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Suppression of nucleon resonances in the total photoabsorption on nuclei

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Abstract

We analyze the recent nuclear-photoabsorption data which have shown the disappearance on nuclei of the resonances higher than the Δ. We discuss how the Fermi motion, the collision broadening, and the Pauli blocking might distort the shape of nucleon resonances in the nuclear medium, and show that the Fermi motion and the collision broadening play the major role. From a fit to the photoabsorption data on uranium we have estimated the total cross sections for scattering of the resonances $P_{33}(1232)$, $D_{13}(1520)$, $F_{15}(1680)$, and $D_{33}(1700)$ on nucleons. The estimate for the P_{33}-resonance cross section is in fair agreement with the value derived from the $NN \to \Delta N$ cross sections, while those for the other resonances are high although do not violate the unitarity limits.

1. Introduction

Recently the total nuclear-photoabsorption cross section on several nuclei has been measured at Frascati [1–4] and at Mainz [5,6], respectively, with a 0.2–1.2 GeV and 0.05–0.8 GeV tagged-photon beams, to study the excitation of baryon resonances in nuclei. These measurements have revealed an interesting result: the $D_{13}(1520)$ and $F_{15}(1680)$ resonances, clearly seen in the photoabsorption measurements on hydrogen and deuteron [7,8], disappear on nuclei, and the cross section on a bound nucleon is damped in comparison with the one on a free nucleon. On

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the contrary, the $P_{33}(1232)$ resonance is only slightly distorted, in agreement with the previous data available in the literature [9,10].

There are several effects that may distort the shape of the baryon resonances in the nuclear medium, specifically: (i) The Fermi motion, which produces a smearing of the resonance peak by suppressing its maximum and increasing its width. (ii) The propagation and interaction of the resonances in the nuclear medium, which affect the parameters (mass and width) of the resonances. The relevant collision shift and broadening depend on the nuclear density as well as on the internal structure of the resonance. (iii) The Pauli blocking, which reduces the resonance width by decreasing the space phase available for the resonance decay products. (iv) The partial deconfinement of quarks in nuclei, which can reduce the strength of the nucleon resonances.

The nuclear photoabsorption in the Δ-resonance region was discussed in the past (see e.g. Refs. [11–13] and references therein). It was shown that the Δ–hole model can describe the data within a 10% accuracy, accounting for the decreasing and broadening of the cross section through the propagation of the Δ-isobar in the nucleus.

In this paper we concentrate our attention on the region of the resonances $D_{13}(1520)$ and $F_{15}(1680)$ and show that the Fermi motion and especially the collision broadening are the main effects which distort their behavior in the nuclear medium. Moreover, from a fit to the nuclear-photoabsorption data on ^{238}U [1] we estimate the total cross sections for scattering of the resonances $P_{33}(1232)$, $D_{13}(1520)$, $F_{15}(1680)$ and $D_{33}(1700)$ on nucleons. For the Δ-resonance, the obtained estimate is in fair agreement with the experimental data on the $NN \to \Delta N$ cross section, while for the other resonances the estimates are higher but do not violate the unitarity limits. We discuss also how the results for the collision cross sections change if one considers a 10% damping of the strengths of all the resonances apart from the Δ.

The organization of the paper is the following: In Section 2 we discuss the effects which can distort the shapes of the nucleon resonances in the nuclear medium. In Section 3 we fit the photoabsorption data and find the N*N scattering cross sections. In Section 4 we discuss the results for the N*N scattering cross sections and compare them with the unitarity limits and the available experimental data. Finally, Section 5 contains the conclusions.

2. The total photoabsorption cross section

2.1. *The forward Compton-scattering amplitude and the Fermi motion*

The total photoabsorption cross section can be expressed through the imaginary part of the forward Compton-scattering amplitude $\mathscr{T}(\omega, \vartheta = 0)$ (see Fig. 1):

$$\frac{\sigma_{\gamma A}^{\text{tot}}}{A} = -\frac{1}{2m\omega} \text{Im}\ \bar{\mathscr{T}}(\omega, \vartheta = 0), \tag{1}$$

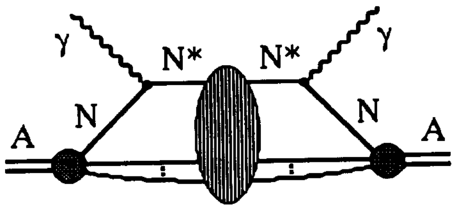

Fig. 1. Forward Compton scattering on nuclei.

where ω is the photon energy, m is the nucleon mass, $\mathscr{F} = \Sigma_i[(Z/A)\mathscr{F}_{i,\gamma p} + (N/A)\mathscr{F}_{i,\gamma n}]$ + background, with the index i referring to different baryon resonances N*. The upper line in Eq. (1) means averaging on the distribution of nucleon momenta in a nucleus (the effect of the Fermi motion).

If one assumes that the resonances formed inside a nucleus do not interact with the surrounding nucleons (see the graph shown in Fig. 2) the corresponding amplitude takes the form

$$\mathscr{F}_{\gamma N}(\omega, \theta = 0) = \int \frac{\mathrm{d}\boldsymbol{p}}{(2\pi)^3} |\varphi(\boldsymbol{p})|^2 \frac{g_{\gamma NN^*}^2}{(k+p)^2 - M^2 + iM\Gamma_0}, \tag{2}$$

where $k = (\omega, \boldsymbol{k})$ and $p = (E, \boldsymbol{p})$ are the photon and nucleon four-momenta, M and Γ_0 are the resonance mass and width in the vacuum, $\varphi(\boldsymbol{p})$ is the momentum distribution of nucleons in the nucleus, which we shall describe here by the Fermi gas model. The vertex function $g_{\gamma NN^*}$ is related to the partial width $\Gamma_{\gamma N}$ as follows:

$$\Gamma_{\gamma N} = \frac{g_{\gamma NN^*}^2}{16\pi} \frac{M^2 - m^2}{M^3}.$$

The Fermi motion results in the broadening of the resonance peak and the suppression of its value at the maximum. The suppression factor S_F is a decreasing function of $x = 2\omega p_F/M\Gamma_0$, where p_F is the Fermi momentum, as shown in Fig. 3. However, because the area under the resonance curve is not altered, the effective width will be by a factor $1/S_F$ larger than the original one.

