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COMPTON SCATTERING ON ^{208}Pb

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A B S T R A C T

In this paper we briefly review the formalism of the nuclear Compton scattering in the frame of the low-energy theorems. We treat the resonant terms of the amplitude, having corrective intermediate nuclear states, as a superposition of Lorentz lines with energy, width and strength fixed by the photo-absorption experiments. The gauge terms are evaluated starting from a simple, but realistic, nuclear Hamiltonian. Dynamical nucleon-nucleon correlations are consistently taken into account, beyond those imposed by the Pauli principle. The theoretical predictions are compared with the data of elastic diffusion of photons from ^{208}Pb.

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1. - INTRODUCTION

The experiment of elastic scattering of photons on ^{208}Pb by Ziegler et al. [1], and the proposal of using the synchrotron radiation of LEP [2] for the same purpose, advocate further studies of the photonuclear elastic (Compton) and inelastic (Raman) scattering. Accordingly, in this paper we further pursue the analysis of Ref. 1), paying special attention to the rôle played in the scattering matrix by the gauge terms, which are treated in the frame of the low-energy theorems (LET) as first suggested by Christillin and Rosa-Clot [3].

These terms not only restore the gauge invariance of the scattering amplitude, at least in the limit of vanishing energy, but are also essential, as shown in Ref. 1), for accounting for the experiment. Furthermore, with a gauge invariant theory it becomes possible, at least in principle, to extract from the photon-scattering data :

a) the enhancement factor κ_D of the dipole sum rule (till now measured only through photo-absorption experiments) ; this possibility stems from the condition imposed by the Thomson limit on the resonant and gauge terms of the scattering amplitude ;

b) the nuclear charge distribution Fourier transform, since the photon waves are not distorted by the nuclear Coulomb field ;

c) the spatial distribution of the mesonic matter within the nuclear medium.

All these items are considered in the present paper, using a non-relativistic nuclear Hamiltonian to ensure gauge invariance. In Section 2 we discuss the formalism of the photon-nucleus scattering and in Section 3 we outline our phenomenological model for the amplitude, which is compared with the experiment in Section 4. In the concluding Section we discuss how to improve upon the present treatment.

2. - THE PHOTON-NUCLEUS SCATTERING MATRIX

Although this topic has been dealt with a number of times in the literature [2)-5)], we briefly review the pertinent formalism here for sake of completeness.

In lowest order perturbation theory the photon-nucleus scattering matrix reads

$$S^{(2)}_{fi} = -\frac{2\pi}{\Omega\hbar\sqrt{\omega\omega'}}\epsilon_\mu\epsilon'_\nu\iint d^4x\, d^4y\, e^{ik\cdot x}e^{-ik'\cdot y}\langle f|T[j^\mu(x)j^\nu(y)]|i\rangle \tag{2.1}$$

where $x_\mu \equiv (\vec{x}, ct)$, $k_\mu \equiv (\vec{k}, \omega/c)$ and $k'_\nu \equiv (\vec{k}', \omega'/c)$ are the four momenta of the initial and the final photons, with polarization ϵ_μ and ϵ'_ν respectively, T is the time-ordered product and Ω the normalization volume. The fully conserved four-vector nuclear current $j_\mu(\vec{r},t) = (\vec{J}/c, \rho)$ obeys the continuity equation

$$\frac{\partial}{\partial x_\mu} j_\mu(\vec{r},t) = 0. \tag{2.2}$$

With standard techniques (2.1) can be rewritten as follows

$$S_{fi}^{(2)} = -\frac{2\pi\hbar c^2}{\Omega\sqrt{\omega\omega'}} \, 2\pi i \, \delta(E_o + \hbar\omega - E_f - \hbar\omega') \epsilon_\mu \epsilon'_\nu \, T^{\mu\nu} \tag{2.3a}$$

with

$$T^{\mu\nu} = \int d\vec{x} \, d\vec{y} \; e^{i\vec{k}\cdot\vec{x}} \, e^{-i\vec{k}'\cdot\vec{y}}$$

$$\sum_n \left\{ \frac{\langle f | j^\nu(0,\vec{y}) | n \rangle \langle n | j^\mu(0,\vec{x}) | i \rangle}{E_o - E_n + \hbar\omega} + \frac{\langle f | j^\mu(0,\vec{x}) | n \rangle \langle n | j^\nu(0,\vec{y}) | i \rangle}{E_o - E_n - \hbar\omega'} \right\} \tag{2.3b}$$

where the energies E_n of the intermediate nuclear excited states $|n\rangle$ appear explicitly in the denominator. Note that $T^{\mu\nu}$ has the dimensions of a length.

As clearly shown in Ref. 3), the perturbative scattering matrix (2.1) is <u>not</u> gauge invariant ; indeed with some algebra one gets

$$k_\mu k'_\nu \iint d^4x \, d^4y \; e^{-ik\cdot x} \, e^{ik'\cdot y} \langle f | T[j^\mu(x) j^\nu(y)] | i \rangle$$

$$= \frac{i}{\hbar} \iint d^4x \, d^4y \; e^{-ik\cdot x} \, e^{ik'\cdot y} \, \delta(x_o - y_o) \langle f | [\rho(y), [\hat{H}, \rho(x)]] | i \rangle \neq 0. \tag{2.4}$$

The above expression allows the determination of the longitudinal components of the gauge terms through the equation

$$k_\mu k'_\nu \left(T^{\mu\nu} + S^{\mu\nu} \right) = 0 \tag{2.5}$$

which works, "stricto sensu", only in the limit $\omega \to 0$, $\omega' \to 0$ (LET).

For the following non-relativistic nuclear Hamiltonian with two-body forces having an attractive Majorana component

$$\hat{H} = \sum_{i=1}^{A} \frac{p_i^2}{2m} + \sum_{i<j}^{A} {}_1 \left[V_W(\vec{r}_i - \vec{r}_j) + V_M(\vec{r}_i - \vec{r}_j) P_{ij} \right] \tag{2.6}$$

(P_{ij} is the space-exchange operator) one gets from the kinetic energy

$$S_{\mu\nu}^{kin} = -\delta_{\mu\nu} r_o \langle 0 | \sum_{i=1}^{A} \left[\frac{1}{2} + t_z(i)\right] e^{-i(\vec{k}-\vec{k}')\cdot\vec{r}_i} | 0 \rangle = -\delta_{\mu\nu} r_o Z F_{el}(\vec{q}) \tag{2.7}$$

and from the Majorana force

$$S_{\mu\nu}^{exch} = r_o \frac{mc^2}{(\hbar c)^2} \langle 0 | \sum_{i<j=1}^{A} \left[t_z(i)-t_z(j)\right]^2 (r_i-r_j)_\mu (r_i-r_j)_\nu \, e^{-i\vec{q}\cdot(\vec{r}_i+\vec{r}_j)/2} \times j_o\left\{\frac{\vec{k}}{2}\cdot(\vec{r}_i-\vec{r}_j)\right\} j_o\left\{\frac{\vec{k}'}{2}\cdot(\vec{r}_i-\vec{r}_j)\right\} V_M(\vec{r}_i-\vec{r}_j) P_{ij} | 0 \rangle . \tag{2.8}$$

In the above, $r_o = e^2/mc^2$, $\vec{q} = \vec{k}-\vec{k}'$, $F_{e\ell}(\vec{q})$ is the nuclear charge form factor and j_o the zeroth-order spherical Bessel functions.

