Multistep compound nuclear reactions with giant resonances as doorway states

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The characteristics of multistep compound nuclear reactions with the giant resonances as doorway states are discussed in this Brief Report. A cross-section formula describing the above nuclear reactions is derived from the Feshbach-Kerman-Koonin theory with a proper modification of its basic assumptions. Parametrization of this formula leads to an extension of the conventional exciton model. As an application, the derived formula is used to calculate the cross sections of the photonuclear reactions exciting the giant dipole resonances in different nuclei, and the results are in good agreement with experimental data.

In recent decades, a large body of experimental data has shown that giant resonances (GR's) exist in a vast region of the periodic table ranging from light to heavy nuclei and that they play an important role in many nuclear reactions.[1-4] It seems urgent to develop a nuclear reaction model incorporating GR's as an essential ingredient. Such a subject has been attracting much attention.[5-7]

It is the purpose of this Brief Report to provide a model, that can describe the multistep compound nuclear (MCN) reactions including GR's as doorway states. The model has been obtained from a proper extension of the Feshback-Kerman-Koonin (FKK) theory and describes the nuclear reaction processes as three major steps: the semidirect process of the GR's excitation and decay, the compound nuclear preequilibrium emission and the compound nuclear (CN) equilibrium emission.

The giant resonances are coherent nuclear excitations with strong collective properties and complex microscopic structures. They are strongly coupled to other nuclear states, especially to the CN states. Their excitation energies are so high that they could induce the multistep processes and lead to the preequilibrium emissions. Because of the above characteristics, the FKK (Ref. 8) theory, after having been properly modified, is thought to be adequate to describe the MCN reactions with GR's as doorway states. The FKK theory assumes that all the CN states (within closed channels) have a random-phase (RP) distribution as an ensemble average is concerned. Since the GR consists of coherent excitations of many particle-hole states with strong phase correlations, the RP approximation is not applicable to it. In order to use the FKK theory to describe the nuclear reactions studied in this report, the basic assumptions of the FKK are properly modified: (1) The RP assumption is applicable to all CN states except the GR; (2) the GR is dominant in the doorway states (the $n=1$ configuration subspace). From the above assumptions, we can derive a cross section formula that describes MCN reactions with GR as an isolated doorway state.

According to FKK, the fluctuation cross section of the MCN is

$$\sigma_{fi}^{(\mathrm{fl})}=\frac{4\pi^3}{k^2}\langle|T_{fi}^{(\mathrm{fl})}|^2\rangle \ . \tag{1}$$

The fluctuation transition matrix element $T_{fi}^{(\mathrm{fl})}$ is a sum, i.e.,

$$T_{fi}^{(\mathrm{fl})}=\sum_{n=1}^{r} T_{fi}^{(n)} \ , \tag{2}$$

(r denotes the equilibrium state) where $T_{fi}^{(n)}$, the transition matrix element in the nth subspace, is defined as

$$T_{fi}^{(n)}=\langle\psi_f^{(-)}|V_{pn}\frac{1}{E-h_{QQ}}V_{1p}|\psi_i^{(+)}\rangle \ . \tag{3}$$

By using the assumption (1), we have

$$\langle T_{fi}^{(n)}\rangle=\langle T_{fi}^{(n)*}\rangle=0, \quad (\text{if } n\geq 2)$$

and

$$\langle|T_{fi}^{(\mathrm{fl})}|^2\rangle=\langle\sum_{nn'}T_{fi}^{(n)}T_{fi}^{(n')*}\rangle$$

$$=\sum_{n=1}^{r}\langle|T_{fi}^{(n)}|^2\rangle \ .$$

The above equation means that even though GR is included as a doorway state in the MCN processes, it is still possible to express the cross section $\sigma_{fi}^{(\mathrm{fl})}$ as a sum of $\sigma_{fi}^{(n)}$ in different subspaces

$$\sigma_{fi}^{(\mathrm{fl})}=\frac{4\pi^3}{k^2}\sum_{n=1}^{r}\langle|T_{fi}^{(n)}|^2\rangle=\sum_{n=1}^{r}\sigma_{fi}^{(n)} \ . \tag{4}$$

