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Nucleon polarizabilities: Model study

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I study a simple model "nucleon," which can only absorb electric dipole radiation. This nucleon will scatter both electrically and magnetically, if its mass is of the same order of magnitude as its excitation energy. Therefore it has a static electric polarizability α as well as a susceptibility β, which are of similar size and of opposite sign. This contradicts naive expectations. The reasons for this unexpected behavior will be analyzed in some detail. Furthermore, numerical results are obtained within a constituent quark model. It turns out that the energy dependence of the electric polarizability and magnetic susceptibility is very weak for photon energies up to 100 MeV.

I. INTRODUCTION

Nucleon polarizabilities are intriguing quantities. The experimental finding that the static electric polarizability of the nucleon is considerably larger than the static magnetic susceptibility is surprising. *How can the nucleon, which is predominantly excited magnetically, respond stronger to a static electric field than a static magnetic field?* This is the question at the root of some recent discussion[1] and the motivation for the model calculations in this paper. If I convince the reader that the observed values for the nucleon polarizabilities are not overly surprising at all, then the main purpose of the present study is fulfilled. Additionally, another important question related to Compton scattering off nuclei and particles will be analyzed; namely, *how important are exchange effects?* It will be shown that terms, which have been proposed as a measure of exchange effects,[2] cancel out to *all* orders in the photon energy.

To this end I will work through a model for nucleon Compton scattering in an attempt to highlight as clearly as possible the essential dynamical features without undue formal complexity. The model considered is not necessarily "realistic" but instructive. The model ground state is a $S^{\pi}=\frac{1}{2}^{+}$ state, and it has only one $\frac{1}{2}^{-}$ excited state apart from the $N\bar{N}$ continuum. That means that such a system can only absorb electric dipole radiation below $N\bar{N}$ threshold. Despite this, I will show that such a system will scatter both electrically and magnetically; that is, it has a static electric polarizability and a magnetic susceptibility. The susceptibility is of similar size as the electric polarizability and of opposite sign. The polarizability and susceptibility are related to the total photoabsorption cross section via an energy-weighted sum rule. I will discuss this sum rule in the above-mentioned model. Numerical studies will be done in the framework of the constituent quark model for photon energies up to about 100 MeV. This is the energy region where modern experiments are done in order to determine the (static) electric polarizability and magnetic susceptibility.[3]

The calculations in this paper are closely related to previous work.[4,5] However, in these papers the nucleon is not discussed specifically. Furthermore, they are much more formal, so that the essential physics may be obscured.

II. SIMPLE MODEL FOR NUCLEON COMPTON SCATTERING

Let me start off by considering a "nucleon" ($\frac{1}{2}^{+}$) with mass ε_{+} and an excitation spectrum as depicted in Fig. 1. There is a discrete $\frac{1}{2}^{-}$ excited state with mass ε_{-} and a continuum corresponding to $N\bar{N}$ production. In this simple model I will study essential features of "nucleon" Compton scattering. In this paper I only consider photon energies k much less than the $N\bar{N}$ threshold. That means that photoabsorption only proceeds via the $\frac{1}{2}^{-}$ discrete state, and a nonrelativistic approach will be sufficient. Of course, the $N\bar{N}$ continuum will contribute to Compton scattering, as will be studied in detail below. The first step in our development will be to demonstrate that the model "nucleon," which only absorbs electric dipole radiation, will scatter both electrically and magnetically, and magnetic dipole scattering may be as important as electric scattering, depending on certain kinematical conditions.

As usual, we separate the Compton amplitude into a "resonance" term R and a "two-photon" or "seagull" amplitude B:

$$T_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})=R_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})+B_{\lambda'\lambda}(\mathbf{k}',\mathbf{k}) \ . \qquad (1)$$

The two-photon amplitude arises in a nonrelativistic theory to describe scattering via the $N\bar{N}$ continuum. Since the assumptions made above guarantee that $N\bar{N}$ will only be produced virtually, the two-photon amplitude is real. Scattering via the $\frac{1}{2}^{-}$ excited state is described by R in the present formulation. It is, however, important to include B since otherwise the model calculation would not be gauge invariant, and the low-energy theorem would be violated.

The resonance amplitude R is determined by the matrix elements of the current operator between the nucleon ground state and the $\frac{1}{2}^{-}$ and $\frac{1}{2}^{+}$ states as intermediate

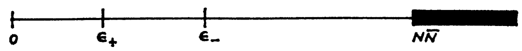

FIG. 1. The model nucleon spectrum.

states. The current operator of the (nonrelativistic) nucleon is given by

$$\mathbf{j}(\mathbf{x})=\tilde{\mathbf{j}}(\mathbf{x})+\frac{1}{2\varepsilon_+}\{\mathbf{P},\tilde{\rho}(\mathbf{x})\}\ , \tag{2}$$

which is correct to order $1/\varepsilon$. The tilde indicates that the operators are intrinsic operators in the nucleon rest frame. The first term is referred to as intrinsic current and the second as recoil current. The total center-of-mass momentum is denoted by $\mathbf{P}$.

The interaction of the intrinsic current with the electric dipole field $\mathbf{E}_\lambda$ and magnetic dipole field $\mathbf{M}_\lambda$ of a real photon with polarization λ and momentum $\mathbf{k}$ ($\boldsymbol{\varepsilon}_\lambda\cdot\mathbf{k}=0$) is given by

$$\int d^3x\ \tilde{\mathbf{j}}(\mathbf{x})\cdot\mathbf{E}_\lambda(\mathbf{x},k)=\frac{1}{k}[\tilde{H},\boldsymbol{\varepsilon}_\lambda\cdot\mathbf{C}(k)]+\boldsymbol{\varepsilon}_\lambda\cdot\mathbf{A}(k)\ , \tag{3a}$$

$$\int d^3x\ \tilde{\mathbf{j}}(\mathbf{x})\cdot\mathbf{M}_\lambda(\mathbf{x},k)=(\boldsymbol{\varepsilon}_\lambda\times\hat{\mathbf{k}})\cdot\mathbf{M}(k)\ , \tag{3b}$$

with

$$\mathbf{C}(k)=3i\int d^3x\ \tilde{\rho}(\mathbf{x})\hat{\mathbf{x}}j_1(kx)\ , \tag{4a}$$

$$\mathbf{A}(k)=\tfrac{9}{2}\int d^3x\ j_2(kx)[\tilde{\mathbf{j}}(\mathbf{x})\cdot\hat{\mathbf{x}}\hat{\mathbf{x}}-\tfrac{1}{3}\tilde{\mathbf{j}}(\mathbf{x})]\ , \tag{4b}$$

$$\mathbf{M}(k)=\frac{3i}{2}\int d^3x\ \hat{\mathbf{x}}\times\tilde{\mathbf{j}}(\mathbf{x})j_1(kx)\ . \tag{4c}$$

