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SINGLE AND DOUBLE POMERON EXCHANGE: FACTORIZATION BREAKING

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It is argued that the factorization breaking observed recently in the diffractive and non-diffractive dijet production at FERMILAB can be naturally explained if diffractive dissociation is considered to be a consequence of a non-uniform absorption of the incident particle wave.

1. Recent measurements of production of two large transverse momentum jets in the central rapidity region at Fermilab energies [1] revealed strong breaking of the Regge Factorization between the "no-gap", "single-gap" and "two-gap" cross-sections. If the presence of the rapidity gap is interpreted as the colour singlet exchange and its absence as the octet exchange, factorization implies that

$$\frac{\sigma_2}{\sigma_1} = \frac{\sigma_1}{\sigma_0} \tag{1}$$

where σ_0, σ_1 and σ_2 denote cross-sections for "no-gap", "single-gap" and "two-gap" events, respectively. Experiment shows that the ratio on the L.H.S of (1) is about 5 times larger than that on the R.H.S.

In the present note I report the results we have obtained on this subject together with Robi Peschanski [2]. We argue that the origin of this effect is basically the same as that observed in measurements of diffractive (virtual) photon-induced and hadron-induced processes [4], i.e., rescattering (absorption) in the final state of the hadron-induced diffractive dissociation [5][a].

The absorptive interpretation of diffractive dissociation was formulated by Good and Walker [6] and we are following this approach. The first step is to expand the incident state into a superposition of "diffractive" (eigen)states

$$|h>= c_1|\psi_1> + c_2|\psi_2> + ... \tag{2}$$

which (by definition) are eigenstates of absorption:

$$T|\psi_i>= \lambda_i|\psi_i> + \omega_i|\phi_i> \tag{3}$$

where $T = 1 - S$ is the absorption operator and thus λ_i is the absorption coefficient[b] of the diffractive state $|\psi_i>$. $|\phi_i>$ describes the inelastic states which are at the origin of absorption. These are multiparticle states which do not show any rapidity

[a] Similar conclusions were reached in [3] where also detailed numerical estimates are given.
[b] At high energy λ_i is dominantly real and positive.

gap[c]. Unitarity of the S-matrix implies

$$|1-\lambda_i|^2 = 1-|\omega_i|^2 = 1-\sigma_i^{nondiff}. \tag{4}$$

When averaged over all diffractive states this relation gives

$$2Re[\lambda] = \sigma^{nondiff}+\sigma^{diff}; \quad |1-\lambda|^2 = 1-\sigma^{inel} \tag{5}$$

where we have denoted $<\lambda>$ by λ and used the relation $<|\lambda|^2>=\sigma^{diff}$ [6].

The physical content of the expansion (2) depends, of course, on the physical meaning one ascribes to the "diffractive" states $|\psi_i>$. Following [5] we take them as states with a fixed number and transverse position of partons [8,9].

2. Let us recall briefly the argument given in [5] where the (single) diffractive production of two large transverse momentum jets was discussed. Since we are dealing with *hard* diffraction, it is reasonable to assume that the expansion (2) is quasi-diagonal. Denoting by ϵ the probability of fluctuation of the incident state $|\psi>$ into the final state $|\psi'>$ one can show [7] that the inelastic transition can be (to the first order in ϵ) written as

$$<\psi'|T|\psi> = \epsilon\,(<\psi'|T|\psi'> - <\psi|T|\psi>) \tag{6}$$

i.e. it is proportional to the *difference of the elastic amplitudes in the final and in the initial state.*

For the transition of γ^* into two jets we have

$$|\gamma^*> = |0> + \epsilon|D^*>; \quad |JJ> = -\bar{\epsilon}|0> + |D^*> \tag{7}$$

where $|D^*>$ represents the parton state consisting of a hard dipole and a bunch of soft partons, while $|0>$ is the state with no partons whatsoever.

For the incident proton we have

$$|p> = |g> + \epsilon|g+D^*>; \quad |p+JJ> = -\bar{\epsilon}|g> + |g+D^*> \tag{8}$$

where $|g>$ is the parton state which represents proton (up to order of ϵ, of course).