Inserting the complete set of states in the 1 subspace and using the assumption (2), one can derive the following results with the same method as used in the FKK:

$$\sigma_{fi}^{(\mathrm{fl})}(E)=\sigma_G^{(i)}(E)\left[\frac{\Gamma_G^{(f)}}{\Gamma_G}+\sum_{n=2}^{r}\frac{\Gamma_n^{(f)}}{\Gamma_n}\left[\prod_{l=2}^{n-1}\frac{\Gamma_l^{\downarrow}}{\Gamma_l}\right]\frac{\Gamma_G^{\downarrow}}{\Gamma_G}\right] \ , \tag{5}$$

with

$$\Gamma_G^{(f)}=2\pi\langle|\langle\psi_f^{(-)}|V_{p1}|\mathrm{GR}\rangle|^2\rangle\ ,$$
$$\Gamma_G^{\downarrow}=\frac{2\pi}{D_2}\langle|\langle\tilde{2}\alpha|V_{2,1}|\mathrm{GR}\rangle|^2\rangle\ ,$$
$$\Gamma_G=\Gamma_G^{\downarrow}+\sum_x\Gamma_G^{(x)}\ , \tag{6}$$
$$\Gamma_G^{(i)}=2\pi\langle|\langle\mathrm{GR}|V_{1p}|\psi_i^{(+)}\rangle|^2\rangle\ ,$$
$$\sigma_G^{(i)}(E)=\frac{2\pi}{k^2}\frac{\Gamma_G\Gamma_G^{(i)}}{2\pi[(E-E_G)^2+\Gamma_G^2/4)}\ ,$$

where $\Gamma_G^{(f)}$ and $\Gamma_G^{\downarrow}$ are the escape and spreading widths of the GR, respectively, Γ_G the total width, $\Gamma_G^{(i)}$ the excitation width of the GR, $\sigma_G^{(i)}$ the excitation cross section of the GR from the initial channel $|i\rangle$, and E_G is the energy position of the GR peak. The other quantities have the same definitions as those in FKK.

Equation (5) is the central result of the present model, which describes the nuclear reactions as the following processes illustrated in Fig. 1.

It should be pointed out that the three quantities $\sigma_G^{(i)}$, $\Gamma_G^{(f)}$, Γ_G, and $\Gamma_G^{\downarrow}/\Gamma_G$ can be calculated by using quantum mechanics because they are related to the excitation and decay of the GR, while the branching ratios $\Gamma_n^{(f)}/\Gamma_n$ and $\Gamma_n^{\downarrow}/\Gamma_n$ can be treated by statistical theory (exciton model).[9]

It is known that FKK can be either computed microscopically or parametrized. Its parametrized form is equivalent to the conventional exciton model (EXM) with the following substitutions

$$\frac{\Gamma_n^{\downarrow}}{\Gamma_n}=\frac{\lambda_+^{(n)}}{\lambda_+(n)+L_n},\quad \frac{\Gamma_n^{(f)}}{\Gamma_n}=\frac{L_n^{(f)}}{\lambda_+^{(n)}+L_n}\quad(\text{for } n<n_r)\ . \tag{7}$$

By using Eq. (7), Eq. (5) can be written as

$$\sigma_{fi}^{(\mathrm{fl})}(E)=\sigma_G^{(i)}(E)\left[\frac{\Gamma_G^{(f)}(n_0)}{\Gamma_G(n_0)}+\sum_{\substack{n=n_0+2\\(\Delta n=2)}}^{n_r-2}\frac{L_n(f)}{\lambda_+(n)+L_n}\times\left[\prod_{\substack{l=n_0+2\\(\Delta l=2)}}^{n-2}\frac{\lambda_+(l)}{\lambda_+(l)+L_l}\right]\times\frac{\Gamma_G^{\downarrow}(n_0)}{\Gamma_G(n_0)}+Q(f)D(n_0,n_r)\right]\ , \tag{8}$$

where the exciton number $n=p+h$, p and h are the particle number and hole number, respectively, and

$$Q(f)=\frac{\Gamma(f)}{\Gamma(\gamma)+\sum_x\Gamma(x)}\ , \tag{9}$$

$$D(n_0,n_r)=\prod_{\substack{l=n_0+2\\(\Delta l=2)}}^{n_r}\frac{\lambda_+(l)}{\lambda_+(l)+L_l}\frac{\Gamma_G^{\downarrow}(n_0)}{\Gamma_G(n_0)}\ . \tag{10}$$