Unlike (2.7), (2.8) does <u>not</u> depend only on $\vec{q}$, but on $\vec{k}$ and $\vec{k}'$ separately and is not diagonal in the tensor indices μ and ν. However,

$$(\vec{\epsilon}\cdot\vec{r}_{ij})(\vec{\epsilon}'\cdot\vec{r}_{ij}) = \vec{\epsilon}\cdot\vec{\epsilon}' r_{ij}^2 - (\vec{\epsilon}\times\vec{r}_{ij})\cdot(\vec{\epsilon}'\times\vec{r}_{ij}) = \vec{\epsilon}\cdot\vec{\epsilon}' \frac{r_{ij}^2}{3} \tag{2.9}$$

for a spherical nucleus *). Furthermore, for photon energies not exceeding the pionic mass, the product of the two j_o is very close to unity since the range of $|\vec{r}_i-\vec{r}_j|$ is limited by the NN interaction, whereas the centre-of-mass co-ordinate varies all over the nucleus. Thus, for a spherical system, (2.8) simplifies to

$$S_{\mu\nu}^{exch} = \delta_{\mu\nu} \frac{r_o}{3} \frac{mc^2}{(\hbar c)^2} \langle 0 | \sum_{i<j}^{A} \left[t_z(i)-t_z(j)\right]^2 (\vec{r}_i-\vec{r}_j)^2 e^{-i\vec{q}\cdot(\vec{r}_i+\vec{r}_j)/2} \times V_M(\vec{r}_i-\vec{r}_j) P_{ij} | 0 \rangle \tag{2.10}$$

which, like $S_{\mu\nu}^{kin}$, depends only on the momentum transfer $\vec{q}$. It is worth noticing that whereas (2.3a) is in general a complex quantity, (2.7) and (2.10) are real, as they correspond to the equal time contribution to the scattering matrix (2.1).

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*) We note that the factor $\frac{1}{3}$ is not included in the formulae of Ref. 3), which may lead to some confusion. After completing this paper, Christillin pointed out to us that he had corrected this factor in Ref. 6).

3. - THE MODEL

According to the discussion of the previous section we split the photo-nucleus scattering amplitude into a resonant and a gauge term

$$\mathcal{F}_{fi}^{total}(E,\theta) = T_{fi}^{res}(E,\theta) + S_{fi}^{gauge}(E,\theta) \tag{3.1}$$

For simplicity we avoid a microscopic calculation and, following Ziegler [1], we assume the resonant term in the forward direction to be well represented by a superposition of Lorentz lines

$$T_{fi}^{res}(E,\theta=0) = r_0 \frac{NZ}{A} \sum_{n,L} \left\{ D_{n,L} E^2 \frac{(E_{n,L}^2 - E^2) + i\Gamma_{n,L} E}{(E_{n,L}^2 - E^2)^2 + \Gamma_{n,L}^2 E^2} + \frac{D_{n,L}\,\delta_{L,1}}{[1+(q/\mu_\pi)^2]^2} \right\} \tag{3.2}$$

with the energy $E_{n,L}$, the width $\Gamma_{n,L}$ and the strength $D_{n,L}$ of each resonance to be fixed by the analysis of the absorption experiments. In (3.2) the index L runs over the multipolarities, whereas n runs over the nuclear resonances of a given multipole character.

The requirements of causality and crossing symmetry are met by the scattering amplitude (3.1). Also worth noticing is that for $E \to 0$ only the second term in the right-hand side of (3.2) survives : it arises partly from intermediate states associated with the nuclear centre-of-mass motion ("dipole spurious state"). Unlike in Ref. 1), where this term is kept constant, we allow for a phenomenological q dependence of it which, at finite scattering angles and high photon energies, should account for the non-vanishing contribution of the mesonic exchange currents (see graphs of Fig. 1). This point will be further discussed later on.

Finally we remind ourselves that from (3.2), via the optical theorem, it follows easily

$$\int_0^\infty \sigma_a(E)\,dE = 2\pi^2\hbar c\, r_0 \frac{NZ}{A} \sum_{n,L} D_{n,L} \tag{3.3}$$

$\sigma_a(E)$ being the "total" photo-absorption cross-section. Thus the amount of $\sum_{n,L} D_{n,L}$ exceeding the unity is directly related to the contribution of the exchange forces to the process.

Let us now consider the gauge terms. We utilize the charge distribution experimentally obtained, in a wide range of momenta, from the elastic scattering of electrons on ^{208}Pb [7]. For practical calculations it is convenient to reproduce it with a Fermi distribution

$$\varrho(\vec{r}) = \frac{\varrho_0}{1+\exp\{(r-\mathcal{R}_p)/c_p\}} \tag{3.4}$$

where $\mathcal{R}_p = 6.49$ fm and $c_p = 0.54$ fm are the radius and surface thickness of the protons' charge distribution. The corresponding elastic form factor turns out to be (almost exactly)

$$F_{e\ell}(q) = \varrho_0 \frac{4\pi}{3}\mathcal{R}_p^3 \frac{\pi c_p q}{\sinh(\pi c_p q)} \frac{3}{q\mathcal{R}_p}\left\{\frac{\sin(q\mathcal{R}_p)}{(q\mathcal{R}_p)^2}\frac{\pi c_p q}{\mathrm{tgh}(\pi c_p q)} - \frac{\cos(q\mathcal{R}_p)}{q\mathcal{R}_p}\right\} \tag{3.5}$$

where, with the choice $F_{e\ell}(0) = 1$,

$$\varrho_0 = \left\{\frac{4\pi}{3}\mathcal{R}_p^3\left[1+\left(\frac{\pi c_p}{\mathcal{R}_p}\right)^2\right]\right\}^{-1} \tag{3.6}$$

