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ABSTRACT:- The several roles of multipole giant resonances in heavy-ion reactions are discussed. In particular, the modifications in the effective ion-ion potential due to the virtual excitation of giant resonances at low energies, are considered and estimated for several systems. Real excitation of giant resonances in heavy-ion reactions at intermediate energies are then discussed and their importance in the approach phase of deeply inelastic processes is emphasized. Several demonstrative examples are given.

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I. INTRODUCTION

In this talk I shall discuss the role of giant resonances in heavy-ion reactions. As is well known the giant resonances have been excited by several probes, most notable of these are electrons, photons, pions, protons, α, etc.. In the case of electrons and photons probes, although the reaction mechanism is quite simple (electromagnetic), the cross section however is very small. A larger cross-section is obtained with hadron probes, with a price paid: a more complicated reaction mechanism. The question of the nature of the background seen in the spectrum of hadron-induced nuclear reactions is still a rather subtle one.

The usual interpretation given to the background is based on the contribution of multistep processes. It is suggested that multi-phonon excitations (of several multi-polarities) results in a spectrum composed of strongly overlapping broad peaks that would show up as a practically structureless background below the isolated, rather prominent, peaks attached to one-photon states (giant resonances). It is well-known that in heavy-ion reactions, the multi-step processes are the rule rather than the exception. One would therefore expect that HI-induced reactions leading to GR excitation exhibit more complicated spectrum, whose back-ground is of an even more subtle nature than that of the reactions induced by simple hadronic probes.

Nevertheless, the excitation of GR in heavy ions has, in the past few years, been shown to play an important role, especially in the mechanism responsible for the large energy loss encountered in deeply inelastic processes. These processes occur at c.m. energies of the order of 2-3 E_B, where E_B is the height of the Coulomb barrier. For very heavy ions, these processes almost completely exhaust the total reaction cross section at these energies. They represent events where the two heavy ions emerge from the reaction after having lost practically all their kinetic energy into intrinsic excitation energies.

Aside from their importance in deeply heavy-ion reactions, the GR's may have a role in the effective ion-ion interaction appropriate for the description of heavy-ion elastic and quasi-elastic scattering at lower energies. At these energies, the GR's enter into the picture in the form of polarization components that should be added to the bare interaction.

This paper is organized as follows. In Section II, the polarization effects of GR's at low energies are discussed in the context of elastic scattering and the optical potential. In Section III a brief account of the experimental evidence in support of the direct excitation of GR's in intermediate energy heavy-ion inelastic scattering is given. In Section IV, the relevance of these excitations in deeply inelastic heavy-ion collisions is discussed and, finally, in Section V, several concluding remarks are made.

II. EFFECT OF GIANT RESONANCES ON THE HEAVY ION OPTICAL POTENTIAL

At low energies, the heavy ion elastic scattering angular distribution is characterized by several features; at small angles the cross-section oscillates about the Rutherford value

($\frac{\sigma_{e\ell}}{\sigma_R}(0) = 1$) and eventually, at intermediate angles, drops rather rapidly to small values, indicating clearly the presence of a phenomenon reminiscent of Fresnel diffraction. The only explicit information about the nuclear structure seems to be associated with the size of the system contained in the grazing angular momentum extracted from the value of the angle $\theta_{1/4}$, at which $\frac{\sigma_{e\ell}}{\sigma_R} = 0.25$ (quarter-point-recipe).

Nuclear structure effects in elastic scattering become more conspicuous at higher energies and larger angles, manifesting themselves through characteristic oscillations in $\frac{\sigma_{e\ell}}{\sigma_R}$ at angles larger than the grazing angle $\theta_{1/4}$. The oscillations arise from interference effects related to specific nuclear structure aspects of the heavy-ion system (coupled channels). Some of these coupled channels effects may, in some cases, become important even at small angles. This is the case seen in the elastic scattering of deformed heavy ions, where multiple Coulomb excitations are important. The strong Coulomb coupling of the elastic channel to the collective inelastic channels in these systems is seen to result in an effective long-range component in the ion-ion potential. This component is found to be predominantly absorptive in the case of coupling to low-lying collective states. The reason is that these states, being low-lying, are excited with such high probabilities and very small energy losses that a treatment based on the sudden approximation is quite appropriate. The reason being that the reaction time is much shorter that the vibrational period. This basically implies a loss of flux from the elastic channel with very little change in the effective real interaction. In contrast, giant resonances, being high-lying states, have a much smaller vibrational period. The system, therefore, reacts adiabatically, resulting in a change in the effective real interaction which becomes more important at sub-barrier energies.

To see this more quantitatively we consider below the amplitude for the excitation of a vibrational state of multipolarity λ , and excitation energy ΔE_λ , using first order time-dependent theory of Coulomb excitation[1].

$$a_\lambda(\infty) = \frac{1}{i\hbar}\int_{-\infty}^{\infty} \langle n|V(r(t))|0\rangle\, e^{i\Delta E_\lambda t/\hbar}\, dt \qquad (1)$$

The interaction, $V(r(t))$ has a matrix element $\langle n|V(r)|0\rangle$ proportional to $r(t)^{-\lambda-1}$. Thus we have for $a_\lambda(\infty)$

$$a_\lambda(\infty) = C\int_{-\infty}^{\infty} \frac{\exp[i\frac{\Delta E_\lambda t}{\hbar}]}{[r(t)]^{\lambda+1}}\, dt \qquad (2)$$

