more so for bound unstable nuclei. The mechanism of breakup is assumed to consist of elongating the projectile, through the action of the interaction, which eventually leads to the production of two or more fragments. This interaction is composed of a short range, nuclear piece and a longer ranged electromagnetic one. A debate has been going on in the literature concerning the way the nuclear part of the breakup cross section depends on the mass of the target nucleus which supplies the interaction. In most references [1, 2, 3], it is assumed that the dependence goes as the cubic root of the mass number. In reference [4], however, it is claimed that this dependence is more like linear! In a recent paper [5], through a careful Continuum Discretized Coupled Channels (CDCC) calculation, the former dependence ($A^{1/3}$) has been established, which corroborates the contention that the nuclear breakup cross section should follow the prediction of the Serber model [6].

It is interesting to compare the numerical CDCC calculation alluded to above with those of simpler analytical models. Specifically, the Austern-Blair adiabatic theory for inelastic scattering comes to mind. If one assumes that the breakup proceeds through a so-called exit doorway [7, 8, 9], then the process can be treated as an inelastic excitation. The idea of exit doorway has been used in the case of the influence of breakup on fusion [7, 8] and in the excitation of giant resonances [9] with success. In a very recent paper [10], a comparison of a preliminary CDCC calculation for the breakup cross section of the system $^{4}$He+$^{27}$Al at low bombarding energies with a simple formula derived by using the Austern-Blair model, showed that such an idea is quite reasonable and encouraged us to pursue the matter further. We do this in the present paper, where we fully develop the Austern-Blair model for the nuclear elastic breakup reaction cross section assumed to proceed through the excitation of an exit doorway [7, 8, 9].

The exit doorway concept has been used in the development of reaction theories involving the the excitation of a doorway in the final state, in contrast to the conventional cases where such resonances are populated in the entrance channel [11]. In the breakup reactions of halo nuclei one may envisage that the process proceeds through the breakup doorway dipole, quadrupole etc.) into the continuum. As such, the detailed description of the exclusive reaction, where the final channels are specified, will necessarily contain the full information about the exit doorway ( its energy, width etc.). This is the case that was encountered in the theory of the excitation of multiple giant resonances [8] and of the influence of the pygmy resonance on the fusion of halo nuclei [7]. In the current paper we will be content with the inclusive quantity of the integrated breakup cross section and the only reference to the exit doorway is made implicitly as a final state that has to be populated for breakup to occur.

The full Hamiltonian which describes the colliding ions can be written as

$$H = H_{0} + F$$

(1)

where $H_{0} = h_{0} + K + V = h_{0} + H^{(0)}$ is diagonal in open channel space, $h_{0}$ is the intrinsic part that describes the structure of the projectile and the target nuclei, $K$ is the kinetic energy operator and $V$ is the optical potential which contains the complex nuclear plus the Coulomb parts. The operator $F$ describes the coupling among the open channels.

The intrinsic Hamiltonian $h_{0}$, which for simplicity is taken here to represent the excitable projectile nucleus with the target considered structureless, is now written as:

$$h_{0} = |\phi_{0} > E_{0} < \phi_{0}| + |d > E_{d} < d| + \sum_{i} |i > E_{i} < i| + \sum_{i} |d > \Delta_{i} < i| + |i > \Delta_{i}^{*} < d| + \sum_{ij} |i > \Omega_{ij} < j| + cc$$

(2)

The first three terms on the RHS above refer to the ground, exit doorway and discretized continuum states,
The doorway) the elastic channels coupling to the breakup
continuum-continuum coupling. Clearly the need to the continuum-continuum coupling
enters Coupled Channels (CDCC) intrinsic Hamiltonian

\[ h_0 = |\phi_0 > E_0 < \phi_0| + \sum_i |i > E_i < i| + \sum i |i > \Delta_i < i| + \sum ij |i > \Omega_{ij} < j| + cc \] (3)

The exit doorway modulated CDCC Hamiltonian, Eq. (2), is our subject of study here. A full development of this new CDCC will be left for a future work. Here we concentrate our effort on understanding the consequence of reaching the breakup continuum from the entrance channel only through the exit doorway \(|d >\). For this purpose we ignore the last term in Eq. (2) and remind ourselves that, whereas \(|\phi_0 >\) and \(|i >\) are eigenstates of \(h_0\), \(|d >\) is not.