Note that $\mathbf{C}(k)$ only depends on the charge density $\tilde{\rho}$. This term arises through the use of the continuity equation

$$\nabla\cdot\tilde{\mathbf{j}}=-i[\tilde{H},\tilde{\rho}]\ , \tag{5}$$

with $\tilde{H}$ the intrinsic Hamiltonian. The separation of the electric multipole into a Coulomb and current multipole in (3a) is known as "Siegert's theorem."

I now briefly investigate the recoil current matrix elements. Naively, one would expect that a $\frac{1}{2}^+$ and a $\frac{1}{2}^-$ state can only be connected by an electric excitation. However, the recoil current generates a magnetic contribution, since a *moving* electric dipole generates a current and therefore a magnetic dipole component

$$\mathbf{M}_R(k)=\frac{3i}{4\varepsilon_+}\int d^3x\ \hat{\mathbf{x}}\times(\mathbf{P}_++\mathbf{P}_-)\times\langle\tfrac{1}{2}^+|\tilde{\rho}(\mathbf{x})j_1(kx)|\tfrac{1}{2}^-\rangle$$
$$=-\frac{1}{4\varepsilon_+}(\mathbf{P}_++\mathbf{P}_-)\times\langle\tfrac{1}{2}^+|\mathbf{C}(k)|\tfrac{1}{2}^-\rangle\ . \tag{6}$$

In single-photon processes such as photoabsorption or electron scattering one can always disregard this magnetic piece by choosing an appropriate reference frame. This is not possible for Compton scattering, as we shall see explicitly below. Moreover, this magnetic "recoil" contribution, which is small for heavy nuclei, is not negligible in the case of the nucleon.

Furthermore, for Compton scattering one also gets ground-state contributions, which never occur in single-photon processes for reasons of energy and momentum conservation, namely,

$$\langle\tfrac{1}{2}^+|\mathbf{M}(k)|\tfrac{1}{2}^+\rangle=\left\langle\tfrac{1}{2}^+\left|\frac{3i}{2}\int d^3x\ \hat{\mathbf{x}}\times\tilde{\mathbf{j}}(\mathbf{x})j_1(kx)\right|\tfrac{1}{2}^+\right\rangle\ , \tag{7}$$

and from the recoil current

$$\frac{1}{2\varepsilon_+}\boldsymbol{\varepsilon}_\lambda\cdot(\mathbf{P}_++\mathbf{P}_-)\left\langle\tfrac{1}{2}^+\left|\int d^3x\ \tilde{\rho}(\mathbf{x})j_0(kx)\right|\tfrac{1}{2}^+\right\rangle$$
$$=\frac{1}{2\varepsilon_+}\boldsymbol{\varepsilon}_\lambda\cdot(\mathbf{P}_++\mathbf{P}_-)G(k)\ . \tag{8}$$

Let me reiterate the physics: In the nucleon rest frame our model nucleon can absorb electric dipole radiation only below $N\bar{N}$ threshold. However, it will scatter magnetically for three reasons: recoil currents, ground-state contributions, and two-photon contributions (to be discussed below). For Compton scattering, it is not possible to choose a frame, where the "recoil current" can be eliminated.

I now write down the resonance amplitude in the photon-"nucleon" c.m. frame (i.e., $\mathbf{P}_+=-\mathbf{k}$):

$$R_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})=\frac{1}{\varepsilon_+-\varepsilon_-+a(k)}\left\langle\tfrac{1}{2}^+\left|\left[\tilde{H},\frac{1}{k}\boldsymbol{\varepsilon}_{\lambda'}^*\cdot\mathbf{C}(k)\right]+\boldsymbol{\varepsilon}_{\lambda'}^*\cdot\mathbf{A}(k)\right|\tfrac{1}{2}^-\right\rangle\left\langle\tfrac{1}{2}^-\left|\left[\tilde{H},\frac{1}{k}\boldsymbol{\varepsilon}_\lambda\cdot\mathbf{C}(k)\right]+\boldsymbol{\varepsilon}_\lambda\cdot\mathbf{A}(k)\right|\tfrac{1}{2}^+\right\rangle$$
$$+\frac{1}{a(k)}\langle\tfrac{1}{2}^+|(\boldsymbol{\varepsilon}_{\lambda'}^*\times\hat{\mathbf{k}}')\cdot\mathbf{M}(k)|\tfrac{1}{2}^+\rangle\langle\tfrac{1}{2}^+|(\boldsymbol{\varepsilon}_\lambda\times\hat{\mathbf{k}})\cdot\mathbf{M}(k)|\tfrac{1}{2}^+\rangle$$
$$+\frac{1}{\varepsilon_+-\varepsilon_--b(\mathbf{k}',\mathbf{k})}\left\langle\tfrac{1}{2}^+\left|\left[\tilde{H},\frac{1}{k}\boldsymbol{\varepsilon}_\lambda\cdot\mathbf{C}(k)\right]+\boldsymbol{\varepsilon}_\lambda\cdot\mathbf{A}(k)+\frac{\boldsymbol{\varepsilon}_\lambda\cdot\mathbf{k}'}{\varepsilon_+}\hat{\mathbf{k}}\cdot\mathbf{C}(k)\right|\tfrac{1}{2}^-\right\rangle$$
$$\times\left\langle\tfrac{1}{2}^-\left|\left[\tilde{H},\frac{1}{k}\boldsymbol{\varepsilon}_{\lambda'}^*\cdot\mathbf{C}(k)\right]+\boldsymbol{\varepsilon}_{\lambda'}^*\cdot\mathbf{A}(k)-\frac{\boldsymbol{\varepsilon}_{\lambda'}^*\cdot\mathbf{k}}{\varepsilon_+}\hat{\mathbf{k}}'\cdot\mathbf{C}(k)\right|\tfrac{1}{2}^+\right\rangle$$
$$-\frac{1}{c(\mathbf{k}',\mathbf{k})}\left\langle\tfrac{1}{2}^+\left|(\boldsymbol{\varepsilon}_\lambda\times\hat{\mathbf{k}})\cdot\mathbf{M}(k)+\frac{\boldsymbol{\varepsilon}_\lambda\cdot\mathbf{k}'}{\varepsilon_+}G(k)\right|\tfrac{1}{2}^+\right\rangle\left\langle\tfrac{1}{2}^+\left|(\boldsymbol{\varepsilon}_{\lambda'}^*\times\hat{\mathbf{k}}')\cdot\mathbf{M}(k)-\frac{\boldsymbol{\varepsilon}_{\lambda'}^*\cdot\mathbf{k}}{\varepsilon_+}G(k)\right|\tfrac{1}{2}^+\right\rangle\ . \tag{9}$$