The largest contribution to the integral in Eq. (12) comes from the vicinity of the classical turning point, $r_{tp} \equiv r(0)$. Thus by expanding $r(t) \simeq r_{tp} + \frac{1}{2}\ddot{r}_{tp}t^2$, and keeping lowest order terms in t , one obtains the simple estimate

$$a_\lambda(\infty) = C\,\frac{\pi}{r_{tp}^{\lambda+1}}\sqrt{\frac{2}{\lambda+1}\frac{r_{tp}}{\ddot{r}_{tp}}}\,\exp\left[-\sqrt{\frac{2}{\lambda+1}\frac{r_{tp}}{\ddot{r}_{tp}}}\,\omega_\lambda\right] \qquad (3)$$

where C is a constant, $\ddot{r}_{tp}$ is the radial acceleration at the classical turning point and $\omega_\lambda \equiv \frac{\Delta E_\lambda}{\hbar}$. Introducing the average collision time, $\tau_{coll} \equiv \sqrt{\frac{r_{tp}}{|\ddot{r}_{tp}|}}$, we have finally[2]

$$a_\lambda(\infty) = C\,\frac{\pi}{r_{tp}^{\lambda+1}}\sqrt{\frac{2}{\lambda+1}}\,\tau_{coll}\,\exp\left[-\sqrt{\frac{2}{\lambda+1}}\,\omega_\lambda\tau_{coll}\right] \qquad (4)$$

Eq. (4) clearly shows that the quantity that decides upon the strength of the transition is $\omega_\lambda\tau_{coll}$; which has the following simple relation to the relevant physical variables

$$\omega_\lambda\tau_{coll} = \sqrt{2}\left(\frac{\Delta E_\lambda}{E_{c.m.}}\right)\eta \qquad (5)$$

where η is the Sommerfeld parameter , $\eta = \frac{z_p z_T e^2}{\hbar v}$, v being the asymptotic relative velocity.

For heavy ion ($z_p z_T >> 1$) reactions at low energies, $\eta >> 1$, therefore one expects $\omega_\lambda\tau_{coll} >> 1$ for giant resonance excitation and $\omega_\lambda\tau_{coll} << 1$ for excitation of low-lying collective states. This indicates that $a_\lambda^{GR}(\infty) << 1$ and accordingly very little amount of flux is lost from the elastic channel. On the other hand, $a_\lambda(\infty) \sim 1$ for the excitation of low-lying states, which reflects the need to incorporate into the elastic channel optical potential, the resulting, absorptive long-range component[3].

The fact that $a_\lambda^{GR}(\infty) << 1$ for heavy ion scattering at low energies implies that the system follows adiabatically the motion. This in turn implies that a way to account for the GR polarization in the optical potential is simply to minimize the multidimensional potential energy surface with respect to the deformation parameters.

To be specific, we consider the polarization effect due to the coupling to the giant quadrupole mode[2]. Then the potential energy is taken to be

$$V = \frac{Z_p Z_t e^2}{r} + \frac{(Z_p Q_t + Z_t Q_p)e^2}{2r^3} + \frac{1}{2}C_p\beta_p^2 + \frac{1}{2}C_t\beta_t^2 \qquad (6)$$

where Q_i is the intrinsic quadrupole moment of the projectile (p) or target (t) nucleus, given by

$$Q_i = \frac{3 Z_i \beta_2(i) R_{o,i}^2}{(5\pi)^{1/2}} \qquad (7)$$

$\beta_2(i)$ is the quadrupole deformation parameter, $R_{o,i}$ the radius and C_i is the spring constant of the assumed harmonic vibrator. Differentiating Eq.(7), first with respect to β_t and then with respect to β_p, setting the resulting first derivative of V equal to zero at each case, we obtain

$$V(r) = \frac{Z_p Z_t e^2}{r} - \left(\frac{R_{o,p}^4}{C_p} + \frac{R_{o,t}^4}{C_t}\right)\left(\frac{9 Z_t^2 Z_p^2 e^4}{40\pi}\right)\frac{1}{r^6} \qquad (8)$$

Using for the spring constants the value 14.8 MeV, obtained by assuming that the isoscalar giant quadrupole resonance at an excitation energy $\Delta E = \frac{60}{A^{1/3}}$ MeV, exhausts the energy-weighted sum rule, we finally obtain the following estimate for the polarization correction (second term on RHS of Eq. (8))

$$\Delta V_{GQR} = -\frac{5}{4}(0.0208)(A_p^{1/3} + A_t^{1/3}) Z_t^2 Z_p^2 / r^6 \ (\mathrm{Mev}) \qquad (9)$$

The factor $\frac{5}{4}$ $(= 1 + \frac{1}{4})$ comes from the approximate inclusion of

the isovector quadrupole resonance $(\Delta E^{T=1}_{\lambda=2} \simeq 2\ \Delta E^{T=0}_{\lambda=2})$ which gives the factor 1/4.

Similar procedure may be followed for the obtention of ΔV due to the isovector dipole resonance (GDR). One obtains

$$\Delta V_{GDR} = -6.7 \times 10^{-3}\left[\frac{N_p}{Z_p A_p^{1/3}} + \frac{N_t}{Z_t A_t^{1/3}}\right] Z_p^2 Z_t^2 / r^4 \ (\mathrm{Mev}) \qquad (10)$$

where $N \equiv A - z$, is the number of neutrons. In deriving Eq.(10), the Goldhaber-Teller model is used together with the empirical excitation energy value of $\Delta E_{GDR} = 80\ A^{-1/3}$ MeV.