$Q(x)$ denotes the evaporation probability of the x particle. The widths $\Gamma(\gamma)$ and $\Gamma(x)$ are defined by

$$\Gamma(x)=\frac{1}{2\pi\rho(A,U)}\int_{V_{\mathrm{Coul.}}}^{U-B_x}dE_x\frac{4\mu_x}{\pi\hbar^2}E_x\sigma(E_x)\rho(A-1,U')\ , \tag{11}$$

$$\Gamma(\gamma)=\frac{1}{2\pi\rho(A,U)}\int_0^U dE_\gamma\frac{2k_\gamma^2}{\pi}\sigma_G^{(i)}(E_\gamma)\rho(A,U-E_\gamma)\ . \tag{12}$$

Equation (8) can be considered as an extended EXM including the GR as a doorway state. It is a mixture of quantum-mechanical and statistical descriptions. $\lambda_+(n)$ is the transition probability from n to $n+2$ exciton states and $L_n(x)$ is the emission probability of the x particle from n-exciton states. Both λ_+ and L_n should be calculated by EXM and are defined as

$$\lambda_+(n)=\frac{2\pi}{\hbar^2}\overline{|M|^2}\frac{gU^2}{2(n+1)}=X_{dk}\left[\frac{g}{A}\right]^3\frac{U}{n+1}\ , \tag{13}$$

$$L_n(x)=\int_{V_{\mathrm{Coul.}}}^{U-B_x}dE_x W(n,E_x)\ , \tag{14}$$

$$L_n=\sum_x L_n(x)\ , \tag{15}$$

with

$$g=6A/(7.5\pi^2)\ ,$$
$$W(n,E_x)=C_x\frac{2S_x+1}{\pi^2\hbar^3}\mu_x E_x\sigma(E_x)\times\frac{\omega(A-1,U',n-1,p-1)}{\omega(A,U,n,p)}\ , \tag{16}$$

and

$$\omega(A,U,n,p)=\frac{g(gU)^{n-1}}{p!h!(n-1)!}\ , \tag{17}$$

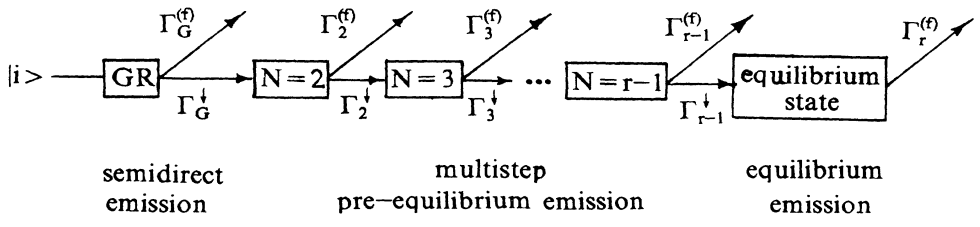

FIG. 1. Reaction processes.

