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Interfering Doorway States and Giant Resonances II: Cross Section Pattern

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Abstract

The mixing of the doorway components of a giant resonance due to the interaction via the common decay channels influences significantly the distribution of the multipole strength and the energy spectrum of the decay products of the giant resonance. The photoemission turns out to be most sensitive to the overlapping of the doorway states. At high excitation energies, the interference between the doorway states leads to a restructuring towards lower energies and apparent quenching of the dipole strength.

1 Introduction

In [1] we investigated analytically as well as numerically the interference of the doorway components of a giant resonance. The main result is the following: In the energy domain of a giant resonance, the interplay of two different types of collectivity inherent in the underlying doorway resonance states plays an important role. According to their origin, they are called internal and external collectivity, respectively. The role of the external collectivity becomes especially important when two or more doorway components of the giant resonance overlap. The interference between these states gives rise to an appreciable redistribution of the dipole strength and shifts it towards lower energies.

In this paper, we study the cross section pattern in order to see the consequences of the interplay of the two types of collectivity in measurabble values. Of special interest are the transition strengths when the interaction via the energy continuum is strong.

To this purpose, we describe in sects. 2 and 3 analytically the cross section pattern observed in different decay channels. The photoemission turns out to be especially sensitive to the degree of overlapping of the doorway states. In sect. 4, we discuss the interaction of the doorway states with the background states which leads to the internal damping of the collective exitation. We show in sect. 5 some numerical results obtained in the same model (without damping), but with the restrictions removed which were introduced into the analytical investigation. The numerical calculations confirm the main features of the interference between the different types of doorway states as they follow from the analytical study. Finally, we summarize the results in sect. 6 and draw some conclusions. Of interest is, above all, the apparent loss of the collective dipole strength at high excitation energy.

All symbols used in this paper are the same as in [1]. We cite to an equation in [1] by writing its number in brackets with the upper index [1], e.g. (2.1)[1] means eq. (2.1) in paper [1].

2 Transition Amplitudes and Partial Transition Strengths

The matrix of the transition amplitudes is

$$\hat{\mathcal{T}}(E) = \mathcal{A}^T \mathcal{G}(E) \mathcal{A} \tag{2.1}$$

(see eqs. (4.1)[1] and (4.3)[1]). To calculate it, one needs the Green's matrix

$$\mathcal{G}(E) = \frac{1}{E - \mathcal{H}} \tag{2.2}$$

(eq. (2.8)[1]) describing the evolution of the intermediate unstable system excited in reactions. In the doorway basis the $(k+1) \times (k+1)$ block $\mathcal{G}^{(dw)}(E)$ of this matrix is the only one which really has to be calculated. The influence of the trapped states is included in the self-energy matrix which contains the coupling between the doorway and trapped states. It manifests itself, as has already been mentioned in [1], in the fine structure variations of the transition amplitudes in the energy region of the unperturbed parental levels. Neglecting this fine structure, one reduces the problem to the calculation of the Green's matrix of the doorway effective Hamiltonian

$$\mathcal{H}^{(dw)} = \begin{pmatrix} \mathcal{H}^{(coll)} & \chi^T \\ \chi & \tilde{\mathcal{H}} \end{pmatrix} \tag{2.3}$$

(eq. (5.23)[1]). We omit also the hermitian part of the coupling

$$\chi = \left(\mathbf{v}^{(0)}\ \mathbf{v}^{(1)}\right) - \frac{i}{2}\left(\mathbf{0}\ \mathbf{w}\right) \tag{2.4}$$

(eq. (5.26)[1]) between the two types of doorway states as discussed in subsection 5.2 of [1].

Further transformations going in close analogy with those described below formula (4.13)[1] lead to the following results:

$$\mathcal{P}(E) = \mathbf{D}^T\, \mathcal{G}^{(coll)}(E)\, \mathbf{D} = \mathbf{D}^2\, \frac{E - \varepsilon_0 + \frac{i}{2}\sin^2\Theta\, \omega(E)}{\Lambda(E)} \tag{2.5}$$

$$T^{cc'}(E) = T^{cc'}_{coll}(E) + \tilde{T}^{cc'}(E) \tag{2.6}$$

where

$$T^{cc'}_{coll}(E) = \left(A^c_1 - \frac{i}{2}\, q^c(E)\right)\left(A^{c'}_1 - \frac{i}{2}\, q^{c'}(E)\right)\frac{E - \varepsilon_0 - \sin^2\Theta\, \mathbf{D}^2}{\Lambda(E)} \tag{2.7}$$

and

$$\tilde{T}^{cc'}(E) = \sum_\alpha \frac{A^c_\alpha\, A^{c'}_\alpha}{E - \tilde{\mathcal{E}}_\alpha}\, . \tag{2.8}$$

It is worthy noting that the collective parts of the transition amplitudes vanish at the energy

$$E_v = \varepsilon_0 + \sin^2\Theta\, \mathbf{D}^2\, . \tag{2.9}$$

The components of the transition vectors $\mathbf{A}^c$ in the doorway basis are defined by eq. (5.10)[1]. The 2×2 collective block

$$\mathcal{G}^{(coll)}(E) = \frac{1}{E - \mathcal{H}^{(coll)} - \mathcal{Q}(E)} =$$

$$\frac{1}{\Lambda(E)}\begin{pmatrix} E - \varepsilon_0 - \cos^2\Theta\, \mathbf{D}^2 + \frac{i}{2}\,\omega(E) & \sin\Theta\cos\Theta\, \mathbf{D}^2 \\ \sin\Theta\cos\Theta\, \mathbf{D}^2 & E - \varepsilon_0 - \sin^2\Theta\, \mathbf{D}^2 \end{pmatrix} \tag{2.10}$$

with the function $\Lambda(E)$ given by

$$\Lambda(\mathcal{E}) \equiv (\mathcal{E} - \varepsilon_0)\ (\mathcal{E} - \varepsilon_{coll}) + \frac{i}{2}\omega(\mathcal{E})\ \left(\mathcal{E} - \varepsilon_0 - \sin^2\Theta\, \mathbf{D}^2\right) = 0 \tag{2.11}$$

(eq. (5.40)[1]) extends the formula

$$G_{coll}(E) = \frac{1}{E - \varepsilon_0 - \mathbf{D}^2 - \mathrm{h}^T\, \frac{1}{E - \bar{H}}\, \mathrm{h}} \tag{2.12}$$

(eq. (4.14)[1]) for the Green's function of the internal collective vibration to the consideration of decaying collective modes.

