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Odderon Contribution in U-matrix Method

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Abstract

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The behaviour of the differences of the proton-proton and proton-antiproton cross-sections is studied with the help of the U-matrix method in the presence of the C-odd partner of pomeron.

Аннотация

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В методе U-матрицы исследуется поведение разностей протон-протонных и протон-антипротонных сечений при наличии вклада нечетного партнера померона.

1 Introduction

One of the most important results obtained in the expirements on hadron-hadron scattering is the discovery of the continuing growth of the total cross-sections in the high-energy region . The data recieved at IHEP , CERN-ISR , FNAL and SppS-colliders suggests that the maximal allowed growth of the total cross-sections , saturating the Froissart bound , is possible :

$$\sigma_{tot} \sim \log^2 s, \tag{1}$$

at $s \to \infty$.

Besides that , the measurement of the differential cross-sections has shown that at $-t = Q^2 \simeq 1 \div 2$ the differential $p\bar{p}$-cross-section is by an order of magnitude higher than the proton-proton one and has the form of a "shoulder " instead of a dip . The recent experiments at SppS reveal that the geometrical scailing is violated: the ratio of the elastic to the total cross - sections σ_{el}/σ_{tot} , which remained constant for a long time , increases from 0.175 at $\sqrt{s} = 53$ to 0.215 at $\sqrt{s} = 900$. Due to this fact the modern models take into account not only pomeron , which leads to the asymptotic equality of $pp-$ and $p\bar{p}$-cross-sections , but also its crossing-odd partner named "odderon" . Since the last one dominates in the region of high t , its presence must lead to different dependence of particle-particle and particle-antiparticle reactions on the scatteing angle ; moreover , it may provide the constancy and even the growth of the difference of the total $pp-$ and $p\bar{p}-$cross-sections $\Delta\sigma_{tot} = \upsilon_{tot} - \sigma^a_{tot}$. For instance , in the scenarios of the "maximal" odderon it grows with the maximal rapidity allowed by the general theory of scattering :

$$\Delta\sigma_{tot} \sim \log s. \tag{2}$$

Lately the interest in the models providing the maximal growth of $\Delta\sigma_{tot}$ has been considerably renewed since they predict the increase of the ratio of the real to the imaginary part of the forward amplitude $\rho(s,0) = \frac{ReF(s,0)}{ImF(s,0)}$ — the effect recently reported by UA4 Collaboration [1] .

However , it has been argued in [2] that the account of the amplitude unitarity in the framework of the eikonal parametrization permits to achieve at most constant , but not growing , $\Delta\sigma_{tot}$. This constancy is obtained only when the leading pomeron and odderon trajectories are degenerate . If the pomeron lies above the odderon , $\Delta\sigma_{tot}$ vanishes as s^{β} , $\beta < 0$. Certainly , it is interesting to find out if it is possible to achieve the maximal odderon in another parametrizations of the scattering amplitude . For example , the approach using the generalized reaction matrix preserves the unitarity of the amplitude [3] and is very convenient for the explicit treatment of our problem . In this work we shall calculate the total and elastic cross-sections with the help of the U-matrix formalism ; after that we shall make a comparison with the results of the eikonal representation . It will be shown that both approaches give the coinciding predictions about the behaviour of the total cross-sections at $s \to \infty$. However , the U-matrix leads to a constant difference of the elastic cross-sections $\Delta\sigma_{el}$ while the eikonal $\Delta\sigma_{el}$ and $\Delta\sigma_{inel}$ grow with the same rate (as $\log s$) , but have opposite signs , and their contributions vanish.