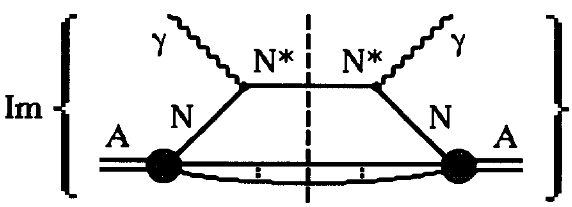

Fig. 2. Imaginary part of the forward Compton-scattering amplitude contributing to the nuclear-photoabsorption cross section. The baryon resonances N* being formed do not interact with other nucleons in the nucleus.

2.2. Propagation of resonances in nuclear medium

The collision broadening of resonances in nuclei was discussed earlier by Bugg [14] who started from the Wigner–Weisskopf approximation [15] for the time development of the resonance and considered its propagation in a uniform nuclear medium. He also considered the effect of the collision broadening in the πd, Kd total cross sections and in the ρ-meson production by pion beam.

For the sake of completeness we present here a different but equivalent derivation of the collision-broadening effect valid in the case when the Compton wavelength of the resonance is much smaller than the nuclear radius. We shall apply this result to the description of the photoabsorption data and discuss the validity of the used approach.

2.2.1. An elementary approach

A resonance created in the nuclear matter interacts with the surrounding nucleons. Let us consider a simple model of a rather fast resonance propagating through the infinite nuclear matter. We assume also that the resonance is rather narrow (to treat it as an elementary particle) and the density of the medium is rather low (to avoid non-linear effects in the density). The intensity at a time t of a resonance beam propagating with a velocity v through the vacuum is described by the well-known expression:

$$|\Psi(t)|^2 = |\Psi(0)|^2 \exp(-\Gamma_0 t/\gamma), \tag{3}$$

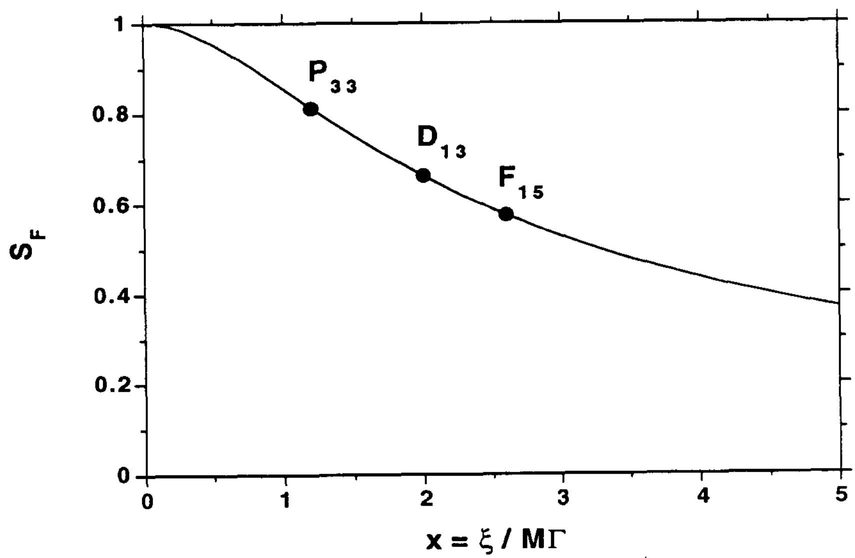

Fig. 3. Suppression factor S_F due to the Fermi motion of nucleons versus x. The dots on the curve show the positions of the resonances $P_{33}(1232)$, $D_{13}(1520)$, and $F_{15}(1680)$.

where γ is the Lorentz factor. Therefore, the intensity of the resonance beam after a path $l = vt$ in the vacuum is decreased by the factor $\exp(-\Gamma_0 l/\gamma v)$.

In the nuclear medium, the intensity of the resonance beam is further decreased by the scattering or absorption processes. Therefore, even in the case $\Gamma_0 \to 0$, the number of the initial particles in the beam decreases with the path l as $\exp(-\rho\sigma^* l)$, where σ^* is the total cross section of the resonance–nucleon interaction, and ρ is the density of the medium. If both effects are present one gets

$$|\Psi(t)|^2 = |\Psi(0)|^2 \exp(-\Gamma t/\gamma),$$

where $\Gamma = \Gamma_0 + \Gamma^*$ and the effect of the collision broadening is

$$\Gamma^* = \rho\sigma^* v\gamma, \tag{4}$$

which is in agreement with the result found by Bugg [14].

Therefore, the propagation of a resonance in the nuclear medium results in the decreasing of its lifetime or, according to the uncertainty principle, in the increasing of its width. The order of magnitude of the width broadening, Γ^*, can be estimated to be 135 MeV for a nuclear density $\rho = 0.17$ fm^{-3} and $\sigma^* v\gamma = 40$ mb.

For a better justification of Eq. (3) we shall consider below the quantum-mechanical approach.

2.2.2. *Infinite nuclear matter: quantum-mechanical approach*

The standard approach to describe the resonance in quantum mechanics is to consider it as a state with the complex energy eigenvalue: $\Psi(t, r) = \Psi(r)\exp(iE^* t)$, where $E^* = M + k^2/2M - \frac{1}{2}i\Gamma_0$, with M and Γ_0 the mass and the width of the resonance in the vacuum. When the resonance propagates through the nuclear medium its wave function can be found by solving the Schrödinger equation:

$$i\frac{\mathrm{d}\Psi(t,r)}{\mathrm{d}t} = \left(M - i\frac{\Gamma_0}{2} - \frac{\nabla^2}{2M} + V_{\mathrm{opt}}(r)\right)\Psi(t,r), \tag{5}$$

where the optical pseudopotential $V_{\mathrm{opt}}(r)$ can be expressed through the forward N*N scattering amplitude, $f_{\mathrm{N^*N}}(0)$, and the density of the nuclear matter (see Fig. 4):

$$V_{\mathrm{opt}}(r) = -\frac{2\pi}{M} f_{\mathrm{N^*N}}(0)\rho(r). \tag{6}$$

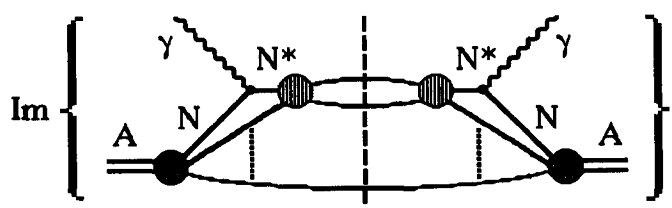

Fig. 4. Imaginary part of the forward Compton-scattering amplitude contributing to the total nuclear-photoabsorption cross section. The baryon resonances N* being formed scatter inelastically on other nucleons in the nucleus.