The amplitudes (2.8) are sums of independent Breit-Wigner terms and contain themselves no interference effects. Indeed, all A^c_α are real and, as one can easily check with the help of eqs. (5.35)[1] – (5.37)[1],

$$\sum_c (A^c_\alpha)^2 = \tilde{\gamma}^\alpha\, . \tag{2.13}$$

All interference effects are included in the collective part (2.7). In particular, the interference of the doorway resonances belonging to the two different types is described by the functions

$$q(E) \equiv -4Q_{11}(E) = \sum_\alpha \frac{w^{(\alpha)2}}{E - \tilde{\mathcal{E}}_\alpha} \tag{2.14}$$

(eq. (5.39)[1]), and

$$q^c(E) = \sum_\alpha \frac{w^{(\alpha)} A^c_\alpha}{E - \tilde{\mathcal{E}}_\alpha} . \tag{2.15}$$

In the collective part (2.7) the dependence on the channel indices c, c' has the desirable factorized form but the factors are generally complex and energy dependent. Therefore, contrary to the case of isolated resonances, the locations of the maxima in the cross sections are not connected with the positions and the residues of the poles of the K- or T- matrices in any simple way. If however the collective resonances do not overlap too strongly all the functions $q(E)$ vary slowly within the energy region of the maximum arising from the giant resonance state and can approximately be considered as some complex constants.

The residues of the elastic reaction amplitudes are expressed in terms of the complex energies of the doorway resonances as

$$ResT^{cc}(\mathcal{E}_{dw}) =$$

$$\left(A^c_1 - \frac{i}{2} q^c(\mathcal{E}_{dw})\right)^2 \left[1 + \frac{1}{4}\sin^2 2\Theta \frac{\mathrm{D}^4}{\left(\mathcal{E}_{dw} - \varepsilon_0 - \sin^2\Theta\, \mathrm{D}^2\right)^2} + \frac{1}{4} q'(\mathcal{E}_{dw})\right]^{-1} . \tag{2.16}$$

In contrast to the simple real residues

$$\Gamma^c_{coll} = (A^c_d)^2 \tag{2.17}$$

and

$$\Gamma^c_r = \gamma^r_\perp \left(\xi_{\perp c}^{(r)}\right)^2 , \qquad (r = 1, 2, ..., k) \tag{2.18}$$

(eqs. (4.34)[1], (4.35)[1]) of the K-matrix, they are complex. They carry information, hidden in the quantities q^c, on the transition vectors $\mathbf{A}^c$ of all the overlapping doorway states.

The above formulae simplify appreciably if one neglects the coupling between the two types of doorway states. In such an approximation the energy dependence of the collective part

$$\sigma^c_{coll}(E) = \frac{1}{2\pi} \left(A^c_1\right)^2 \langle\gamma\rangle \frac{(E - E_v)^2}{(E - \varepsilon_0)^2 (E - \varepsilon_{coll})^2 + \frac{1}{4}\langle\gamma\rangle^2 (E - E_v)^2} \tag{2.19}$$

of the strength

$$\sigma^c(E) = -\frac{1}{\pi} \mathrm{Im} T^{cc}(E) = \sigma^c_{coll}(E) + \tilde{\sigma}^c(E) \tag{2.20}$$

of the transition into a particular decay channel c turns out to have the same universal form as in the single–channel model of ref. [2]. This part reveals two equally high maxima

$$\sigma^c_{coll}(\varepsilon_0) = \sigma^c_{coll}(\varepsilon_{coll}) = \frac{2}{\pi} \frac{\left(A^c_1\right)^2}{\langle\gamma\rangle} \tag{2.21}$$

at the poles of the K–matrix

$$\hat{K}(E) = \frac{\hat{\mathbf{A}}_d^T \hat{\mathbf{A}}_d}{E - \varepsilon_{coll}} + \frac{\hat{X}_\perp}{E - \varepsilon_0} \,. \tag{2.22}$$

(eq. (4.32)[1]). Between them, the transition strength (2.19) drops to zero at the point $E = E_v$, eq. (2.9). Therefore, the relative values of all KPW of the internal collective state with the energy ε_{coll} can be easily found as

$$\frac{\Gamma^c_{coll}}{\Gamma^{c'}_{coll}} = \frac{(A^c_d)^2}{(A^{c'}_d)^2} = \frac{(A^c_1)^2}{(A^{c'}_1)^2} = \frac{\sigma^c_{coll}(\varepsilon_{coll})}{\sigma^{c'}_{coll}(\varepsilon_{coll})} \,. \tag{2.23}$$

For any value of the mixing parameter λ, these widths satisfy the sum rule

$$\sum_c \Gamma^c_{coll} = \cos^2\Theta \, \langle\gamma\rangle \,. \tag{2.24}$$

The remaining part $\sin^2\Theta\,\langle\gamma\rangle$ belongs to the collective state at the energy ε_0 where the noncoherent contribution $\tilde{\sigma}^c(\varepsilon_0)$ is also large.

Therefore, the parameters of the K–matrix are directly extracted from the collective part of the cross sections σ^c in the adopted two-level approximation. Nevertheless, only in the limit $\lambda \ll 2$, when the collective states become isolated and the equations

$$\mathcal{E}_0 = \varepsilon_0 - \frac{i}{2}\sin^2\Theta\,\langle\gamma\rangle \,, \qquad \mathcal{E}_1 \equiv \mathcal{E}_{gr} = \varepsilon_0 + \mathbf{D}^2 - \frac{i}{2}\cos^2\Theta\,\langle\gamma\rangle \,, \tag{2.25}$$

(see (5.51)[1]) are valid, the quantity

$$\frac{(A^c_1)^2}{\langle\gamma\rangle} = \frac{(A^c_d)^2}{\langle\gamma\rangle\cos^2\Theta} = \frac{\Gamma^c_{coll}}{\langle\gamma\rangle\cos^2\Theta} \Rightarrow \frac{\Gamma^c_{gr}}{\Gamma_{gr}} \tag{2.26}$$

(see eqs. (4.34)[1], (3.7)[1]) coincides with the branching ratio $\mathcal{B}^c_{coll}$ obtained from

$$\sigma^c(E_{dw}) = \frac{2}{\pi}\frac{\Gamma^c_{dw}}{\Gamma_{dw}} \equiv \frac{2}{\pi}\,\mathcal{B}^c_{dw} \tag{2.27}$$

(see eq. (2.5)[1]). For finite values of λ the branching ratio differs also from the ratio

$$\frac{(A^c_d)^2}{\Gamma_{1;1/2}} \tag{2.28}$$

where

$$\Gamma_{1;1/2} = \frac{1}{2}\left[1 + \frac{1}{2}\left(\sqrt{1 + \frac{4}{\lambda^2} + \frac{4}{\lambda}\cos 2\Theta} - \sqrt{1 + \frac{4}{\lambda^2} - \frac{4}{\lambda}\cos 2\Theta}\right)\right]\langle\gamma\rangle \tag{2.29}$$

is the width on the half height of the right collective peak. The analogous width of the left peak is equal to

$$\Gamma_{0;1/2} = \frac{1}{2}\left[1 - \frac{1}{2}\left(\sqrt{1 + \frac{4}{\lambda^2} + \frac{4}{\lambda}\cos 2\Theta} - \sqrt{1 + \frac{4}{\lambda^2} - \frac{4}{\lambda}\cos 2\Theta}\right)\right]\langle\gamma\rangle \,. \tag{2.30}$$

The sum

$$\Gamma_{0;1/2} + \Gamma_{1;1/2} = \Gamma_{dw=0} + \Gamma_{dw=1} = \langle\gamma\rangle \tag{2.31}$$

of both of them is however independent of the values of the parameters λ and Θ and equal to the total width of both collective doorway resonances.