Here I have used

$$a(k)=k+\frac{k^2}{2\varepsilon_+},\quad b(\mathbf{k}',\mathbf{k})=k+\frac{(\mathbf{k}'+\mathbf{k})^2}{2\varepsilon_-}-\frac{k^2}{2\varepsilon_+},\quad c(\mathbf{k}',\mathbf{k})=k+\frac{k^2}{2\varepsilon_+}+\frac{\mathbf{k}'\cdot\mathbf{k}}{\varepsilon_+}\ . \tag{10}$$

The first two terms in this expression correspond to Fig. 2(a) (direct process) and the last two to Fig. 2(b) (exchange process). Note that due to the choice of our reference frame (i.e., photon nucleon c.m. frame) the recoil current enters only in the exchange process. It is now useful to rewrite the above expression in the following way:

$$R_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})=\sum_{i=1}^{6}R^{(i)}_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})\ . \tag{11}$$

The first term contains a double commutator (DC) and two commutators

$$R^{(1)}_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})=\left\langle \tfrac{1}{2}^+\left|\left[\frac{1}{k}\varepsilon^*_{\lambda'}\cdot\mathbf{C}(k),[\tilde{H},\frac{1}{k}\varepsilon_{\lambda'}\cdot\mathbf{C}(k)]\right]\right|\tfrac{1}{2}^+\right\rangle+\left\langle \tfrac{1}{2}^+\left|\left[\frac{1}{k}\varepsilon^*_{\lambda'}\cdot\mathbf{C}(k),\varepsilon_{\lambda'}\cdot\mathbf{A}(k)\right]\right|\tfrac{1}{2}^+\right\rangle$$
$$-\left\langle \tfrac{1}{2}^+\left|\left[\varepsilon^*_{\lambda'}\cdot\mathbf{A}(k),\frac{1}{k}\varepsilon_{\lambda}\cdot\mathbf{C}(k)\right]\right|\tfrac{1}{2}^+\right\rangle\ . \tag{12}$$

These three terms cancel exactly against corresponding terms arising in the two-photon amplitude.

$$R^{(2)}_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})=\varepsilon^*_{\lambda'}\cdot\varepsilon_{\lambda}\frac{\varepsilon_+-\varepsilon_-}{3k^2}\left[\frac{a(k)}{\varepsilon_+-\varepsilon_-+a(k)}-\frac{b(\mathbf{k}',\mathbf{k})}{\varepsilon_+-\varepsilon_--b(\mathbf{k}',\mathbf{k})}\right]\langle\tfrac{1}{2}^+|\mathbf{C}^2(k)|\tfrac{1}{2}^+\rangle\ , \tag{13a}$$

$$R^{(3)}_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})=\varepsilon^*_{\lambda'}\cdot\varepsilon_{\lambda}\frac{1}{3k}\left[\frac{a(k)}{\varepsilon_+-\varepsilon_-+a(k)}-\frac{b(\mathbf{k}',\mathbf{k})}{\varepsilon_+-\varepsilon_--b(\mathbf{k}',\mathbf{k})}\right]\langle\tfrac{1}{2}^+|[\mathbf{A}(k);\mathbf{C}(k)]|\tfrac{1}{2}^+\rangle\ , \tag{13b}$$

$$R^{(4)}_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})=\varepsilon^*_{\lambda'}\cdot\varepsilon_{\lambda}\frac{1}{3}\left[\frac{1}{\varepsilon_+-\varepsilon_-+a(k)}+\frac{1}{\varepsilon_+-\varepsilon_--b(\mathbf{k}',\mathbf{k})}\right]\langle\tfrac{1}{2}^+|\mathbf{A}^2(k)|\tfrac{1}{2}^+\rangle\ , \tag{13c}$$

$$R^{(5)}_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})=-\frac{1}{3}\frac{\varepsilon^*_{\lambda'}\cdot\mathbf{k}\,\varepsilon_{\lambda}\cdot\mathbf{k}'}{\varepsilon_+-\varepsilon_--b(\mathbf{k}',\mathbf{k})}\left[2\frac{\varepsilon_+-\varepsilon_-}{\varepsilon_+k^2}\langle\tfrac{1}{2}^+|\mathbf{C}^2(k)|\tfrac{1}{2}^+\rangle+\frac{1}{\varepsilon_+k}\langle\tfrac{1}{2}^+|[\mathbf{A}(k);\mathbf{C}(k)]|\tfrac{1}{2}^+\rangle\right]\ , \tag{13d}$$

$$R^{(6)}_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})=(\varepsilon^*_{\lambda'}\times\hat{\mathbf{k}}')(\varepsilon_{\lambda}\times\hat{\mathbf{k}})\frac{1}{3}\frac{\mathbf{k}'\cdot\mathbf{k}}{\varepsilon_+k^2}\langle\tfrac{1}{2}^+|\mathbf{M}^2(k)|\tfrac{1}{2}^+\rangle\ . \tag{13e}$$

To achieve this structure, I have used that

$$\frac{\varepsilon_+-\varepsilon_-}{\varepsilon_+-\varepsilon_-+a}=1-\frac{a}{\varepsilon_+-\varepsilon_-+a},\quad \frac{\varepsilon_+-\varepsilon_-}{\varepsilon_+-\varepsilon_--b}=1+\frac{b}{\varepsilon_+-\varepsilon_--b} \tag{14}$$

on several occasions. In the ground-state contribution $R^{(6)}$, I have only kept terms to order $1/\varepsilon_+$.

