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ON ELECTROMAGNETIC EFFECTS IN DIFFRACTION SCATTERING OF HADRONS

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1. INTRODUCTION

When constructing the theory of diffraction electromagnetic effects are usually neglected [1,2]. However, it is well known that electromagnetic process cross-sections increase with energy (e.g.[3,4,5]). The purpose of this paper is to ascertain the order of magnitude of electromagnetic corrections. We restrict ourselves to the electromagnetic corrections to the vacuum pole intercept and to the triple-Pomeron coupling constant.

The electromagnetic shift of the vacuum pole due to self-energy part, shown in Fig. 1,

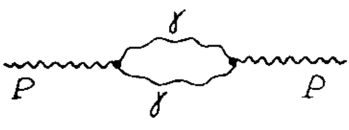


Fig. 1

has been discussed in Ref. [5]. If the input Pomeron has $\alpha_P(0) = 1 + \Delta^h$ ($\alpha_P(t)$ is the Pomeron trajectory), then the pair of poles arise due to the lack of photon reggeization: $\alpha_\pm(0) = 1 + \Delta^h \pm \Delta_\pm^{\gamma\gamma}$, where

$$\Delta_\pm^{\gamma\gamma} = -\frac{\Delta^h}{2} \pm \sqrt{\left(\frac{\Delta^h}{2}\right)^2 + \Sigma^{\gamma\gamma}} \qquad (1)$$

$$\Sigma^{\gamma\gamma} = \frac{1}{4\pi^3 R^2}\,\frac{(\sigma_{tot}^{\gamma N})^2}{\sigma_{tot}^{NN}} \,. \qquad (2)$$

Here $\sigma_{tot}^{\gamma N} \simeq 0.1$ mb , $\sigma_{tot}^{NN} \simeq 40$ mb – are the total cross-sections of nucleon photoabsorbtion and of NN- interaction, respectively; $R^{-2} \simeq 0.5\ (GeV/c)^2$ is the parameter characterising the t-dependence of the deep inelastic scattering cross-section.

It is clear from (1) that $\alpha_+(0) > 1$, $\alpha_-(0) < 1$ independently of the sign of Δ^h. The vacuum pole shift due to interference of electromagnetic and strong interactions is also possible. The self-energy part of such type is pictured in Fig. 2.

Fig. 2

V denotes vector mesons ρ, ω . The quantity $\Delta^{\gamma V}$ has been calculated in Ref. [6,7]. In the present paper integration over

energy corresponding to the γ-V loop, is separated into two parts. Unlike in Ref. [7] the main contribution to $\Delta^{\gamma V}$ comes from the region where the V-exchange is nonreggeized. It is found that $\Delta^{\gamma V} \sim 10^{-3}$.

In Section 2 the radiative correction to triple-Pomeron coupling due to the diagram of Fig. 3 is estimated

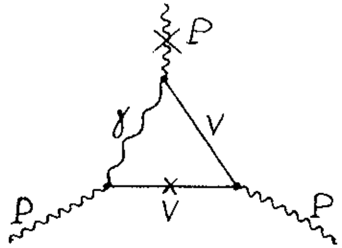


Fig. 3

The cross in Fig. 3 means that the particle V is on the mass shell. The corresponding correction to the triple-Pomeron vertex g_{PPP} is simply connected with the pole shift: *)

$$g_{PPP}^{\gamma V} = \Delta^{\gamma V}\sigma_{tot}^{VN} \Big/ \sqrt{\sigma_{tot}^{NN}} \qquad (3)$$

In Section 3 the calculated quantities are compared with experimental data. It is also shown that in cases of strong [2] and weak [1] Pomeron coupling strong interactions do not shadow electromagnetic corrections.

1. Calculation of $\Delta^{\gamma V}$

The total hadron cross-section in the first order in $\Delta^{\gamma V}$ depends on the energy in the following way

*) We normalize g_{PPP} as in Ref. [8]

$$\sigma_{tot} = \sigma^{h}_{tot} + \sigma^{int}_{tot}$$
$$\sigma^{int}_{tot} = \Delta^{\gamma V}\sigma^{h}_{tot}\,\ln\left(\frac{s}{M^2}\right), \tag{4}$$

where $M^2 \sim 4\ (\mathrm{GeV})^2$ is the minimum permissible energy squared for Pomeron exchange. The first term in (4) σ^{h}_{tot} is determined by strong interaction, the second one σ^{int}_{tot} arises due to interference of strong and electromagnetic interactions (see, e.g., diagram at Fig. 4).