In the same two-level approximation, the residues (2.16) at the poles $\mathcal{E}_{dw=0,1}$ can be presented in a very simple form

$$ResT^{cc}(\mathcal{E}_{dw}) = \frac{(A_1^c)^2}{\langle\gamma\rangle}\Gamma_{dw}\frac{\mathcal{E}_{dw} - \mathcal{E}^*_{dw'}}{\mathcal{E}_{dw} - \mathcal{E}_{dw'}} \,. \tag{2.32}$$

Correspondingly, the TPW of the collective states look as

$$\Gamma^c_{dw} = \frac{(A_1^c)^2}{\langle\gamma\rangle}\Gamma_{dw}\sqrt{\frac{1 + [\tan\delta_{dw}(E_{dw'}) - \tan\delta_{dw'}(E_{dw})]^2}{1 + [\tan\delta_{dw}(E_{dw'}) + \tan\delta_{dw'}(E_{dw})]^2}} \,. \tag{2.33}$$

Here the scattering phase of the resonance dw taken at the energy of the resonance dw' is defined by the standard relation

$$\tan\delta_{dw}(E_{dw'}) = -\frac{1}{2}\frac{\Gamma_{dw}}{E_{dw'} - E_{dw}} \,. \tag{2.34}$$

These phases vanish when the resonances are isolated. The last factor on the r.h.s. of eq. (2.33) is just the matrix element U_{dw} of the Bell-Steinberger nonorthogonality matrix

$$\Gamma_{dw} = \frac{1}{U_{dw}}\sum_c |A^c_{dw}|^2 \tag{2.35}$$

(see eq. (2.30)[1]). Using eqs. (5.43)[1] – (5.45)[1], one can present the latter factor explicitly in terms of the mixing parameters Θ and λ,

$$U_{dw=0,1} = \frac{1}{\sqrt{2}}\left[1 + \frac{1 + \frac{1}{4}\lambda^2}{\sqrt{\left(1 - \frac{1}{4}\lambda^2\right)^2 + \lambda^2\cos^2 2\Theta}}\right]^{\frac{1}{2}} \,. \tag{2.36}$$

In both limiting cases, $\lambda \ll 2$ and $\lambda \gg 2$, this factor goes to unity. The factor is maximal in the intermediate region of $\lambda \approx 2$. In particular, for $\lambda = 2$

$$U_{0,1} = \begin{cases} \frac{1}{\sqrt{1-\tan^2\Theta}}\,; & 0 < \Theta < \frac{\pi}{4} \\ \frac{1}{\sqrt{1-\cot^2\Theta}}\,; & \frac{\pi}{4} < \Theta < \frac{\pi}{2} \,. \end{cases} \tag{2.37}$$

The quantity (2.37) becomes infinite for $\Theta = \frac{\pi}{4}$ as mentioned in [1].

As has been shown in ref. [3], the energy spectrum of the decay products of an arbitrary two-level unstable system can generally be expressed in terms of the energies and partial widths of the resonances and one additional real mixing parameter which satisfies a sum rule following from the Bell-Steinberger relation

$$\hat{\mathbf{A}}^*_{dw} \cdot \hat{\mathbf{A}}_{dw'} = i\,U_{dw\,dw'}\,(\mathcal{E}_{dw'} - \mathcal{E}^*_{dw}) \tag{2.38}$$

(eq. (2.29)[1]). The situation is even simpler in our quasi single-channel case (see the remark below eq. (2.19)) where the latter parameter is easily found explicitly [3] as a function of only the complex resonance energies $\mathcal{E}_{0,1}$. The resulting expression is remarkably simple,

$$\sigma^c_{coll}(E) = \frac{2}{\pi}\frac{(A^c_1)^2}{\langle\gamma\rangle}\sin^2\left[\delta_0(E)+\delta_1(E)\right] = \sigma^c_{coll}(\varepsilon_{coll})\sin^2\left[\delta_0(E)+\delta_1(E)\right] \ . \tag{2.39}$$

This yields the relation

$$\sigma^c_{coll}(E_{dw}) = \sigma^c_{coll}(\varepsilon_{coll})\cos^2\delta_{dw'}(E_{dw}) \tag{2.40}$$

between the values of the cross section at the resonance energies E_{dw} and its maximal value. One can easily convince oneself with the help of eqs. (5.44)[1], (5.45)[1] that both phases $\delta_{dw'}(E_{dw})$ drop to zero when $\lambda \ll 2$ and the resonances are isolated. However, in the opposite case of $\lambda \gg 2$ only the phase $\delta_0(E_1)$ of the narrow resonance is small. The other phase, $\delta_1(E_0)$, belonging to the level with the large width $\sim\langle\gamma\rangle$ is close to $\pi/2$. The cross section (2.19) has a narrow dip at the energy $E = E_v$ of the state $dw = 0$. In the limit of very large λ the narrow state decouples and gets invisible in the particle decay cross sections. At the same time, the interference of the state $dw = 1$ with the other doorway states becomes important. Then the decay cross section pattern can generally exhibit a more rich picture caused by the interference of the k similar doorway states.

3 Photoemission

The process of photoemission by the collective states turns out to be most sensitive to their interference. To take the electromagnetic radiation into account, one has only to add to the antihermitian part $-\frac{i}{2}AA^T \equiv -\frac{i}{2}W$ of the effective Hamiltonian $\mathcal{H}$ (eq. (3.2)[1]) the new term

$$W_{el} = -\frac{i}{2}\alpha_{el}\,\mathbf{D}\mathbf{D}^T \tag{3.1}$$

describing the radiation of the same multipolarity as the internal coupling vector $\mathbf{D}$. Therefore, the corresponding external coupling amplitude

$$\mathbf{A}^{(rad)} = \sqrt{\alpha_{el}}\,\mathbf{D} \tag{3.2}$$

is proportional to this vector with the constant α_{el} characterizing the strength of the electromagnetic interaction.

The elastic matrix element of the K-matrix in the photo-channel is equal to $\alpha_{el}P(E)$ (see eq. (4.8)[1]). The radiation KPW are therefore proportional to the dipole strengths

$$f^r = \left(\mathbf{d}\cdot\mathbf{\Phi}^{(r)}\right)^2 \tag{3.3}$$

(eq. (4.18)[1]), and

$$\Gamma^{(rad)}_r = \alpha_{el}\,\mathbf{D}^2 f^r \ . \tag{3.4}$$

One can immediately see from

$$f^1 = 1-\kappa^2 \ , \qquad f^r \sim \frac{\kappa^2}{N-1} \quad (r\neq 1) \tag{3.5}$$

(eq. (4.27)[1]) that, in the limit of small κ, the internal collective state appropriates the main part of the total radiation width $\alpha_{el}\mathbf{D}^2$. When $\kappa \to 0$, only the pole at the energy ε_{coll} survives in the radiation K-matrix element.

