THE FLAVOURING OF THE POMERON AND REGGEON FIELD THEORY WITH THRESHOLDS

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The flavoured pomeron is introduced into reggeon field theory to show that, in the absence of a secondary vacuum trajectory P′, the critical solution is in fair agreement with experimental data.

A recent detailed model [1] for non-diffractive renormalization of the pomeron intercept due to the production of strangeness, baryon number and other "flavours", has been able to reproduce available experimental data. On the other hand it has been found [2,3] that the introduction of rapidity thresholds for the manifestation of pomeron effects leads to a regularization of reggeon field theory (RFT) graphs, making it possible to evaluate their relevance – related to diffractive and absorptive corrections – at finite rapidities.

The purpose of this note is to introduce a simplified model for the flavoured pomeron into RFT, identifying the above mentioned thresholds with those for flavour production as suggested in refs. [1] and [4]. It is possible in this way to examine the consistency of the critical RFT solution [5] with experimental data. It turns out that the critical pomeron, without a secondary P′ trajectory, correctly reproduces pp total cross section (except for cosmic ray data, where the prediction is too low), $\bar{p}$ multiplicity, new data for pp → pX inclusive spectrum [6] and forward real to imaginary ratio for pp scattering (again with low predictions in the ISR range). While the general agreement with a critical pomeron is good, because of the mentioned discrepancies and several simplifying assumptions of the model, a slightly supercritical solution [7] of RFT can also be accepted.

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A simple model for the pomeron contribution to pp scattering may be written through the inverse propagator as

$$\mathrm{i}\Gamma(E,k^2) = \mathrm{i}\hat{\Gamma}(E,k^2) + \delta - \Sigma(E,k^2)\,, \qquad (1)$$

related to the elastic amplitude by

$$A(Y,t) = \frac{\beta(t)}{2\pi\mathrm{i}}\int_{c-\mathrm{i}\infty}^{c+\mathrm{i}\infty} d(-E)\,\frac{\mathrm{i}}{\Gamma(E,k^2)}\,\mathrm{e}^{(1-E)Y}\,. \qquad (2)$$

As usual, the 'energy' and 'momentum' of RFT are defined in terms of angular momentum and momentum transfer $E = 1 - J$, $k^2 = -t$. The rapidity is taken as $Y = \ln s$ and

$$\mathrm{i}\hat{\Gamma}(E,k^2) = (E - \hat{\Delta} - \alpha_0' k^2)\exp[-Y_0(E - \hat{\Delta} - \alpha_0' k^2)]\,, \qquad (3)$$

which corresponds to an 'embryonic' pomeron pole with intercept $\hat{\alpha} = 1 - \hat{\Delta}$ and slope α_0' for $Y > Y_0$. The counterterm δ can be nonperturbatively calculated for the critical RFT solution in terms of the previous parameters and the triple-pomeron coupling r_0. The 'self-mass' correction Σ, when evaluated through the renormalization group techniques, gives a relevant contribution to A only at high rapidity values [3,4]. Therefore, even though the complete expression eq. (1) corresponds to a very complicated J-plane structure, its effect will be equivalent, after Laplace transforming, to the following more intuitive view. For $Y_0 < Y < 2Y_0$ the relevant propagator is the embryonic pomeron one, eq. (3), which gives a contribution to the

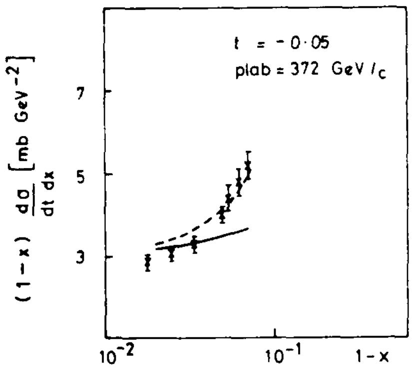

Fig. 1. Inclusive distribution pp → pX. Solid line is triple-embryonic pomeron coupling eq. (10). Dashed line includes ω and π exchange. Data are from ref. [6].