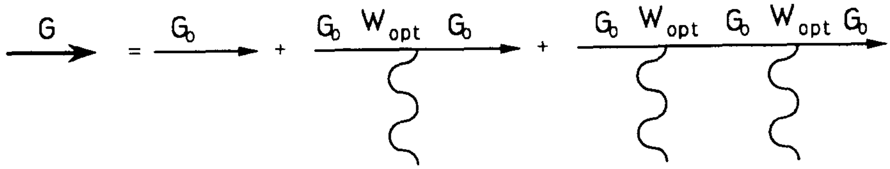

Fig. 5. The diagrams describing the relativistic resonance propagator in the nuclear medium.

For a constant density $\rho(r)=\rho_0$ the non-relativistic propagator of the resonance, $G_{\text{n.r.}}$, can be written as

$$G_{\text{n.r.}}=\frac{1}{E-(M+\delta M)+\frac{1}{2}i(\Gamma_0+\Gamma^*)}, \tag{7}$$

where the two extra terms in the denominator, δM and Γ^*, describe the shift of the resonance mass and the increase of its width (the effect of the collision broadening),

$$\delta M=-\frac{2\pi}{M}\text{Re } f_{\text{N*N}}(0)\ \rho_0, \tag{8}$$

$$\Gamma^*=\frac{4\pi}{M}\text{Im } f_{\text{N*N}}(0)\ \rho_0. \tag{9}$$

Using the optical theorem, Im $f_{\text{N*N}}(0)=(k/4\pi)\sigma^*$, we find $\Gamma^*=\rho\sigma_0^*v\gamma$, which is in agreement with Eq. (4).

2.2.3. *Relativistic generalization*

The relativistic propagator of the resonance G can easily be constructed using the diagram technique. Summing up the series of the diagrams depicted in Fig. 5 we can write the following expression for the propagator:

$$G=G_0+G_0W_{\text{opt}}G_0+G_0W_{\text{opt}}G_0W_{\text{opt}}G_0+\cdots,$$

where $G_0^{-1}=s-M^2+iM\Gamma_0$, and the relativistic optical pseudopotential has the form $W_{\text{opt}}=-4\pi f_{\text{N*N}}(0)\rho$. In the case of constant density $\rho=\rho_0$ this series can easily be summed up:

$$G=\left(G_0^{-1}-W_{\text{opt}}^{(0)}\right)^{-1}, \tag{10}$$

and the mass and width corrections are given by the expressions

$$\delta M^2=-4\pi\ \text{Re } f_{\text{N*N}}(0)\ \rho_0\approx 2M\delta M,$$

$$\Gamma^*=\frac{4\pi}{M}\text{Im } f_{\text{N*N}}(0)\ \rho_0=\rho_0\sigma^*v\gamma,$$

which are equivalent to Eqs. (8) and (9).

2.2.4. *Finite nuclear size and average density approximation*

In principle the exact resonance propagator should be used to take into account the effects of a nuclear medium in case of finite nuclei. One can also think that for medium and heavy nuclei the average-density approximation (ADA) for the resonance propagator should be good:

$$G_{\mathrm{av}} = G_0 + G_0 W_{\mathrm{av}} G_0 + G_0 W_{\mathrm{av}} G_0 W_{\mathrm{av}} G_0 + \cdots = \left(G_0^{-1} - W_{\mathrm{av}}\right)^{-1}, \tag{11}$$

where $W_{\mathrm{av}} = -4\pi f_{\mathrm{N^*N}}(0)\rho_{\mathrm{av}}$, and ρ_{av} is the averaged nuclear density.

The derivation of ADA and the comparison of this approximation with the particle–hole model [11–13] are discussed elsewhere [16]. Here, it is useful to stress that ADA can be justified in the eikonal approximation when the wavelength of the resonance is much smaller than the nuclear size. In this case a resonance formed by photons in a nuclear medium propagates along a straight line with a fixed impact parameter. Then the amplitude is integrated over all the impact parameters weighted by the nuclear-thickness function. The average momenta of the $\Delta(1232)$ and N*(1520) formed by photons inside the nuclei are about 300 and 760 MeV/c, respectively, which means that the necessary condition for validity of ADA is satisfied even for $\Delta(1232)$. In the case of heavier resonances ADA is expected to work even better.

Taking into account that in our case the resonances are formed by photons in all of the nuclear volume and that their mean free path even for decay does not exceed 0.5 fm, we can conclude that at least for heavy nuclei we can safely substitute in Eq. (11) $\rho_{\mathrm{av}} = \rho_0$, where ρ_0 can be taken from the standard parametrization of the nuclear density.

2.2.5. *Rescattering of the resonance decay products*

A resonance formed by a photon (see Fig. 1) may decay before it collides with nucleons inside a nuclear medium. Then its decay products may rescatter and this effect may give some corrections. Let us demonstrate that these corrections are small in the case of the photoabsorption processes.

At first we consider the formation and decay of the Δ-resonance where the only πN decay channel is dominant. The time development of the state vector of the resonance in its rest frame can be written as the superposition of two components:

$$|\Psi\rangle = \exp(-iMt)\left\{\exp(-\tfrac{1}{2}\Gamma_0 t)\,|\mathrm{N}^*\rangle + \left[1 - \exp(-\Gamma_0 t)\right]^{1/2} |\pi\mathrm{N}\rangle\right\}. \tag{12}$$

Because the Δ-resonance decays mainly inside the nucleus, the pion and the nucleon may rescatter after its decay. It is convenient to discuss those rescattering effects in the framework of the multiple-scattering theory and distinguish the coherent and noncoherent parts of the photoabsorption cross section.

In the case of the noncoherent cross section the rescattering of the particles in the final state does not change the value of the cross section. For the rescattering and absorption of pions this is the consequence of unitarity which can be applied when the wavelength of the pion is much less than the nuclear radius (see e.g. Ref. [17]). For the rescattering of nucleons this is the consequence of the completeness condition for the eigenstates of the hamiltonian for the system of A nucleons in the final state.

In the case of the coherent photoproduction of π^0-mesons the cross section is affected by the rescattering of pions which produces nuclear shadowing. However, the coherent photoproduction of π^0-mesons on nuclei is strongly suppressed. Indeed the coherent-photoproduction amplitude is proportional to $\boldsymbol{q}\cdot(\boldsymbol{k}\wedge\boldsymbol{e})$, where $\boldsymbol{q}$ is the momentum transfer and $\boldsymbol{e}$ is the photon-polarization vector, and is equal to zero at zero angle. Therefore, the contribution of the coherent-photoproduction cross section is equal to zero for infinite nuclear matter and is negligibly small for heavy nuclei. On the contrary, if a Δ-resonance is formed by a pion beam the coherent rescattering of pions is very important and this effect produces an extra distortion of the Δ-peak in the total cross section of pion–nucleus interactions (a discussion of these effects in the framework of the Δ–hole model is given in Refs. [11–13]).