Similarly, the photoelastic scattering amplitude can easily be obtained from the function $\mathcal{P}(E)$. One has to substitute only $\mathbf{D}^2$ by $(1-\frac{i}{2}\,\alpha_{el})\,\mathbf{D}^2$ when calculating the collective Green's matrix (2.10). In our two-level approximation, this leads to the result

$$\sigma^{(rad)}(E) = \frac{1}{2\pi}\,\alpha_{el}\,\mathbf{D}^2\,\langle\gamma\rangle\times$$

$$\frac{(E-\varepsilon_0)^2\,(\cos^2\Theta + \alpha_{el}/\lambda) + \frac{1}{4}\,\alpha_{el}\mathbf{D}^2\,\langle\gamma\rangle\sin^4\Theta}{\left[(E-\varepsilon_0)(E-\varepsilon_{coll}) - \frac{1}{4}\,\alpha_{el}\mathbf{D}^2\,\langle\gamma\rangle\sin^2\Theta\right]^2 + \frac{1}{4}\,\langle\gamma\rangle^2\left[(1+\alpha_{el}/\lambda)(E-\varepsilon_0) - \sin^2\Theta\,\mathbf{D}^2\right]^2}\,. \tag{3.6}$$

For small values of the parameter λ, the principal maximum of the photoemission strength lies at the energy ε_{coll}. Near this point the expression (3.6) reduces to the standard Breit-Wigner formula

$$\sigma^c(E) = \sigma_0\,\frac{1}{2\pi}\frac{\Gamma_{dw}}{(E-E_{dw})^2 + \frac{1}{4}\Gamma_{dw}^2}\,\Gamma_{dw}^c \tag{3.7}$$

(eq. (2.4)[1]) with

$$\Gamma_{gr}^{(rad)} = \alpha_{el}\,\mathbf{D}^2\,, \qquad \Gamma_{gr} = \langle\gamma\rangle\cos^2\Theta + \alpha_{el}\,\mathbf{D}^2 \tag{3.8}$$

describing the radiation from the isolated giant resonance. The radiation branching ratio $\mathcal{B}^{(rad)}$ decreases when the parameter λ grows. For large values of λ the main maximum is displaced to the point $E = E_v$ where the particle cross sections have the minimum due to the narrow collective state. The energy dependence is of the Breit-Wigner shape again when the radiation and total widths become equal to

$$\Gamma_0^{(rad)} = \alpha_{el}\,\mathbf{D}^2\sin^2\Theta\,, \qquad \Gamma_0 = \frac{1}{\lambda^2}\,\langle\gamma\rangle\sin^2 2\Theta + \alpha_{el}\,\mathbf{D}^2\sin^2\Theta\,. \tag{3.9}$$

This peak contains only the part $\sin^2\Theta$ of the total transition strength and is naturally ascribed to the collective state with the energy $E_0 = \varepsilon_0 + \sin^2\Theta\,\mathbf{D}^2 = E_v$ (see eq. (5.53)[1]) which acquires the dipole strength $f^0 = \sin^2\Theta$, (eq. (5.54)[1]), due to the external interaction. The nucleon width of this state diminishes and the radiation branching ratio increases together with λ. Therefore, the radiation appears again as a narrow line near the centroid of the broad resonance visible only in the particle decay channels. The radiation from the broader collective resonance is suppressed and manifests itself only as a long low tail which stretchs towards higher energies. If λ is very large the narrow line becomes the only manifestation of the giant resonance in the photoemission.

In the most interesting intermediate domain of parameters $\alpha_{el} \ll \lambda \ll 1/\alpha_{el}$ the photoemission strength is

$$\sigma^{(rad)}(E) \approx -\frac{\alpha_{el}}{2\pi}\,\mathrm{Im}\,\mathcal{P}(E)\,. \tag{3.10}$$

The interference of the radiation from the two resonances becomes strongest when $\lambda \approx 2$. The frequency spectrum of the radiation is broad in this case, its characteristic width is $\sim \mathbf{D}^2$ and the radiation intensity in the maximum is small. Generally, the shape of the spectrum is not Lorenzian when $\lambda \approx 2$.

4 Spreading Width

We now briefly discuss the role of the spreading width. As in [2], we suggest that the doorway states couple effectively to $N_{bg} \gg N_{dw}$ compound states which lie in the energy domain of the GR and have no direct access to the continuum. We also assume that the coupling matrix elements $V_{dw\,bg}$ are random Gaussian variables with zero mean value. Then, in the limit $N_{bg} \longrightarrow \infty$ the doorway Green's function changes as $\mathcal{G}^{(dw)}(E) \longrightarrow \mathcal{G}^{(dw)}(E - \Delta + \frac{i}{2}\Gamma^{\downarrow})$ [2] where Δ and $\Gamma^{\downarrow}$ are the energy shift and spreading width respectively. Neglecting their possible slow energy dependence in the whole domain of the GR, we can fully incorporate the hermitian shift Δ (which is in fact small due to statistical reasons) into the mean position ε_0. The only effect of the interaction with the background states is then the additional shift of the poles of the transition amplitudes along the imaginary direction in the complex energy plane. Note that the integral sum rule

$$\int_{-\infty}^{\infty} dE\, \sigma^c(E) = (\mathbf{A}^c)^2 = \sum_{\tau} \Gamma^c_{\tau} \tag{4.1}$$

(eq. (2.36)[1]) survives the transformations made.

We will not present here the rather cumbersome general expressions. Confining ourselves for the sake of simplicity to the two-level approximation, the shift considered does not influence the relations (5.45)[1] – (5.49)[1] together with (5.44)[1], (5.45)[1]. We suggest further that the displacement $\mathbf{D}^2$ is smaller than both the escape and spreading widths. Then, the transition strength corresponding to the particle emission in a channel c acquires the Breit-Wigner shape with the centroid $E_{gr} = \cos^2\Theta\,\mathbf{D}^2$ and the total width $\Gamma_{gr} = \langle\gamma\rangle + \Gamma^{\downarrow}$.