The full doorway-modulated CDCC equations can be obtained as follows. The full Schrodinger equation of the colliding system is,

\[ [E - (H_0 + F)]|\psi > = 0 \] (4)

which when projecting onto the different channels gives:

\[ (E - E_0 - H_0^{(0)})|\psi_i^+ > = \Xi_0|\psi_i^+ > \] (5)

\[ (E - E_i - H_i^{(0)})|\psi_i^+ > = F_{i0}|\psi_0^+ > \] (6)

We now invoke the exit-doorway hypothesis,

\[ F_{i0} = F_{0i}\alpha_{id}^* \quad F_{i0} = F_{0i}\alpha_{id}^* \] (7)

The overlaps \(\alpha_{di}\) and \(\alpha_{id}^*\) and can be easily obtained from Eq.(2)(without the last term). \[\alpha_{di}|^2 = (\Gamma^d_i/2\pi)/[(E_i - E_d)^2 + (\Gamma^d_i/2)^2] \] (8)

where \(\Gamma^d_i\) the exit-doorway spreading width describing its average coupling to the continuum states of the projectile, is related to the \(\Delta_i\) factors through,

\[ \Gamma^d_i = 2\pi|\Delta_i|^2\rho \] (9)

where \(|\Delta_i|^2\) is an average value and \(\rho\) is the average density of discretized continuum states in the vicinity of \(d\). Clearly the need to the continuum-continuum coupling terms would be very important if exclusive cross sections are to be calculated, since through them (and through the doorway) the elastic channels coupling to the breakup channel continuum can be fully accounted for. Including the C-C coupling term, would result in a more complicated expression for \(|\alpha_{di}|^2\) than that of Eq\[\Box\]

Equations (5) and (6) can be recast into the following, after setting \(E_0 = 0\) and \(F_{ij} = 0\),

\[ (E - H^{(0)}_0)|\psi_0^+ > = \alpha_{di}F_{0i}|\psi_i^+ > \] (10)

\[ (E - E_i - H_i^{(0)})|\psi_i^+ > = \alpha_{id}^*F_{i0}|\psi_0^+ > \] (11)

The breakup cross section, within the exit doorway model then becomes,

\[ \sigma_{bu} = \sum i \frac{k_i}{k_0}|\alpha_{di}|^2 < |\psi_i^+ > |F_{i0}|\psi_0^+ > |^2 \] (12)

\[ \approx \frac{k_d}{k_0} < |\psi_i^+ > |F_{i0}|\psi_0^+ > |^2 \] (13)

Where the sum over \(i\) has been performed by appropriate contour integration over \(E_i\). Note that the Q-value in \(\psi_d^-\) is complex owing to the non-zero width of the exit doorway whose energy is \(E_d - i\Gamma_d^2/2\). A simple way to see how the complex Q-value arises is to eliminate \(\psi_i^+\) in Eq.(5) in favor of \(\psi_0^+\) by employing Eq.(6), which gives \(\psi_0^+ = \frac{1}{E_i - E_d - H_d^{(0)} + i\epsilon}F_{i0}|\psi_0^+ >\). With this Eq.(5) becomes

\[ (E - E_0 - H_0^{(0)} - \sum_i F_{i0}) \frac{1}{E_i - E_d - H_d^{(0)} + i\epsilon}F_{i0}|\psi_0^+ > = 0 \] with the exit doorway hypothesis, the polarization potential contribution, \(\sum_i F_{i0} \frac{1}{E_i - E_d - H_d^{(0)} + i\epsilon}F_{i0}\) becomes

\[ \sum i F_{i0} \frac{\Gamma_d^2/2\pi}{(E_i - E_d)^2 + (\Gamma_d^2/2)^2} \frac{1}{E_i - E_d - H_d^{(0)} + i\epsilon}F_{i0} \approx \frac{F_{0d}}{E - (E_d - i\Gamma_d^2/2) - H_d^{(0)} + i\epsilon}F_{0d} \]

This suggests defining the exit-doorway scattering wave function by setting \(H_i^{(0)} = H_d^{(0)}\) such that Eqs.(10) and (11) become:

\[ (E - H^{(0)}_0)|\psi_0^+ > = F_{0d}|\psi_d^+ > \] (14)

\[ (E - (E_d - i\Gamma_d^2/2) - H_d^{(0)})|\psi_d^+ > = F_{0d}|\psi_0^+ > \] (15)

The “inelastic” cross-section is thus given by Eq.(13) above with the aforementioned proviso that the Q-value of the excited state is complex. The width of this Q-value is a measure of the continuum contribution to the coupling.