It is clear from Eqs. (9) and (10) that the giant resonance polarization is represented by an attractive, real, long-range component in the ion-ion potential. Though this component is small (e.g. in $^{40}Ar + {}^{60}Gd$, $V_{GQR} \cong -0.2$ MeV at $r \sim R_B$, see Table 1 where ΔV is calculated for several HI systems), it may be quite important in those heavy-ion processes that are sensitive to the tail of the nuclear potential (quasi-elastic reactions populating discrete states). The GR polarization component may also be important in heavy-ion fusion at sub-barrier energies where barrier penetration is the dominant mechanism. A slight lowering of the barrier height, that would arise from the inclusion of ΔV_{GR}, may result in a significant increase in the fusion cross-section over the value obtained from simple one-dimensional barrier penetration calculation.

It is of interest to compare the purely real polarization potential discussed above with the predominantly imaginary dynamical polarization potential arising from the coupling of the elastic channel to low-lying vibrational or rotational states. Owing to the small value of $\omega_\lambda \tau_{coll}$ in this last case, a sudden-limit treatment of the excitation process is adequate. The resulting dynamical polarization potential, has the following simple form[3] (setting the energy loss equal to zero)

$$\Delta V_{DPP}(r) = -i\left[\frac{a_\ell}{r^3} + \frac{b_\ell}{r^4} + \frac{c_\ell}{r^5}\right] \tag{11}$$

where the coefficients a_ℓ, b_ℓ and c_ℓ are complicated functions of the center of mass energy, E , the reduced transition probabilities, the Sommerfeld parameter, η , and the orbital angular momentum. The numerical value of $|\Delta V_{DPP}|$ is comparable to that of ΔV_{GQR} .

Whereas the consideration of ΔV_{GQR} results in a slight change of the real ion-ion potential, and a subsequent slight deviation of the classical trajectory and the deflection function from the Rutherford ones, the effect of $\Delta V_{DPP}(r)$ on the elastic scattering cross section is much more drastic owing to its absorptive nature. The inclusion of $\Delta V_{DPP}(r)$ in the calculation of $\frac{d\sigma_{e\ell}}{d\Omega}$ results in the following simple form for the cross section, valid at sub-barrier energies

$$\frac{d\sigma}{d\Omega} = \frac{d\sigma_R}{d\Omega} \exp\left[-\sum_{i=1}^{3} F_i(\ell,\eta,E)\, f_i(\theta)\right] \tag{12}$$

where F_i are functions related to the coefficients a_ℓ, b_ℓ and c_ℓ and $f_i(\theta)$ are simple function of the center of mass angle. The presence of the exponential factor in Eq. (12) results in a significant damping in the elastic cross section that starts at small angles and extends all the way to back angles (see Ref. (3) for more details).

It should be stressed that the real nature of the giant resonance polarization potential V_{GQR} and V_{GDR} is a direct consequence of the adiabatic nature of the excitation process encountered at the sub-barrier energies we have been considering so far. At higher energies the excitation of giant resonances is believed to be responsible for a significant part of the total reaction cross section especially in cases where deeply inelastic reactions are found to dominate over fusion. This implies that if one were to construct e.g. V_{GQR} to be used for the calculations of the heavy-ion elastic scattering cross section at intermediate energies, one would find it to contain a strong absorptive component. In the following section, we shall discuss the experimental situation of GR excitations in heavy ion collisions at intermediate energies.

III. GIANT RESONANCES EXCITED IN HEAVY-ION REACTIONS AT INTERMEDIATE ENERGIES

As one can see from Eqs. (4) and (5), as the energy is increased, the factor in the exponential $\tau_{coll}\omega_\lambda$, becomes smaller, even for the excitation of giant resonances, rendering the amplitude $a_\lambda(\infty)$, for these excitations, appreciable. Note, however, that at the higher energies, considered here $(E_{CM} > 2E_B)$, it is the nuclear excitation (short-ranged) rather than the Coulomb excitation (long-ranged) of the giant resonances which is dominant. This has an immediate consequence of restricting the excitation process to the surface region of the nucleus. A result of this localization is the characteristic Fraunhofer pattern in the inelastic distribution. Employing the usual terminology of heavy-ion physics, the above pattern in the angular distribution results from the interference between the near- and far-side components of the amplitude. One immediate consequence of this interference phenomena is the Blair phase-rule[4], which says that the oscillations of the angular distributions of inelastic transitions to odd multipolarities are in phase with those of elastic scattering, whereas even multipolarities are out of phase.

As an example we show in figure 1 the data on the light-heavy-ion system $^{12}C + ^{27}Al$ at $E_{^{12}C} = 82$ MeV studied by Betts et al.[5]. The Coulomb barrier, E_B, of this system is about 22 MeV, so we are talking about $\frac{E_{CM}}{E_B} \simeq 3.0$, i.e. center of mass energies well above the barrier. The angular distributions do show the Fraunhofer pattern, with the Blair rule approximately satisfied. One also sees clearly the large background in the spectrum, which is customarily attributed to multiple excitations of several GR modes. As has already been discussed earlier, this interpretation is not the only possibility. Multiple excitation of incoherent modes might certainly contribute as well.