For heavier resonances the $\pi\pi$N decay channel may also be important. In this case we have to consider also the possibility of the coherent photoproduction of a dipion in the 1^- state. The amplitude of this process is equal to zero at zero angle. However, this cross section is strongly suppressed by the nuclear form factor because of the rather high value of the longitudinal momentum transfer in the considered energy interval (see also the discussion at the end of Section 3.3).

2.2.6. Medium effects on the Δ propagation in the potential model

In many papers (see e.g. Refs. [11–13,18]) the medium effects on the Δ propagation in nuclei at the resonance energies were treated in the framework of the potential model where the interaction of the Δ-resonance with the nucleus is described by a complex phenomenological potential [18]

$$V_\Delta(E, r) = V_0(E)\frac{\rho(r)}{\rho(0)} + 2\boldsymbol{l}_\Delta\cdot\boldsymbol{s}_\Delta V_{1s}(E). \tag{13}$$

The central part, V_0, was taken proportional to the nuclear density $\rho(\boldsymbol{r})$, and the spin–orbit term, V_{ls}, ($\boldsymbol{l}_\Delta$ and $\boldsymbol{s}_\Delta$ are the orbital angular momentum and the spin operators of the propagating Δ) was introduced by analogy with the nucleon–nucleus optical potential. For pion energies between 100 and 250 MeV it was found that the central potential V_0 does not depend on the energy and its imaginary and real parts are: Im $V_0 = -40$ MeV and Re $V_0 = -30$ MeV. The spin–orbit term, V_{ls}, is considerably smaller and is needed only for the description of $\mathrm{d}\sigma/\mathrm{d}\Omega$ at large center-of-mass angles, $\vartheta_{\mathrm{c.m.}} > 60°$.

The potential (13) with the above parameters was used by Koch et al. [11] for the description of the photoabsorption data on ^{16}O and ^{12}C in the framework of the Δ–hole model. Several improvements were later made by Carrasco and Oset [13], who used a more complex model with π and ρ exchanges for the ΔN scattering amplitude.

The real and imaginary parts of the optical potential (13) produce the shift of the mass δM and the broadening Γ^* of the resonance. Here we will use the pseudopotential (6). The imaginary and real parts of the optical pseudopotential are proportional to the imaginary and real parts of the forward N*N scattering amplitude. In Section 3, we consider the shift of the mass, δM, and the broadening of the width, Γ^*, as free parameters of the fit and use them to determine the real

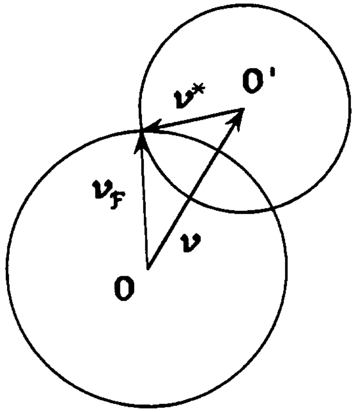

Fig. 6. The illustration of the geometrical meaning of B_F for the case $v > v_F$. The decay of the resonance with a nucleon in the final state is allowed when the velocity $\boldsymbol{v} + \boldsymbol{v}^*$ of the nucleon is outside the sphere formed by the vector $\boldsymbol{v}_F$.

and imaginary parts of the forward N^*N scattering amplitude. Moreover, in the case of Δ we compare our results with the predictions of the Δ-hole model.

2.3. Pauli blocking

Another correction to the resonance width in the nuclear medium is due to the Pauli blocking, which arises because of the occupation of the momentum space $p < p_F$ by other nucleons. This effect leads to an increase of the resonance lifetime in the nuclear medium and to a decrease of its width. If only the pion–nucleon decay mode is dominant, the blocking factor in the non-relativistic approximation can be written as

$$B_F = \begin{cases} 1 & \text{when} \quad |v - v^*| \geqslant v_F, \\ 0 & \text{when} \quad v + v^* \leqslant v_F, \\ \dfrac{(v + v^*)^2 - v_F^2}{4vv^*} & \text{when} \quad |v - v^*| \leqslant v_F \leqslant v + v^*, \end{cases} \tag{14}$$

where v is the velocity of the resonance, $\boldsymbol{v}^*$ is the velocity of the nucleon in the resonance rest frame, and $v_F = p_F/m$. The geometrical meaning of B_F is illustrated in Fig. 6 for the case when $v > v_F$. The decay of the resonance is allowed when the velocity of the nucleon $(\boldsymbol{v} + \boldsymbol{v}^*)$ in the rest frame of the nucleus is out of the sphere formed by the vector $\boldsymbol{v}_F$. Therefore, B_F can be identified with the solid angle formed in the resonance rest frame by the allowed nucleon directions:

$$B_F = 1 - \int_{\vartheta' < \vartheta} \frac{d\Omega'}{4\pi} = \tfrac{1}{2}(1 + \cos\vartheta), \tag{15}$$

where $\cos\vartheta = (v^2 + v^{*2} - v_F^2)/2vv^*$.

Therefore, the width of the resonance in the nuclear medium is given by

$$\Gamma = \Gamma_0 B_F + \Gamma^*. \tag{16}$$

We get, in particular, $B_F = 0.68$ for the Δ and observe that, neglecting the collision broadening, the Δ-width in the nuclear medium would be $\Gamma_0 B_F = 82$ MeV only.

It should also be taken into account that the higher-mass resonances have considerable probabilities of decaying with more than one pion in the final state. For example, the resonance N*(1520) has branching ratios (50–60)% and (40–50)% into the $N\pi$ and $N\pi\pi$ channels, respectively. The last branching ratio is distributed with the probabilities (15–30)% and (10–25)% between the $\Delta\pi$ and $N\rho$ states. For this resonance the effect of the Pauli blocking is about 10%. For the other resonances the effect is even smaller.