Due to the simple structure of our "model" nucleon spectrum I could reduce everything to ground-state matrix elements. This is not possible in general, but all essential features for our considerations are nicely contained in the above expressions. In particular, all recoil corrections are systematically included. While on first sight $R^{(2)}-R^{(4)}$ look purely electric dipole, and $R^{(6)}$ magnetic dipole, actually all expressions are electric and magnetic dipole as well as electric quadrupole due to the fact the $b(\mathbf{k}',\mathbf{k})$ is still depending on both $\mathbf{k}'$ and $\mathbf{k}$. We will separate out the electric and magnetic dipoles later.

It is now important to isolate the DC and the two commutators Eq. (12) also in the two-photon amplitude (TPA) using the explicit expression of the electric dipole field

$$\mathbf{E}_{\lambda}(\mathbf{x},k)=\frac{3}{k}\nabla_x[j_1(kx)\varepsilon_{\lambda}\cdot\hat{\mathbf{x}}]+\tfrac{9}{2}j_2(kx)(\hat{\mathbf{x}}\varepsilon_{\lambda}\cdot\hat{\mathbf{x}}-\tfrac{1}{3}\varepsilon_{\lambda})=\nabla c_{\lambda}(\mathbf{x},k)+\mathbf{d}_{\lambda}(\mathbf{x},k)\ . \tag{15}$$

The contribution of the two-photon amplitude to scattering of electric dipole radiation is then given by

$$\left\langle\tfrac{1}{2}^+\left|\int d^3x'\int d^3x\sum_{l'l}E_{\lambda'l'}(\mathbf{x},k)\tilde{B}_{l'l}(\mathbf{x}',\mathbf{x})E_{\lambda l}(\mathbf{x},k)\right|\tfrac{1}{2}^+\right\rangle=\sum_i B^{(i)}_{\lambda'\lambda}(k',k) \tag{16}$$

in terms of the intrinsic two-photon operator $\tilde{B}_{l'l}(\mathbf{x}',\mathbf{x})$. Inserting Eq. (15) into Eq. (16) yields

$$B^{(1)}_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})=-R^{(1)}_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})\ , \tag{17a}$$

$$B^{(2)}_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})=\frac{1}{\varepsilon_+}\left\langle \tfrac{1}{2}^+ \left| \int d^3x' \int d^3x\, c_{\lambda'}(\mathbf{x}',k)\nabla_{x'}\tilde{\rho}(\mathbf{x}')\cdot\nabla_x\tilde{\rho}(\mathbf{x})c_\lambda(\mathbf{x},k) - c_{\lambda'}(\mathbf{x}',\mathbf{k})\nabla_{x'}\tilde{\rho}(\mathbf{x}')\cdot\mathbf{d}_\lambda(\mathbf{x},k)\tilde{\rho}(\mathbf{x}) - \tilde{\rho}(\mathbf{x}')\mathbf{d}_{\lambda'}(\mathbf{x}',k)\cdot\nabla_x\tilde{\rho}(\mathbf{x})c_\lambda(\mathbf{x},k) \right| \tfrac{1}{2}^+ \right\rangle$$
$$=\boldsymbol{\varepsilon}^*_{\lambda'}\cdot\boldsymbol{\varepsilon}_\lambda \frac{1}{\varepsilon_+}\langle \tfrac{1}{2}^+ | b^{(2)}(k) | \tfrac{1}{2}^+\rangle + \frac{1}{\varepsilon_+}\langle \tfrac{1}{2}^+ | (\boldsymbol{\varepsilon}^*_{\lambda'}\times\boldsymbol{\varepsilon}_\lambda)\cdot\boldsymbol{\sigma}\bar{b}^{(2)}(k) | \tfrac{1}{2}^+\rangle \ , \tag{17b}$$

$$B^{(3)}_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})=\left\langle \tfrac{1}{2}^+ \left| \int d^3x' \int d^3x \sum_{l'l} d_{\lambda' l'}(\mathbf{x},k)\tilde{B}_{l'l}(\mathbf{x}',\mathbf{x})d_{\lambda l}(\mathbf{x},k) \right| \tfrac{1}{2}^+ \right\rangle$$
$$=\boldsymbol{\varepsilon}^*_{\lambda'}\cdot\boldsymbol{\varepsilon}_\lambda \langle \tfrac{1}{2}^+ | b^{(3)}(k) | \tfrac{1}{2}^+\rangle + \langle \tfrac{1}{2}^+ | (\boldsymbol{\varepsilon}^*_{\lambda'}\times\boldsymbol{\varepsilon}_\lambda)\cdot\boldsymbol{\sigma}\bar{b}^{(3)}(k) | \tfrac{1}{2}^+\rangle \ . \tag{17c}$$

Here, I have used the important gauge condition for the intrinsic two-photon operators[4,5]

$$\sum_{l'} \frac{\partial}{\partial x'_{l'}}\tilde{B}_{l'l}(\mathbf{x}',\mathbf{x})=i[\tilde{\rho}(\mathbf{x}'),\tilde{\mathbf{j}}_l(\mathbf{x})]+\frac{1}{\varepsilon_+}\frac{\partial}{\partial x'_l}\tilde{\rho}(\mathbf{x}')\tilde{\rho}(\mathbf{x}) \tag{18}$$

as well as the continuity equation (5). Note that $B^{(1)}_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})$ is cancelled exactly by $R^{(1)}_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})$. As we shall see later $B^{(2)}_{\lambda'\lambda}$ is the only term of the whole amplitude which is finite at zero-photon energy.