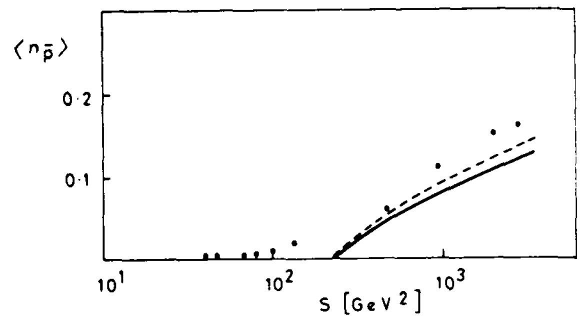

Fig. 2. Average $\bar{p}$ multiplicity. Solid line corresponds to the critical value $\delta = 0.125$ and dashed line to the supercritical case $\delta = 0.14$. Data are from ref. [17].

total cross-section combination

$$\sigma_{tot} = [\sigma_{tot}(\bar{p}p) + \sigma_{tot}(pp)]/2 = A(Y, t=0)/s\ ,$$

of the form

$$\sigma_{tot} = \beta(0)\, e^{-\hat{\Delta} Y}\ , \quad Y_0 < Y < 2Y_0\ . \tag{4}$$

For $2Y_0 < Y < \bar{Y}$, where $\bar{Y}$ is a high-rapidity value above which contributions coming from Σ are important, the effective inverse propagator is

$$\mathrm{i}\Gamma_0(E, k^2) = \mathrm{i}\hat{\Gamma}(E, k^2) + \delta\ , \tag{5}$$

corresponding to a real pole with intercept α_0 (bare pomeron from the point of view of RFT) and numerically not too important complex poles [1]. Through the Laplace transform one obtains

$$\sigma_{tot} = \beta(0)\, e^{-\hat{\Delta} Y}[1 + \delta(Y - 2Y_0) + \tfrac{1}{2}\delta^2 (Y - 3Y_0)^2 \theta(Y - 3Y_0) + \ldots] \tag{6a}$$

$$\approx \beta' e^{(\alpha_0 - 1)Y}\ , \quad 2Y_0 < Y < \bar{Y}\ . \tag{6b}$$

Finally, for $Y > \bar{Y}$, the critical solution is well approximated by the asymptotic behaviour

$$\sigma_{tot} = a Y^{-\gamma}\ , \quad Y > \bar{Y}\ . \tag{7}$$

We adopt here the Chew–Rosenzweig scheme [8] in which f does not exist as a trajectory (P′) independent from the pomeron. This choice is done to emphasize the rise of σ_{tot} due to the bare pomeron α_0. The inclusion of the decreasing P′ contribution requires a comparatively higher value of α_0 which would clearly be in the supercritical region [9].

The effective renormalization $\alpha_0 - \hat{\alpha}$ which appears at $Y > 2Y_0$ can be physically related to a threshold effect for the production of one, two etc. baryon–antibaryon pairs in correspondence to the terms of eq. (6a). The fact that this threshold is twice the one, Y_0, for the manifestation of the embryonic pomeron is a simplification of the model. Indeed Y_0 is a sort of average between the low threshold for a π ladder and the higher one for $K\bar{K}$ production. The oversimplification of the low-energy thresholds does not allow the application of the model as it stands to πp and Kp scattering; more parameters are needed to include these cases [1,10].

Starting our analysis of data we fix $2Y_0$ in correspondence with the minimum of σ_{tot}, i.e. $Y_0 \simeq 2.7$. Consequently, to reproduce σ_{tot} for $Y < 5.4$ with the embryonic pole eq. (4) we must choose $\hat{\alpha} \simeq 0.93$ and $\beta(0) \simeq 57.5$ mb.

