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THE DYNAMICAL ORIGIN
OF THE RIGHT SIGNATURE FIXED POLE IN COMPTON SCATTERING ?

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A B S T R A C T

We find a new sum rule for the residue of the $J = 0$ fixed pole in Compton scattering. The sum rule applies to any target and links the residue to a calculable set of contributions from the resonance region. We discuss proton, neutron, deuteron and pion targets and the possibility of discriminating from previous models. The relationship with the parton model is noted and the feasibility of measuring pion photoabsorption cross-sections is discussed.

The possibility of having other than moving singularities in both real and virtual Compton scattering has been known for a long time ¹⁾. The elementarity of the photon allows for a variety of singularities at sense and nonsense points ^{2),3)}, some of these poles being additive and some multiplicative. A classical example is the fixed pole that restores the otherwise vanishing Pomeron coupling to on shell photons in the imaginary part of the forward spin averaged Compton amplitude.

In this paper we focus our attention on the right signature $J = 0$ additive fixed pole which contributes a piece of the form c/s to the real part of the asymptotic forward spin averaged Compton amplitude.

Because of analyticity one can immediately write down a sum rule that expresses C as an integral over the difference between the physical photo-absorption cross-section $[\sigma_T(\nu)]$ and the extrapolation to all energies of its asymptotic form $["\sigma_T(\nu)"]$. The sum rule reads

$$C = -\frac{\alpha Q^2}{M} - \frac{1}{2\pi^2} \int d\nu [\sigma_T(\nu) - "\sigma_T(\nu)"] \quad (1)$$

where Q is the charge of the target. A derivation can be found in Refs. 5), 6) and 7).

The interesting theoretical problem is to find an independent way to determine C and compare it with that evaluated from Equation (1). Our purpose is to provide such an equation and discuss it, as well as to discuss other proposals and their experimental status.

First let us discuss the models in which the current interacts with constituents locally, since these are the only known models to provide a mechanism for the $J = 0$ fixed pole ⁷⁾⁻¹⁰⁾. We believe that the sum rule we discuss may clarify the role of elementary constituents within the nucleon.

In the parton model, as an example of the family of constituent models, we see that $\sigma_T(\gamma p)$ receives contributions from two classes of diagrams on Figs. 1a and 1b. Diagram 1a contains vector meson poles in the photon legs and their contribution to $\sigma_T(\gamma p)$ will be like a typical hadronic process, e.g., $\sigma(\pi p)$. Furthermore, these diagrams can contain resonant enhancements in

the s channel. Diagrams 1b and 1c are characterized ^{7),8)} by the elementary photon interaction, do not contain vector meson poles in the photon legs, have no obvious counterpart in hadronic processes, have no s channel resonant enhancements and generate the $J = 0$ fixed pole behaviour in the Compton amplitude. Specifically, it is the two photon contact interaction on Fig. 1c that gives rise to the fixed pole ^{*)}. Our present approach is to concentrate on the channels that build $\sigma_{\text{Tot}}(\gamma p)$ and ask how the $J = 0$ fixed pole is built from this point of view. This is unrelated to the parton model but again there appears to be an intimate connection between the fixed pole and elementary photon couplings in certain channels. Specifically these latter are contact terms, not the two photon contact terms of the parton model (Fig. 1c), but the familiar contact interactions of the electric Born model (Fig. 2b).

NUCLEON COMPTON SCATTERING

In photon initiated reactions on nucleon targets it is well known that in many two-body and quasi-two-body reactions an electric Born type model is an excellent representation of the amplitude. On the other hand, the presence of the well-known resonances, including the nucleon, is also well established, and so by duality one expects normal Regge trajectories to contribute as well. These ideas seem confirmed if one plots (see Fig. 3) the photon cross-section and the corresponding hadron ($I = 1$) cross-section. Except for a normalization factor they agree almost everywhere. The difference can be totally attributed to the integral over all contact terms induced by the electric Born model ¹¹⁾. Hence one is lead to the following equation

$$\sigma_{\gamma N} = k \sigma_{\pi N} + \sigma_{\text{CONTACT}} \quad (2)$$

where k is the normalization constant (of order α) and the second piece is the contact interaction cross-section generated because of gauge invariance. For reasons that are explained below we shall neglect interference terms.