2 General properties of the U-matrix

In the high-energy region the scattering amplitude F(s,t) can be written with the help of the generalized reaction matrix (U-matrix) , arising from the single-time formulation of the two-body problem in QFT [4] . The essential feature of this theory is the introduction of the quasipotential , which is the analogue of the usual quantum-mechanics potential for the relativistic case . Unlike the Bethe-Salpeter equation , the choice of boundary conditions using the quasipotential can be made in a rather simple way ; moreover , due to the non-conservation of the number of the particles in quantum field theory , the effective potential is complex , and its imaginary part must have definite sign to satisfy the untarity conditions . There exists a simple relation between the analytical properties of the quasipotential and those of the scattering amplitude , which allows first to make the analytical continuation from $s-$ to $t-$channel for the quasipotential and after that go over to the analytical continuation of the amplitudes without any extra restrictions on the behaviour of Regge trajectories .

According to [5, 6] , in the centre-of-mass system we can write the following singletime equation:

$$F(\vec{p},\vec{q}) = U(\vec{p},\vec{q}) + \frac{i\pi\rho(s)}{8}\int d\Omega_{\vec{k}} U(\vec{p},\vec{k})F(\vec{k},\vec{q}) \ , \tag{3}$$

where $F(\vec{p},\vec{q})$ is the scattering amplitude , $p = k = q$, $\rho(s) = \sqrt{\frac{s-4m^2}{s}}$. The U-matrix , just like the amplitude , contains all the nessesary information about the scattering process . It can be shown that non-Hermitian part of the U-matrix appears due to the presence of inelastic channels ; if we introduce the symbolic notations , in which Eq.3 is written as

$$F = U + \mathrm{i}UDF, \tag{4}$$

and the amplitude unitarity condition as

$$F - F^{+} = 2\mathrm{i}F^{+}DF + 2\mathrm{i}\mathrm{H} \tag{5}$$

(here H is the inelastic contribution) , then

$$U - U^{+} = 2\mathrm{i}(1 + \mathrm{i}U^{+}D)\mathrm{H}(1 - \mathrm{i}DU). \tag{6}$$

Passing over to the impact parameter space we get

$$F(s,t) = \frac{s}{\pi^2}\int_0^\infty bdb \frac{U(b,s)}{1 - iU(b,s)} J_0(\mathrm{b}\sqrt{-\mathrm{t}}) \tag{7}$$

and

$$\sigma_{inel}(s,t) = 8\pi\int_0^\infty bdb \frac{Im\mathrm{U}(\mathrm{b},\mathrm{s})}{|1 - \mathrm{i}\mathrm{U}(\mathrm{b},\mathrm{s})|^2} J_0(\mathrm{b}\sqrt{-\mathrm{t}}), \tag{8}$$

where U(b,s) is the Fourier-Bessel transform for U(s,t) :

$$U(b,s) = \frac{\pi^2}{s}\int_0^\infty \sqrt{-t}d\sqrt{-t}U(s,t)J_0(\mathrm{b}\sqrt{-\mathrm{t}}). \tag{9}$$

To close this section let us say a few words about the description of Regge poles in this approach . Suppose that not F(s,t) but U(s,t) is proportional to $s^{\alpha}(t)$, where $\alpha(t)$ is the pole trajectory in the complex angular momentum plane . A crossing-even (crossing-odd) pole $\alpha_{\pm}(t)$ is thus related to the matrix

$$U_{\pm}(s,t) = -g_{\pm}\xi_{\pm}(\mathrm{t})\mathrm{s}^{\alpha_{\pm}(\mathrm{t})} , \tag{10}$$

where $g_{\pm} > 0$ are real constants and $\xi_{\pm}(t)$ are signature factors . Such choice makes the conditions of the amplitude unitarity explicit just from the beginning and helps to avoid procedures of the additional unitarisation for superhigh energies . If it is nessesary to take account of several Regge poles , the resulting U-matrix must be the sum of the U-matrices for the individual poles .