Turning to $S_{\mu\nu}^{exch}$, note the identity

$$[t_z(i)-t_z(j)]^2 = \left[\frac{1}{2}+t_z(i)\right]\left[\frac{1}{2}-t_z(j)\right] + (i\leftrightarrow j) \tag{3.7}$$

which entails the vanishing of the exchange matrix element in (2.10) and explicitly displays the appearance of proton and neutron densities. To evaluate (2.10) one would need a proper ground state wave function of ^{208}Pb correlated by a realistic two-body force ; we limit ourselves to the Brueckner-Hartree-Fock ground state of infinite nuclear matter and to a double square well two-body potential

$$V(1,2) = \vartheta(|\vec{r}_1-\vec{r}_2|)(1-\alpha+\alpha P_{12}) \tag{3.8}$$

with

$$\vartheta(r) = \begin{cases} U_0\ , & \alpha = 0 \quad \text{for } 0 \leq r < a \\ -V_0\ , & \alpha = 1/2 \quad \text{for } a \leq r \leq b \\ 0 & \text{for } r > b\,. \end{cases}$$

We hope to account, at least partially, for the finite size effects with the phenomenological protons and neutrons Fermi distributions (3.4). Some algebra then leads to the following expression

$$S_{\mu\nu}^{exch} = \delta_{\mu\nu}\frac{r_0}{3}\frac{mc^2}{(\hbar c)^2}\iint d\vec{r}_1\, d\vec{r}_2\, e^{-i\vec{q}\cdot(\vec{r}_1+\vec{r}_2)/2}\, r_{12}^2\, \varrho_p(\vec{r}_1)\varrho_n(\vec{r}_2)F(r_{12})V_M(r_{12})$$

$$= -\delta_{\mu\nu}\, r_0\, Z\, F_{ex}(\vec{q}) \tag{3.9}$$

which defines the exchange form factor $F_{ex}(\vec{q})$. With the force (3.8), the above six-dimensional integral can be reduced to the following three-dimensional one

$$S_{\mu\nu}^{exch} = -\delta_{\mu\nu} \frac{r_o}{3} \frac{mc^2}{(\hbar c)^2} NZ \varrho_o^p \varrho_o^n V_o \frac{(2\pi)^2}{q} \int_0^\infty dR\, R \sin(qR) \int_a^b dr\, r^4 F(r)$$
$$\times \int_{-1}^{+1} dx \left\{1+\exp\left[\left(\sqrt{R^2+r^2/4+Rrx} - \mathcal{R}_p\right)/c_p\right]\right\}^{-1} \left\{1+\exp\left[\left(\sqrt{R^2+\frac{r^2}{4}-Rrx} - \mathcal{R}_n\right)/c_n\right]\right\}^{-1} \tag{3.10}$$

where $\mathcal{R}_p$, c_p, $\mathcal{R}_n$ and c_n are the radii and the surface thicknesses of the proton and neutron Fermi distributions, respectively. If the two-body wave function is correlated only through the Pauli principle, then

$$F(r) = \left[\frac{3}{k_F r} j_1(k_F r)\right]^2 \tag{3.11}$$

where k_F is the Fermi momentum and j_1 the spherical Bessel function of order one. Otherwise,

$$F(r) = \frac{24}{k_F^3} \sum_{\ell=0}^{\infty} (-1)^\ell (2\ell+1) \int_0^{k_F} dk\, k^2 \left[1-\frac{3}{2}\left(\frac{k}{k_F}\right)+\frac{1}{2}\left(\frac{k}{k_F}\right)^3\right] j_\ell(kr) \frac{u_\ell(k,r)}{kr} \tag{3.12}$$

where $u_\ell(k,r)$ is the radial two-body Bethe-Goldstone wave function for the two-body potential (3.8) [8].

The natural question now arises of whether our scattering amplitude (3.1) has the correct $E \to 0$ behaviour. Indeed, in this limit one has to recover the Thomson amplitude

$$\lim_{E\to 0} \mathcal{F}_{fi}(E,\theta=0) = r_o\left(\frac{NZ}{A} - Z\right) = -r_o \frac{Z^2}{A} \tag{3.13}$$

where the resonant and the gauge contributions are separated. In our case ($|\vec{k}| = |\vec{k}'| = E/\hbar c$)

$$\lim_{E\to 0} \mathcal{F}_{fi}(E,\theta=0) =$$
$$= r_o\left\{\frac{NZ}{A}\sum_{n,L} D_{n,L}\,\delta_{L,1} - Z F_{e\ell}(0) - Z F_{ex}(0)\right\} \tag{3.14}$$

Therefore since

$$\sum_n D_{n,1} = 1 + K_D \tag{3.15}$$

(3.13) is recovered from (3.14) if

$$\lim_{q \to 0} \frac{A}{N} F_{ex}(\vec{q}) = K_D \tag{3.16}$$

which can be regarded either as an independent method to evaluate the enhancement factor of the electric dipole sum rule or as a consistency test between the resonant and the gauge terms of the photon-nucleus scattering amplitude.

4. - RESULTS

As shown by Ziegler et al. [1], the imaginary part of the scattering amplitude (3.1) in the forward direction accounts satisfactorily for the photo-absorption experiment of Bergère et al. [9] on ^{208}Pb when the four dipole states and the two quadrupole states of Table 2 in Ref. 1) are considered. The first three dipole excitations correspond to the familiar giant dipole resonance and eventually to its first harmonic and the fourth one simulates with its large width the substantial photo-absorption above the dipole collective mode, which is generally interpreted in terms of a quasi-deuteron mechanism [10].

Beyond the isovector quadrupole resonance of ^{208}Pb we have also considered the one at 16 MeV, whose electric quadrupole character has not yet been established. In Fig. 2 the experimental total photo-absorption cross-section in ^{208}Pb is seen to be quite nicely reproduced with the nuclear resonances taken into account.

Let us now consider the scattering cross-section. Its expression

$$\frac{d\sigma}{d\Omega}(E,\theta) = \left|\mathcal{F}_{fi}^{E1}(E,\theta)\right|^2 \frac{1+\cos^2\theta}{2} + \left|\mathcal{F}_{fi}^{E2}(E)\right|^2 \frac{1-3\cos^2\theta+4\cos^4\theta}{2} + $$
$$+ 2\cos^3\theta \left\{ \mathrm{Re}\, \mathcal{F}_{fi}^{E1}(E,\theta)\, \mathrm{Re}\, \mathcal{F}_{fi}^{E2}(E) + \mathrm{Im}\, \mathcal{F}_{fi}^{E1}(E,\theta)\, \mathrm{Im}\, \mathcal{F}_{fi}^{E2}(E) \right\} \tag{4.1}$$

is easily obtained by combining the forward scattering amplitude (3.1) with standard angular factors [11]. Note that the gauge contribution is purely dipole, but it introduces an additional angular dependence through the momentum transfer $q = 2(\omega/c) \sin(\theta/2)$. This also happens for the non-Lorentzian component (4.2) of the resonant amplitude.