^{*)} For spinless partons. For $\frac{1}{2}$ spin partons in the infinite momentum frame an effective seagull diagram arises, see Ref. 8).

In passing we note already a parallel with the parton model classification. The $k\sigma_{\pi N}$ corresponds to the "hadronic photon" Fig. 1a for which the s channel resonance enhancements generate the t channel Regge exchanges. The contact terms are direct elementary photon couplings, as in Fig. 1b, which have no s channel resonance enhancements and as we shall now show appear intimately connected with the $J = 0$ fixed pole.

Inserting the cross-section Eq. (2) into the sum rule Eq. (1) we find our sum rule

$$C = -\frac{i}{2\pi^2} \int_0^\infty d\nu \sigma_{\text{cont}}(\nu) \quad (3)$$

This equation is theoretically appealing. First of all both sides are purely electromagnetic. Second, since we know the explicit channels that have contact terms it can be evaluated and compared with the value of C fixed by dispersion relations.

Let us discuss all possible processes from the beginning.

For the proton case (the best known) the way in which $\sigma_{\text{Tot}}(\gamma p)$ is built from the exclusive channels is experimentally known¹²⁾ and the important contact terms have been calculated¹¹⁾. Beyond the resonance region these terms are absorbed^{*)} ensuring rapid convergence of the sum rule Eq. (3) as it is also clear from Fig. 3. A good example is $\gamma N \rightarrow A_2 \Delta$ that contributes very little already (about $2\mu\text{b}$) to the photoabsorption cross-section.

Typical channels contributing to the sum rule are¹¹⁾

$$\begin{aligned} \gamma N &\rightarrow \eta N \\ \gamma N &\rightarrow \eta \Delta \\ \gamma N &\rightarrow \pi \Delta \\ &\dots \end{aligned} \quad (4)$$

*) The absorption follows the prescription of Ref. 12) for $\gamma N \rightarrow \pi \Delta$.

The presence of these terms seems well established. They explain near their thresholds a large number of details of these processes including angular distributions, spin correlations and energy dependence. Most important their over-all size is fixed. As seen in Fig. 3 all together they yield a sizeable contribution to the proton photoabsorption cross-section. Of course these contact terms are gauge dependent and so the Born terms must also be included, however, these latter are negligible ¹³⁾. Further since these terms are real they interfere weakly with nearby resonances ; in any event these latter are also small compared to the contact terms ^{13),14)}.

Upon inserting the magnitude of the contact terms into the sum rule Eq. (3) one finds

$$C = - 3\mu b. GeV$$

This equals, in our opinion accidentally, the value of the Thomson limit on protons. Since the gauge terms are generated by isovector photons the contact pieces are equal for protons and neutrons except for small corrections due to interference. Thus we expect $C_p = C_n$. The deuteron, because of the weak coupling, will have $C_D = 2C_p$.

We discuss now the experimental situation concerning baryon Compton scattering. Unfortunately Eq. (1) is very sensitive to the relative amounts of Pomeron and tensor trajectories in " σ_T " and especially to the intercept of the latter. In fact the low energy extrapolation of " σ_T " can be changed substantially by changing arbitrarily the intercepts over a wide range of values, without altering the quality of the fits for the high energy tail ^{7),15)}. The natural assumption, supported from Chew-Frautschi plots is to assume the intercept to be $\frac{1}{2}$ ⁵⁾. In this case one gets for the proton target a value that seems to be $- 3\mu b GeV$. When one turns to the neutron the problem is even more involved. To disentangle the contribution of f and A_2 one must have good measurements of the difference between total proton and neutron cross-sections at high energy. Present experiments cannot yield this information. This is best illustrated in Ref. 16) where two different fits to this difference lead to very different conclusions concerning C_n . In fact if one defines the high energy region by the condition $E \geq 2GeV$, the fit chooses a remarkably low value for the ratio f/A_2 , giving a small neutron fixed pole ¹⁷⁾. However, if one takes a more conservative point of view and takes $E \geq 4 GeV$, one gets