3. Results

3.1. Description of the proton and neutron data

In order to describe the nuclear-photoabsorption data and isolate the medium effect, we first determined the resonance masses and widths by fitting the proton- and neutron-photoabsorption data [7,8] to a sum of eight Breit–Wigner resonances plus a smooth background. The resonances used in the fit are the $P_{33}(1232)$, $P_{11}(1440)$, $D_{13}(1520)$, $S_{11}(1535)$, $D_{15}(1675)$, $F_{15}(1680)$, $D_{33}(1700)$, and $F_{37}(1950)$, because they give the main contributions to the total γN cross section below 1.2 GeV. The contribution of the ignored resonances is lower than 20 μb, and was included into the background. The resonance contributions were parametrized in the usual Breit–Wigner form:

$$\sigma^i(s) = \frac{\sigma^i_{\text{res}} M^2 \Gamma^2}{(s - M^2)^2 + M^2 \Gamma^2}, \tag{17}$$

where $i = $ p,n. We required that the masses and widths of the resonances agree with their experimental values taken from the Particle Data Book [19]. The ratio $\sigma^{\text{n}}_{\text{res}} / \sigma^{\text{p}}_{\text{res}}$ was also fixed according to the experimental data [23] in the case of $I = \frac{1}{2}$ resonances, and it was assumed equal to 1 in the case of the $I = \frac{3}{2}$ resonances. The values of $\sigma^{\text{p}}_{\text{res}}$ were free parameters of the fit.

The background was taken to be of the form

$$\sigma^i_{\text{B}} = \frac{a_{1i}(\omega - \omega_0) + a_{2i}(\omega - \omega_0)^{1/2}}{c_i + \omega - \omega_0}, \tag{18}$$

which has the correct Regge asymptotic behavior. Here again i stays for p or n. The values a_{1i} and a_{2i} were fixed by matching with the results of the high-energy data [19]. They are different for protons and neutrons, specifically: $a_{1\text{p}} = 91$ μb/GeV, $a_{2\text{p}} = 71$ μb/GeV, $a_{1\text{n}} = 87$ μb/GeV, and $a_{2\text{n}} = 65$ μb/GeV. The value

Table 1
The resonance masses and widths determined from our fit (pn) to the proton- and neutron-photoabsorption data. The experimental data [19] and the results obtained in Ref. [7] from the fit to the proton (p) and neutron (n) data are also given

State	M(MeV)				Γ(MeV)			
	Exp	p	n	pn	Exp	p	n	pn
P_{11}	≈ 1440 1430–1470	1425	1440	1430	≈ 350 250–450	90	110	115
D_{13}	≈ 1520 1515–1530	1515	1515	1505	≈ 120 110–135	135	150	140
S_{11}	≈ 1535 1520–1555	1595	1570	1560	≈ 150 100–250	50	100	255
D_{15}	≈ 1675 1670–1685			1680	≈ 150 140–180			145
F_{15}	≈ 1680 1675–1690	1680	1695	1670	≈ 130 120–140	130	185	115
P_{33}	≈ 1232 1230–1234	1235	1250	1230	≈ 120 115–125	130	185	120
D_{33}	≈ 1700 1650–1750			1710	≈ 300 200–400			220
F_{37}	≈ 1950 1940–1960	1950	1960	1960	≈ 300 290–350	425	420	230

ω_0 corresponds to the pion threshold. The values c_i were free parameters in the fit: their values were found to be $c_p = 0.42$ GeV and $c_n = 0.47$ GeV.

We used the neutron cross sections extracted by Armstrong et al. [7,8] from the proton and deuteron data, after having corrected for the Fermi-motion effect.

In Table 1 we give the masses and widths of the resonances determined from the simultaneous fit to the proton and neutron data [7,8] and compare these values with those given in the Particle Data Book [19] and by Armstrong et al. [7,8].

The proton- and neutron-photoabsorption cross sections are shown in Figs. 7a and 7b. Three resonance structures are clearly seen in the energy dependence of the proton data: they are mainly due to the contributions of the $P_{33}(1232)$, $D_{13}(1520)$, and $F_{15}(1680)$ resonances. The strengths of the first two resonances in the neutron channel are almost the same as in the proton one, while the coupling of the third resonance to the neutron channel is weaker. In the figures there are also shown the fitting curves (solid-line curves) and the contribution of the different resonances and the background: the variance of the fit was found to be 1.9σ. In Refs. [7,8] the same data were fitted separately with the variance 1.4σ for the proton and 1.9σ for the neutron using more parameters to describe the background. In that case, the masses and widths of some resonances resulted different in the fits to the proton and deuteron data. Moreover, they were also different from the Particle Data Book values [19]. On the contrary, the values obtained from our fit are in good agreement with the parameters of the resonances given in the Particle Data Book.

In Table 2 we present the photoabsorption cross sections on the proton and neutron, $\sigma_{\text{res}}^{\text{p}}$ and $\sigma_{\text{res}}^{\text{n}}$ listed together with the experimental data [19]. In the case of the proton the comparison is also made with Ref. [7].

3.2. The description of the uranium data

Then, we fitted the ^{238}U photofission data [1] using the same resonance parameters obtained from the fit to the proton and neutron data. We used the following expression for the forward Compton-scattering amplitude:

$$\bar{\mathscr{T}}(\omega, \vartheta = 0) = \int \frac{\mathrm{d}\boldsymbol{p}}{(2\pi)^3} |\vartheta(\boldsymbol{p})|^2 g_{\gamma\mathrm{NN}^*}^2 G, \tag{19}$$

where $G = (G_0^{-1} - W_{\text{av}})^{-1} = [(p+k)^2 - (M + \delta M)^2 + iM\Gamma]^{-1}$, and Γ is given by Eq. (16).

As said above, in our approach the effects of the nuclear medium for each resonance were described by the two parameters δM and Γ^*. To decrease the total number of parameters we put $\delta M = 0$ for all the resonances except for the Δ(1232). In fact, in the nuclear medium the Δ is much narrower than the other resonances and even a small shift of its mass contributes very significantly to the χ^2 value. This is due to the very high strength of the Δ-resonance whose cross section is about 450 μb. The other resonances are much more broadened by the Fermi motion, while the Pauli-blocking effect is less important. Therefore, their effective widths in the nuclear medium are much larger than the one for the Δ, and comparatively small shifts of their masses affect the fit negligibly. It is not surprising that the imaginary parts of the mass operator of the resonances receive large shifts, while the real parts do not. In the framework of the pseudooptical potential, this can be interpreted to mean that the real part of the N*N forward-

Table 2
The photoabsorption cross sections on the proton and neutron as determined from our fit (pn), together with the experimental values [19] and the proton cross sections from Ref. [7]: In the last four columns the velocity of the resonances in the nucleus, v, the Pauli blocking factor, B_{F}, the collision broadening, Γ^*, and the product of the interaction cross sections, σ^*, and velocity of the resonances are given