The evolution of the γ-strength $\overline{\sigma^{(rad)}(E)}$, when the escape width $\langle\gamma\rangle$ changes from values smaller than $\Gamma^{\downarrow}$ to larger ones, is appreciably richer. The strength transforms smoothly from

$$\overline{\sigma^{(rad)}(E)} = \frac{1}{2\pi}\,\alpha_{el}\mathbf{D}^2\,\frac{\Gamma^{\downarrow}}{(E-\mathbf{D}^2)^2 + \frac{1}{4}\,(\Gamma^{\downarrow})^2} \tag{4.2}$$

for $\langle\gamma\rangle \ll \Gamma^{\downarrow}$ to

$$\overline{\sigma^{(rad)}(E)} = \frac{1}{2\pi}\,\alpha_{el}\sin^2\Theta\,\mathbf{D}^2\,\frac{\Gamma^{\downarrow}}{(E-\sin^2\Theta\mathbf{D}^2)^2 + \frac{1}{4}\,(\Gamma^{\downarrow})^2} \tag{4.3}$$

in the opposite limit $\langle\gamma\rangle \gg \Gamma^{\downarrow}$. In the intermediate region, the maximum monotonously decreases and moves towards lower energies. The shape of the radiation spectrum is not Lorentzian when both widths are of comparable values. Eq. (4.3) implies in particular the loss of an appriciable part ($= \cos^2\Theta$) of the transition strength, since the contribution of the broader collective state which is described by a right long tail in Fig. 3(d) (see next section) is invisible in eq. (4.3). It is worthy noting that the width of the γ-spectrum in our schematic model is always determined mainly by the spreading width. The escape width $\langle\gamma\rangle$ drops out not only from eq. (4.2) but also from eq. (4.3). This is due to the fact that the radiating state becomes almost trapped.

Therefore, all results obtained are valid only if the spreading width $\Gamma^{\downarrow}$ is noticeably less than the escape width of the GR. In hot nuclei in which we are interested, this condition seems to be fulfilled: According to experimental data [4, 5], the escape width is growing with excitation energy whereas the spreading width saturates.

5 Numerical Results and Discussion

The behaviour of the dipole strengths, energies and widths of the interfering resonance states is reflected in the cross section pattern as shown above analytically by using the two-level approximation. Below, we show the results of numerical investigations performed under less restrictive assumptions. The calculations are performed with the same 10 levels and 3 channels as in [1]. Damping is not taken into account, i.e. the results are true only for $\langle\gamma\rangle \gg \Gamma^{\downarrow}$ (see the discussion in Sect. 4).

In Figs. 1 to 3 we show the energy dependence of the transition strengths into particle and photo channels for the three values of the overlap parameter $\lambda = 0.1$, 2 and 5. As in the figures in [1], the energy E is measured in units of $\mathbf{D}^2$. Due to the strong interference, the pattern is noticeably different in the different final channels. One nicely sees the shift of the maximum at the higher energy towards lower energies which is predicted by the two-level approximation. Moreover, the fragmentation of the maximum at the lower energy into a number of resonances can be seen which, of course, disappears in the limit of degenerate unperturbed levels e_n. At last, the growing restructuring of the dipole strength with increasing external coupling in favor of the lower-lying components is seen in Figs. 1(d) to 3(d). For example the summed strength above $E > 0$ amounts to 99%, 87% and 85% in the case of the degenerate unperturbed spectrum (dashed lines). As to the maximum value of the transition strength into the photo channel at the higher energy, it drops down by a factor of more than 10 when λ increases from 0.1 to 2, while a narrow high peak appears in agreement with the analytical consideration at lower energy when λ becomes large.

The elastic and photo-nuclear reaction cross sections are shown in Figs. 4 and 5. They are calculated for the same three values $\lambda = 0.1, 2$ and 5 as the transition strengths in Figs. 1 to 3. Both the shift of the dipole resonance to lower energies and the loss of its dipole strength are seen very clearly also in these values.

Thus, we have the following picture. Provided that the coupling (2.4) is omitted, the two doorway states of the two-level approximation fully exhaust the total dipole strength so that only they can radiate γ-rays. The radiation pattern determined by the two doorway states turns out to be very sensitive to their degree of overlapping: as long as the energy displacement of one of them is appreciably larger than the sum of the particle escape widths (i.e. $\lambda \ll 1$) only one of them radiates. If, however, they overlap ($\lambda \sim 1$ or $\lambda > 1$) the interference leads to a strong redistribution of the dipole strength as well as the escape width between the two states. When the degree of overlapping exceeds some critical value the escape width of one of the states starts to decrease (dynamical trapping effect). This effect is governed by the avoided crossing of two resonances described in detail in [6].

If the coupling (2.4) is taken into account, then the mixing of all the doorway states leads to an additional restructuring of the total dipole strength in favour of the low-lying components. In any case, the nearly trapped state acquires an appreciable dipole strength and therefore would, in the absence of any internal damping, radiate a narrow electromagnetic line in the vicinity of the centroid of the broad bump observed in the particle channels. The other broad states, which also possess noticeable dipole strengths, contribute mostly to a long radiation tail stretched towards larger energies. They manifest themselves mainly in the particle channels. The internal damping enlarges the width of the γ-line from the trapped state by $\Gamma^{\downarrow}$ and masks the tail. This fact results in a seeming loss of a part of the radiation transition strength $\overline{\sigma^{(rad)}(E)}$.

Both the loss of the dipole strength and the shift of a part of it towards lower energies are discussed at present in connection with experimental results obtained for collective excitations in hot nuclei (see e.g. the Proceedings of the Gull Lake Nuclear Physics Conference on Giant resonances, 1993, [7]). The γ-ray multiplicity from the decay of giant dipole resonances is shown experimentally to increase with the excitation energy, as long as it is not too high, in agreement with the 100% sum-rule strength. At higher energies, however, its saturation signals the quenching of the multiplicity and the existence of a limiting excitation energy for observation of GDR through its γ emission. The different existing theoretical approaches can only partly explain the experimental situation observed [8].

The results obtained in the present paper point to a new mechanism which could possibly shed an additional light on the problem. The main feature of this interference pattern, described well in the two-level approximation, is caused by the avoided resonance crossing of two states decaying into the same open channels. The visible bulk of the GDR γ emission originates from a specific state with dynamically reduced particle escape width but large dipole moment (the trapped collective state) while the emission from the broader state is suppressed being spread over a wide energy range.

6 Summary

On the basis of a phenomenological schematic model we investigated the interferences between the different doorway components of a giant multipole resonance and considered its influence onto the cross section pattern. The photoemission turns out to be especially sensitive to the degree of overlapping of the collective doorway states. The interference between them leads, at a certain critical value of the external coupling, to a strong redistribution of the dipole strength in favour of the low-lying components.

The internal damping due to the coupling of the doorway states to the background of complicated states masks or even smears out the effects of the interference when the corresponding spreading width exceeds the total escape widths of the doorway components. In very hot nuclei, however, it is possible that the spreading width is smaller than the escape width [4, 5]. If so, the interference manifests itself, in particular, in an apparent quenching of the dipole strength of GR. The saturation of the γ multiplicity observed experimentally at about 250 MeV excitation energy in heavy nuclei [7, 8] may be therefore, at least partly, caused by the interference phenomena discussed in the present paper. Further investigations of this interesting question are necessary.