At this point we comment on the inclusion of the continuum-continuum coupling, namely the last term in Eq.(3). In this situation the amplitudes \(\alpha_{di}\) are obtained by matrix diagonalization and, among other things, the resulting overlap probability \(|\alpha_{di}|^2\), deviates from the
Breit-Wigner form of Eq.(8). A possible form which may incorporate some of the c-c effects is a Lorentzian:

\[
|\alpha_{dl}|^2 = \frac{2}{\pi} \frac{\Gamma_d^2 E_d^2}{(E_d^2 - E_a^2) + \Gamma_d^2 E_d^2}
\]

The above form results in an equation for \(\psi_d^{(\pm)}\) with a modified form factor which depends on the position and width of the exit doorway \(\tilde{V}_{d0} \approx f(E_d, \Gamma_d^1)V_{d0}\) where \(f(E_d, \Gamma_d^1)\) is generally complex. Accordingly the cross-section would be: \(\sigma = |f(E_d, \Gamma_d^1)|^2 \sigma_{DWBA}\). In the limiting case of \(\Gamma_d \ll E_d\), the factor \(f(E_d, \Gamma_d^1)\) is approximately given by \(\sigma \approx (1 + \frac{\Gamma_d^1}{2E_d})\sigma_{DWBA}\). In the case of coupling to the breakup continuum considered here, the other limit, \(\Gamma_d \gg E_d\) is more appropriate, as \(E_d\) is roughly given by the Q-value of the breakup (\(\leq 1\)MeV) while \(\Gamma_d^1\) measures the extent in continuum excitation the discretization is performed (\(\approx 10\) MeV). The function \(f(E_d, \Gamma_d^1)\) can be calculated in such a situation, but we leave this for a future investigation. The important point we are making here, is that a DWBA calculation with complex excitation energy in the final state, and with a form factor of the type \(f(E_d, \Gamma_d^1)V_{d0}\), should be an adequate candidate to treat the elastic breakup process.

In the following we take the exit doorway to be excited states of different multipoles and use the Austern-Blair sudden/adiabatic theory \[12, 13\]. We employ the Distorted Wave Born Approximation for \(\psi_0^{(+)}\) and \(\psi_d^{(-)}\).

The elastic breakup cross section and its dependence on the target mass can be analysed within the Distorted Wave Born Approximation (DWBA). If we treat the breakup as an inelastic multipole process, the amplitude \(T_{LM} = \langle \psi_d^{(-)} | F_{d0} | \psi_0^{(+)} \rangle\) would look like:

\[
T_{LM} = \sum_{l_f} (2l_f + 1)^{1/2} \int |l_f|<l_f L;00{|l_f|} > R_{l_f,l_i}(k_f, k_i) \left( e^{i\sigma_{l_f}(k_f)} + e^{-i\sigma_{l_i}(k_i)} \right) Y_{l,-M}(\theta, \phi). \quad (16)
\]

The unpolarized cross section of the dipole transition is then obtained from the expression

\[
d\sigma_L/d\Omega = \frac{\mu}{(2\pi\hbar^2)^2} \frac{k_f}{k_i} \sum_{M=-L}^{M=L} |T_{LM}|^2 \quad (17)
\]

The radial integrals \(R_{l_f,l_i}(k_f, k_i)\) for pure nuclear excitation are given by,

\[
R_{l_f,l_i}(k_f, k_i) = (4\pi/k_fk_i) \int dr f_{l_f}(k_f, r) F_L(r) f_{l_i}(k_i, r)
\]

where the form factors \(F_L(r)\) are given by the following expressions for the monopole, \(L = 0\), dipole, \(L = 1\) and quadrupole excitations \(L = 2\) \[14\].

\[
F_0(r) = -\delta_0^{(N)}(3V(r) + r \frac{dV(r)}{dr}) \quad (19)
\]

\[
F_1(r) = -\delta_1^{(N)}(\frac{3}{2}\frac{\Delta R_P}{R_P}) \left( \frac{dV(r)}{dr} + \frac{R}{3} \frac{d^2V(r)}{dr^2} \right) \quad (20)
\]

\[
F_2(r) = -\delta_2^{(N)} \frac{dV(r)}{dr} \quad (21)
\]

with \(\Delta R_P = R_n - R_p\) being the difference between the rms radii of the neutron and proton distributions of the projectile, and \(V(r)\) is the elastic scattering channel optical potential. The quantities \(R_n\) and \(R_p\) can be extracted from the analysis of Refs. \[16, 18\]. In Ref. \[18\] a power expansion in \(\Delta R_P\) was employed in the analysis of \(\alpha\)-inelastic scattering from neutron skin nuclei.