In calculating (4.1) we have tentatively used $\mathcal{R}_n = \mathcal{R}_p$ and $c_n = c_p$ for the radius and surface thickness of the neutron distribution, which is more uncertain than that of the proton.

Our, admittedly crude, two-body Majorana potential yields the correct binding energy, saturation density and compression modulus of infinite nuclear matter in a lowest order G matrix calculation [12] with the following values for the parameters : $V_o = 53$ MeV, $a = 0.57$ fm and $b = 2.1$ fm. The photo-absorption and scattering data require, however, a stronger Majorana attraction. Indeed, the correct Thomson limit is recovered with $V_o = 70$ MeV when only the Pauli correlations are taken into account.

The experimental differential cross-section $d\sigma/d\Omega$ at $\theta = 60^{\circ}$ and $\theta = 150^{\circ}$ and the ratio $d\sigma(150^{\circ})/d\sigma(60^{\circ})$ are displayed in Figs. 3, 4 and 5 together with the results of the present analysis, both with the Pauli and the dynamical correlation function. In the first case, the Thomson limit is exactly recovered and the data are seen to be nicely reproduced. On the other hand, the G matrix correlations included in $S^{exch}_{\mu\nu}$ are essentially induced by the attractive intermediate range nucleon-nucleon force. As a consequence their rôle is mostly felt at low energy, where they sensibly reduce the cross-section. In Fig. 6 the dynamical correlation function (3.12) is compared with the Pauli one.

We stress the crucial rôle of the q dependence in

$$g(q) = \frac{N}{A} \frac{\sum_n D_{n,1}}{[1+(q/\mu_\pi)^2]^2} \tag{4.2}$$

[see the last term in the right-hand side of (3.2)] in order to achieve agreement with the experiment. Also, in the high frequency regime, the relevance of (4.2) should be recognized. Indeed there our scattering amplitude (3.1), while in agreement with the causality requirement [13]

$$\mathcal{F}^{total}_{fi}(E\to\infty, \theta=0) = S^{gauge}_{fi}(E\to\infty, \theta=0) = -r_o Z\,[1+F_{ex}(0)] \tag{4.3}$$

in the forward direction, is not vanishing at finite θ. This reflects the contribution, for $E \to \infty$, of the contact pionic term [graph e) of Fig. 1] to the scattering amplitude.

In Fig. 7 we display the behaviour of the elastic and exchange form factors versus q, together with g(q). It is immediately apparent that a strong cancellation takes place between the gauge terms and g(q) itself. This feature was also present in the approach of Ziegler et al. [1].

One is thus forced to conclude that, although the Compton scattering might in principle yield precious information on κ_D, the proton/neutron distribution and the pionic one, namely

$$\rho_{exch}(r) = NZ\rho_0^p\rho_0^n \frac{mc^2}{(\hbar c)^2} V_0 \frac{2\pi}{3}\int_a^b dr' r'^4 F(r')$$
$$\times \int_{-1}^{+1} dx \left\{1+\exp\left[\left(\sqrt{r^2+r'^2/4+rr'x} - R_p\right)/c_p\right]\right\}^{-1} \times$$
$$\times \left\{1+\exp\left[\left(\sqrt{r^2+r'^2/4-rr'x} - R_n\right)/c_n\right]\right\}^{-1}, \qquad (4.4)$$

In practice, this could be difficult to achieve owing to the delicate interplay between the resonant and gauge terms of the amplitude. So we believe that any conclusion on a possible diffraction pattern associated with the pionic matter distribution should be postponed until a truly microscopic description of the various components in the scattering amplitude is available.

5. - CONCLUDING REMARKS

In this paper we have analyzed the Compton scattering data on ^{208}Pb utilizing the LET expressions for the gauge terms in the photon-nucleus scattering amplitude in addition to the model of Ziegler et al. for the resonant ones. A schematic nuclear Hamiltonian has been employed, with a Majorana space-exchange interaction. We have explored the rôle of the dynamical correlations among nucleons in the potential gauge term. They appear to play a minor rôle, except at very small frequencies, where they substantially reduce the cross-sections.

In order to achieve an agreement with the experiment, we have been forced to let the constant term in the resonant amplitude of Ziegler become q dependent, admittedly with an arbitrary functional form. However, we consider this procedure reasonable also because, in so doing, our high frequency scattering amplitude at finite angles is non-vanishing.

We also note that the use of LET, in the energy region we are interested in, is not warranted in principle. Indeed, Arenhövel [14] pointed out that a proper calculation should avoid LET, but rather attempt to compute directly the meson exchange current (MEC) contributions. This amounts to evaluating the five diagrams of Fig. 1, which he has shown [14] to yield, in the limit of vanishing frequency, the gauge exchange term (3.9) as given by LET, if the pion exchange is considered.

On the other hand, from Fig. 1 one infers that only diagram e) is associated with a contact scattering of the photon out of the pions, thus being the dominant one at high frequency and related to the $E \to \infty$ limit of $S_{\mu\nu}^{exch}$. However, in this energy regime our $F_{exch}(q)$ is non-vanishing only in the forward direction. Thus, as previously stated, we have been forced to phenomenologically account for the effects of the diagram e) at $\theta \neq 0$ with the function $g(q)$.

In conclusion, the Compton scattering is indeed a powerful tool for getting meaningful information on the nuclear collective response [15] and on meson-related nuclear phenomena, but a proper interpretation of their interplay can only be obtained with a microscopic theory for the whole scattering amplitude.

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FIGURE CAPTIONS

Figure 1 Meson exchange contributions (MEC) to the photon-nucleus scattering amplitude.

Figure 2 The experimental photo-absorption cross-section in ^{208}Pb, together with the predictions of the present model. The experimental points are taken from Ref. 7).