3 Pomeron and odderon in the U-matrix method

For a long time , just till the early 70's , the collider data seemed to demonstrate the tendency of the particle-particle and particle-antiparticle total cross-sections to the same constant limit satisfying the well-known Pomeranchuk theorem . Nevertheless the results obtained at Serpukhov and later at CERN-ISR clearly indicated that in the new domain of energies the total cross-sections began to increase again ; soon it was found that the differential $pp-$and $p\bar{p}-$cross-sections as well had different behaviour . Since such growth does not depend on the sort of the colliding particles , it is natural to suppose that in the asymptotical region the main contribution is provided by the Pomeranchuk pole , or pomeron ; moreover , the trustworthy description of σ^{p}_{tot} and σ^{a}_{tot} must also include its crossing-odd partner , or odderon (see early references in [7]) .

From the point of view of QCD the pomeron-odderon structure of the amplitude directly arises from the hypothesis of the"gluon dominance" [8] . The asymptotical interaction must occur due to the exchange by some number of gluons in a colourless state , because the contributions from the quark exchanges fall down like a power of s. The reggeization of crossing-even exchanges leads to the pomeron and of crossing-odd ones — to the odderon . At the same time the introduction of odderon does not contradict the basic QFT principles , and *a priori* nothing forbids the maximal growth of basic QFT principles , and *a priori* nothing forbids the maximal growth of $\Delta\sigma_{tot}$, namely as $\log s$.

The scenarios with such saturating increase of $\Delta\sigma_{tot}$ are named " the models of the maximal odderon " : for instance , in [9] the form of the amplitude , chosen on the base of motto "Strong interactions must be maximally strong" , provides both the saturation of $\Delta\sigma_{tot}$ and the increase of $\rho(s,0)$. On the other hand there exists a number of models with the "minimal" odderon [10, 11] , which predict $\Delta\sigma_{tot} \to 0$; in one of these approaches [12] it is even possible to verify the value of $\rho(s,0)$ and show that the special parametrization of the experimental cross-sections gives the reasonable values of ρ [0.136 instead of 0.24] . Certainly , $\Delta\sigma_{tot}$ in reality may also tend to some constant . Let us now see , which alternative is realized in the U-matrix treatment of the amplitude .

Suppose that the main contributions into the high-energy amplitude are made by one pomeron and one odderon . Of course , there arises a question if the really existing singularities have such simple structure or not . For instance , in the perturbative QCD using the "leading log" approximation , pomeron and odderon appear as two- and three-gluon bound states ; in this case pomeron must consist at least of two Regge poles [13] . Nevertheless this calculation implies that the odderon trajectory lies above the pomeron one , which is incompatible with the amplitude unitarity . Besides that , the presence of several Regge poles does not differ , in principle , (as a consequence of the U-matrix additivity) from the case of the non-coinciding pomeron and odderon and has no influence on the qualitative behaviour of the leading term in the asymptotical expansion : it is still determined only by leading poles .

If the pomeron and odderon trajectories are linear and correspondingly are written as $\alpha_\pm(t) = 1 + \Delta_\pm + \alpha'_\pm(t)$, where $\Delta_\pm > 0, \alpha'_\pm > 0$, by the direct calculation of Eq. 9 we get

$$U_+(b,s) = \frac{\lambda_+}{B_+(s)}\bar{s}^{\Delta_+}\exp(-\frac{b^2}{B_+(s)}),$$
$$U_-(b,s) = \frac{i\lambda_-}{B_-(s)}\bar{s}^{\Delta_-}\exp(-\frac{b^2}{B_-(s)}). \quad (11)$$

Here $\lambda_+ = \frac{2\pi^2 g_+}{\sin\frac{\pi}{2}(1+\Delta_+)}, \lambda_- = \frac{2\pi^2 g_-}{\cos\frac{\pi}{2}(1+\Delta_-)}, \bar{s} = s\exp(-i\frac{\pi}{2}), B_\pm(s) = 4\alpha'_\pm \log\bar{s}$.

One can note that the logarythmical behaviour of $B_\pm(s)$ provides the saturation of the Froissart bound : it follows from Eq. 11 that the effective radius of the interaction $b_{max} \sim \log s$ and $\sigma_{tot} \sim b^2_{max} \sim \log^2 s$.