In the adiabatic limit, \(k_i = k_f = k\), and for large orbital angular momenta, \(l_f = l_i = l\), the radial integral can then be evaluated in closed form following the procedure of Austern and Blair \[12, 13\]. For the dipole and quadrupole cases we have

\[
R_{1,l}^{(1)}(k) = -i\delta_1^{(N)}(\pi \hbar^2/\mu) \frac{3}{2} \frac{\Delta R_P}{R_P} \times \left( \frac{dS_1^{(N)}(k)}{dl} + \frac{R}{3} \frac{d^2S_1^{(N)}(k)}{dl^2} \right) \quad (22)
\]

\[
R_{1,l}^{(2)}(k) = -i\delta_2^{(N)}(\pi \hbar^2/\mu) \frac{dS_2^{(N)}(k)}{dl} \quad (23)
\]

where \(\delta_1^{(N)}\) is the nuclear deformation length given by \(\delta_1^{(N)} = \beta^{(N)}_L R_P\) with \(\beta^{(N)}_L\) being the nuclear deformation parameter and \(R_P\) is the radius of the excited projectile.

The above expression for the radial integrals can be associated with the nuclear elastic breakup radial integral. Thus we can obtain analytical expression for the integrated nuclear breakup cross section by simply integrating the cross section formula, eq.(2). In performing this calculation the angular momentum coupling coefficients are evaluated exactly and the sum over \(l_i\) can be performed by using the Coulomb phase shifts both as functions of \(l_f \equiv l\). The amplitude of eq. (13) is given now by

\[
T_{LM} = i\sqrt{2} \sum_{l=0}^{\infty} (2l+1)^{1/2} R_{l_f,l_i}^{(k)}(k)e^{2i\sigma_{l_i}(k)}Y_{l,-M}(\theta, \phi) \quad (24)
\]

with the condition that \(T_{LM} = 0\) if \(L + M\) is odd. The integrated pure nuclear breakup cross section containing dipole and quadrupole contributions then becomes the following

\[
\sigma = \left[ (\delta_1^{(N)})^2 + (\delta_2^{(N)})^2 \right] \sum_{l=0}^{\infty} (2l+1) \left( \frac{dS^{(N)}(k)}{dl} \right)^2 \quad (25)
\]
where terms proportional to the second derivative of $S^{(N)}_l(k)$ have been dropped.

A simple estimate of the above formula can be made by approximating the sum in $l$ by an integral in $\lambda = l + 1/2$:

$$
\sum_{l=0}^{\infty} (2l + 1) \frac{dS^{(N)}_l(k)}{dl}^2 \rightarrow \int_{0}^{\infty} 2\lambda \frac{dS^{(N)}_l(k)}{d\lambda}^2 d\lambda = I.
$$

(26)

Assuming a real nuclear S-matrix which depends on $\lambda$ through $[1 + \exp(\lambda - \Lambda)/\Delta]^{-1}$ then the derivative of $S$ would peak around the grazing angular momentum $\Lambda$ with a width given by $\Delta$. The integral (26) is then obtained as: $I = \frac{1}{\sqrt{\Delta}}$ for $\Lambda/\Delta \gg 1$. Using $\Delta = kR$, $\Lambda = kR$, with $R = r_0(A^1/3 + A^{1/3})$ and $a$ being the diffuseness of the optical potential we find the simple formula for $\sigma$.

$$
\sigma = c\left[(\delta^{(N)}_1)^2 \left(\frac{3}{2}\right)^2 \frac{\Delta R_P}{R_P^2} - (\delta^{(N)}_2)^2 \right] R/3a
$$

(27)

where $c$ is a constant normalization factor which depends among other things on the exit doorway nature of the excited state exemplified by the factor $f(E_d, \Gamma, 1_d)$. It is clear that $\sigma$ depends linearly on the radius of the target and, more importantly on the square of the nuclear dipole and quadrupole deformation lengths. Thus, the $A^{1/3}$ dependence is established.