$f/A_2 \approx 5$. This new ratio gives $C_p \approx C_n$. Hence this is an open question. For the deuteron one can make an independent fit and its fixed pole can be shown to lie between the following limits ¹⁸⁾

$$-6.6 \leq C_D \leq -2.2 \quad \text{in } \mu b \cdot GeV$$

These bounds are compatible with all models except that of Ref. 10). All of these results are put together in the Table.

The final question concerning baryons is the behaviour of the fixed pole with varying photon mass. The sum rule for arbitrary q^2 (fixed) reads ^{19),20)}

$$B(q^2) + \frac{2M}{q^2} \int d\nu \left[\nu W_2(\nu, q^2) - \nu W_2(\nu, q^2) \right] = \frac{R(q^2)}{q^2} \quad (5)$$

where $B(q^2)$ is the Born term and $R(q^2)$ is the fixed pole residue.

In many models $R(q^2)$ is predicted to be linear in q^2 ^{8),9),10),20)}. In our model the behaviour of $R(q^2)$ is dependent upon the behaviour of the contact terms with q^2 and is a priori unknown. At least from the form of our sum rule Eq. (3), we would expect $R(q^2)$ to have the same sign for all q^2 .

Is this compatible with the data? Naïvely one has the impression ¹⁹⁾ that the fixed pole changes sign between photoproduction and the deep inelastic regime but this overlooks the fact that in proceeding from ν (at $q^2 = 0$) to ω a dimension is lost and so there is a question of what sets the scale ²⁰⁾. Indeed, as noted in Ref. 20), present data do not rule out a polynomial residue especially if $F_2(\omega) \ll 0.35$, i.e., diffraction is small in the deep inelastic. This is perhaps not unreasonable since several clues from exclusive and semi-inclusive channels suggest that non-diffractive mechanisms may play an important role at large q^2 ^{*)}.

*) We have in mind the works of Ref. 21), 22), the observation that ρ diffractive production seems to die out of $\sigma_{Tot}(\gamma p)$ as q^2 increases, the π^+/π^- ratio in the current fragmentation region $\neq 1$, and $\nu W_2^{\gamma n}/\nu W_2^{\gamma p} \lesssim 0.9$ even at $W \approx 25$. See the reviews in Ref. 23).

Further, as noted in Ref. 20), in a region where the data scale then the fixed pole residue would be linear in q^2 . This was extended in Ref. 21) where it was noted that the success of the Bloom-Gilman sum rule over a range of q^2 is indirect support for such a fixed pole behaviour in that q^2 range. By use of the variable ω_w it seems that the global scaling can be extended to $q^2 = 0$ (22). If the data indeed scale in this fashion, i.e.,

$$F(\omega_w) = \frac{\omega'}{\omega_w} F(\omega_w) \approx \frac{q^2 + a^2}{q^2} F(\omega_w) \quad (6)$$

where

$$\omega_w = \frac{2Mv + b^2}{q^2 + a^2} ; \quad \omega' = \frac{2Mv + b^2}{q^2}$$

then the residue of the fixed pole is given by

$$\int d\omega_w [F(\omega_w) - "F(\omega_w)"] = \frac{R(q^2)}{q^2} \quad (7)$$

for all q^2 , i.e., is linear in q^2 .