The $pp-$ and $p\bar{p}-$ forward amplitudes look like

$$F^{p(a)}(s,0) = \frac{s}{\pi^2}\int bdb\frac{U_+(b,s) \pm U_-(b,s)}{1-i(U_+(b,s) \pm U_-(b,s))}. \tag{12}$$

We can as well denote

$$F^+(s,0) = F^p(s,0) + F^a(s,0) = \frac{2s}{\pi^2}\int bdb\frac{U_+^2 - U_-^2 + U_+}{[1+U_+]^2 - U_-^2}, \tag{13}$$

$$F^-(s,0) = F^p(s,0) - F^a(s,0) = \frac{2s}{\pi^2}\int bdb\frac{U_-}{[1+U_+]^2 - U_-^2} \tag{14}$$

and obtain the equations

$$\sigma_{tot} = \sigma_{tot}^p + \sigma_{tot}^a = \frac{8\pi^3}{s} ImF^+(s,0) \tag{15}$$

$$\Delta\sigma_{tot} = \sigma_{tot}^p - \sigma_{tot}^a = \frac{8\pi^3}{s} ImF^-(s,0). \tag{16}$$

The unitarity of $F^{p(a)}(s,0)$ leads to the inequalities $\Delta_+ \geq \Delta_-$, $\alpha'_+ \geq \alpha'_-$. In this case the behaviour of the intergrands in Eq. 13 and Eq. 14 at fixed b and growing s is different : it tends to a constant in the first equation and infinitely vanishes in the second one . This gives some reason to suspect that $\Delta\sigma_{tot}$ does not grow in the "asymptopia" ; let us prove that for a special choice of the parameters $\alpha'_+ = \alpha'_- = \alpha', \Delta_+ = \Delta_- = \Delta, B_\pm(s) = B(s)$.

4 The calculation of $\Delta\sigma_{tot}$

It can be shown that $\Delta\sigma_{tot}$ grows more rapidly if the trajectories are degenerate rather if they do not coincide ; moreover , it is argued in [15] , that this degeneracy may appear in QCD as a result of the quark confinement . Then

$$U^{p(a)}(b,s) = \frac{\lambda_+ \pm i\lambda_-}{B(s)}\bar{s}^\Delta \exp(-\frac{b^2}{B(s)}). \tag{17}$$

From Eq. 12 we come to

$$F^{p(a)}(s,0) = \frac{s}{\pi^2}\int_0^\infty bdb\frac{\lambda_+ \pm i\lambda_-}{B(s)}\bar{s}^\Delta \exp(-\frac{b^2}{B(s)})\frac{1}{1 - i\frac{\lambda_+ \pm i\lambda_-}{B(s)}\bar{s}^\Delta \exp(-\frac{b^2}{B(s)})}. \tag{18}$$

After the standard substitution this integral is easily calculated and we get

$$\sigma_{tot}^{p(a)} = 4\pi Re\Big\{ \mathrm{B(s)} \log(1 + \frac{\lambda_+ \pm i\lambda_-}{\mathrm{B(s)}} \bar{s}^{\Delta})\Big\}. \tag{19}$$

Denote

$$\rho_\lambda = \sqrt{\lambda_+^2 + \lambda_-^2}\ , \theta_0 = \frac{\pi}{2}(\Delta + 1) - arctg\frac{\pi}{2\log s} \rightarrow \frac{\pi}{2}(\Delta+1)$$

as $s \rightarrow \infty$. Now we can pick out the growing term $\log(\frac{\lambda_+ \pm i\lambda_-}{B(s)}\bar{s}^{\Delta})$ and expand the remaining summand into the series of $\left(\frac{B(s)}{\lambda_+ \pm i\lambda_-}\bar{s}^{-\Delta}\right)$. After the calculation of the real part we find that