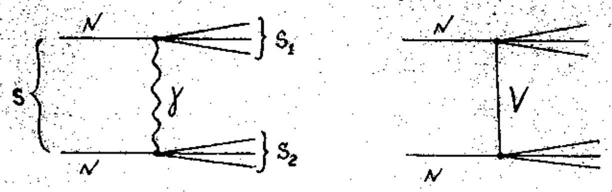


Fig. 4

The differential cross-section of NN-interaction with the production of two showers has the form/6/

$$\frac{d^3\sigma^{int}}{dq^2 ds_1 ds_2} = \frac{1}{(2\pi)^3 S^2}\,\frac{1}{q^2(q^2+m_V^2)}\sum_V \mathrm{Im}\,T^{\gamma N\to VN}_{\mu\nu}(s_1,t,q^2)\Big|_{t=0}\ \mathrm{Im}\,T^{\gamma N\to VN}_{\mu\nu}(s_2,t,q^2)\Big|_{t=0}, \tag{5}$$

where $-q^2$ is the momentum transfer squared, s_1, s_2 are the shower effective masses squared; $T^{\gamma N\to VN}_{\mu\nu}(s_i, t, q^2)$ is the vector meson photoproduction amplitude; t is the momentum transfer squares from γ to V.

We use the vector dominance model and neglect the absorption of the longitudinal vector meson and γ-quantum to obtain:

$$\mathrm{Im}\,T^{\gamma N\to VN}_{\mu\nu}(s_1,t,q^2)\Big|_{t=0} = \frac{s_1}{1+q^2/m_V^2}\left(16\pi\,\frac{d\sigma^{\gamma N\to VN}}{dt}\Big|_{t=0}\right)^{1/2}\cdot$$
$$\cdot\left(g_{\mu\nu} - \frac{P_\mu q_\nu + P_\nu q_\mu}{Pq} + q^2\,\frac{P_\mu P_\nu}{(Pq)^2}\right)\ . \tag{6}$$

Here $d\sigma^{\gamma N\to VN}/dt\,\big|_{t=0}$ is the cross-section of vector meson photoproduction on the nucleon having the momentum P. Using $q_m^2 = S_1 S_2/S$ (the minimum value of q^2) instead of S_1, S_2 and putting (6) into (5) we can perform integration over S_1.

$$\frac{d^2\sigma^{int}}{dq^2 dq_m^2} = \frac{4}{\pi}\sum_V \frac{m_V^4}{(q^2+m_V^2)^3}\,\frac{d\sigma^{\gamma N\to VN}}{dt}\Big|_{t=0}\left(2\frac{q^2}{q_m^2} + \frac{q_m^2}{q^2} - 2\right)\ln\left(\frac{s}{M^2}\right) \tag{7}$$

At small q_m^2 values the energy corresponding to V-exchange is large and we should account for the reggeization of V. (From the point of view of the complex momentum theory the contribution of this region is essentially smaller than the above discussed). So we substitute the usual vector meson propagator by

$$\left(\frac{m_V^2}{q_m^2}\right)^{\alpha_V(q^2)-1}\frac{\pi}{2}\,\alpha_V'\ \mathrm{tg}\left(\frac{\pi\alpha_V(q^2)}{2}\right)\alpha_V(q^2)\ , \tag{8}$$

where $\alpha_V(q^2) = \alpha_V(0) - \alpha_V' q^2$ (V - a reggeon trajectory), $\alpha_V(0) = 1/2$; $\alpha_V' \simeq 1\ (\mathrm{GeV/c})^{-2}$. The factor $\alpha_V(+q^2)$ in (8) is expected on the basis of exchange degeneracy of ρ and A_2 trajectories (e.g. /9/).

At large q^2 the factor $tg(\pi\alpha_v/2)$ contains poles, that should be compensated in some way. Instead of this we limit the integration region to $q^2 \lesssim m_v^2$, in accordance with multiperipheral kinematics. Dividing the integration region over q_m^2 into two intervals we obtain:

$$\sigma_{tot}^{int} = \frac{4}{\pi^2}\sum_V m_v^2 \left.\frac{d\sigma^{\gamma N\to VN}}{dt}\right|_{t=0} (I_1 + I_2)\ln\left(\frac{s}{M^2}\right) \qquad (9)$$

I_1 is the contribution from the region, where V is not reggeized

$$I_1 = \int_{y_0}^{\infty} dy \int_y^{\infty} \frac{dx}{(1+x)^3}\left(2\frac{x}{y} + \frac{y}{x} - 2\right), \qquad (10)$$

where $x = q^2/m^2$, $y = q_m^2/m^2$. The integration over x, y in (10) is enlarged to infinity due to fast convergence of the integrand.

$$I_2 = \int_0^{y_0} dy \int_y^1 \frac{dx}{(1+x)^2}\, y^{1-\alpha_v(xm_v^2)}\left(2\frac{x}{y} + \frac{y}{x} - 2\right)\frac{\pi}{2}\alpha_v' m_v^2\, tg\left(\frac{\pi\,\alpha_v(xm_v^2)}{2}\right) \qquad (11)$$

is the integral over the region where V is reggeized.