Another attempt to observe GR in heavy ion reactions was made by Buenerd et al.[6]. These authors looked at ^{12}C inelastic scattering on ^{90}Zr and ^{208}Pb at $E_{^{12}C} \simeq 120$ MeV, and ^{14}N on ^{40}C, ^{90}Zr, ^{197}Au, ^{208}Pb and ^{209}Bi at $E_{^{14}N} = 161$ MeV. Fig. (2) shows the spectrum and angular distributions at $E_x = 11.0$ MeV and 2.61 MeV of ^{208}Pb (^{14}N, $^{14}N'$). One notices that in the present case of a heavy target, the GQR

exciation cross section shows a distribution in angles which is more of a Coulomb-nuclear interference type (near-side phenomenon, no Blair phase-rule) than a Fraunhofer type. This is primarily a consequence of the larger Coulomb interaction ($\eta >> 1$) and excitation than in the $^{12}C + ^{27}Al$ system mentioned earlier.

Similar features in the angular distribution as the ones mentioned above may be seen in the data of Doll et al.[7], Fig. (3). These authors measured the inelastic scattering of ^{16}O ($E_{Lab} = 315$ MeV) on ^{208}Pb and ^{12}C, in the Q-value region corresponding to the excitation of the GQR in the targets. Again one sees clearly the near-side dominance in $^{16}O + ^{208}Pb$ and near-far interference in $^{16}O + ^{12}C$ (for comparison, $\eta(^{16}O + ^{12}C) = 1.7$ and $\eta(^{16}O + ^{208}Pb) = 23.3$, both evaluated at $E_{^{16}O} = 315$ MeV).

In all of the above cases, several GR's with different multipolarities were identified. It should be clear from the above discussion that the disentangling of nuclear structure information of the GR's (e.g. damping width), using heavy-ion inelastic scattering is more complicated than in light-ion-induced reaction. This is principally a consequence of the much more complicated nature of the background in the former. However, the observation of these resonances in HI is important in so far as the interpretation of the nature of deeply inelastic reactions (DIC), is concerned. The relevance of our discussion in this section to DIC becomes clear when one recognizes that the energies at which GR have been populated in HI reactions ($E_{CM} > 3E_B$) correspond closely to those at which the DIC cross-section becomes a major part of the HI total reaction cross section.

IV. GIANT RESONANCES IN DEEPLY INELASTIC REACTIONS OF HEAVY IONS

The usual scenario of heavy-ion reactions is that at CM energies, E_{CM} , close to the Coulomb barrier, E_B , the total reaction cross section, σ_R , is almost completely exhausted by complete fusion, σ_F . This last process is characterized by a complete amalgamation of the two ions to form a compound nucleus. All relevant degrees of freedom reach their equilibrium values in this process usually viewed as statistical. At center-of-mass energies $E_{CM} > 2E_B$, σ_F starts deviating appreciably from σ_R , and at higher energies, $\frac{\sigma_F}{\sigma_R} << 1$. The usual interpretation of this phenomenon is the occurence, at these energies, of a process intermediate, in complexity, between fusion and qausi-elastic reactions[8].

This new mechanism of the heavy-ion reaction, usually referred to as deeply inelastic collision, involves a partial equilibrium of the system, and might be accordingly considered as the HI-analogue of light-ion-induced pre-equilibrium reactions. The difference between the two reactions is, however, quite clear. Figure (4) shows a typical spectrum of a HI-induced reaction at intermediate energies. For comparison we show in the inset a typical spectrum of a light-ion-induced reaction. The rather wide hump centered at the exit channel Coulomb barrier is the DIC component.

The picture employed to describe DIC is that the two ions after reaching the interaction zone, suffer a large amount of energy loss, exchange many nucleons and emerge as two highly excited, deformed, fragments. They do not fuse because not all the degrees of freedom reach their equilibrium values.

Several macroscopic variables are employed in the description of the evolution of the system. In particular, owing to the classical nature of relative motion, the detection angle (deflection function) is used as a measure of the reaction time. The longer the reaction time is, the greater would be the energy loss. This last observation has a clear consequence on the angular distribution. Depending on the final fragment masses and the total energy loss the angular distribution shows a marked evolution from side-peaked (for final fragments equal to initial ones) to forward peaked. Fig. (5) shows the spectra of heavy particle products and angular distributions in $^{40}Ar + ^{232}Th$ at E_{CM} = 330 MeV.

Other macroscopic variables relevant for the description of DIC are the mass, charge, orbital angular momentum and charge-to-mass ratio. It is now evident that the charge-to-mass ratio equilibrates quite fast ($\tau_{z/A} \approx 1.3 \times 10^{-22}$ s) , followed by the energy ($\tau_E \approx 3 \times 10^{-22}$ s) , the orbital angular momentum ($\tau_\ell \approx 15 \times 10^{-22}$ s), and finally the charge or mass ($\tau_z \approx 60 \times 10^{-22}$ s). Actually in most DIC events the charge or mass never reaches equilibrium[9].

The equilibration process associated with, z/A z and A , in DIC, is nicely described by a diffusion equation[10]. The important point one discovers from this description is that the width, Γ , of the distribution $P(x,t)$ of a given macroscopic variable (x = z/A , z or A), satisfies the usual diffusion relation $\Gamma \propto t^{1/2}$. With the association of t with the deflection angle (see earlier discussion), one reaches the conclusion that the larger the deflection angle is, the larger the width of the distribution ($\Gamma^2 \propto \theta$) . Larger angles, in our case, may be obtained by simply allowing the system to scatter to negative angles. The "angle" variable (which is connected with the nuclear reaction time) is measured from a reference angle, θ_g , related to the position in angle-space where quasi-elastic processes (small energy loss) reach their maximum values. A very nice demonstration of the above picture is seen in fig. (6) which represents the width of the charge distribution in the system $^{40}Ar + ^{232}Th$.