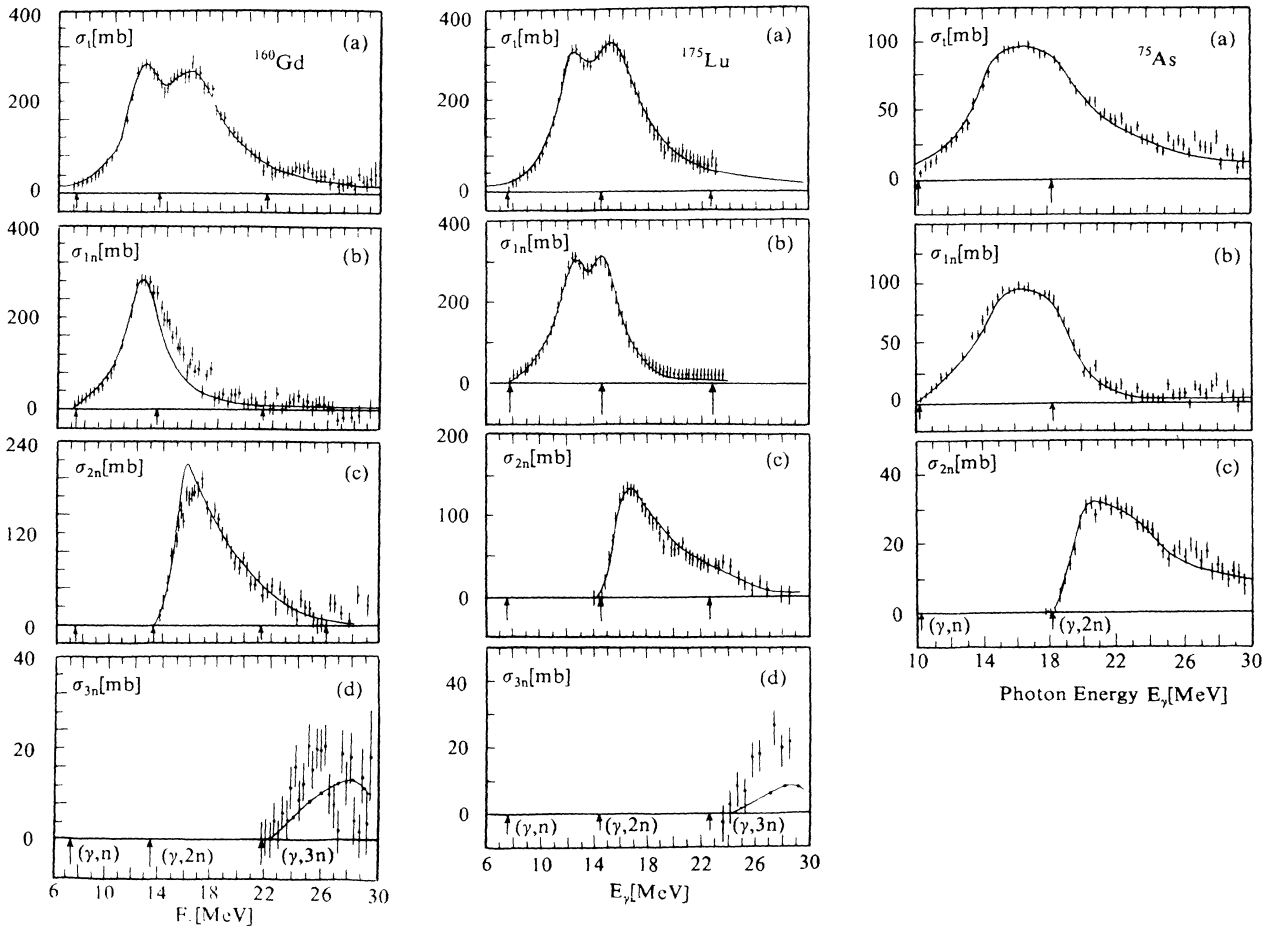

FIG. 2. Photoneutron cross sections for ^{208}Pb, ^{186}W, ^{181}Ta, ^{175}Lu, ^{160}Gd, and ^{75}As: (a) total; (b) single; (c) double; (d) triple.

where $C_n=(n+1)/n$ (for neutron), $C_p=(n-1)/n$ (for proton), $U'=E-E_x-B_x$, $|M|^2$ the square average of the matrix element of two-body interactions, ω the state number, and $W(n,E_x)$ the emission probability of the x particle at the emission energy E_x from the n-exciton states. A is the mass number, U the excitation energy, g is the single-particle level density, and $\sigma(E_x)$ the inverse cross section. B_x and μ_x are the separation energy and reduced mass of x particle, respectively. $x_{dk}=4.77286d_k$, where d_k is an adjustable parameter and are taken to be 570 MeV[3]. All these quantities are defined in Refs. 10–14.