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Figure Captions

Fig.1
The transition strengths into particle (a,b,c) and photo (d) channels for $\lambda = 0.1$ and the electromagnetic interaction strength $\alpha_{el} = 0.01$. The resonance states are the same as in Fig. 2 in [1]. The dashed lines correspond to the case of parental levels fully degenerated ($\Delta_e = 0$).

Fig.2
The same as in Fig.1 but for $\lambda = 2$.

Fig.3
The same as in Fig.1 but for $\lambda = 5$. Note the different E scale in (d).

Fig.4
The elastic cross section for $\lambda = 0.1$ (a), $\lambda = 2$ (b), and $\lambda = 5$ (c). The resonance states are the same as in Fig. 2 in [1]. Note the different E scale in (a).

Fig.5
The photo-nuclear cross section for $\lambda = 0.1$ (a), $\lambda = 2$ (b), and $\lambda = 5$ (c). The resonance states are the same as in Fig. 2 in [1]. Note the different E scale in (c).

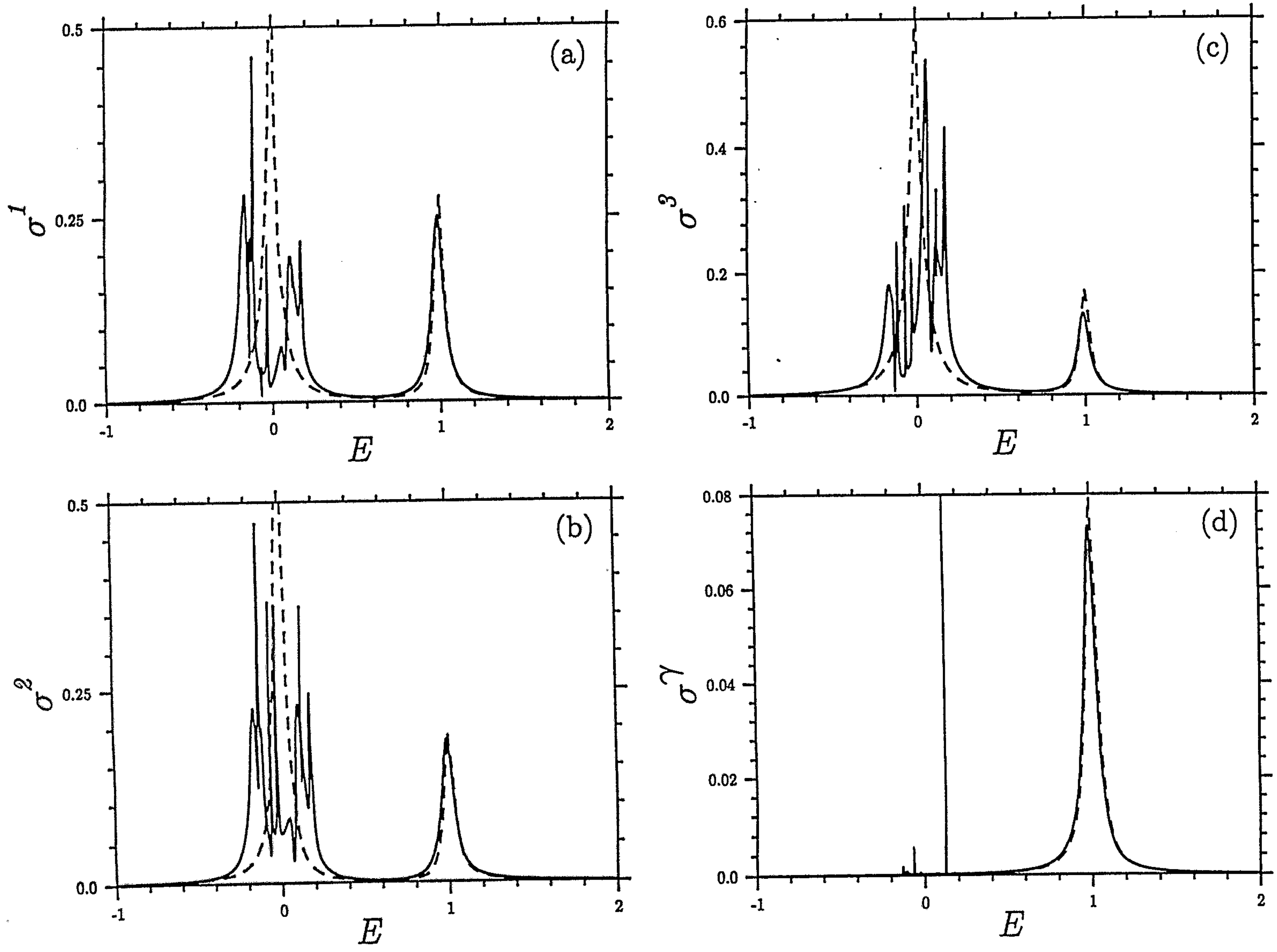
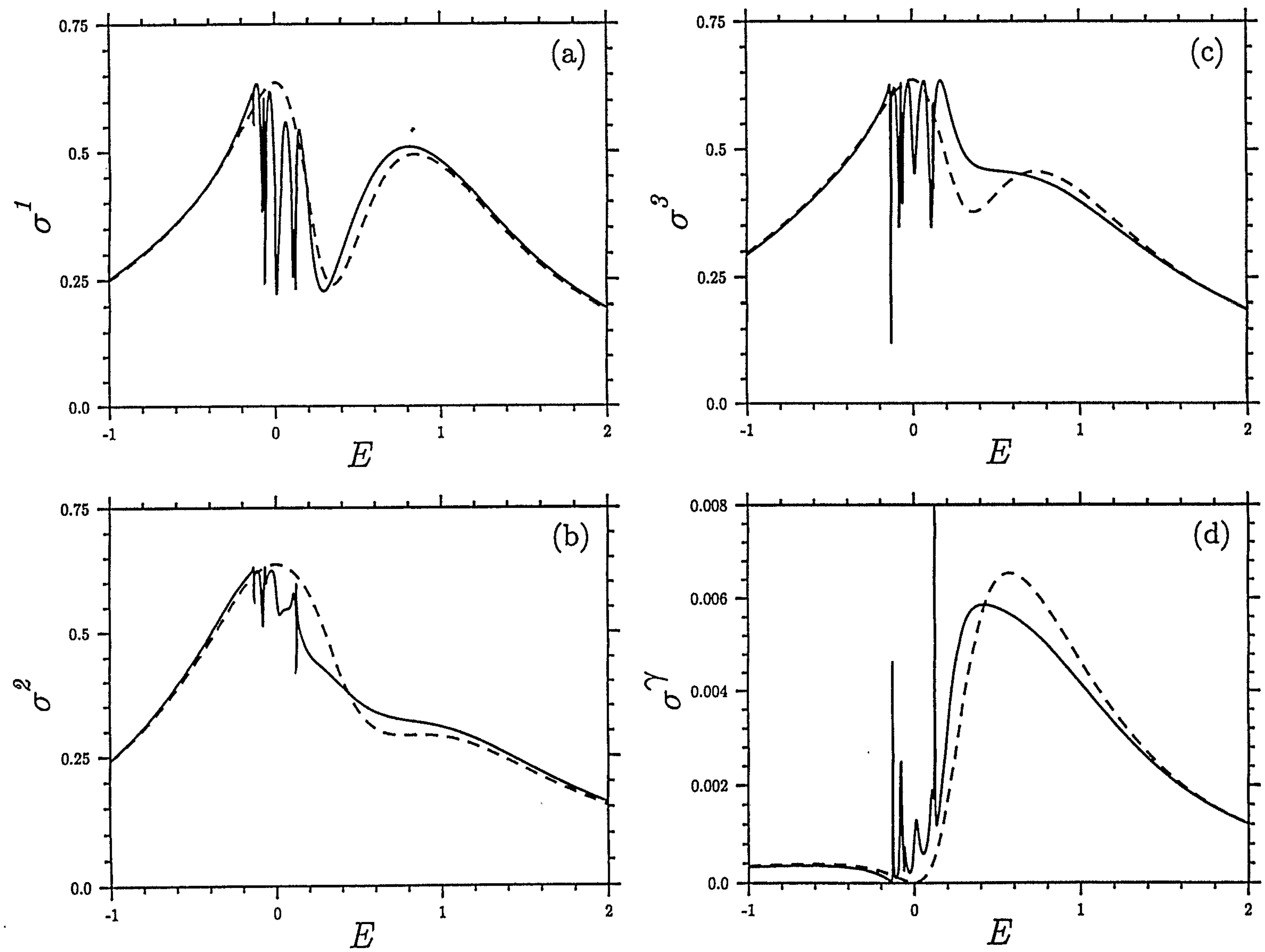