In the calculation to follow we use the cluster model to calculate deformation lengths for the different multi-polarities $[19, 20, 21]$. This model assumes that the projectile is composed of two clusters, a core of mass and charge $a_c$ and $z_c$ and a “valence” particle with $a_b$ and $z_b$. The separation energy is denoted by $Q$, the Q-value of the breakup. Calling the spectroscopic factor of finding the cluster configuration in the ground state of the projectile, $S$, one obtains the following expression for the distribution of $B(E\lambda)$ in the excitation energy $E_x$ [19, 21].

$$
\frac{dB(E\lambda)}{dE_x} = SN_b^2 \frac{2^{\lambda-1}}{\pi^2} (\lambda!)^2(2\lambda + 1) \left(\frac{\hbar^2}{\mu_b}\right) \times
\frac{Q^{1/2}(E_x - Q)^{\lambda+1/2}}{E_x^{\lambda+2}} \times
\frac{(Z_b A_c^\lambda + (-1)^\lambda Z_c A_b^\lambda)}{A_p^\lambda} e^{2}\frac{h^2}{r}
$$

(28)

where $N_0$ is normalization factor which takes into account the finite range, $r_0$ of the $c + b$ potential. The latter is assumed to be such as to give a Yukawa type wave function at large distances, $\psi_{bc}(r) = N_0 \sqrt{K/(2\pi)} e^{-K}\times e^{-K_{ro}/r}$ with $K = \sqrt{2\mu_{bc}Q/\hbar^2}$ and $N_0 = \frac{\sqrt{K_{ro}}}{\sqrt{2\mu_{bc}Q}}$. It is easy to obtain $B(E\lambda)$ by simply integrating of Eq. (28) and employing the expression:

$$
\int_{0}^{\infty} \frac{y^{\lambda+1/2}}{(y + 1)^{2\lambda+2}} dy = \frac{\left((-2\lambda+3)\pi\right)}{(2\lambda + 1)\sin[(\lambda + 3/2)\pi]} \prod_{k=1}^{2\lambda+1} \left[(\lambda + 3/2 - k)\right]
$$

(29)

FIG. 1: CDCC calculations for the nuclear breakup (dots) compared to the results of Eq. (27). See text for details.

We get for the cluster-model deformation lengths $\delta^1_1$ and $\delta^2_2$ the following:

$$
(\delta^{(N)}_1)^2 = \left(\frac{2\pi}{3} \frac{A_p}{Z_p N_p}\right)^2 \frac{B(E1)}{e^2} = \frac{N_0^2 S}{3Z_p N_p} 2 \frac{3}{16\pi} \frac{h^2}{\mu_{bc}} (A_z Z_b - A_b Z_c)^2 1 Q^2
$$

(30)

$$
(\delta^{(N)}_2)^2 = \left(\frac{4\pi}{3Z_p R_P}\right)^2 \frac{B(E2)}{e^2} = \frac{N_0^2 S}{3Z_p R_P} 5 \frac{2}{32\pi} \frac{h^2}{\mu_{bc}} (A_z^2 Z_b^2 + A_b^2 Z_c^2)^2 1 Q^2
$$

(31)

where $p(= b + c)$ refers to the projectile.

For our three nuclei discussed here, we have $^{11}$Be = $^{10}$Be + $n$, $^{8}$B = $^{7}$Be + $p$ and $^{7}$Be = $^{4}$He + $^{3}$He, which define their cluster character, with the corresponding breakup Q-values, 0.504 MeV, 0.137 MeV and 1.587 MeV. The factor $N_0^2 S$ could be related to the Asymptotic Normalization Coefficient (ANC) of the bound state wave function and is taken as a parameter to be adjusted so as to account for the experimentally known $B(E\lambda)$.

Simple estimate of $\frac{\Delta R_P}{R_P}$ can be obtained from [14, 15] who gave $\Delta R_P \approx \frac{|N_0^{1/3} - Z_b^{1/3}|}{A_P^{1/3}}$. In Table 1 we present the results of the deformation lengths obtained using the cluster model of formulas [30] and [31]. For $^{11}$Be, we used $B(E1) = 1.05 \pm 0.06$ eV-fm$^2$ [1] and we get $(\delta^{(N)}_1)^2 = 0.71 \pm 0.04$ fm$^2$. Further, $(\delta^{(N)}_2)^2 = 1.27 \pm 0.25$ fm$^2$ from the same reference.
TABLE I: Deformation lengths for the $^7\text{Be}$, $^8\text{B}$ and $^{11}\text{Be}$ projectiles. The deformation lengths for $^7\text{Be}$ and $^8\text{B}$ have been calculated using formulas (30) and (31) using $N_0 S = 1$. For the $^{11}\text{Be}$ the $\delta_1$ and $\delta_2$ are the values from Ref.[1]

| Projectile | $(\delta_1(N))^2$ (fm$^2$) | $(\delta_2(N))^2$ (fm$^2$) | $(\Delta R^2)_c$ | $c$ |
|------------|-----------------------------|-----------------------------|-----------------|-----|
| $^7\text{Be}$ | 0.41 | 1.54 | 0.00576 | 1.37 |
| $^8\text{B}$ | 1.62 | 3.24 | 0.0179 | 0.307 |
| $^{11}\text{Be}$ | 0.84(2) | 1.27(25) | 0.0214 | 2.52 |