Using (2), (3) and (6) we obtain

$$\begin{aligned} <jj|T|\gamma^*> &= \epsilon <D^*|T|D^*>; \\ <p+JJ|T|p> &= \epsilon\,(<g+D^*|T|g+D^*> - <g|T|g>) \end{aligned} \tag{9}$$

To calculate the difference of the elastic amplitudes in (9) we assume that the two objects in the final state scatter independently. This gives

$$<g+D^*|T|g+D^*> - <g|T|g> = <D^*|T|D^*>(1- <g|T|g>) \tag{10}$$

When this is introduced in (9) one immediately sees that the amplitude for the process with incident protons is reduced by the factor $(1-<g|T|g>) = (1-\lambda_p)$, where λ_p is the elastic *pp* amplitude. It was shown in [5] that this factor does account for the observed factorization breaking.

3. To discuss factorization breaking at FERMILAB we consider the transitions $p\to p+C$ and $p\to p+C^*$ where (C) denotes the centrally produced system

[c]The first term in (3) is usually interpreted as "Pomeron exchange" where no colour is exchanged between $|\psi_i>$ and the target. The second term would then describe all colour exchanges and thus $|\omega_i|^2$ represents the probability of this colour-exchange interaction.

(at rapidity y) of two large transverse momentum jets and a "background" of the soft particles nearby and $(C^*) = (C)$ + additional particles which fill the rapidity between p and (C).

Now we have to write down the expansion into the diffractive states. Since we are discussing *hard* diffraction, we again assume that the expansion is quasi-diagonal and write

$$\begin{aligned} |p> &= |g> + \epsilon_P|g+D> + \epsilon|g+D^*> \\ |p+C> &= -\bar{\epsilon}_P|g> + |g+D> + \epsilon'|g+D^*> \\ |p+C^*> &= -\bar{\epsilon}|g> - \bar{\epsilon}'|g+D> + |g+D^*> \end{aligned} \tag{11}$$

where ϵ, ϵ_P and ϵ' are small (and thus will be kept only up to first order)[d]. As before, $|g>$ is a superposition of partonic states representing proton. (D) is a superposition of partonic states consisting of a *hard* (i.e. small in transverse space) dipole at rapidity y and a number of soft partons with rapidities nearby (this system eventually decays predominantly into two large transverse momentum jets and the "background" of soft particles). $(D^*) = (D)$ + soft partons at rapidities between the projectile and (D).

Using (3), (11) and (6) we thus obtain

$$\begin{aligned} <p|T|p> \equiv \lambda_p = \lambda_g; \quad <p+C|T|p> &= \epsilon_P\ (\lambda_{(g+D)} - \lambda_g) \\ <p+C^*|T|p> &= \epsilon\ (\lambda_{(g+D^*)} - \lambda_g), \end{aligned} \tag{12}$$

where all quantities on the R.H.S. are to be understood as superpositions of the corresponding quantities defined for truely diffractive parton states.

Following [5] we now assume that $\lambda_{(g+D)}$ and $\lambda_{(g+D^*)}$ can be calculated assuming independent scattering of their components. When inserted into (12), this leads to

$$\begin{aligned} <p+C|T|p> &= \epsilon_P\lambda_D(1-\lambda_g) \\ <p+C^*|T|p> &= \epsilon\lambda_{D^*}(1-\lambda_g), \end{aligned} \tag{13}$$

To obtain the cross-sections for one-gap and two-gap events, we have to square the corresponding amplitudes and sum over the final hadronic states. The result is

$$\begin{aligned} \sigma_2(\vec{s},y) &= |\epsilon_P(\vec{s},y)|^2\sigma_D^{diff}(\vec{b}-\vec{s},Y-y)\left[1-\lambda_p(\vec{b},Y)\right]^2 \\ \sigma_1(\vec{s},y) &= |\epsilon(\vec{s},y)|^2\sigma_{D^*}^{diff}(\vec{b}-\vec{s},Y-y)\left[1-\lambda_p(\vec{b},Y)\right]^2 \end{aligned} \tag{14}$$

We have used here the general formulae of Good and Walker [c.f. (4) and (5)]:

$$[\lambda_g(\vec{b},Y)]_{av} = \lambda_p; \quad \left[|\lambda_D|^2\right]_{av} = \sigma_D^{diff}; \quad \left[|\lambda_{D^*}|^2\right]_{av} = \sigma_{D^*}^{diff} \tag{15}$$

where σ^{diff} is the diffractive cross-section and λ_p is the elastic *pp* amplitude.