The magnetic dipole contribution can also be written down easily in terms of the two-photon operator,

$$\left\langle \tfrac{1}{2}^+ \left| \int d^3x' \int d^3x \sum_{l'l} M_{\lambda' l'}(\mathbf{x}',k)\tilde{B}_{l'l}(\mathbf{x}',\mathbf{x})M_{\lambda l}(\mathbf{x},k) \right| \tfrac{1}{2}^+ \right\rangle$$
$$=-\tfrac{3}{4}(\boldsymbol{\varepsilon}_{\lambda'}\times\hat{\mathbf{k}}')\cdot(\boldsymbol{\varepsilon}_\lambda\times\hat{k})\sum_{\substack{s'l'\\ slm}}\varepsilon_{s'l'm}\varepsilon_{slm}\left\langle \tfrac{1}{2}^+ \left| \int d^3x' \int d^3x\, j_1(kx')\hat{\mathbf{x}}'_s, B_{l'l}(\mathbf{x}',\mathbf{x})\hat{\mathbf{x}}_s j_1(kx) \right| \tfrac{1}{2}^+ \right\rangle$$
$$=(\boldsymbol{\varepsilon}_{\lambda'}\times\hat{\mathbf{k}}')\cdot(\boldsymbol{\varepsilon}_\lambda\times\hat{\mathbf{k}})\langle \tfrac{1}{2}^+ | b^{(4)}(k) | \tfrac{1}{2}^+\rangle \ , \tag{19}$$

but obviously one needs an explicit model for $B_{l'l}$ in order to calculate this expression. That will be followed up later. As was pointed out already, this magnetic contribution arises due to the $N\bar{N}$ continuum in our model. Of course, there exist also higher multipoles, which will not be discussed explicitly here.

We will now write down the complete scattering amplitude for our model system, but will only consider $E1E1$ and $M1M1$ multipoles. To separate out the multipoles from the resonance term we will approximate

$$b(\mathbf{k}',\mathbf{k})\approx k^2+\frac{k^2}{\varepsilon_-}-\frac{k^2}{2\varepsilon_+}=b(k) \tag{20}$$

in all denominators in (13a)–(13e). The complete expression $b(\mathbf{k}',\mathbf{k})$ will be kept in the numerators. This is an excellent approximation for moderate photon energies.

The complete expression now reads

$$T_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})=T^{E1}_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})+T^{M1}_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})+(\text{other multipoles}) \ , \tag{21}$$

$$T^{E1}_{\lambda'\lambda}(\mathbf{k}',\mathbf{k})=\boldsymbol{\varepsilon}^*_{\lambda'}\cdot\boldsymbol{\varepsilon}_\lambda t^{E1}(k), \quad T^{M1}_{\lambda'\lambda}=(\boldsymbol{\varepsilon}^*_{\lambda'}\times\hat{\mathbf{k}})\cdot(\boldsymbol{\varepsilon}_\lambda\times\hat{\mathbf{k}})t^{M1}(k) \ , \tag{21a}$$

$$\begin{aligned} t^{E1}(k)=&\frac{\varepsilon_+-\varepsilon_-}{3k^2}\left[\frac{a(k)}{\varepsilon_+-\varepsilon_-+a(k)}-\frac{b(k)}{\varepsilon_+-\varepsilon_--b(k)}\right]\langle \tfrac{1}{2}^+ | \mathbf{C}^2(k) | \tfrac{1}{2}^+\rangle \\ &+\frac{1}{3k}\left[\frac{a(k)}{\varepsilon_+-\varepsilon_-+a(k)}-\frac{b(k)}{\varepsilon_+-\varepsilon_--b(k)}\right]\langle \tfrac{1}{2}^+ | [\,\mathbf{A}(k);\mathbf{C}(k)] | \tfrac{1}{2}^+\rangle \\ &+\frac{1}{3}\left[\frac{1}{\varepsilon_+-\varepsilon_-+a(k)}+\frac{1}{\varepsilon_+-\varepsilon_--b(k)}\right]\langle \tfrac{1}{2}^+ | \mathbf{A}^2(k) | \tfrac{1}{2}^+\rangle+\frac{1}{6\varepsilon_-}\langle \tfrac{1}{2}^+ | \mathbf{M}^2(k) | \tfrac{1}{2}^+\rangle \\ &-\frac{1}{9}\frac{1}{\varepsilon_+-\varepsilon_--b(k)}\left[2\frac{\varepsilon_+-\varepsilon_-}{\varepsilon_+}\langle \tfrac{1}{2}^+ | \mathbf{C}^2(k) | \tfrac{1}{2}^+\rangle-\frac{k}{\varepsilon_+}\langle \tfrac{1}{2}^+ | [\,\mathbf{A}(k);\mathbf{C}(k)] | \tfrac{1}{2}^+\rangle\right] \\ &+\frac{1}{\varepsilon_+}\langle \tfrac{1}{2}^+ | b^{(2)}(k) | \tfrac{1}{2}^+\rangle+\langle \tfrac{1}{2}^+ | b^{(3)}(k) | \tfrac{1}{2}^+\rangle \ , \end{aligned} \tag{21b}$$

$$t^{M1}(k)=-\frac{1}{6\varepsilon_-}\frac{\varepsilon_+-\varepsilon_-}{\varepsilon_+-\varepsilon_--b(k)}\langle\tfrac{1}{2}^+|\mathbf{C}(k^2)|\tfrac{1}{2}^+\rangle-\frac{k}{6\varepsilon_-}\frac{1}{\varepsilon_+-\varepsilon_--b(k)}\langle\tfrac{1}{2}^+|[\,\mathbf{A}(k);\mathbf{C}(k)]|\tfrac{1}{2}^+\rangle$$
$$-\frac{1}{6\varepsilon_+}\frac{1}{\varepsilon_+-\varepsilon_--b(k)}\{2(\varepsilon_+-\varepsilon_-)\langle\tfrac{1}{2}^+|\mathbf{C}^2(k)|\tfrac{1}{2}^+\rangle+k\langle\tfrac{1}{2}^+|[\,\mathbf{A}(k);\mathbf{C}(k)]|\tfrac{1}{2}^+\rangle\}+\langle\tfrac{1}{2}^+|b^{(4)}(k)|\tfrac{1}{2}^+\rangle\ . \tag{21c}$$

Equations (21b) and (21c) are the main results of this paper. These expressions are rather involved despite the extreme simplicity of the model "nucleon" spectrum. This stems from the fact that we have kept recoil corrections exactly up to order k/ε_+, which is important for the nucleon.