As a test of our model, Eq. (8) is used to describe photonuclear reactions exciting the giant dipole resonances (GDR's) in various nuclei. Since our present objective is only to test our model in a tentative manner, those quan-

TABLE I. The neutron escape widths $\Gamma_G^{(n)}$ and the spread widths $\Gamma_G^{\downarrow}$ of GDR at the excitation energy $E_G=(E_{G1}+E_{G2})/2$. $\Gamma_G^{(n)}/\Gamma_G$ and $\Gamma_G^{\downarrow}/\Gamma_G$ are calculated by using Eqs. (18) and (19), and the total widths Γ_G for the deformed nuclei assume the average values $\Gamma_G=(\sigma_{01}^{(G)}\Gamma_{G1}+\sigma_{02}^{(G)}\Gamma_{G2})/(\sigma_{01}^{(G)}+\sigma_{02}^{(G)})$.

	^{75}As	^{160}Gd	^{175}Lu	^{181}Ta	^{186}W	^{208}Pb $\beta=1$	^{208}Pb $\beta=10$
$\sigma_0^{(G)}$ (mb)	41 77	215 233	217 287	296 341	211 334	645	
E_G (MeV)	14.89 17.61	12.23 15.96	12.32 15.47	12.30 15.23	12.59 14.88	13.63	
Γ_{G1} (MeV)	3.64	2.77	2.57	2.43	2.29		
Γ_{G2} (MeV)	7.26	5.28	4.70	4.48	5.18		
Γ_G (MeV)	6.00	4.08	3.78	3.53	4.06	3.94	
$\Gamma_G^{(n)}/\Gamma_G$	0.0475	0.0558	0.0376	0.0454	0.0529	0.0477	0.3338
$\Gamma_G^{(n)}$ (MeV)	0.29	0.23	0.14	0.16	0.21	0.19	1.32
$\Gamma_G^{\downarrow}/\Gamma_G$	0.9521	0.9442	0.9624	0.9546	0.9471	0.9523	0.661
$\Gamma_G^{\downarrow}$ (MeV)	5.71	3.85	3.64	3.37	3.85	3.75	2.62