4. The cross-section for one-gap events, σ_1, given in (14), can be also calculated as a product of (i) the probability of fluctuation of the target into $(g+D)$, i.e. $|\epsilon_P|^2$, (ii) probability of *non-diffractive* interaction of (D) with the projectile, i.e. $\sigma_D^{nondiff}$ and (iii) probability that no inelastic interaction of the final antiproton

[d] The relations between expansion coefficients follow from orthonormality of the states.

with the projectile took place, i.e. $1-\sigma^{inel}=(1-\lambda_{\bar{p}})^2$ (c.f. (5)). Thus we can write

$$\sigma_1(\vec{s},y) == \left|\epsilon_P(\vec{b}-\vec{s},Y-y)\right|^2 \sigma_D^{nondiff}(\vec{s},y)\left[1-\lambda_p(\vec{b},Y)\right]^2 \tag{16}$$

Comparing (16) and (14) we have

$$|\epsilon_P(\vec{b}-\vec{s},Y-y)|^2\sigma_D^{nondiff}(\vec{s},y)=|\epsilon(\vec{s},y)|^2\sigma_{D^*}^{diff}(\vec{b}-\vec{s},Y-y) \tag{17}$$

This equation has a simple physical meaning. It says that the probability of a fluctuation of the projectile into the projectile+gap+D configuration is proportional to the diffractive interaction in the projectile-D system, whereas the probability of fluctuation into no-gap configuration is proportional to the non-diffractive interaction in this system. It thus provides the interpretation, in the Good-Walker language, of the "Pomeron-exchange" mechanism for gap creation.

5. To discuss the no-gap events we have to go beyong the Good-Walker argument, as the latter refers only to diffractive interactions. To this end we observe that the probability of a no-gap event to occur can be written as the product of (i) the probability of fluctuation of the projectile into $|g+D^*>$, i.e. $|\epsilon|^2$, and (ii) the probability of *nondiffractive* interaction of the central system (D^*) with the target, i.e. $\sigma_{D^*}^{nondiff}$:

$$\sigma_0(\vec{s},y)=|\epsilon(\vec{s},y)|^2\sigma_{D^*}^{nondiff}(\vec{b}-\vec{s},Y-y). \tag{18}$$

Using (14), (16), (18) and (17) we thus obtain

$$\begin{aligned}
\sigma_2(\vec{s},y) &= |\epsilon_P(\vec{b}-\vec{s},Y-y)|^2\sigma_D^{diff}(\vec{s},y)\left[1-\lambda_p(\vec{b},Y)\right]^2\\
\sigma_1(\vec{s},y) &= \left|\epsilon_P(\vec{b}-\vec{s},Y-y)\right|^2\sigma_D^{nondiff}(\vec{s},y)\left[1-\lambda_p(\vec{b},Y)\right]^2;\\
\sigma_0(\vec{s},y) &= |\epsilon_P(\vec{b}-\vec{s},Y-y)|^2\sigma_D^{nondiff}(\vec{s},y)\frac{\sigma_{D^*}^{nondiff}(\vec{b}-\vec{s},Y-y)}{\sigma_{D^*}^{diff}(\vec{b}-\vec{s},Y-y)}
\end{aligned} \tag{19}$$

where we have taken advantage of the symmetry (between projectile and target) when writing the expressions for σ_2 and σ_0.

These equations show explicitely that the factorization formula (1) is violated:

$$\frac{\sigma_1(\vec{s},y)}{\sigma_0(\vec{s},y)}=\frac{\sigma_{D^*}^{diff}(\vec{b}-\vec{s},Y-y)}{\sigma_{D^*}^{nondiff}(\vec{b}-\vec{s},Y-y)}[1-\lambda_p(\vec{b},Y)]^2;\ \ \frac{\sigma_2(\vec{s},y)}{\sigma_1(\vec{s},y)}=\frac{\sigma_D^{diff}(\vec{s},y)}{\sigma_D^{nondiff}(\vec{s},y)} \tag{20}$$

One sees that the factorization violating factor $[1-\lambda_p(\vec{b},Y)]^2$ is basically the same as that which was shown in [5] to be reponsible for violation of factorization between single diffractive processes at HERA and at FERMILAB[e].

6. In conclusion, we have shown that the factorization breaking observed in difractive and non-diffractive two jet production at FERMILAB has the same origin as that observed in diffractive two-jet production at HERA and FERMILAB. The

[e] One should note, however, that its effects are sensitive to the details of the hard process in question (mostly to its impact parameter dependence [5]) and thus may be different in different processes [3].

phenomenon can be naturally explained if diffractive dissociation is understood as a consequence of the inelastic interactions ("absorption") of the incident beam.

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