Several comments are in order. (i) The DC in (12) and (16) cancels exactly to all orders in k. (ii) Despite the extreme simplicity of the nucleon excitation spectrum assumed here, which only admits electric dipole radiation to be absorbed, we get a magnetic contribution which consists of a recoil contribution, a ground-state contribution, and a two-photon contribution [cf. (21c)]. The recoil contribution depends on the charge and current densities. The ground-state contribution requires a knowledge of the current, and the magnetic two-photon contribution requires explicit knowledge of the two-photon amplitude. One can calculate t^{E1} and t^{M1} with an ansatz for these operators and a knowledge of the nucleon ground-state wave function. This will be addressed in Sec. V. (iii) I have demonstrated explicitly the cancellation of the double commutator arising in both the resonance amplitude and the two-photon amplitude. This cancellation is important for the low-energy theorem to hold. But even more important is to realize that there is no physical significance to be given to the DC as has been done in many papers concerned with the experimental evaluation of data.[2] The same DC arises in the photonuclear Thomas-Reiche-Kuhn (TRK) sum rule. While the DC evaluated with the kinetic-energy operator alone gives the classical TRK results, the DC evaluated with the *NN* interaction has been interpreted as a "measure" of non-nucleonic degrees of freedom. A similar interpretation has been attempted for the DC arising in the TPA. However, as demonstrated here and more formally in earlier work,[5] the TPA DC just cancels exactly against a corresponding term in R (to all orders in k), and it is therefore not appropriate to attempt a measurement of exchange effects via this term.

III. LOW-ENERGY BEHAVIOR AND STATIC POLARIZABILITIES

Let me now explicitly study the low-energy behavior (up to order k^2) of the electric and magnetic dipole amplitudes. First I note that for $k\to 0$

$$\mathbf{C}(k)\to ik\mathbf{D},\quad \mathbf{M}(k)\to ik\boldsymbol{\mu}\ ,$$

$$\beta^{(2)}\to\left[\int d^3x\,\tilde{\rho}(\mathbf{x})\right]^2-\frac{k^2}{3}\int d^3x\,\tilde{\rho}(x)\int d^3x'\tilde{\rho}(\mathbf{x}')x'^2\ ,$$

$$\beta^{(4)}\to k^2\hat{\chi}_D\ ,$$

$$\hat{\chi}_D=-\tfrac{1}{12}\sum_{s'l'slm}\varepsilon_{s'l'm}\varepsilon_{slm}\int d^3x'\int d^3x\,\mathbf{x}_{s'}B_{l'l}(\mathbf{x}',\mathbf{x})\mathbf{x}_s\ ,$$

where $\mathbf{D}$ is the (static) electric dipole operator, and $\boldsymbol{\mu}$ the magnetic moment operator. Secondly, one now has to carefully take into account the k dependence of the factors multiplying the matrix elements in Eqs. (21b) and (21c).

$$t^{E1}(k)=\frac{e^2}{\varepsilon_+}-\frac{e^2}{3\varepsilon_+}\langle r^2\rangle k^2+\frac{2}{3}\left[\frac{1}{\varepsilon_+-\varepsilon_-}+\frac{1}{2\varepsilon_+}-\frac{1}{2\varepsilon_-}\right]\langle\tfrac{1}{2}^+|\mathbf{D}^2|\tfrac{1}{2}^+\rangle k^2-\frac{1}{6\varepsilon_-}\langle\tfrac{1}{2}^+|\hat{\boldsymbol{\mu}}^2|\tfrac{1}{2}^+\rangle k^2\ , \tag{22a}$$

$$t^{M1}(k)=\left[\frac{1}{6\varepsilon_-}+\frac{1}{3\varepsilon_+}\right]\langle\tfrac{1}{2}^+|\mathbf{D}^2|\tfrac{1}{2}^+\rangle k^2+\langle\tfrac{1}{2}^+|\hat{\chi}_D|\tfrac{1}{2}^+\rangle k^2\ . \tag{22b}$$

Expressions (22a) and (22b) differ somewhat from the usually quoted ones, since one usually approximates further than I did here [e.g., $\varepsilon_-\approx\varepsilon_+$ in (22b)].

The static electric polarizability α and the static magnetic susceptibility β are defined as

$$\alpha=\lim_{k\to 0}\alpha(k)=-\lim_{k\to 0}[t^{E1}(k)-t^{E1}(0)]/k^2\ , \tag{23a}$$

$$\beta=\lim_{k\to 0}\beta(k)=-\lim_{k\to 0}t^{M1}(k)/k^2\ , \tag{23b}$$

and we obtain from (22a) and (22b)

$$\alpha=\left[-\frac{2}{3}\frac{1}{\varepsilon_+-\varepsilon_-}-\frac{1}{3}\frac{\varepsilon_+-\varepsilon_-}{\varepsilon_+\varepsilon_-}\right]\langle\tfrac{1}{2}^+|\mathbf{D}^2|\tfrac{1}{2}^+\rangle+\frac{e^2}{3\varepsilon_+}\langle r^2\rangle+\frac{1}{6\varepsilon_-}\langle\tfrac{1}{2}^+|\hat{\boldsymbol{\mu}}^2|\tfrac{1}{2}^+\rangle\ , \tag{24a}$$

$$\beta=-\left[\frac{1}{6\varepsilon_-}+\frac{1}{3\varepsilon_+}\right]\langle\tfrac{1}{2}^+|\mathbf{D}^2|\tfrac{1}{2}^+\rangle-\langle\tfrac{1}{2}^+|\hat{\chi}_D|\tfrac{1}{2}^+\rangle\ . \tag{24b}$$

Of course, our model system does not have a paramagnetic susceptibility, since there is only a $\frac{1}{2}^-$ excited state. As

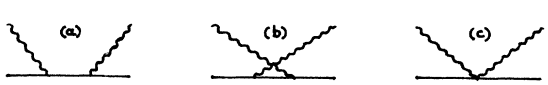

FIG. 2. The Compton scattering amplitude: (a) direct graph, (b) exchange graph, and (c) two-photon or "seagull" amplitude.

a matter of fact, this is the only essential difference between a "real" nucleon and the model nucleon. In our model β obviously arises entirely due to recoil and two-photon contributions. For nuclei β would be small, for nucleons it is substantial.

I now would like to carefully discuss this result for different cases. Assume first that $\varepsilon_+ \gg \varepsilon_- - \varepsilon_+$. This case is actually realized in nuclei; i.e., the mass of the system is much larger than its first excitation energy. Then

$$\alpha \approx \frac{2}{3} \frac{\langle \frac{1}{2}^+ | D^2 | \frac{1}{2}^+ \rangle}{\varepsilon_- - \varepsilon_+}, \quad \beta \approx 0 , \tag{25}$$

if the diamagnetic susceptibility is small as can be safely assumed. The above result corresponds to what one naively expects: A system which absorbs $E1$ radiation also has a sizable electric polarizability and the magnetic response is zero.