Therefore it does seem possible that, to some approximation, the fixed pole residue $R(q^2)/q^2$ is constant or at most slowly varying. Which of these two occurrences obtain is dependent, in our picture, upon the q^2 dependence of the contact terms. As pointed out by Elitzur (24) the excitation form factors are consistent with a universal form that is a function of the ratio q^2/s where s is the square of the centre-of-mass energy. If the contact terms also have such a behaviour then the ω_w scaling of the total cross-section will be exact and the fixed pole will be linear in q^2 due to the scaling contact terms. However, it seems possible that the contact terms might die faster than "scaling", i.e., $\gamma N \rightarrow MB$ dies like q^2/M_B^2 , hence faster than q^2/s when the M_B invariant mass squared is $s > M_B^2$. Indeed comparison of data on $\sigma(\gamma p \rightarrow \Delta^{++} \pi^-)$ in the vicinity of the second resonance suggests that the contact term dies more rapidly than the resonance form factor in this region (25). This is probably already evident in the rapid disappearance

with q^2 of the prominent photoproduction shoulder on the low mass side of the second resonance, this shoulder arising from the opening of the $\pi\Delta$ channel which is in turn generated by a contact term (see Fig. 3). If this interpretation is correct then we expect some non-scaling behaviour to be present in σ_{Tot} , the contact terms to die fast in q^2 and as a result the fixed pole to tend to zero in the large q^2 limit. Experimental verification of such a correlation between the q^2 dependence of the contact terms and the fixed pole would be most interesting.

COMPTON SCATTERING ON PIONS

The Harari-Pagels sum rule reads ²⁶⁾

$$-\frac{\alpha}{m_\pi} = \frac{1}{2\pi^2} \int_{\nu_0}^{\infty} d\nu [\sigma_{\gamma\pi^+} - \sigma_{\gamma\pi^-}] + C_\pi \quad (8)$$

where C_π is the residue of a $J = 0$ fixed pole that may be present in the Compton amplitude. Harari and Pagels (HP) showed that if one assumes $C_\pi = 0$ then resonance saturation by ω gives a good result.

Under similar assumptions to those employed in the derivation of the HP sum rule a superconvergence sum rule can be discussed for $\rho\pi$ scattering ²⁷⁾

$$0 = \frac{1}{2\pi^2} \int_{\nu_0^{\rho}}^{\infty} d\nu [\sigma_{\rho\pi^+} - \sigma_{\rho\pi^0}] \quad (9)$$

where ν_0^{ρ} is the threshold appropriate to $\rho\pi$.

Given the HP conclusion that $C_\pi = 0$ then it has been argued ¹⁰⁾ that the two sum rules are inconsistent with vector dominance and further that VD requires $C = -\alpha/m$. We also believe that the two sum rules are not simply related by VD but do not accept the claim that $C = -\alpha/m$.

The essential reason for this is the schizophrenic nature of the photon coupling, a feature common to both the parton model and the present approach, namely that it has sometimes a hadronic character and sometimes manifests an elementary "pointlike" coupling. For example in the parton model

$\sigma(\gamma\pi)$ receives contributions from the hadronic photon (Fig. 1a), for which VD has some success, and also the elementary photon (Fig. 1b). We hence obtain two independent sum rules, one for the pointlike and one for the hadronic photon. Specifically, the superconvergence one Eq. (9) is an equation for the hadronic photon

$$-\frac{1}{2\pi^2} \int d\nu \left[\sigma_{\gamma\pi^+} - \sigma_{\gamma\pi^0} \right]_{\text{hadronic}} = 0 \quad (10)$$

and so combining with Eq. (8) we have a sum rule relating the pointlike photon coupling to the fixed pole residue ^{*)}

$$C_\pi = -\frac{1}{2\pi^2} \int_{\nu_0}^{\infty} d\nu \left[\sigma_{\gamma\pi^+} - \sigma_{\gamma\pi^0} \right]_{\text{pointlike}} \quad (11)$$

The picture is satisfyingly self-consistent in that one would in general expect $C_\pi \neq 0$ arising from the seagull terms in the pointlike interaction topologies (see Figs. 1b and 1c) and this is now related to an integral over this same class of pointlike contributions to the cross-section. If VD were exact then the latter would all vanish and so $C_\pi = 0$ in Eq. (11). Conversely it is the failure of VD, expressed by the elementarity of the photon couplings, that generates C_π .