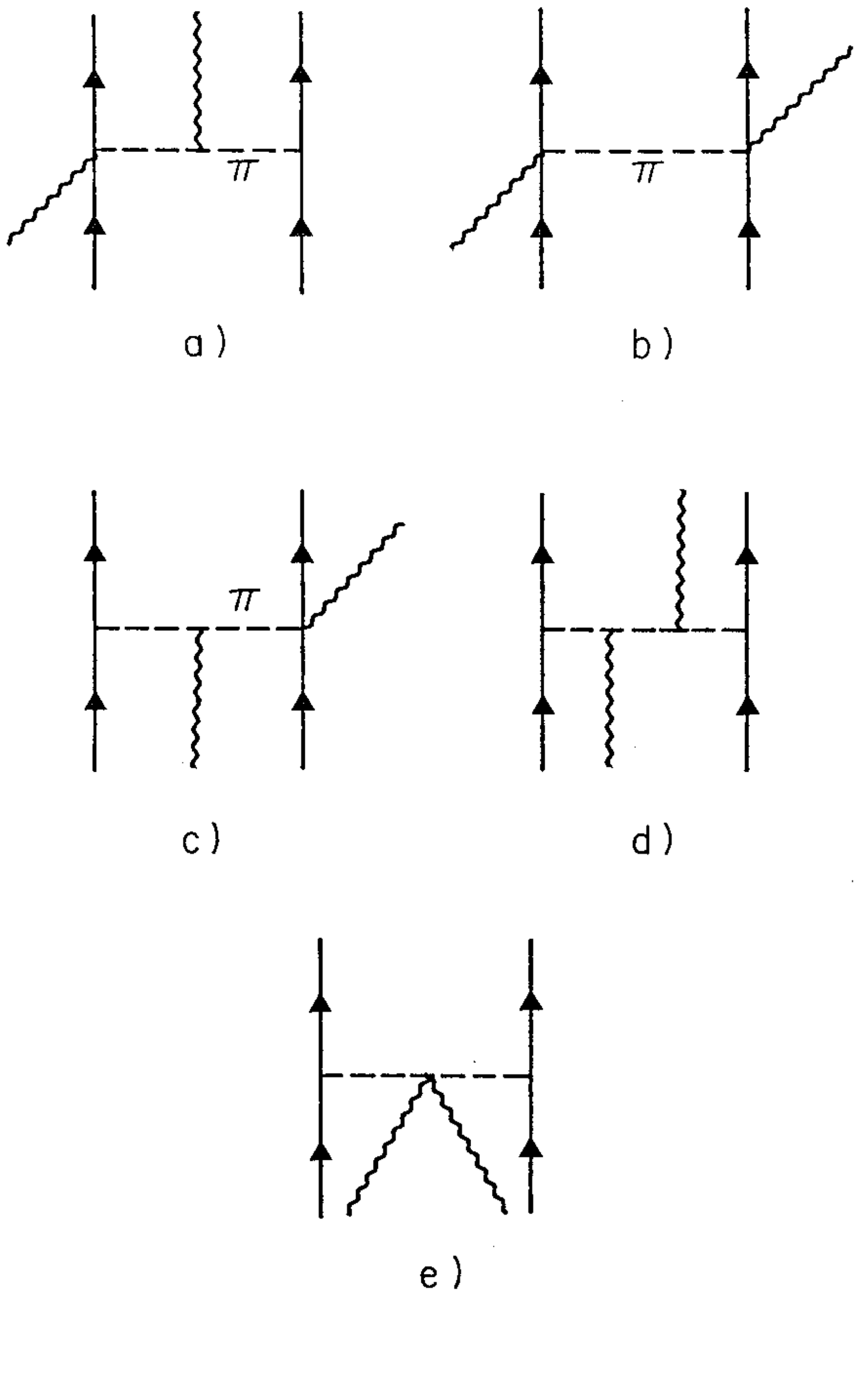
Figure 3 The Compton scattering cross-section on ^{208}Pb at $\theta = 60^{\circ}$. The continuous line is obtained with the inclusion of dynamical correlations into $S_{\mu\nu}^{exch}$, whereas the dashed line stems from the Pauli correlations only. The arrows indicate the exact Thomson limit. The experimental points are taken from Ref. 1).

Figure 4 The same as in Fig. 3, at $\theta = 150^{\circ}$.

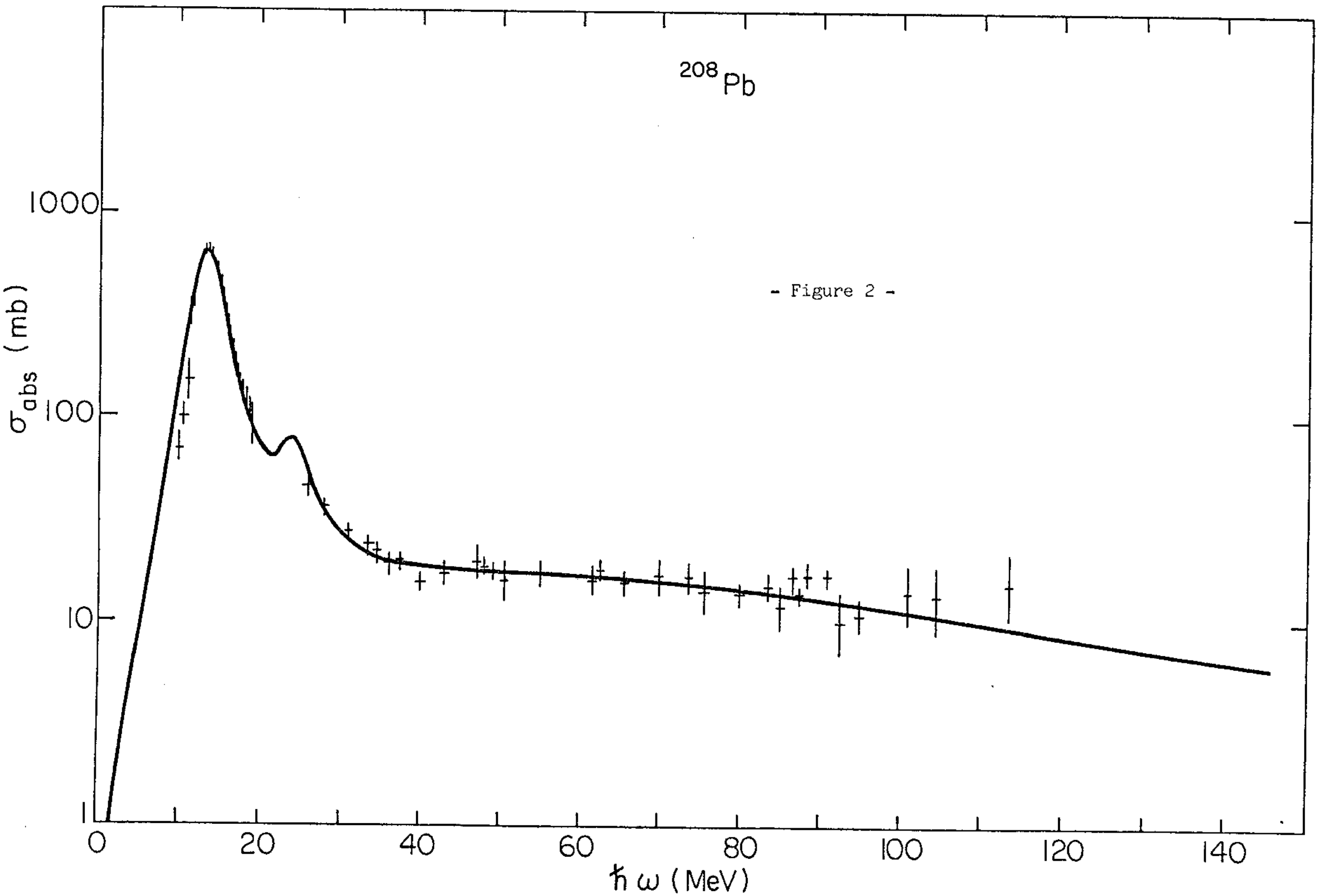
Figure 5 The ratio $d\sigma(150^{\circ})/d\sigma(60^{\circ})$. The continuous and dashed lines refer to the same situations as in Fig. 3.