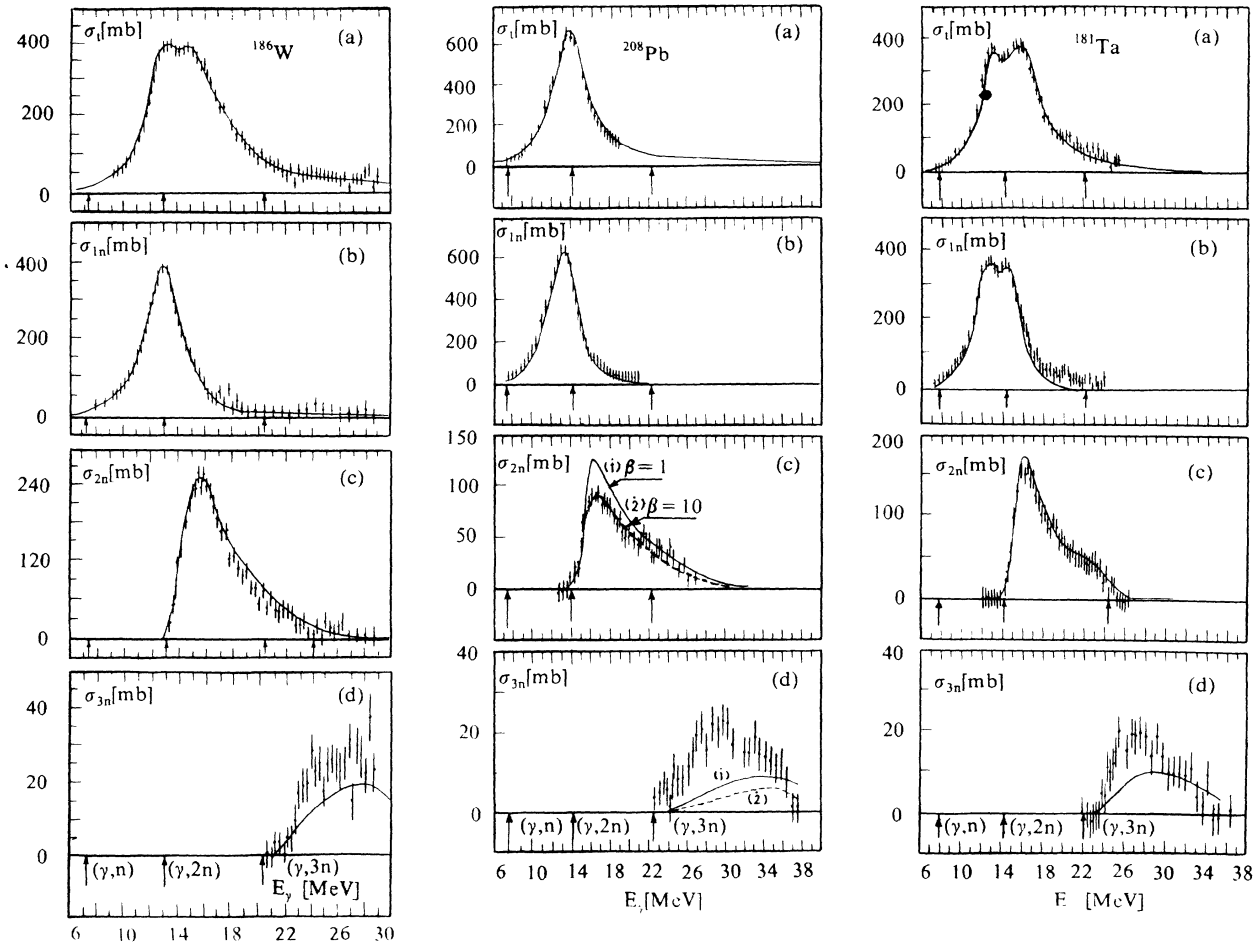

FIG. 2. (*Continued*).

tities related to GR will not be calculated from a more delicate theory (such as the RPA), but rather be further approximated.

Since the GR state is populated dominantly in 1p-1h subspace, we assume

$$\Gamma_G^{\downarrow}(n_0)=a\Gamma_{n_0}^{\downarrow},\ \Gamma_G^{(f)}(n_0)=b(f)\Gamma_{n_0}^{(f)},\ (n_0=1p+1h)\ , \tag{18}$$

where a and $b(f)$ are ratios of the GR widths to the average of the single-particle widths.

By virtue of Eqs. (18) and (7), Eq. (8) becomes

$$\sigma_{fi}^{(\mathrm{fl})}(E_\gamma)=\sigma_G^{(i)}(E_\gamma)\left\{\frac{\beta(f)L_{n_0}(f)}{\lambda_+(n_0)+\sum_x \beta(x)L_{n_0}(x)}+\sum_{\substack{n=n_0+2\\(\Delta n=2)}}^{n_r-2}\frac{L_n(f)}{\lambda_+^{(n)}+L_n}D(n-2,n_0)+Q(f)D(n_r,n_0)\right\}\ , \tag{19}$$

with

$$D(n,n_0)=\left[\prod_{\substack{l=n_0+2\\(\Delta l=2)}}^{n}\frac{\lambda_+(l)}{\lambda_+(l)+L_l}\right]\times\frac{\lambda_+(n_0)}{\lambda_+(n_0)+\sum_x \beta(x)L_{n_0}(x)}\ , \tag{20}$$

where $\beta(x)=b(x)/a$ is an adjustable parameter indicating the relative strength of the particle emission from the GR. The GDR excitation cross section can be very well parametrized by the Lorentzian form

$$\sigma_G^{(i)}(E_\gamma)=\sum_{j=1}^{m}\sigma_{0_j}^{(G)}\frac{E_\gamma^2E_{Gj}^2}{(E_\gamma^2-E_{Gj}^2)^2+E_\gamma^2\Gamma_{Gj}^2}\ ,\qquad m=\begin{cases}1 \text{ for spherical}\\ 2 \text{ for deformed}\ .\end{cases} \tag{21}$$

In our calculation, the parameters of optical potential are taken from Ref. 15 and the energy level density, GDR parameters, pairing and shell corrections from Refs. 1, 16, and 17.