In figure 1 we compare our results, Eq.(27) using $a = 0.65$ fm with the CDCC calculation of Ref.[5] at $E_{lab} = 200$ MeV.A. Clearly we underestimate the CDCC calculation. The reason resides in the neglect, in our model, of the higher-order channel coupling terms alluded to above. We also show in figure 1 the comparison with the CDCC calculation for $^8\text{B}$ and $^7\text{Be}$ using the formula(27). Clearly the scaling law is better obeyed in the “normal” nucleus $^7\text{Be}$ as has already been discussed in [5]. The value of the normalization $c$ is close to unity for the “normal” nuclei $^7\text{Be}$ For the $^8\text{B}$ the normalization is very close the average value of the Asymptotic Normalization Coefficient $\text{ANC} = 0.45$ measured in ref.[22]. For the $^{11}\text{Be}$ a higher normalization is obtained probably due to higher order effects which are not accounted by our DWBA description.

In conclusion, we derived an expression for the nuclear breakup cross-section using the Austern-Blair theory. The obtained cross-section exhibits the scaling law and should serve to supply a simple mean for an estimate of the nuclear breakup contribution. The expression found should be contrasted with the purely geometric Serber-like expression $\sigma = 2\pi Ra$.[3].

Acknowledgements

This work is supported in part by Fundação de Amparo à Pesquisa do Estado de São Paulo (FAPESP) and the Conselho Nacional de Desenvolvimento Científico e Tecnológico (CNPq). M.S.H is the Martin Gutzwiller Fellow, 2007/2008.

[1] N. Fukuda et al. Phys. Rev. C70(2004) 054606.
[2] R. Palit et al. Phys. Rev. C68(2003) 0344318.
[3] T. Aumann, Eur. Phys. J. A26(2005) 441.
[4] M. A. Nagarajin, C. H. Dasso, S. M. Lenzi and A. Vitturi, Phys. Lett. B, 503(2001) 65.
[5] M. S. Hussein, R. Lichtenthäler, F. M. Nunes and I. J. Thompson, Phys. Lett., B640(2006) 91.
[6] R. Serber, Phys. Rev. 72(1947) 1008.
[7] M. S. Hussein and A. F. R. de Toledo Piza, Phys. Rev. Lett. 72(1994) 2693.
[8] M. S. Hussein, M.P. Pato and A. F. R. de Toledo Piza, Phys. Rev. C52(1995) 846.
[9] L.F. Canto, A. Romanelli, M. S. Hussein and A. F. R. de Toledo Piza, Phys. Rev. Lett. 72(1994) 2147.
[10] E.J. Benjamin et al., Phys. Lett. B 647(2007) 30.
[11] N. Auerbach and V. Zelevinsky, Nucl. Phys. A781(2007) 67.
[12] N. Austern and J.S. Blair, Ann. Phys. (NY) 33(1965) 15.
[13] W.E. Frahn, Nucl. Phys. A272(1976) 413.
[14] G.R. Satchler, Nucl. Phys. A195(1972) 1.
[15] G.R. Satchler, Nucl. Phys. A472(1987) 215.
[16] G. Alkhazov et al. Phys. Rev. Lett. 78(1997) 2313.
[17] I. Tanihata et al., Phys. Lett. 289(1988) 261.
[18] A. Krasznahorkay et al., Phys. Rev. Lett. 66(1991) 1287.
[19] C.A. Bertulani, G. Baur and M.S. Hussein, Nucl. Phys. A526(1991) 751.
[20] C.A. Bertulani and A. Sustich, Phys. Rev. C46(1992) 2340.
[21] C.A. Bertulani, M.S. Hussein and G. Münzenberg, Physics of Radioactive Beams (Nova Science, New York, 2001)
[22] L. Trach, F. Carstoiu, C.A. Gagliardi and R.E. Tribble Phys. Rev. Lett. 87(2001) 271102-1