Many more evidence[11] has accumulated in support of the above diffusion picture. However, the question of the relaxation process associated with the energy variable remains unsetteled. A simple diffusion picture as the above was found not to work so well for the E-relaxation. This clearly points for the need of treating the evolution of several of the collective variables associated with the E-relaxation process in a completely non-statistical, coherent fashion .

It is here that the GR,s come into the picture. Within the DIC model developed by Broglia et al.[12] the GR,s are explicitly involved in the removal of energy from the radial motion to intrinsic excitations. A set of coupled classical equations of motion that involves as coordinates, the vector that specifies the relative position of the two heavy ions, as well as the usual variables that describe harmonic oscillators representing the different modes of collective vibrations, are solved. In order to guarantee that the energy deposited into the GR,s does not return to the relative motion, these oscillators are rendered damped. The damping widths attached to these oscillators represent, on the average, the fragmentation of the GR due to its coupling to the non-coherent, statistical, (ph), degrees of freedom.

What one usually ends up obtaining from such a model, are average quantities: the energy loss as a function of the impact parameter, the average deflection function, and the

classical cross-section. In order to smooth out the classical singularities in the classical cross section (caustics) and to obtain a measure of the dispersion in energy, Broglia et al.[12], invoke the zero point motion of the harmonic oscillators which generates dispersions in measurable quantities, of a purely quantal nature. This is to be contrasted with the purely statistical nature of the dispersion that results from a transport (diffusion) interpretation of DIC[10].

One should mention that the Copenhagen model for DIC does contain some statistical aspects as well. In order to reproduce the average energy loss, it was found necessary to involve not only the GR,s but also particle transfer treated as a diffusion process. This last mechanism was found to be responsible for as much as 50% of the average energy loss. One would expect, therefore, that a significant part of the widths (dispersion) of the distributions of the observable macroscopic variables, arises from statistical fluctuations related to the diffusion of particles.

It seems clear that both quantal and statistical fluctuations are present in DIC. The quantal effects are mostly operative in the approach phase of the collision process. During this stage GR,s may paly a dominant role as "doorways" to the more complicated p-h configurations. A recent discussion on unifying both statistical and coherent effects in DIC may be found in Ref. (13). Actually the importance of GR,s in the processes involved in DIC has been clearly demonstrated through the rather extensive TDHF calculations that have been performed in the last few years[14].

Though in the discussion we have had so far, we have considered the isoscalar GR,s; recently several suggestions have been made concerning the possible importance of the isovector giant resonances in the charge equilibration process. It is suggested in Ref. 15) that the experimental width of the Z distribution at fixed mass asymmetry might be related to the thermal and quantal fluctuations of the axial component of the isovector E1 mode associated with the intermediate complex. The authors of Ref. 15) suggest further that the rather small widths of the Z-distributions observed in asymmetric systems such as $^{40}Ar + ^{197}Au$, are predominantly thermal in nature, whereas the large widths observed for nearly symmetric systems, e.g., $^{86}Kr + ^{98}Mo$, are quantal.

The above suggestions, though not directly confirmed by experimental findings point clearly for the possible role of the giant dipole resonance in the charge equilibration process in DIC. It should be mentioned, however, that a completely different interpretation of the charge width can be found in Ref. 16) where a simple diffusion picture is invoked together with a proper account of the Pauli blocking effect.

V. CONCLUSIONS

In this talk the role of giant multipole resonances in heavy ion reactions has been discussed. It is seen that at low center of mass energies, the GR,s enter into the picture in the form of polarization effects modifying primarily the real part of the ion-ion potential. At higher energies the excitation of the GR,s become possible. Several illustrative examples were

presented demonstrating the existence of the GR peaks on top of a rather complicated background. The possible importance of the GR excitation in deeply inelastic collisions was then discussed, both in the process of energy dissipation and equlibration, and in determining the width of the charge distribution at fixed mass asymmetry.

In conclusion, the giant resonance excitation in heavy ion reactions should be considered as an integral part of a process involving necessarily many collective and intrinsic degrees of freedom. The interplay between the resulting coherent and statistical responses of the system is a dominant theme of present day research in heavy ion physics.

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System	$-V_{GQR} \cdot r^6$ (MeV.fm^6)	$-V_{GDR} \cdot r^4$ (MeV.fm^4)
$^{40}Ar + {}^{160}Gd$	3.77×10^5	6.239×10^3
$^{40}Ar + {}^{148}Sm$	3.48×10^5	5.71×10^3
$^{16}O + {}^{148}Sm$	4.99×10^4	1.08×10^3
$^{16}O + {}^{208}Pb$	9.45×10^5	1.892×10^3
$^{208}Pb + {}^{208}Pb$	1.39×10^7	2.37×10^5
$^{238}U + {}^{238}U$	2.3×10^7	2.45×10^5

TABLE 1 - Numerical values of the giant resonance polarization potentials for several heavy-ion systems. These potentials should be used in the description of heavy ion scattering at sub-barrier energies (no overlap of the heavy ion densities). A reasonable value of the minimum radius down to which the ΔV's of Eqs. (9) and (10) are valid is $\sim 1.6\ (A_p^{1/3} + A_t^{1/3})$ [fm]. The values of ΔV given above for the ^{16}O-induced reactions should be considered with reservation, since the estimate $\Delta E_1 \sim \frac{80}{A^{1/3}}$ (MeV) for dipole giant resonance excitation energy is really valid for $A > 40$.