State	$\sigma_{\text{res}}^{\text{p}}$ (μb)			$\sigma_{\text{res}}^{\text{n}}$ (μb)		v	B_{F}	Γ^* (MeV)	$\sigma^* v$ (mb)
	Exp	p	pn	Exp	pn				
$P_{11}(1440)$	10–15	50	15	2–10	10	0.40	0.87	150	40 [a]
$D_{13}(1520)$	95–130	160	140	75–115	120	0.45	0.90	315	82
$S_{11}(1535)$	10–20	30	10	15–35	0	0.46	0.90	160	40 [a]
$D_{15}(1675)$	~ 2		5	15–25	5	0.52	0.92	160	40 [a]
$F_{15}(1680)$	45–65	90	45	5–10	0	0.52	0.92	320	80
$P_{33}(1232)$	410–480	390	420	410–480	420	0.27	0.68	75	21
$D_{33}(1700)$	20–45		20	20–45	20	0.53	0.93	335	83
$F_{37}(1950)$	15–30	45	25	15–30	25	0.62	0.95	170	40 [a]

[a] The cross sections for these resonances were kept fixed in the fit (see text).

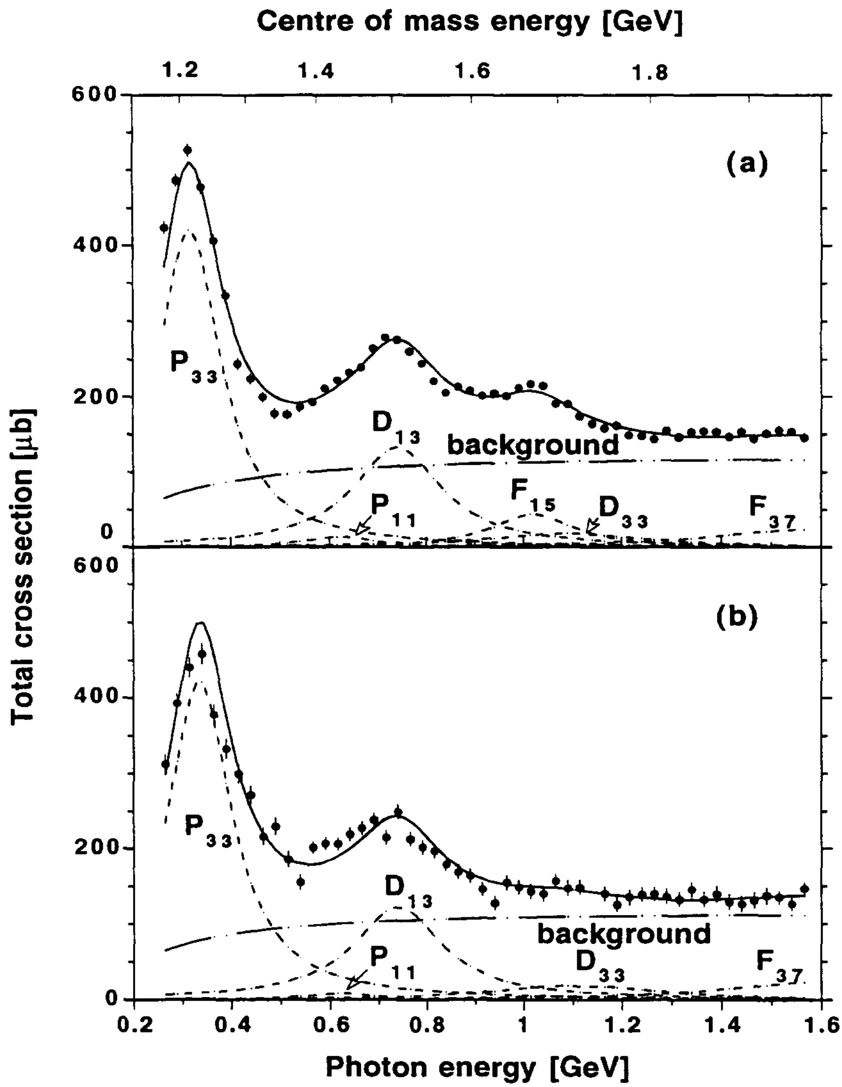

Fig. 7. Smooth fit to the proton (a) and neutron (b) data. The contribution of some resonances and of the background are also shown.

scattering amplitude is small compared to the imaginary part. This really happens in the strong-absorption regime when many different final states are open and the forward-scattering amplitude is almost imaginary. Moreover, we fixed the broadening Γ^* of the resonances $P_{11}(1440)$, $S_{11}(1535)$, $D_{15}(1675)$, and $F_{37}(1950)$ (assuming $\sigma^* v = 40$ mb) because their photon couplings are rather small and the nuclear-photoabsorption cross section is not sensitive to the changes of these parameters.

In conclusion, in the fit to the uranium data we had five free parameters, specifically: the values of Γ^* (or $\sigma^* v$) for the $P_{33}(1232)$, $D_{13}(1520)$, $F_{15}(1680)$, and

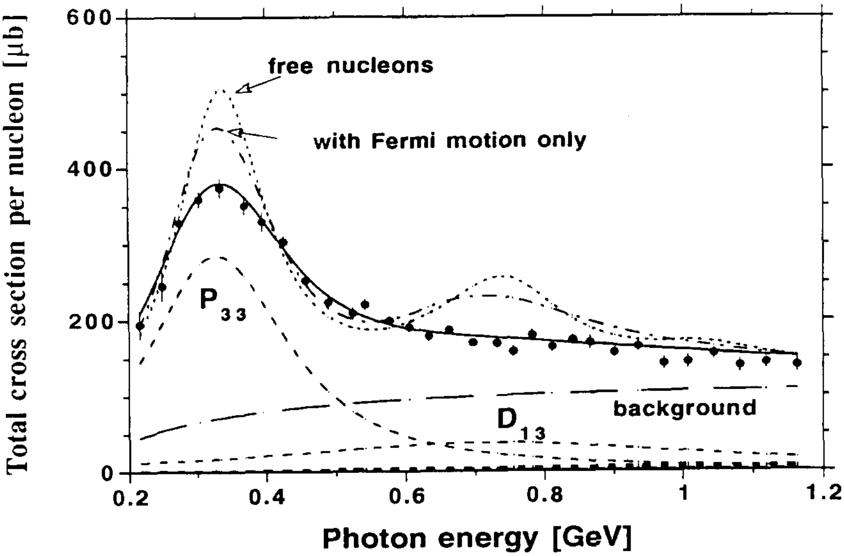

Fig. 8. Smooth fit to the Frascati data for uranium [1] (curve 3). The contribution of some resonances and of the background are also shown. The curve 1 shows the result which could be observed on free nucleons. The curve 2 shows the result which accounts for the Fermi motion only.