If the RFT solution is the critical one, the strength of baryon pair production δ must be related to the triple-pomeron coupling. In fact, from the requirement that $\Gamma(0,0) = 0$, eq. (1) gives

$$\delta = \hat{\Delta} e^{Y_0 \hat{\Delta}} + \int_{-\infty}^{0} \left\{ [1 - Y_0(E' - \hat{\Delta})] e^{Y_0 \hat{\Delta}} - \frac{(1 - Y_0 E')}{Z(E')} \right\} \times e^{-Y_0 E'} \mathrm{d}E'\ , \tag{8}$$

where Z is defined through the normalization condition

$$\left.\frac{\partial \mathrm{i}\Gamma}{\partial E}\right|_{E=-E_N,\ k^2=0} = \frac{1}{Z}(1 + Y_0 E_N)\, e^{Y_0 E_N}\ , \tag{9}$$

and can be worked out by means of the one-loop approximation of the renormalization group functions as done in refs. [3,4]. Thus eq. (8) can be numerically

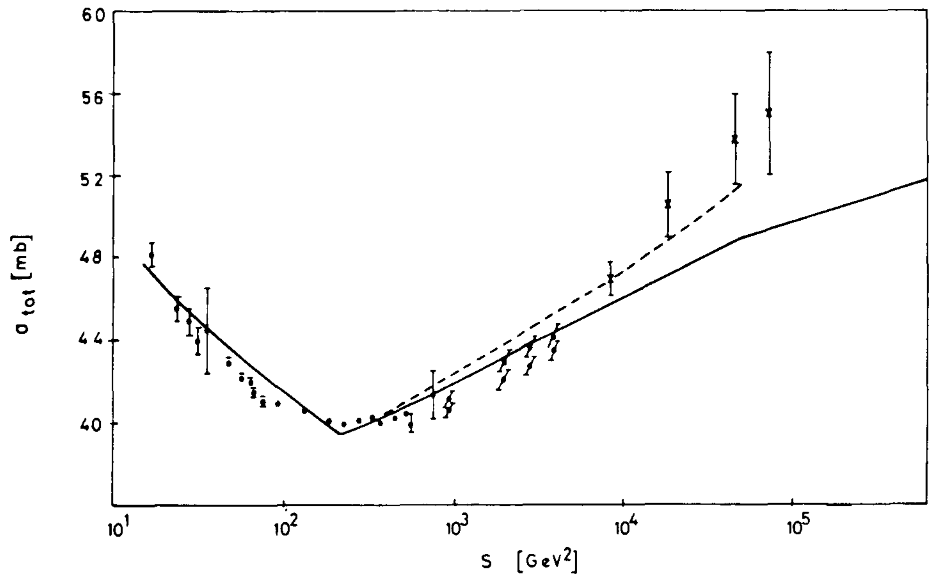

Fig. 3. Total cross section combination $(\sigma_{pp} + \sigma_{\bar{p}p})/2$. Solid and dashed lines correspond to critical (eqs. (4), (6b) and (7)) and supercritical cases respectively. Data points • are from ref. [18] and data points x from ref. [15].

evaluated in terms of the experimental parameters $\alpha_0' \simeq 0.3\ \mathrm{GeV}^{-2}$ and $r_0 = 0.5\ \mathrm{GeV}^{-1}$. Regarding this last controversial coupling, it is found that recent data [6] on pp → pX inclusive cross section for $1-x < 0.05$ are correctly reproduced by the triple-embryonic pomeron formula

$$(1-x)\frac{\mathrm{d}\sigma}{\mathrm{d}t\,\mathrm{d}x} = \frac{\sqrt{2}r_0}{16\pi}\beta(0)^{3/2}\frac{(1-x)^{\hat{\Delta}}}{s^{\hat{\Delta}}} \times \exp\left[5 - 2\alpha_0' \ln(1-x)\right]t\,, \qquad (10)$$

with the above values of the parameters (see fig. 1), which is in agreement with the experimental decrease with increasing p_{lab}. For larger values of $1-x$, corrections due to the exchange of ω and π become relevant. The former is calculated through a formula analogous to eq. (10) (taking into account that $[\sigma_{tot}(\bar{p}p) - \sigma_{tot}(pp)]/2$ is well reproduced by $40.8/e^{0.6Y}$) with $r_{\omega\omega p} \simeq 1\,\mathrm{GeV}^{-1}$ and the latter is estimated according to the parametrization of ref. [11]. The contribution of π-exchange at $1-x \approx 0.07$ is ~20% and explains the experimental difference between pp → pX and dp → dX as was suggested in ref. [12].

