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Photoproduction total cross-sections at very high energies and the Froissart bound

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Abstract
A previously successful model for purely hadronic total cross-sections, based on QCD minijets and soft-gluon resummation, is here applied to the total photoproduction cross section. We find that our model in the γp case predicts a rise with energy stronger than in the $pp/\bar{p}p$ case.

1 Introduction

In this note, we shall describe (and apply to data) a model for the total cross-section [1, 2], based on the ansatz that infrared gluons provide the saturation mechanism in the rise of all total cross-sections (thus obeying the Froissart bound), with the rise calculated through the increasing number of hard collisions between low-x, but perturbative gluons. These collisions produce low p_t partons which hadronize in so called mini-jets. We assume that for $p_t \geq 1 \div 2\ GeV$ the parton-parton cross-section can still be calculated perturbatively and set a minimum p_t cut-off, p_{tmin} in the jet cross-section calculation. To make connection with actual phenomenological inputs, the mini-jet cross-sections are calculated [3, 4] using DGLAP evolved parton densities: for the proton we have used GRV [5], MRST [6] and CTEQ [7], for the photon GRS [8] and CJKL [9]. In our model we use only LO densities, as part of the NLO effects are described by soft gluon resummation and the use of NLO would result in some double counting. Similarly, we have opted for tree level parton-parton cross-sections and one loop α_s. As the c.m. energy increases, with fixed p_{min}, these mini-jet cross-sections increase and their contribution to the total cross-section becomes larger than any observed cross-section, violating unitarity. This has resulted in discarding the mini-jet model. Embedding the mini-jet cross-section in the eikonal representation, restores unitarity, but requires modelling of the matter distribution in the colliding particles via an impact parameter distribution. Convolution of the electromagnetic form factors is frequently used and more fundamental attempts exist in the framework of Reggeon calculus and perturbative QCD. Our model focuses on very soft gluons as the source of a dynamical description of the impact factor and its energy dependence. Thus the name Bloch Nordsieck (BN) underlies the infrared region and its resummation. We shall briefly present this model, show its results for purely proton processes, and then apply it to photoproduction processes.

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2 The Bloch-Nordsieck Model (BN)

Our BN model exhibits fractal behaviours for quantities such as (i) the energy rise of the mini-jet cross-sections for which $\sigma_{jet} \approx s^{\delta}$ with $\delta \approx 0.3$ and (ii) the very low momentum single gluon emission probability which we propose to be proportional to k_t^{-p-1} with $0 < p < 1$, for gluons of transverse momentum k_t .

This model, which was initially developed for purely hadronic total cross-section, incorporates QCD inputs such as parton-parton cross-sections, realistic parton densities, actual kinematics, and soft gluon resummation. We write, for a general process,

$$\sigma_{tot}^{AB} = 2 \int d^2\vec{b}[1 - e^{-n(b,s)/2}] \tag{1}$$

with the imaginary part of the eikonal related to the average number of inelastic collision $n(b, s)$. We isolate all hard perturbatively calculated collisions into

$$n_{hard}(b, s) = A(b, s)\sigma_{jet}(p_{tmin}, s) \tag{2}$$

and phenomenologically determine the remaining collisions which we call $n_{soft}(b, s)$. At present our model is unable to make an *ab initio* calculation of this quantity, and we use a QCD inspired modelling, described in [2].

The impact parameter distribution is obtained from the Fourier transform of the soft gluon transverse momentum, resummed, distribution, namely

$$A(b, s) = \mathcal{N} \int d^2\mathbf{K}_\perp \frac{d^2P(\mathbf{K}_\perp)}{d^2\mathbf{K}_\perp} e^{-i\mathbf{K}_\perp\cdot\mathbf{b}} = \frac{e^{-h(b,q_{max})}}{\int d^2\mathbf{b} e^{-h(b,q_{max})}} \equiv A(b, q_{max}(s)) \tag{3}$$

with

$$h(b, q_{max}(s)) = \int d^3\bar{n}_g(\vec{k})[1-e^{i\mathbf{k}_\perp\cdot\mathbf{b}}] = \frac{16}{3}\int_0^{q_{max}(s)} \frac{dk_t}{k_t}\frac{\alpha_s(k_t^2)}{\pi}\left(\log\frac{2q_{\max}(s)}{k_t}\right)[1 - J_0(k_t b)] \tag{4}$$