$D_{33}(1700)$ resonances and δM for the Δ. In the calculations, we took into account also the Pauli-blocking effect.

The results of the calculations are shown in Fig. 8 together with the uranium data measured at Frascati [1]. In the figure are given: (i) the results of calculations when the nucleus is considered as a system of free nucleons, the resonances $P_{33}(1232)$ and $D_{13}(1520)$ are clearly seen on this curve, while the resonance $F_{15}(1680)$ is much less pronounced due to its weak coupling to the neutron channel, (ii) The result of the calculation which takes into account only the Fermi-motion effect (for the momentum distribution of nucleons in the nucleus we used the Fermi gas model with $p_F = 280$ MeV/c), and (iii) the result of the calculation which accounts also for the nuclear-medium effect. The variance of the fit is 1.3σ; as is shown it reproduces the data satisfactory.

In Table 2 we give the values for Γ^* and $\sigma^* v$ obtained from the fit. The best description of the Δ-peak was found at $\delta M_\Delta = 15$ MeV and $\Gamma^*_\Delta = 75$ MeV (or $\sigma^* v \approx 21$ mb). This means that without the effect of the Fermi motion the effective width of the Δ in the nuclear medium would be about 155 MeV. For the resonances $D_{13}(1520)$, $F_{15}(1680)$, and $D_{33}(1700)$ we found $\Gamma^* = 315$, 320, and 335 MeV, respectively (which corresponds to $\sigma^* v \approx 80$–83 mb for all these resonances). As these values of $\sigma^* v$ seem to be too large we performed a second fit to the uranium-photoabsorption data assuming a 10% damping of the strengths of all these resonances apart from the Δ. (Possible reasons of such a damping are discussed in Section 3.3). In this case the best fit to the data was found for the same value $\sigma^* v = 21$ mb for the Δ, and $\sigma^* v = 65$, 42, and 53 mb for the resonances D_{13}, F_{15}, and D_{33}, respectively. The quality of this second fit was almost the same as for the first one (the variance of the fit is 1.34σ).

3.3. Comparison with other approaches

There are many papers devoted to the analysis of the photoabsorption data in the Δ-resonance region. The data on light and medium nuclei, like ^{4}He, ^{12}C, and ^{16}O, were analyzed in the framework of the Δ–hole model (see e.g. Refs. [11–13]). Using the effective ph and Δh interactions Carrasco and Oset [13] developed a model with a many-body expansion in the number of ph excitations in the nucleus which can be also applied for heavy nuclei. This model could describe the data available on ^{12}C, ^{16}O, and ^{208}Pb below 500 MeV.

It is not easy to make a direct comparison of our results for δM_Δ and Γ_Δ^* with the results of the Δ–hole model for ^{12}C and ^{16}O. Nevertheless, some comparison can be made. In Ref. [11], the parameters of the forward nuclear Compton amplitude for ^{12}C in the one-doorway approximation are presented. The expectation value of the imaginary part of the spreading potential cited in Table IV of Ref. [11] is equal to -25 MeV. This would correspond to $\Gamma_\Delta^* = 50$ MeV, which is somewhat smaller than our value for uranium. Several improvements to the Δ–hole model were made by Carrasco and Oset [13], who described well the old total photoabsorption data on ^{12}C, ^{16}O, and ^{208}Pb (see Figs. 45, 46 of Ref. [13]) in the Δ-region. The value of Γ_Δ^* in that paper is ≈ 60 MeV at $\rho = 0.75\rho_0$. If we assume a linear dependence of Γ_Δ^* on ρ, we shall find quite a reasonable agreement with our result.

As was already mentioned before, it is expected that our model works better for higher resonances than for the Δ. However, there are still several important questions which should be better clarified, specifically:

(i) How big can the damping of the resonance strength because of the quark-deconfinement effect be? Such a damping was predicted by Giannini et al. [20] in the framework of the non-relativistic quark model with an oscillator potential. According to this model the strength of the resonance in nuclei is damped because of the exchange of quarks between overlapping nucleons and the damping coefficient is different for different resonances, specifically the $P_{33}(1232)$ is not damped, while the $D_{13}(1520)$ and $F_{15}(1680)$ are damped by 11% and 23%, respectively. However, it is worth to notice that recently Akulinichev and L'vov [21] have also analyzed the suppression mechanism for the $D_{13}(1520)$ excitation due to overlap of nucleons in the framework of a non-relativistic quark model with a short-range repulsion and conclude that this effect is negligible.

(ii) Are there important corrections due to the shadowing? These corrections might be related to the coherent photoproduction of one and two pions. However, the cross sections of both processes are strongly suppressed for heavy nuclei by the nuclear form factor. As it was said before, the amplitude of the first reaction is equal to zero at forward angles. For the second reaction the forward amplitude is not zero when the pion pair is in the 1^- state. However, the 2π photoproduction cross section on the nucleon becomes significant when the invariant mass of the 2π

system is higher than 500 MeV. This means that the forward coherent photoproduction of 2π on heavy nuclei will be suppressed by the nuclear form factor depending on the minimal momentum transfer $q_{\min} = m_{2\pi}^2/2k_\gamma$, which is equal to 125 MeV at $k_\gamma = 1$ GeV/c.

4. Estimates of the N* N scattering cross sections

The inelastic cross section of the ΔN $\to$ NN interaction can be estimated using the data on the reaction pp $\to$ pnπ^+. The average energy which the Δ-resonance has after its formation in a nucleus corresponds to $\sqrt{s} = 2.19$ GeV or $p_{\text{lab}} = 1.3$ GeV/c. According to Shimizu et al. [22] at this energy the total cross section of the reaction pp $\to$ pnπ^+ is equal to 11.4 mb. It is known that at $\sqrt{s} = 2.32$ and 2.51 GeV the cross section of the reaction pp $\to$ pnπ^+ is dominated by the Δ^{++} production in the intermediate state (see Refs. [23–25]). Therefore it is reasonable to assume that the same dominance takes place at $\sqrt{s} = 2.19$ GeV. At this energy we are very close to the threshold energy $\sqrt{s} = m + M = 2.17$ GeV and can take into account only the ΔN S-wave. Then the cross section of the reaction pp $\to$ pnπ^+ can be written as

$$\frac{\mathrm{d}\sigma(\mathrm{pp} \to \mathrm{pn}\pi^+)}{\mathrm{d}\mu^2} = \frac{1}{64\pi^2 s}\left(\frac{p_2}{p_1}\right)_{\text{c.m.}} \overline{|A(\mathrm{pp} \to \mathrm{p}\Delta^{++})|^2} \times \frac{M_\Delta \Gamma_0}{(\mu^2 - M_\Delta^2) + (M_\Delta \Gamma_0)^2}, \tag{20}$$

where the invariant amplitude squared is summed up over the spin projections of the initial and final particles.