$$\sigma_{tot}^{p(a)} \rightarrow 4\pi\Big\{ 4\alpha'\Big\{ \Delta \log^2 s + \log(\frac{\rho_\lambda}{4\alpha' \log s}) \log s +$$
$$+ \frac{4\alpha'}{\rho_\lambda} \cos(\theta_0 \mp arctg(\frac{\lambda_-}{\lambda_+})) s^{-\Delta} \log^2 s \Big\} - 2\pi\alpha' \Big\{ \theta_0 \mp arctg(\frac{\lambda_-}{\lambda_+}) -$$
$$- \frac{4\alpha'}{\rho_\lambda} \sin(\theta_0 \mp arctg\frac{\lambda_-}{\lambda_+}) s^{-\Delta} \log s \Big\}\Big\}. \tag{20}$$

This , in its turn , implies

$$\sigma_{tot} \simeq 32\pi\alpha'\Delta \log^2 s \tag{21}$$

and

$$\Delta\sigma_{tot} = 16\pi\alpha' arctg\frac{\lambda_-}{\lambda_+} + \frac{\pi}{2}\lambda_-(\frac{16\alpha'}{\rho_\lambda})^2 \cos(\frac{\pi}{2}\Delta) s^{-\Delta} \log^2 s +$$
$$+ \lambda_-(\frac{8\pi\alpha'}{\rho_\lambda})^2 sin(\frac{\pi}{2}\Delta) s^{-\Delta} \log s. \tag{22}$$

As we have expected , the presence of the pomeron with the non-zero intercept Δ leads to the saturating growth of σ_{tot} . However , the odderon provides only constant $\Delta\sigma_{tot}$; we also notice that $\Delta\sigma_{tot} \rightarrow 0$ if $\lambda_- \rightarrow 0$, i.e. the difference of the total cross-sections is really appearing due to the odderon .

5 The calculation of the total elastic cross-sections

From Eq. 8 and Eq. 17 we get the expression

$$\sigma_{inel}(s,0) = 8\pi \int_0^\infty bdbIm\frac{\lambda_+ \pm i\lambda_-}{B(s)}\bar{s}^\Delta \exp(-\frac{b^2}{B(s)})\frac{1}{|1 - i\frac{\lambda_+ \pm i\lambda_-}{B(s)}\bar{s}^\Delta \exp(-\frac{b^2}{B(s)})|^2}. \quad (23)$$

After the substitution $z = z_0 e^{-\frac{b^2}{B}}, z_0 = -i\frac{\lambda_+ \pm i\lambda_-}{B(s)}\bar{s}^\Delta$, Eq. 23 takes the form

$$\sigma_{inel}^{p(a)}(s,0) = 8\pi Re \int_0^{z_0} \frac{B}{2}\frac{dz}{|1+z|^2}. \quad (24)$$

Since the integrand has no singularities in the integration region , Eq. 23 is easily calculated when $s \to \infty$.

Finally we get

$$\sigma_{inel}^{p(a)} = 4\pi\Big\{4\alpha' \log s + \frac{16\alpha'^2}{\rho_\lambda}\Big\{\cos(\theta_0 \mp arctg\frac{\lambda_-}{\lambda_+}) - 2\Big\}s^{-\Delta}\log^2 s - \\ -\frac{8\alpha'^2}{\rho_\lambda}\pi \sin(\theta_0 \mp arctg\frac{\lambda_-}{\lambda_+})s^{-\Delta}\log s\Big\}. \quad (25)$$

The difference $\Delta\sigma_{inel} = \sigma_{inel}^p - \sigma_{inel}^a$ vanishes in the high-energy limit:

$$\Delta\sigma_{inel} = \frac{\pi}{2}\lambda_-(\frac{16\alpha'}{\rho_\lambda})^2\cos(\frac{\pi}{2}\Delta)s^{-\Delta}\log^2 s + \\ +\lambda_-(\frac{8\pi\alpha'}{\rho_\lambda})^2 sin(\frac{\pi}{2}\Delta)s^{-\Delta}\log s; \quad (26)$$

with the help of Eq. 22 we find

$$\Delta\sigma_{el} = \Delta\sigma_{tot} - \Delta\sigma_{inel} = 16\pi\alpha' arctg\frac{\lambda_-}{\lambda_+}. \quad (27)$$

In addition , asymptotically $\sigma_{el} \sim \log^2 s$.