In the above equation, we need to extend the integral to zero momentum gluons, which supply the saturation effect of resummation. One needs then an ansatz for the single soft gluon distribution in Eq. 4, namely for $\alpha_s(k_t^2)$ as $k_t \to 0$. Our model for this behaviour is inspired by the Richardson potential, but in order to have a finite result for the integral in Eq. 4 we use

$$\alpha_s = \frac{12\pi}{33 - 2N_f}\frac{p}{\ln[1 + p(\frac{k_t}{\Lambda})^{2p}]}. \tag{5}$$

This expression gives the asymptotic freedom value for large k_t , and is singular (but integrable) at $k_t = 0$ (for $p < 1$). The closer p is to 1, the more the minijet cross-sections will be quenched at any given energy.

The energy dependence of the impact function $A(b, s)$ is introduced through the upper limit of integration in Eq. 4. As amply discussed in ref. [1], the function q_{max} is obtained through an averaging over the parton densities

$$q_{max}(s) = \sqrt{\frac{s}{2}}\,\frac{\sum\limits_{i,j}\int\frac{dx_1}{x_1}\int\frac{dx_2}{x_2}\int_{z_{min}}^1 dz f_i(x_1)f_j(x_2)\sqrt{x_1x_2}(1-z)}{\sum\limits_{i,j}\int\frac{dx_1}{x_1}\int\frac{dx_2}{x_2}\int_{z_{min}}^1 dz f_i(x_1)f_j(x_2)}, \tag{6}$$

Notice that in our model, the impact parameter distribution depends on the energy and the process under consideration through the parameter $q_{\max}(s)$, which is evaluated using the given parton densities.

The BN model thus described has been applied to proton-proton scattering, obtaining a total cross-section for LHC to be $\sigma(\sqrt{s} = 14\ TeV) = (100 \pm 12)\ mb$, where the error reflects various uncertainties as in the choice of densities, minimum parton p_t cut-off and the IR behaviour of the soft gluon coupling. Our results for proton-proton and proton-antiproton scattering are shown in Fig. 1 with labelling and references defined as in [4].

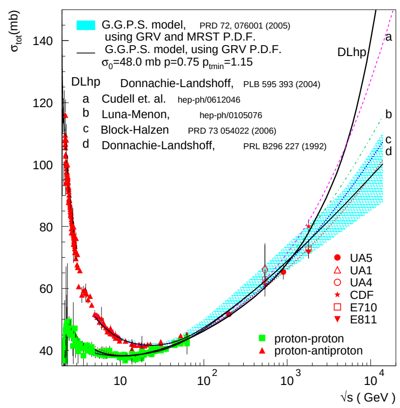

Fig. 1: Data and models for proton-proton and proton-antiproton total cross-section from ref. [4].

3 Photon processes and the total γp cross-section at high energies

Application to photons requires the probability that a photon behaves like a hadron. One possibility is to use Vector Meson Dominance (VMD) in the eikonal representation, as in [10, 11],

$$\sigma_{tot}^{\gamma p} = P_{\gamma \to hadron} \sigma_{tot}^{\gamma_{had} p} = 2 P_{had} \int d^2\vec{b} [1 - e^{-n(b,s)/2}] \tag{7}$$

with $P_{\gamma \to hadron} = 1/240$. As for the proton case, the average number of inelastic collisions, $n(b, s)$, is split between hard collisions calculable as QCD minijets, and a soft part. Hence, the average number of collisions is written as

$$n(b, s) = n_{soft}(b, s) + n_{hard}(b, s) = \frac{2}{3} n_{soft}^{pp}(b, s) + A(b, s)\sigma_{jet}(s)/P_{had} \tag{8}$$

with n_{hard} including all outgoing parton processes with $p_t > p_{tmin}$. The jet cross-sections are calculated using actual photon densities, which themselves give the probability of finding a given quark or gluon in a photon, and thus P_{had} needs to be canceled out in n_{hard}. For the soft part, a good description is obtained with $n_{soft}^{pp}(b,s)$ being the same as in the pp case[4]. The impact function $A(b,s)$ again supplies saturation and is calculated using photon and proton densities in Eq. 6. Once this energy parameter has been calculated, $A(b,s)$ is fully determined. More fundamental attempts to obtain the impact function for photons can be found in ref. [12].