Having gained confidence in the value of r_0, its use in eq. (8) gives $\delta \approx 0.125$. This is probably an upper bound for the critical δ because, even if other data could suggest a higher value of r_0, the triple-flavoured pomeron (smaller than the embryonic coupling) would be required in the calculation of Z. With the above value of δ eq. (5) determines $\alpha_0 \simeq 1.04$ which agrees with the critical bare intercept calculated by several different methods [3,13]. The use of eq. (6a) and the assumption [1] $\langle n_{\bar{p}}\rangle \simeq \frac{1}{2}\langle n_{\bar{B}}\rangle$ determines the $\bar{p}$ multiplicity shown in fig. 2.

We now may predict σ_{tot} for all the experimental range (see fig. 3) using eqs. (6b) and (7). The validity of the asymptotic expression is assumed [3] for $Y \gtrsim 4Y_0$, and the critical exponent is taken as $\gamma = -0.26$ according to the most refined evaluation [14]. It is seen that the prediction is good except for the cosmic-ray range where it is too low compared with data [15].

Notice that for schematic purposes, Σ has been introduced only for $Y > \bar{Y}$. Had we included it from its threshold $3Y_0$, the predicted rise of σ_{tot} would have been slightly smoother [3] (about 0.5 mb less at $Y \sim 4Y_0$). This, and the assumed absence of P', gives a somehow maximum rise for the critical pomeron evaluated in the present framework.

A related problem appears when calculating the forward real-to-imaginary ratio of pp scattering ρ. Using the dispersion relation

$$\mathrm{Re}\,A(s) = (2s^2/\pi)P\int \mathrm{d}s'\,\sigma_{tot}/(s'^2 - s^2)\,, \qquad (11)$$

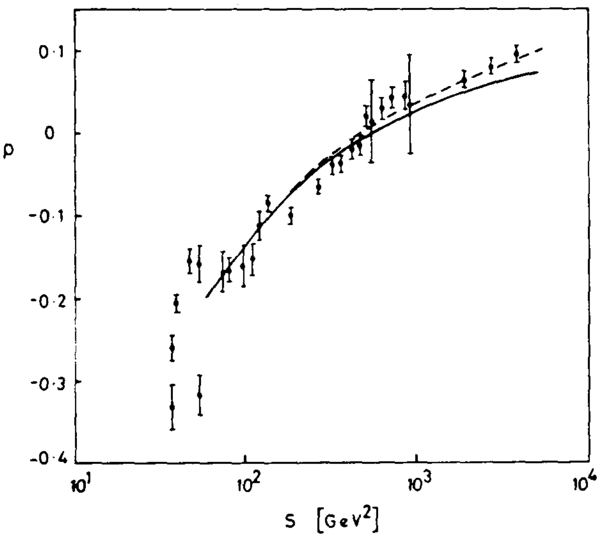

Fig. 4. Forward pp real-to-imaginary ratio $\mathrm{Re}\,A_{pp}/\mathrm{Im}\,A_{pp}|_{t=0}$. Solid and dashed lines correspond to critical and supercritical case respectively, calculated with eq. (11) and ω contribution. Data are from refs. [16] and [19].

with eqs. (4), (6b) and (7) and including the ω contribution, ρ is obtained as shown in fig. 4. Due to the small total cross sections in the cosmic-ray range, too low values of ρ are predicted in comparison with ISR data [16].

These facts suggest a higher value of α_0. Taking $\delta \simeq 0.14$, still compatible with $\langle n_{\bar{p}} \rangle$, and from which the bare pomeron intercept turns out to be $\alpha_0 \simeq 1.05$, one obtains improvement both for σ_{tot} in the cosmic-ray range and for ρ in the ISR region. This value of α_0 corresponds to a supercritical intercept in the present simplified model. It is clear that different assumptions in the calculation may cause a variation of 0.01 in the critical intercept. Moreover, contributions to $A(Y, t)$ of three- and four-pomeron Green functions have been deliberately left out to avoid additional coupling parameters. Even though their rapidity dependence seems to compensate in an eikonal approximation [9], their strongly model-dependent relevance could also account for the difference between the two fits shown in this note.

In conclusion, the present model indicates the fair agreement of a simple flavoured critical pomeron with data, its phenomenological difference with a slightly supercritical theory being a matter of more refined details.

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