Figure 6 The two-particle exchange correlation function $F(r)$ as given by (3.12) (continuous line) and (3.11) (dashed line).

Figure 7 The elastic form factor (3.5) (dot-dashed line) and the exchange form factor $F_{ex}(q)$ defined in (3.9) both with the $F(r)$ (3.11) (dashed line) and (3.12) (continuous line). The function $g(q)$, Eq. (4.2), is also shown (double-dot-dashed line). Note that $F_{e\ell}$ and F_{ex} enter into the scattering amplitude with opposite signs with respect to $g(q)$.

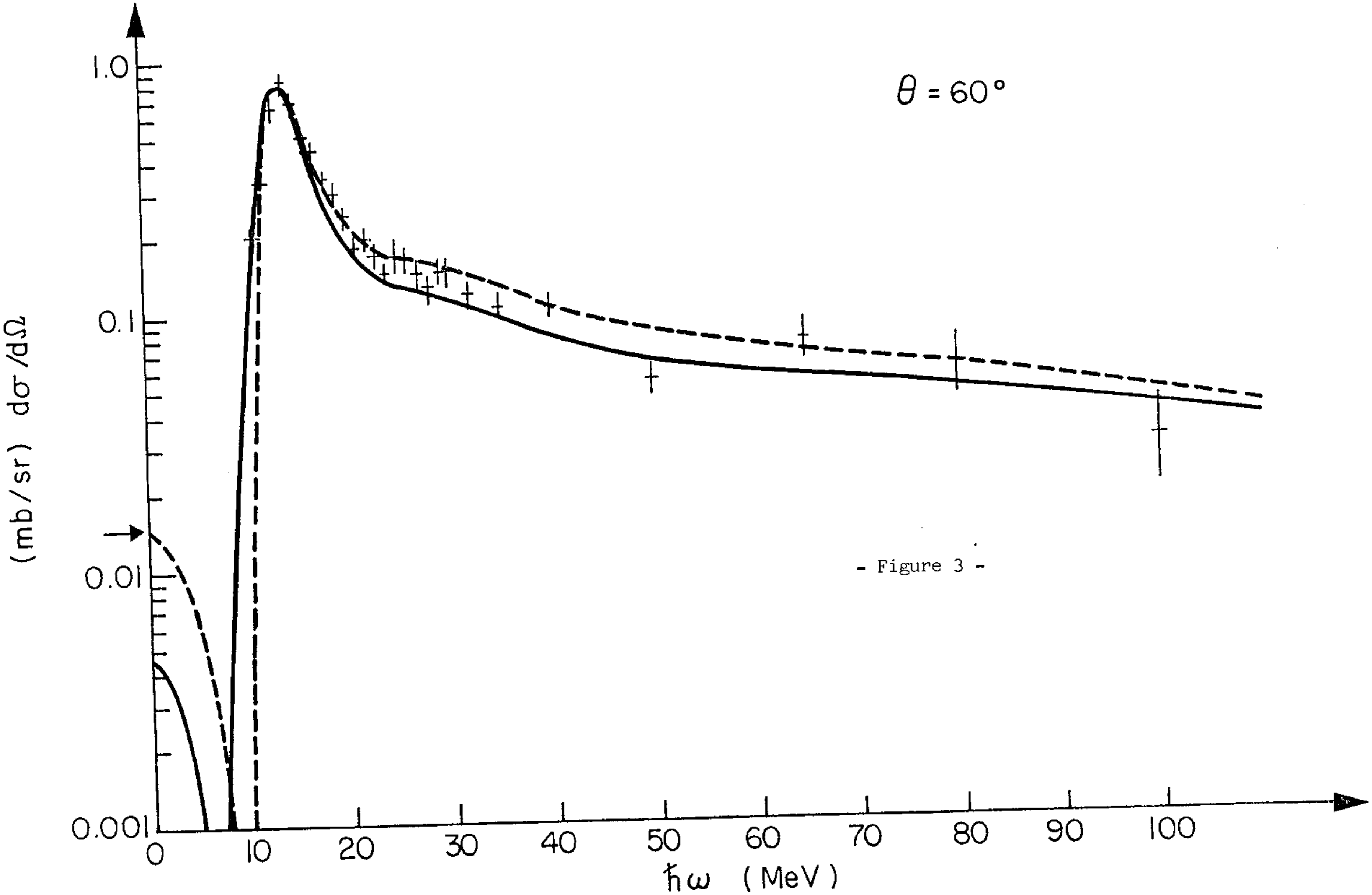


- Figure 1 -

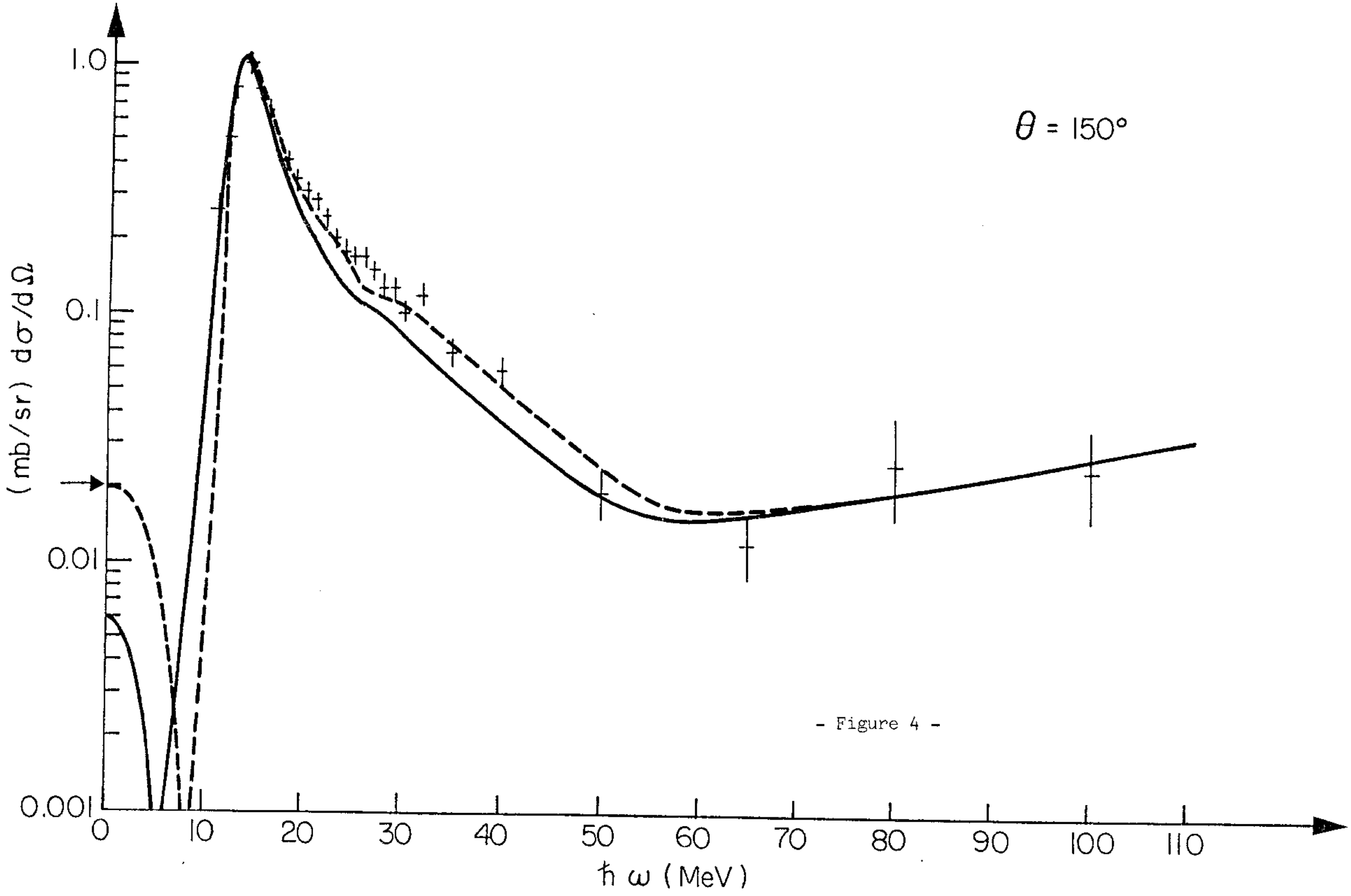


^{208}Pb
σ_{abs} (mb)
1000
100
10
1
0
20
40
60
80
100
120
140
$\hbar\omega$ (MeV)