The sum rule Eq. (11) is of little use until the nature of the pointlike couplings is specified. In our model the pointlike contributions to the cross-sections are the channels containing the contact interactions (Fig. 2) and so Eq. (11) becomes again our sum rule Eq. (3)

$$C_\pi = -\frac{1}{2\pi^2} \int_{\nu_0}^{\infty} d\nu (\sigma_{\gamma\pi^+} - \sigma_{\gamma\pi^0})_{\text{contact}} \quad (12)$$

*) For $\nu_0 < \nu_0^0$ the ω contribution and the Thomson term cancel following HP.

In particular there is the large contribution from the channel which is depicted in Fig. 4.

Hence, contrary to some authors we believe that the Harari-Pagels saturation is correct because the pieces they neglected, contact terms and fixed pole, cancel each other !

However, we predict a large Compton cross-section for the channels that do not contain ω (in particular $\rho^0 \pi^+$), while others, though of some importance as seen in Fig 4, will be absorbed and contribute little.

This leads to our final point. Can Compton pion scattering be measured and how should it be done ? It is interesting to notice that the classical Chew-Low procedure that has been successfully applied for $\pi\pi$ scattering presents problems here. Indeed the presence of contact terms makes the problem of isolating the pion pole hard. An interesting example is the measurement of $\gamma p \rightarrow \Delta^{++} \pi(\rho)$ which shows a large cross-section. One is tempted to assign a large value to A_2 formation in the $\gamma\pi \rightarrow \rho\pi$ subsystem by extrapolation to the pion pole and looking to the relevant region in missing mass. This would be erroneous because of the presence of a contact coupling between $\gamma N A_2 \Delta$ that explains the size of this channel quite naturally.

Hence one must go to a region in energy where the pion pole survives and the contact term is unimportant. This is true at moderate energies since the π exchange contributes to all partial waves and the contact term gets absorbed.

A counter experiment for the reaction

$$\gamma + N \longrightarrow \Delta^{++} + \text{anything}$$

(where the $t_{N\Delta}$ is small and at energies of 10 GeV or larger) in which one detects the recoiling delta will undoubtedly be a candidate to measure the relevant Compton cross-section. A counter experiment seems well within present possibilities. If the cross-section follows Fig. 4, this could be an interesting result per se and could add weight to our model and to the idea that there are direct four particle couplings for photons like $\gamma\rho\pi\pi$, as proposed elsewhere²⁸⁾. Furthermore, the possible connection between these pointlike interactions and those of the parton model is intriguing and perhaps worthy of further study.

	Quark parton model	Born term model	Contact term model	Experiment
Proton	a	- 3.0	- 3.0	-3.0±1.0
Neutron	2/3 a	0	- 3.0	-2.0±2.0
Deuteron	5/3 a	- 1.5	- 6.0	-6.6 -2.2
Pion (charged)	b	-20.0	-10.0	unknown
Pion (neutral)	b	0	- 5.0	unknown

Table caption :

Values of the $J = 0$ fixed pole for different processes in different models and experiment a and b are not predicted.

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FIGURE CAPTIONS

- Fig. 1 : Parton model contributions to the $\sigma_T(\gamma p)$ (1a and 1b) and origin of the $J = 0$ fixed pole in the Compton amplitude (1c).
- Fig. 2 : Contributions to $\sigma_T(\gamma p)$ from exclusive channels. 2a contains s channel enhancements and the photon is typically hadronlike in its behaviour. In 2b a meson (M) baryon (B) system is generated as a result of the elementary photon contact interaction.
- Fig. 3 : Normalized $\sigma_{\gamma p}$ cross-section :
a) to yield the same $\sigma_{\pi N}$ ($I = 1$) except for over-all factor both at the Δ region and asymptotically.
b) the difference between $k\sigma_{\gamma p} - \sigma_{\pi N}$ (see text).
- Fig. 4 : Typical contact term contribution to pion Compton scattering. Full lines are unabsorbed contributions, dotted lines absorbed contributions (see text).

FIG. 3

FIG. 4