As a boundary between these two regions we use $y_0 = 1$ and obtain

$$I_1 = 0.85; \quad I_2 = 0.03. \qquad (12)$$

Substituting these values as well as

$$\left.\frac{d\sigma^{\gamma N\to\rho N}}{dt}\right|_{t=0} \simeq 0{,}1 \text{ mb} \qquad \left.\frac{d\sigma^{\gamma N\to\omega N}}{dt}\right|_{t=0} \simeq 0{,}01 \text{ mb} \quad /10/$$

into (9) we have

$$\sigma_{tot}^{int} = 0{,}024 \ln\left(\frac{s}{M^2}\right) \text{ mb}. \qquad (13)$$

Hence it follows from (4) that vacuum pole shift is

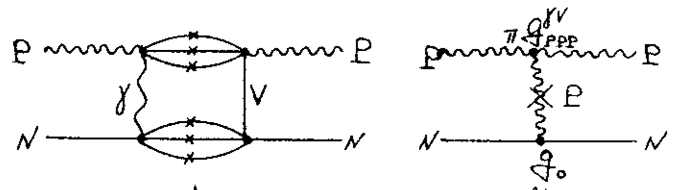
$$\Delta^{\gamma V} \simeq 0{,}6\cdot 10^{-3}. \qquad (14)$$

Note for comparison, that in Refs. /6,7/ the value $\Delta^{\gamma V} = 2.3 \times 10^{-3}$ has been obtained. *)

Let us emphasize that the calculation of $\Delta^{\gamma V}$ contains some uncertainties and should be considered only as a rough estimate. Besides there are lots of electromagnetic contributions to Δ^{int} unaccounted here. Nevertheless one may hope that (4) gives us the right order of magnitude of Δ^{int} as well as vector dominance model permits semiquantitative description of low energy electromagnetic hadron interaction.

2. Radiative corrections to the triple-Pomeron coupling constant

To calculate the contribution of diagram of Fig. 3 we shall examine the Pomeron-nucleon scattering shown at Fig. 5



a) б)

Fig. 5

*) In Ref. /6/ due to a numerical error in getting (27) from (21) this result was 4 times increased.

The calculations here are carried out analogously to NN scattering, see Fig. 4. The difference is in taking into account only low effective masses of the upper shower in the diagram of Fig. 5a. Confining ourselves to the eikonal approximation, we obtain that the contribution to σ_{tot}^{PN} from the diagram of Fig. 5a is connected with σ_{tot}^{int} found in Section 1, in the following way:

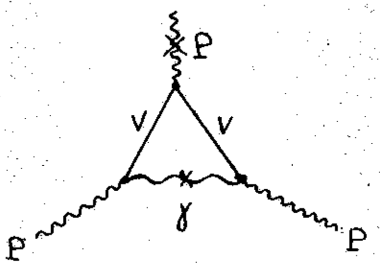
$$\left(\sigma_{tot}^{PN}\right)^{\gamma V} = \pi\,\frac{\sigma_{tot}^{VN}}{\sigma_{tot}^{NN}}\,\frac{\sigma_{tot}^{int}}{\ln\left(\frac{s}{M^2}\right)} \qquad (15)$$

The comparison of (13) with (4) and with Fig. 5b leads us just to (3).

Note that all uncertainties, contained in calculations of PN interaction cross-section, are included in the quantity σ_{tot}^{int}. Putting into (3) $\sigma_{tot}^{VN} \simeq 30$ mb [11] and $\Delta^{\gamma V} \sim 10^{-3}$ we have

$$g_{PPP}^{\gamma V} \simeq 10^{-2}\ (\mathrm{GeV}/c)^{-1}. \qquad (16)$$

The modern experimental data on the reaction $p + p \to p + X$ do not contradict the value $g_{PPP} = 0.25 \div 0.5\ \mathrm{GeV}^{-2}$ (e.g. [8,12]). Thus the correction (16) is about several per cent. of g_{PPP}. There are also some other electromagnetic contributions to g_{PPP}. One of them is shown in Fig. 6.



It is interesting to note that in the case of vector dominance type theory the diagram of Fig. 6 contains rather large essential transfered momenta: $q_\perp^2 \sim m_\rho^2$. One may expect this effect changing the form of the inclusive spectrum in the reaction $p + p \to \gamma + X$ at $q_\perp \geqslant m$, since at low q γ-quanta arise mainly from π^0-decays. This example shows that the γ-quantum spectrum is sensitive to the structure of hadron parton wave function at small distances (i.e., input multiperipheral ladder structure).