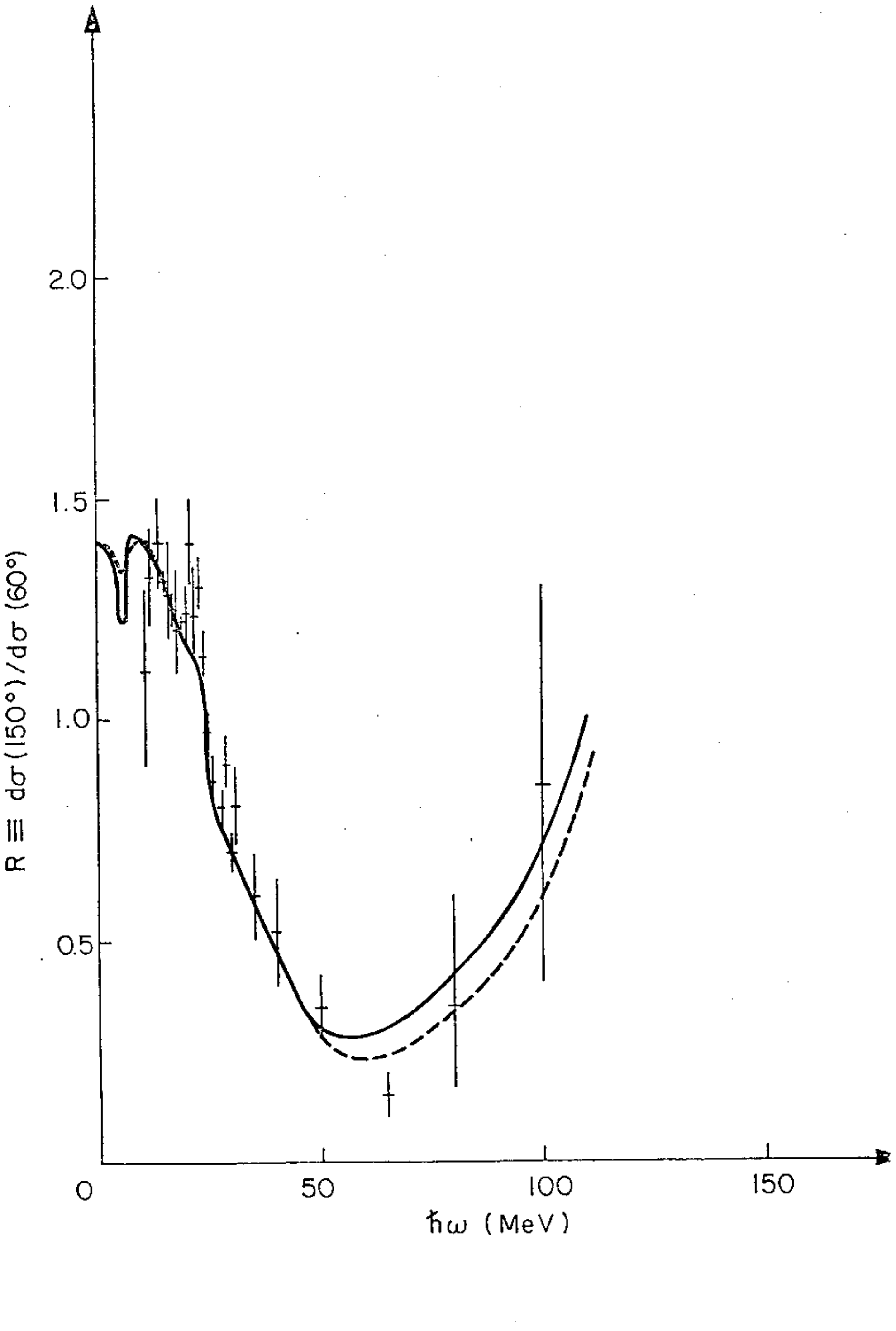
- Figure 2 -



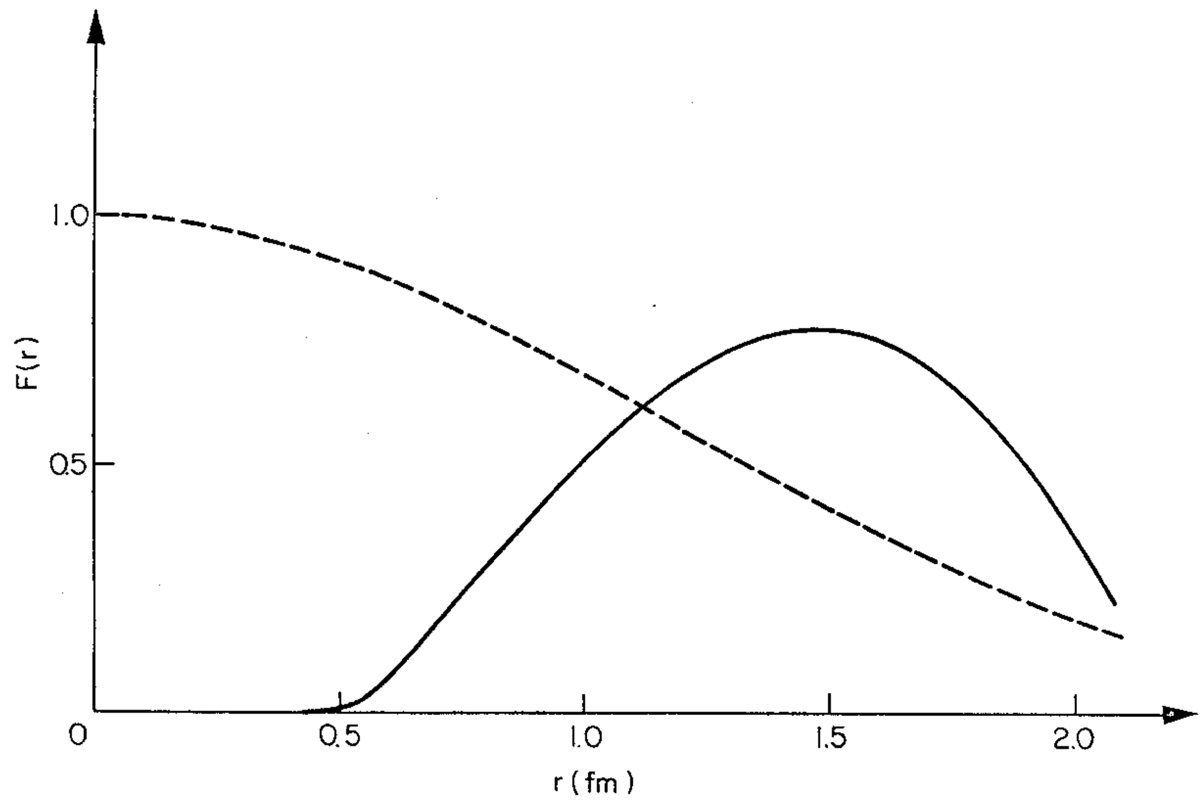
- Figure 3 -


θ = 150°
(mb/sr) dσ/dΩ
1.0
0.1
0.01