The conclusion changes dramatically if $\varepsilon_+ \approx \varepsilon_- - \varepsilon_+ = \varepsilon$ which is the case for nucleons. Then

$$\alpha = \frac{7}{6\varepsilon} \langle \tfrac{1}{2}^+ | \mathbf{D}^2 | \tfrac{1}{2}^+ \rangle , \qquad \beta = -\frac{5}{12\varepsilon} \langle \tfrac{1}{2}^+ | \mathbf{D}^2 | \tfrac{1}{2}^+ \rangle - \chi_D , \tag{26}$$

and now α and β are of the same order of magnitude and of opposite sign. (Sure, for a "real" nucleon there are other terms which do not arise in our simple model.)

The message to be drawn from this investigation is as follows: If the mass of the object we scatter off is large as compared to the first excitation energy, then our naive expectation holds that an object which absorbs electric dipole radiation also responds predominantly electrically in Compton scattering. This is the case for nuclei. If the mass of the object is of the same order of magnitude as the relevant excitation energy, then we cannot neglect recoil, ground-state and two-photon effects and the system responds both magnetically and electrically. This is the case for nucleons. As our model calculation shows, there is then nothing particularly surprising about the actually observed values of α and β for nucleons.

IV. PHOTOABSORPTION CROSS SECTION AND SUM RULES

One can also calculate the total photoabsorption cross section in the above model for energies below $N\bar{N}$ threshold. It is given by

$$\sigma(k) = \frac{4\pi^2}{k} \left| \left\langle \tfrac{1}{2}^+ \left| \frac{(\varepsilon_- - \varepsilon_+)}{k} \boldsymbol{\varepsilon}_\lambda \cdot \mathbf{C}(k) + \boldsymbol{\varepsilon}_\lambda \cdot \mathbf{A}(k) \right| \tfrac{1}{2}^- \right\rangle \right|^2 \delta((\varepsilon_+^2 + k^2)^{1/2} - \varepsilon_- + k) . \tag{27}$$

It is now tempting to check the validity of the sum rule of the nucleon polarizabilities

$$\alpha + \beta = \frac{1}{2\pi^2} \int_0^\infty \frac{\sigma(k)}{k^2} dk , \tag{28}$$

which is easily obtained from the Gell-Mann, Goldberger, and Thirring dispersion relation[6]

$$T_{\lambda\lambda}(k) - T_{\lambda\lambda}(0) = \frac{k^2}{2\pi^2} \int_0^\infty \frac{\sigma(k')}{k'^2 - k^2} dk' \tag{29}$$

by expanding both sides of (29) to order k^2. Furthermore, the following approximate sum rules have been proposed[1]

$$\alpha \approx \frac{1}{2\pi^2} \int_0^\infty \frac{\sigma^{E1}(k')}{k'^2} dk' , \qquad \beta \approx \frac{1}{2\pi^2} \int_0^\infty \frac{\sigma^{M1}(k')}{k'^2} dk' , \tag{30}$$

where σ^{E1} (σ^{M1}) is the electric dipole (magnetic dipole) contribution to the total absorption cross section.

The δ function in (27) requires the integrand in (28) to be evaluated at

$$\bar{k} = \frac{\varepsilon_-^2 - \varepsilon_+^2}{2\varepsilon_-} . \tag{31}$$

If $\varepsilon_- - \varepsilon_+ \ll \varepsilon_+$, i.e., the mass of the system is large as compared to the first excitation energy, then $\bar{k} \simeq \varepsilon_- - \varepsilon_+$. If, furthermore, $\bar{k}R \ll 1$ (R is the size of the system), then one finds from Eqs. (27) and (28)

$$\alpha + \beta = \frac{2}{3} \frac{1}{\varepsilon_- - \varepsilon_+} \langle \tfrac{1}{2}^+ | D^2 | \tfrac{1}{2}^+ \rangle , \tag{32}$$

which is consistent with the result (25) and the sum rules in Eq. (30).

However, if $\varepsilon_+ \approx \varepsilon_- - \varepsilon_+$, Eq. (32) [or, alternatively, Eq. (24a)] is not obtained. This is actually not surprising, since the formula for $\sigma(k)$ [Eq. (27)] is only correct for k smaller than $N\bar{N}$ threshold. Furthermore, $\bar{k}$ is now not small compared to ε_+, so that one cannot use a low-energy expansion for $C(\bar{k})$, which has been used to derive Eq. (32). From this we conclude that

$$\alpha = \frac{1}{2\pi^2} \int_0^\infty \frac{\sigma^{E1}(k)}{k^2} dk \tag{33}$$

is approximately valid for systems with a mass much heavier than the first excitation energy (as, e.g., nuclei), but Eq. (33) is wrong for systems such as the nucleon. The fundamental reason why Eq. (33) fails for nucleons is the fact that partial wave dispersion relations [cf. Eq. (29)] do not strictly hold, since they are not a Lorentz-invariant concept. For nuclei they are valid approximately, but they are not to be used for nucleons. Unfortunately, Eq. (33) has been applied to the nucleon by Hayward in Ref. 1, which has led to completely erroneous conclusions.

V. NUMERICAL ANALYSIS AND DISCUSSION

I now want to calculate $t^{E1}(k)$ and $t^{M1}(k)$ in a simple quark model and see how they behave for $0<k<100$ MeV. Reason for such an analysis is the fact that experiments to measure α and β can only be carried out for $k \approx 50-150$ MeV and extrapolation to $k \to 0$ is necessary.