Fig. 1


(a)
(b)
(c)
(d)
σ^1
σ^2
σ^3
σ^γ
E
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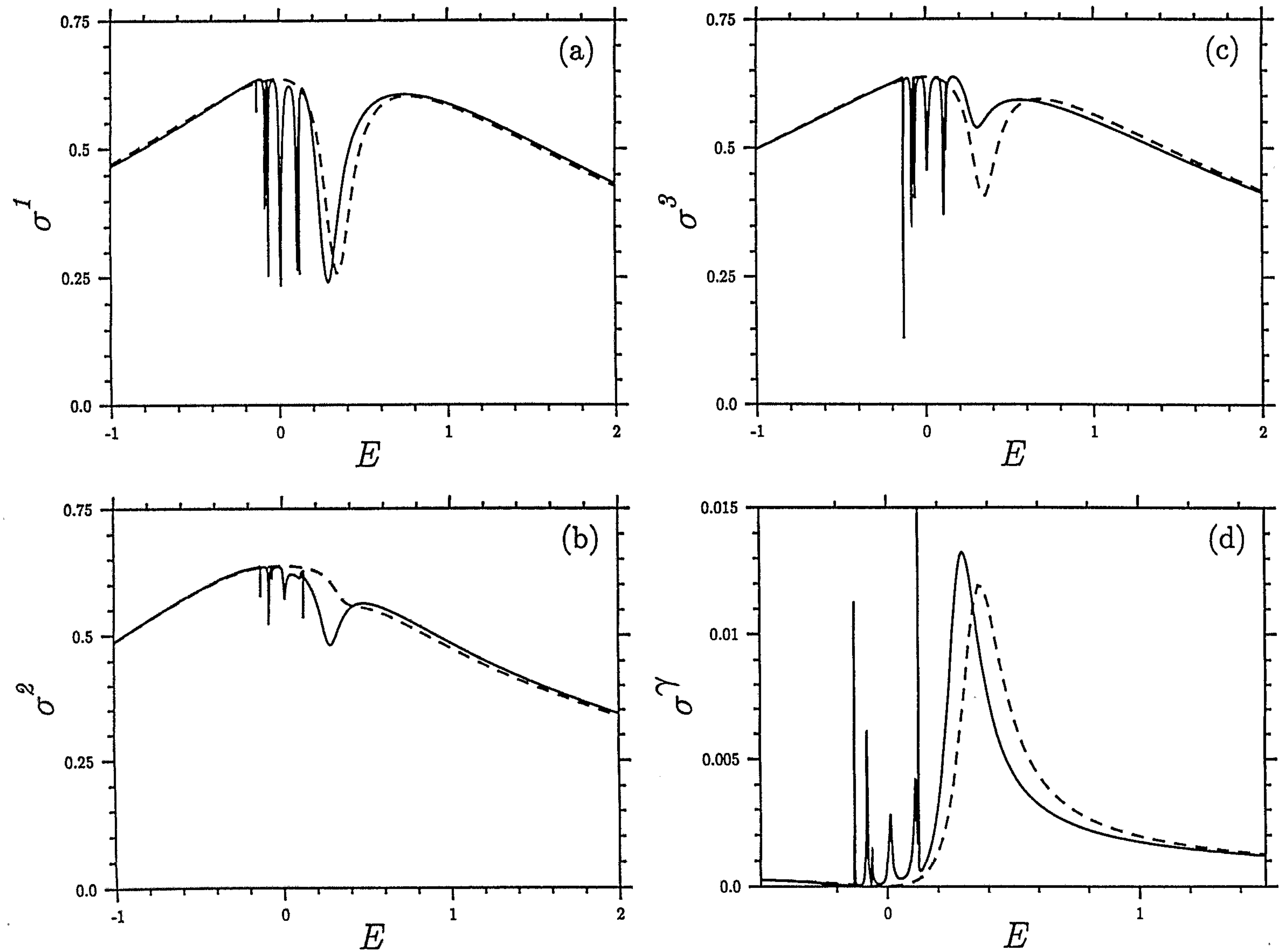
0.25
0.0
0.008
0.006
0.004
0.002
-1
0
1
2

Fig. 2



(a)
(b)
(c)
(d)
σ^1
σ^2
σ^3
σ^γ
E
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0.5
0.25
0.0
0.015
0.01
0.005
-1
0
1
2