This relation implies that the saturation of the Froissart bound arises due to the elastic scattering ; though σ_{inel} continues to rise , we find $\sigma_{el}/\sigma_{tot} \to 1$, which is in accordance with the demands of the general field theory .

6 U-matrix in comparison with eikonal

The cross-sections found in Sec.4,5 were previously obtained with the help of the eikonal approximation [2],[14] . We must emphasize that both methods give coinciding predictions about the behaviour of σ_{tot} and $\Delta\sigma_{tot}$, while the description of the scattering mechanisms is not the same . Though eikonal $\Delta\sigma_{tot}$ also tends to constant when the leading trajectories are degenerate , the inelastic difference $\Delta\sigma_{inel}$ still grows as $\log s$; at the same time $\Delta\sigma_{el}$ increases with the same rapidity but with opposite signs , so their contributions into $\Delta\sigma_{tot}$ compensate each other . In the U-matrix treatment $\Delta\sigma_{inel}$ falls down as a power of s ; that is why even the constant $\Delta\sigma_{tot}$ arises at the expense of the elastic reactions . This fact does not seem surprising since in the eikonal parametrization one obtaines $\sigma_{el}/\sigma_{tot} \to \frac{1}{2}$; at the same time the U-matrix approach always provides $\sigma_{el}/\sigma_{tot} \to 1$. We shall as well observe that , unlike the eikonal picture where all $\Delta\sigma$ depend on $\lambda_\pm$ parameters only through their ratio λ_-/λ_+ , in our case $\Delta\sigma_{inel}$ is wholly determined by λ_- and ρ_λ . For the realistic choice of the intercept $\Delta \approx 0.1 \div 0.2$ the signs of λ_- and $\Delta\sigma_{inel}$ are to be the same , according to Eq. 26 . Moreover it follows from Eq. 22 that the information about the asymptotical signs of $\Delta\sigma_{tot}$ and $d\Delta\sigma_{tot}/ds$ is sufficient to make definite conclusions about the signs of λ_+ and λ_- .

The same speculations can be conducted in the case of non-degenerate trajectories . Certainly , we cannot get the exact expressions for the values concerned since we have to take into account more than one exponential piece each time . However , the upper estimates show that $\Delta\sigma_{tot}$ falls down if the pomeron lies above the odderon ; namely $\Delta\sigma_{tot} \sim s^\beta, \beta < 0$, which is rather compatible with the eikonal predictions.

To conclude we stress that all the results of this paper concern the asymptotical behaviour of the cross-sections . Of course the functions in question may have extrema at finite energies . For example , it can be seen from Eq. 22 that in the used approximation $d\Delta\sigma_{tot}/ds$ has two zeros , which satisfy the quadratic equation for $\log s$. These zeros are $\log s = 0$ and $\log s = \frac{2}{\Delta}$, which for $s_0 \approx 1$ GeV correspond to $s \approx 1$ GeV and $s \approx e^{\frac{2}{\Delta}} \approx e^{10} \div e^{20}$ GeV ; due to this $\Delta\sigma_{tot}$ may temporarily grow at current energies . These speculations are on no account strict but the fact is that even in the one pomeron-one odderon model $\Delta\sigma_{tot}$ may increase within a finite interval , though slower than $\log s$. However , if the collider energies have already reached "asymptopia" (though the experimental values of σ_{el}/σ_{tot} are still suspiciously far from 1) , then the measurement of the elastic scattering may become critical to choose between U-matrix and eikonal parametrizations .

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