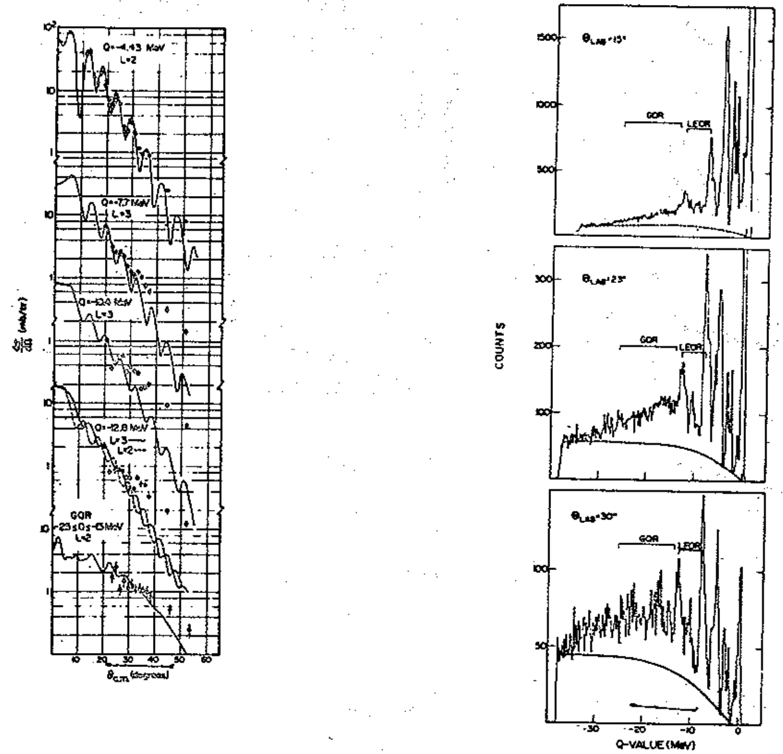

FIGURE 1 - Energy spectra and angular distribution of inelastically scattered ^{12}C on $^{27}A\ell$ at $E_{^{12}C} = 82$ MeV. From Ref. (6).

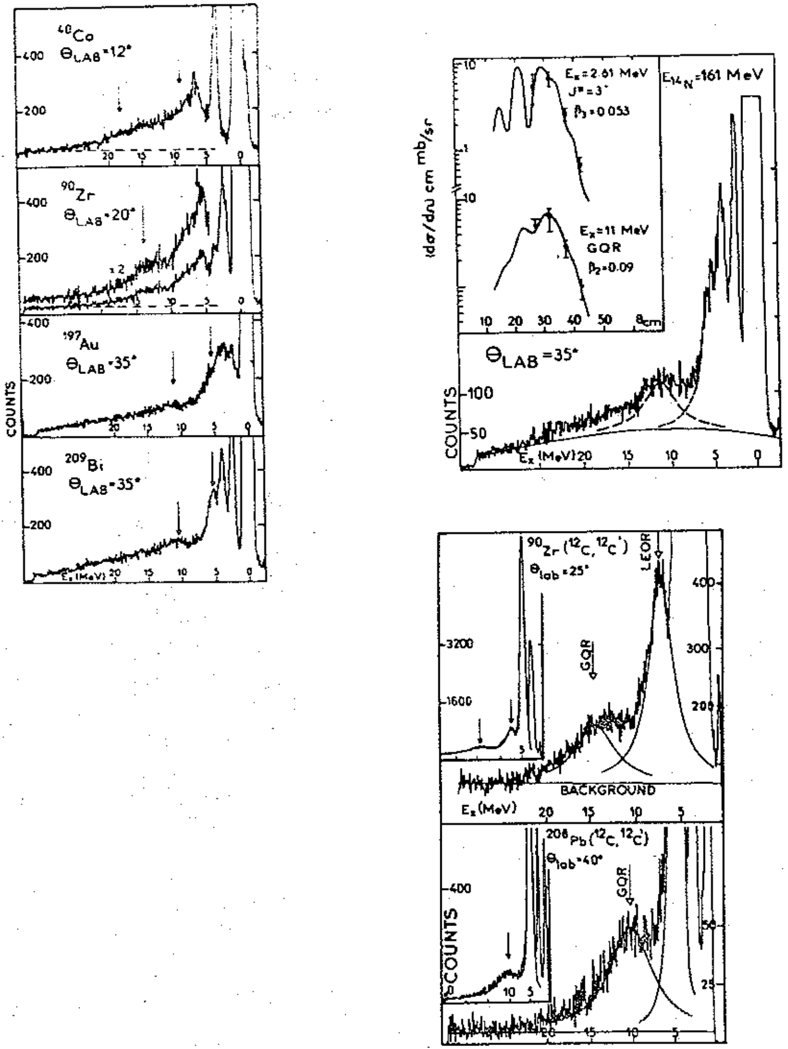

FIGURE 2 - Typical energy spectra of inelastically scattered ^{14}N (top) and ^{12}C (bottom). Also shown are two angular distributions of N in the reaction $^{208}Pb(^{14}N,^{14}N)$. The figures were taken from Ref.(6).