Fig. 3

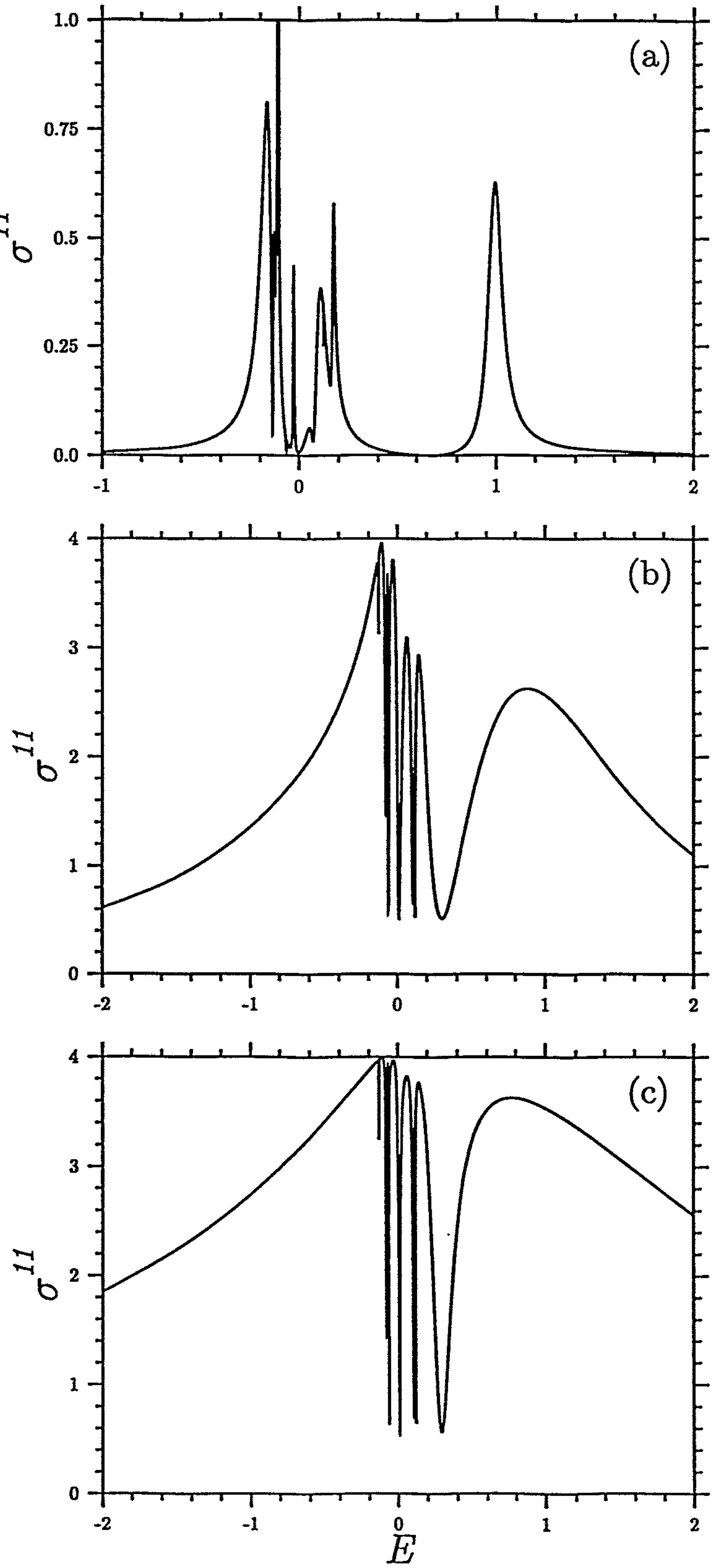

Fig. 4

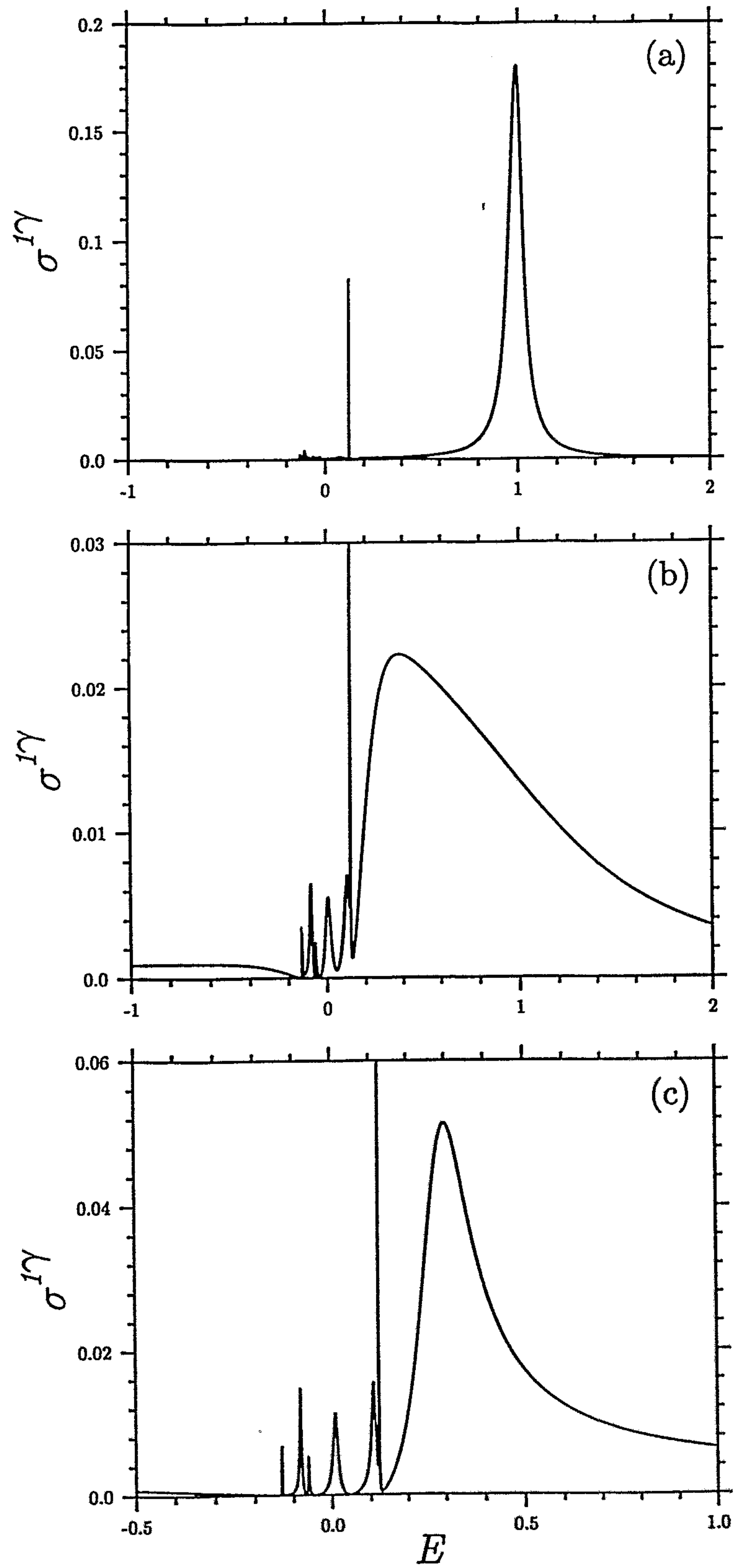
(a)
(b)
(c)
$\sigma^{1\gamma}$
E
0.2
0.15
0.1
0.05
0.0
0.03
0.02
0.01
0.06
0.04
-1
0
1
2
-0.5
0.5
1.0

Fig. 5