As the experimental distributions on the nπ^+ invariant mass are very close to the Breit–Wigner form (see Refs. [23–25]) we can conclude that the invariant amplitude squared,

$$\overline{|A(\mathrm{pp} \to \mathrm{p}\Delta^{++})|^2} = \overline{|A(\mathrm{p}\Delta^{++} \to \mathrm{pp})|^2}, \tag{21}$$

does not depend on μ in the interval $M_\Delta - \Gamma_\Delta \leqslant \mu \leqslant M_\Delta + \Gamma_0$. In this case we can define the quantity.

$$v\sigma(\Delta^{++}\mathrm{p} \to \mathrm{pp}), \tag{22}$$

which is expected to be weakly dependent on the Δ-mass when it is varied inside the Breit–Wigner peak.

Then integrating Eq. (20) over μ^2 we can define $\overline{|A(pp \to p\Delta^{++})|^2}$ and calculate

$$v\sigma(\Delta^{++}\mathrm{p} \to \mathrm{pp}) = 34 \text{ mb}.$$

Using the isotopic invariance (isovector t-channel exchange) we obtain for an average nucleon

$$v\sigma(\Delta^{+}\mathrm{N} \to \mathrm{pN}) = v\sigma(\Delta^{0}\mathrm{N} \to \mathrm{nN}) = 16.5 \text{ mb} \quad (v = 0.266).$$

If we assume that the ratio of the cross sections $\sigma(\Delta N \to N\Delta)$ and $\sigma(\Delta N \to NN)$ can be estimated using the one-pion exchange model, we shall find for the sum of the elastic and charge exchange cross sections the following value: $\upsilon\sigma(\Delta^+ N \to N\Delta) = 12$ mb, at $\sqrt{s} = 2.19$ GeV. Then the estimation for the total cross section of the Δ^+N (or Δ^0N) interaction will be

$$\upsilon\sigma(\Delta^+ N \to pN) + \upsilon\sigma(\Delta^+ N \to N\Delta) = 28.5 \text{ mb},$$

which is in reasonable agreement (within 30%) with the value given in Table 2.

The upper limit of $\upsilon\sigma^*$ for the resonances $D_{13}(1520)$ and $F_{15}(1680)$ can be obtained using the upper limit $\sigma_{max} \approx 20$ mb for the cross sections of the reactions $NN \to ND_{13}$ and $NN \to NF_{15}$ at the energies corresponding to their photoproduction. We get, from the inverse reactions,

$$\upsilon\sigma(ND_{13} \to NN) \leqslant 40 \text{ mb}, \quad \text{at} \quad \sqrt{s} = 2.53 \text{ GeV},$$
$$\upsilon\sigma(NF_{15} \to NN) \leqslant 26 \text{ mb}, \quad \text{at} \quad \sqrt{s} = 2.72 \text{ GeV}.$$

For the scattering on a black sphere the elastic cross section should be the same as the inelastic one, so the upper limits for the total cross sections can be estimated as

$$\upsilon\sigma_{tot}(ND_{13}) \leqslant 80 \text{ mb}, \quad \text{at} \quad \sqrt{s} = 2.53 \text{ GeV},$$
$$\upsilon\sigma_{tot}(NF_{15}) \leqslant 50 \text{ mb}, \quad \text{at} \quad \sqrt{s} = 2.71 \text{ GeV}.$$

In view of the existing uncertainties for the damping factor, the values $\upsilon\sigma^* =$ 68–82 mb for D_{13} and 42–80 mb for F_{15} derived from the fits do not contradict these limits.

5. Conclusions

We have discussed the implications of the recent Frascati photoabsorption data on nuclei [1–4] in the region of the resonances $D_{13}(1520)$ and $F_{15}(1680)$. We have shown that the main effects distorting the resonance behavior in the nuclear medium are the Fermi motion and the collision broadening. Moreover, we have shown that the Fermi motion alone cannot explain the suppression of the resonances and especially of the $D_{13}(1520)$ one. The collision broadening resulted in the most important effect, although about 10% damping of the resonance strength in the nuclei cannot be excluded.

We developed a reasonable but simple model for the description of the collision broadening; taking into account that the average momenta of those resonances are about 0.76 and 1.03 GeV/c we could describe their propagation in the nucleus using the eikonal approximation and introduce the optical pseudopotential for N*–nuclear interaction depending only on the average density of the nucleus and the forward N*N scattering amplitude, which is predominantly imaginary. In this case the collision broadening depends only on one parameter: the total N*N scattering cross section.

By fitting the experimental photonuclear data we found the N^*N total cross sections for the $P_{33}(1232)$, $D_{13}(1520)$, $F_{15}(1680)$, and $D_{33}(1700)$ resonances, which contribute substantially to the total γN cross section below 1.2 GeV. We found the following values:

$$\begin{aligned} &\upsilon\sigma_{tot} = 21 \text{ mb for the } P_{33}(1232) \text{ resonance,} && \text{at } p_{lab} = 0.34 \text{ GeV}/c, \\ &\upsilon\sigma_{tot} = 65\text{–}82 \text{ mb for the } D_{13}(1520) \text{ resonance,} && \text{at } p_{lab} = 0.76 \text{ GeV}/c, \\ &\upsilon\sigma_{tot} = 42\text{–}80 \text{ mb for the } F_{15}(1680) \text{ resonance,} && \text{at } p_{lab} = 1.03 \text{ GeV}/c, \\ &\upsilon\sigma_{tot} = 53\text{–}83 \text{ mb for the } D_{33}(1700) \text{ resonance,} && \text{at } p_{lab} = 1.07 \text{ GeV}/c. \end{aligned}$$

It is worth to notice that all the found N^*N cross sections do not violate the unitarity limits and in the case of ΔN interaction are also consistent with the experimental data on the reaction $NN \to \Delta N$.

For a better understanding of the situation one needs experimental data on different partitionings of the total photoabsorption cross section and over a broader energy region. It would also be important to investigate the suppression of the different resonances which have smaller widths (meson, baryon, hyperon, etc.) in the nuclear medium in experiments of the production type.

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