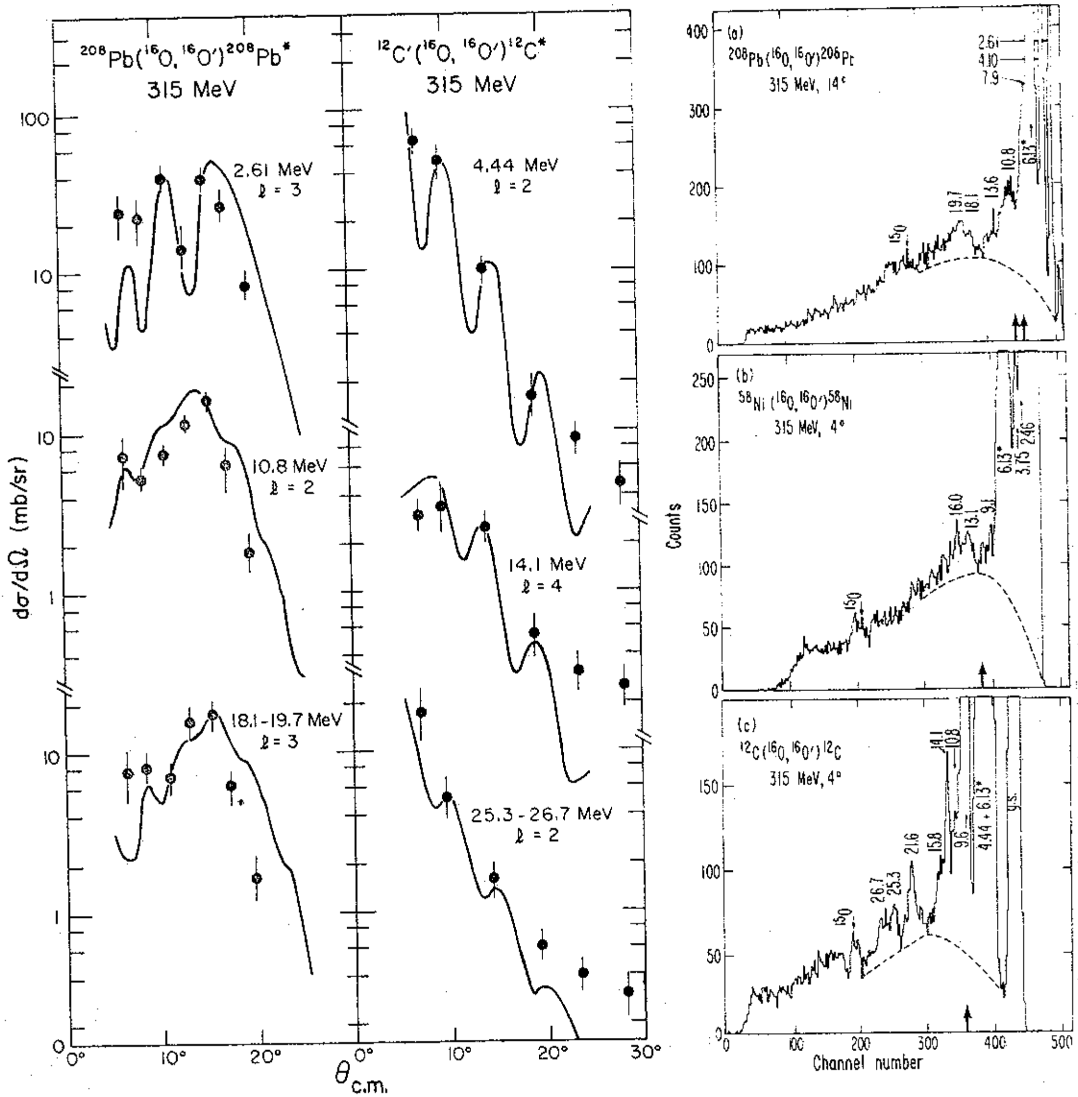

FIGURE 3 - Energy spectra and angular distribution of inelastically scattered ^{16}O on ^{208}Pb and ^{12}C at $E_{^{16}O}$ = 315 MeV. From Ref. (7).