For γp we show in Fig. 2 both the saturation parameter q_{max} plotted as a function of the γp c.m. energy, as well as the resulting impact parameter function at four representative energies. Unlike other models, based on the convolution of the form factors of the colliding particles, the impact function in the BN model is energy dependent, with a shape shrinking with energy. Both

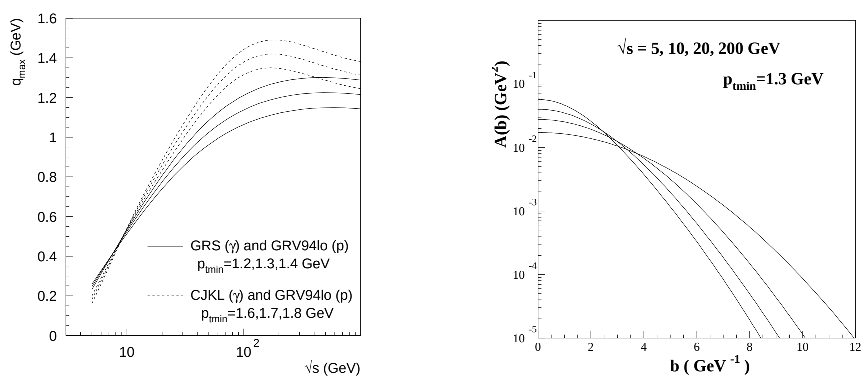

Fig. 2: Left: average maximum tramsverse momentum allowed to single soft gluon emission in γp scattering. Right: Impact parameter distribution for γp scattering at various c.m. energies

the mini-jet cross-section and the impact function (the latter through q_{max}) for γp depend on the set of Parton Density Functions (PDFs). We choose GRV94 for the proton, GRS and CJKL for the photon, input them into the eikonal and compare the results with HERA data [13, 14], including a set of ZEUS BPC data extrapolated from $\gamma^* p$ to $Q_\gamma^2 = 0$ [15]. These results for the BN model are shown in Table 1 for different parameters sets chosen so that all of HERA data are included in a band defined by the last two columns in table 1.

In ref. [16], we have observed that at very high γp energies, our results indicate a faster rise than is the case for proton inspired models. For energies beyond HERA, $q_{\max}(s)$, computed through the photon densities, no longer increases (unlike the proton case) thus blocking saturation earlier than for protons. As a result, since the mini-jet cross-sections keep on increasing, the photon cross-sections, past present accelerator energies, would grow faster than the purely hadronic ones. We have noted [16] that this prediction from our model finds independent support in the fit by Block and Halzen [17] which gives results close to ours in the very high energy region.

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Table 1: Values (in mb) for total cross-section for γp scattering evaluated in the c.m. energy of colliding particles, for different parameter sets

$\sqrt{s}$ GeV	EMM with Form Factors $p_{tmin} = 1.5$ GeV	BN_γ model GRS, p=0.75 $p_{tmin} = 1.2$	BN_γ model CJKL, p=0.8 $p_{tmin} = 1.8$	BN_γ GRS, p=0.75 $p_{tmin} = 1.15$
5	0.116	0.116	0.116	0.116
11.46	0.114	0.115	0.114	0.1155
48.93	0.122	0.130	0.121	0.132
112.14	0.139	0.155	0.140	0.16
478.74	0.238	0.228	0.203	0.236
1097.3	0.352	0.279	0.250	0.289
4684.6	0.635	0.384	0.338	0.395
10736.8	0.829	0.449	0.390	0.461
20000	0.985	0.499	0.429	0.512

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