In order to reveal the influence of GR widths, we first take $\beta(n)=\beta(p)=1$. The results are presented in Fig. 2 in comparison with experimental data. The values of

$\Gamma_G^{\downarrow}/\Gamma_G$, $\Gamma_G^{(n)}/\Gamma_G$, $\Gamma_G^{\downarrow}$ and $\Gamma_G^{(n)}$ are given in Table I. The calculated results are in good agreement with experimental data. We found that the semidirect neutron emission component is about $\sim 3-6\%$, the preequilibrium emission at the second step contributes about $\sim 4\%$ and the total preequilibrium emission (including "semidirect") about $\sim 12\%$. This result indicates that the semidirect and the preequilibrium processes must be taken into account in the photonuclear reactions of GDR's.

However, for ^{208}Pb, ^{181}Ta, and ^{175}Lu, σ_{3n} are smaller, and the positions of the peaks shift toward higher-energy regions. The peak of σ_{2n} is especially higher for ^{208}Pb.

To improve σ_{2n}, β has been adjusted for ^{208}Pb. As $\beta=10$, the σ_{2n} curve is much improved, while σ_{3n} becomes worse. The reason is as follows: a larger β means stronger semidirect particle emissions that suppress equilibrium emissions significantly, and therefore decrease σ_{2n} and σ_{3n}. Including the 2p-2h configurations in the GDR will increase the weight of equilibrium emissions and therefore improve σ_{3n}. This will be discussed in a forthcoming paper.

The angular distribution is not considered in the present report and will be a future topic.

Finally, we conclude with a comparison between our results and those given by Dias *et al.*[5] and Bracco *et al.*[6] In Ref. (5), a formalism for a description of GR decay has been derived from the framework of a multistep compound nuclear model.[18] The decay of GR is described as containing two components: the direct GR emission and the compound nuclear emission. A mixing parameter is introduced. The compound nuclear decay widths are treated by the Hauser-Feshbach formulas, while the direct decay widths can be calculated by the microscopic RPA theory. As for the formulas applied to ^{208}Pb, it has been shown that for GDR, the direct component contributes about 9–17%. In Bracco's papers,[6] the formulas proposed in Ref. 5 are employed to analyze the GMR decay data of ^{208}Pb. For GMR, the branching ratio of direct emission is about 15% calculated from the experimental data, which is consistent with the theoretical calculation within the experimental error. In Ref. (6a), the above formulas are extended to three step processes, and the precompound decay of the 2p-2h doorway of ^{208}Pb into the compound 1p-2h states of ^{207}Pb is also included. Since a precise determination of precompound emission needs accurate measurement and a better knowledge of level density of precompound states, only estimated values of precompound decay widths are given. In this report, we have presented a multistep description of the same processes from the framework of FKK without introducing mixing parameters. In our formulas, the GR decay consists of three types of contributions: the semidirect emission, the preequilibrium emission and the equilibrium emission. The semidirect decay should be treated quantum mechanically (for instance, by the RPA), while the preequilibrium and equilibrium decay widths are calculated by an exciton model (an extension of Hauser-Feshbach formulas). As for the formulas applied to GDR in ^{208}Pb, the first component contributes about 5%, while the first two components contribute about 12%, which is in agreement with the direct component given by Dias *et al.*[5] In short, our formalism is just an extension of the formalism proposed by Dias *et al.* to include multistep precompound emissions, as can be seen in Ref. 6a.

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