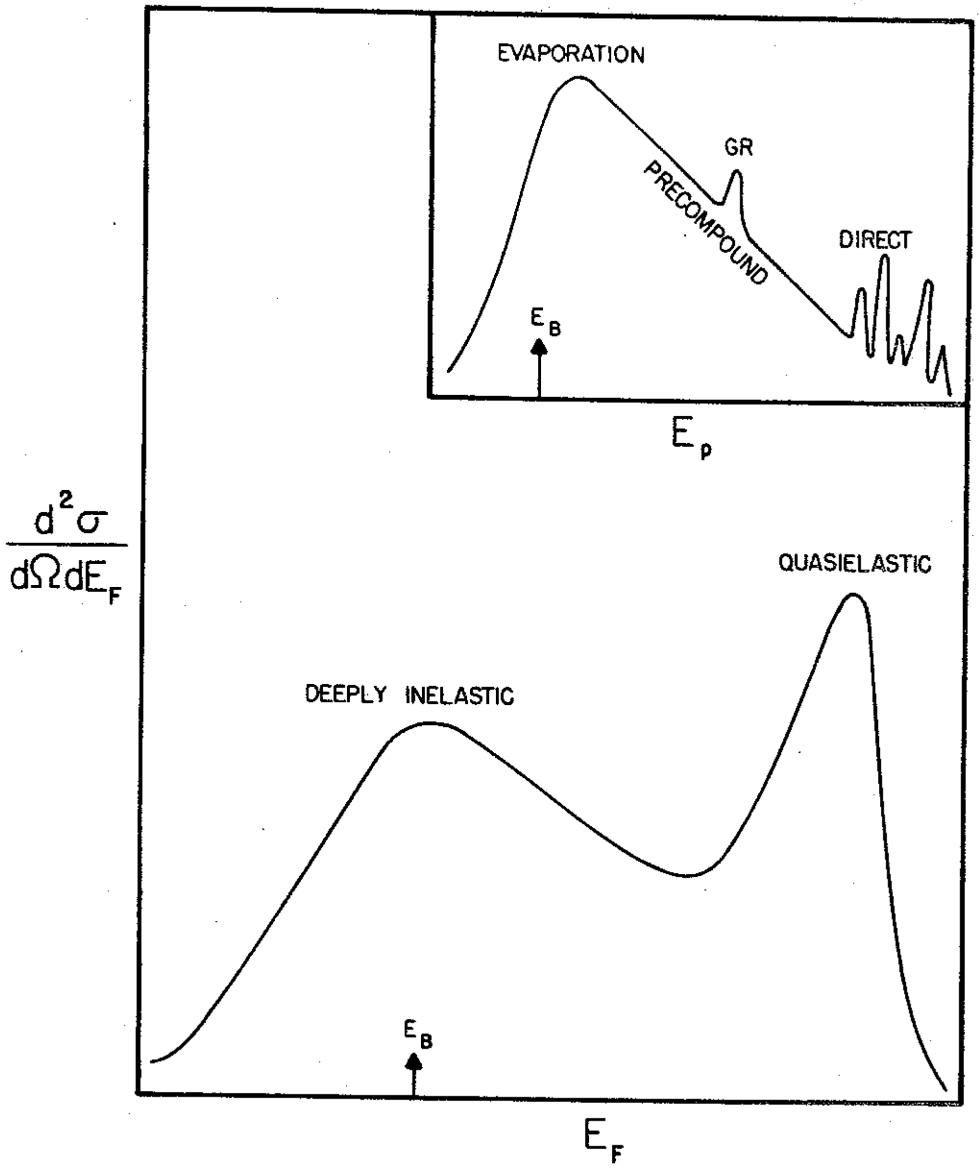

FIGURE 4 - A schematic plot of a typical heavy ion energy spectrum at intermediate energies. The broad peak centered at E_F close to the exit channel Coulomb barrier, corresponds to the DIC component. In the inset we show a typical spectrum of a light ion induced reaction.

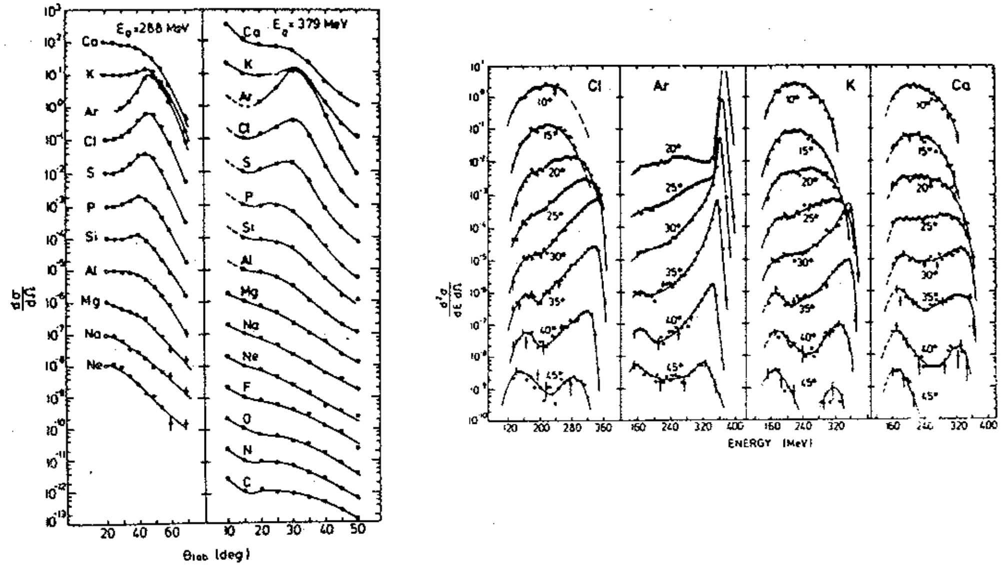

Figure 5 - Energy spectra and angular distributions of several nuclear species originating from the reaction $^{40}Ar + ^{232}Th$. From Ref. (8).

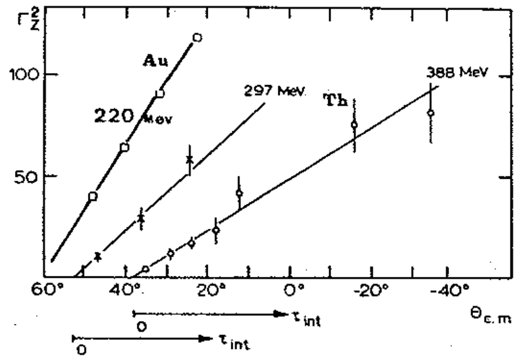

Figure 6 - The width-at-half-maximum of the charge distribution squared, vs. the c.m. angles (i.e. interaction time) for several systems. From Ref. (8).