The space part of the s-state quark wave function is given by

$$\psi(\boldsymbol{\rho},\boldsymbol{\lambda})=\frac{a_0^3}{\pi^{3/2}}e^{-a_0^2(\rho^2+\lambda^2)/2} \tag{34}$$

with relative coordinates $\boldsymbol{\rho}$ and $\boldsymbol{\lambda}$.[7] All three quarks are put into the s state and antisymmetrization is achieved via the color part of the wave function. The harmonic-oscillator constant a_0 is the only parameter of the model. It is now easy to express α and β in terms of this parameter and one obtains for the proton

$$\alpha_p=\left[\frac{2}{3}\frac{1}{\varepsilon_- - \varepsilon_+}+\frac{1}{3}\frac{\varepsilon_- - \varepsilon_+}{\varepsilon_+\varepsilon_-}+\frac{1}{3\varepsilon_+}\right]\frac{e^2}{a_0^2}+\frac{\mu_p^2 e^2}{24\varepsilon_-\varepsilon_+^2}, \tag{35a}$$

$$\beta_p=-\left[\frac{1}{6\varepsilon_-}+\frac{1}{3\varepsilon_+}\right]\frac{e^2}{a_0^2}-\frac{1}{2\varepsilon_+}\frac{e^2}{a_0^2} \tag{35b}$$

and for the neutron

$$\alpha_n=\left[\frac{4}{9}\frac{1}{\varepsilon_- - \varepsilon_+}+\frac{2}{9}\frac{\varepsilon_- - \varepsilon_+}{\varepsilon_+\varepsilon_-}\right]\frac{e^2}{a_0^2}+\frac{\mu_n^2 e^2}{24\varepsilon_-\varepsilon_+^2}, \tag{36a}$$

$$\beta_n=-\left[\frac{1}{9\varepsilon_-}+\frac{2}{9\varepsilon_+}\right]\frac{e^2}{a_0^2}-\frac{1}{9\varepsilon_+}\frac{e^2}{a_0^2}. \tag{36b}$$

Here I have given the quark an anomalous magnetic moment, so that the experimental proton (neutron) magnetic moment μ_p (μ_n) is reproduced. However, the terms $\sim\mu^2$ in Eqs. (35a) and (36a) are completely negligible. Furthermore, I have assumed a two-photon operator

$$B_{l'l}(\mathbf{x}',\mathbf{x})=\delta_{l'l}\sum_q \frac{e_q^2}{M_q}\delta(\mathbf{x}'-\mathbf{r}_q)\delta(\mathbf{x}-\mathbf{r}_q) \tag{37}$$

and the quark mass to be $M_q=\varepsilon_+/3$.

In the harmonic-oscillator constituent quark model $\varepsilon_- - \varepsilon_+$ and a_0 are related by

$$\frac{3a_0^2}{\varepsilon_+}=\varepsilon_- - \varepsilon_+ . \tag{38}$$

For $\varepsilon_- - \varepsilon_+ \approx 500$ MeV this predicts $a_0=2$ fm^{-1} and a root-mean-square radius of the nucleon of 0.5 fm, which is rather small. On the other hand, the experimental root-mean-square radius of 0.86 fm corresponds to $\varepsilon_- - \varepsilon_+ \approx 170$ MeV, which is definitely too small. Consequently, the simplest quark model does not provide us with a consistent scheme which describes the experimentally observed electric dipole excitation energy and the root-mean-square radius simultaneously.

Nevertheless, in the light of the present model calculations I present in Table I results for both cases discussed above as well as a case where I have disregarded relation (38) and chosen $\langle r^2\rangle^{1/2}=0.86$ and $\varepsilon_- - \varepsilon_+ \approx 500$ MeV (ε_+ is chosen to be 940 MeV in all cases). The experimental value[8] for α_p is $(11.6\pm2.4)(10^{-4}\ \mathrm{fm}^3)$, so that case 3 in Table I is not unrealistic. Of course, β_p is much too big since our model does not predict a paramagnetic susceptibility χ_p. A typical value for χ_p is 10 (10^{-4} fm^3) so that β_p in case 3 would be +2 (10^{-4} fm^3), which also agrees nicely with the experimental value $(2.9\pm2.4)\times10^{-4}$ fm^3. This case also agrees well with a recent experimental determination[9] of the neutron electric polarizability ($\alpha_n=11.7^{+4.3}_{-11.7}\times10^{-4}$ fm^3). Certainly, one should not take these numbers too seriously, since the model is much too naive to be quantitative. But in any case it contains the important physics qualitatively. Similar results have been obtained in Refs. 10 and 11.

I would like to point out that the term

$$\frac{1}{3}\frac{\varepsilon_- - \varepsilon_+}{\varepsilon_+\varepsilon_-}\frac{e^2}{a_0^2} \tag{39}$$

in Eq. (35a), which is usually neglected, actually provides an important contribution of about 10% to α_p.

I, furthermore, have evaluated $\alpha(k)$ and $\beta(k)$ as defined in Eqs. (23a) and (23b). It turns out that $\alpha(k)$ and $\beta(k)$ are very weakly energy dependent, so that a measurement of α (100 MeV) instead of α (0 MeV) introduces an error of about 2% only, which is certainly beyond experimental errors. Of course, this is a model-dependent statement.

TABLE I. Static electric polarizability and magnetic susceptibility for proton and neutron calculated from Eqs. (35) and (36).

	case 1	case 2	case 3	units
$\langle r^2\rangle^{1/2}$	0.5	0.86	0.86	fm
$\varepsilon_- - \varepsilon_+$	500	170	500	MeV
α_p	6.4	34.0	14.2	10^{-4} fm^3
β_p	−3.6	−8.2	−8.0	10^{-4} fm^3
α_n	3.3	28	10.3	10^{-4} fm^3
β_n	−1.6	−4.9	−4.6	10^{-4} fm^3

VI. CONCLUSIONS

In the present study I have shown in detail that a system, which only absorbs electric dipole radiation nevertheless scatters electrically and magnetically. The determining factor is the relation between the three different energy scales involved in the Compton scattering process: The photon energy k, the "nucleon" mass ε_+, and the excitation energy $\varepsilon_- - \varepsilon_+$. The naive expectation, that an object that absorbs electrically also predominantly scatters electrically, only holds if $\varepsilon_- - \varepsilon_+ \ll \varepsilon_+$. This case is indeed realized for nuclei. For nucleons we have $\varepsilon_- - \varepsilon_+ \approx \varepsilon_+$ and our naive expectation does not hold. The reason is that in this case recoil currents, ground-state contributions, and two-photon effects become important.

I have also pointed out that the double-commutator term which has been advocated to measure non-nucleonic degrees of freedom in nuclei[2] actually exactly cancels out to all orders in k and can therefore not be determined experimentally. Finally, I have calculated numerical values for α and β in a constituent quark model, which are in qualitative agreement with experiment.

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