0.001
0
10
20
30
40
50
60
70
80
90
100
110
ħ ω (MeV)

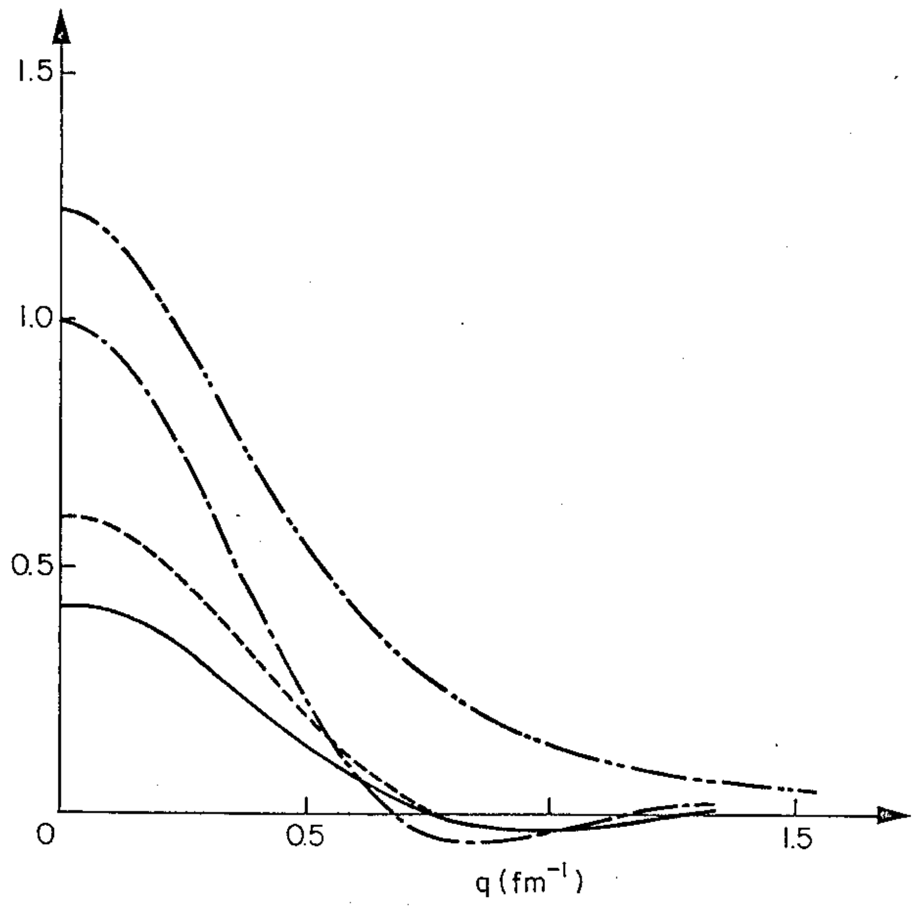
- Figure 4 -



- Figure 5 -



- Figure 6 -



- Figure 7 -