3. Discussion of results

The obtained estimates of radiative corrections to $\alpha_P(0)$ (14) and g_{PPP} (16) show that in the available energy range these effects are small. Nevertheless, it is useful to discuss the role of electromagnetic interactions at asymptotic energies.

1. Strong coupling of Pomerons [2]. This case seems the most interesting from the point of view of possible electromagnetic interaction influence, as in this case there is small parameter, namely the triple-Pomeron coupling constant g_{PPP}^h. To meet the condition $\alpha_P(0) = 1$ after inclusion of triple-Pomeron interaction, there must be the gap in the input Pomeron spectrum:

$$\Delta^h \simeq -\frac{\left(g_{PPP}^h\right)^2}{32\pi\alpha_P'}\ln\frac{\left(g_{PPP}^h\right)^2}{32\pi\alpha_P'} \simeq 1{,}5\cdot 10^{-2}$$

An additional shift of vacuum pole intercept, connected with self-energy part shown in Fig. 1 is

$$\Delta_+^{\gamma\gamma} \simeq \Sigma^{\gamma\gamma}/\Delta^h \simeq 1{,}7\cdot 10^{-4} \qquad (17)$$

$\Delta^{\delta V}+\Delta^{\gamma\gamma}$ is about 10 % of Δ^h. We should emphasize that the problem of including electromagnetic corrections of type Δ^{int} into Δ^h is in principle not solved and must be considered separately. In any case it is not excluded that the increase of the input gap by Δ^{int} is completely compensated due to increasing the triple-Pomeron coupling constant by $g_{ppp}^{\delta V}$. Such compensation will occur if

$$\frac{g^h_{PPP}}{16\pi\alpha'_P}\,\frac{\sigma^{PN}_{tot}}{\sqrt{\sigma^{NN}_{tot}}}\left|\ln\frac{g^h_{PPP}}{16\pi\alpha'_P}\right| = 1. \qquad (18)$$

Condition (18) holds within the accuracy to uncertainties of g^h_{ppp} value. At the same time it follows from (1) that the shift $\Delta^{\delta\gamma}$ cannot be compensated by triple-Pomeron contribution. Therefore at $e^2 \cdot \ln(s/s_0) \ll 1$ the effective $\alpha_P(0) > 1$ due to electromagnetic interaction contributions.

It is interesting that the shadowing connected with vacuum cuts does not change the situation. This statement can be most easily checked in four-dimensional transversed momentum space and then continued over space dimension similar to /2/. Let us consider at first the shadowing of interference contribution Δ^{int}. After inclusion of the triple-Pomeron coupling the electromagnetic correction may increase due to changing the Pomeron Green function. To the first order in Δ^{int} the electromagnetic contribution to the Green function G is

$$\delta G = \Delta^{int} G^2(\ell, k^2)$$
$$G = \frac{\beta(\ell)}{\omega + k^2 R^2(\ell)} \qquad (19)$$

where $l = \ln(1/\omega_m)$; $\omega_m = \max(\omega, k^2)$ (all the notations here and in what follows are taken from Ref. /2/). But taking into account shadowing of type similar to Fig.7

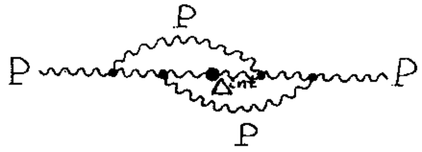


Fig. 7

changes (19), and at $k^2 = 0$ it takes the form

$$\delta G = \frac{\Delta^{int}}{\omega^2} . \qquad (20)$$

Thus the contribution of shadowing is cancelled. The validity of (20) can be easily seen from the fact that shadowing of Δ^{int} coincides with Γ_0' shadowing, for which one has $\beta^2\Gamma'/\Gamma_0'$, where Γ_0' and Γ' are the input and renormalized vertices for particle emission from the Pomeron /2/.

It is clear that the same reasoning is applicable to the double photon exchange as well, i.e., the electromagnetic effects in the case of strong Pomeron coupling lead to increase of total hadron cross-sections.

2. Weak Pomeron coupling /1/. In this case $g^h_{ppp}(0) = 0$. The problem of estimating radiative corrections to g_{ppp} becomes more uncertain, as the condition $g^h_{ppp}(0) = 0$ means the well known fact, that in calculating it one can not use individual diagrams. The situation with the problem of Δ^{int} influence on vacuum pole intercept is similar to the case of strong Po-

meron coupling. It is unclear, whether this contribution must be separated from the strong interaction one. Nevertheless, the pure electromagnetic contribution $\Delta^{\gamma\gamma}$ is of another nature and changes asymptotics.

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