# AN APPLICATION OF THE *U*-MATRIX APPROACH TO HIGH ENERGY PROTON-PROTON SCATTERING

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We consider elastic proton-proton scattering at high energies and small momentum transfer. The expressions for the pp-scattering helicity amplitudes are obtained using the model of pomeron as an $SU(3)$-scalar photon with helicity-non-conserving interaction vertex. On the basis of the *U*-matrix method, we calculate unitarized helicity amplitudes in elastic pp-scattering at high energies and small momentum transfer. The isoscalar formfactor of proton is shown to give the main contribution into helicity-flip amplitudes. The momentum transfer dependence of amplitudes is analyzed. It is shown that the ratio of the single-flip amplitude to the non-flip amplitude $r_5$ equals to a half of the nucleon isoscalar anomalous magnetic moment.

## Introduction

General principles of quantum mechanics require the unitarity of the $S$-matrix. The unitarity imposes a bound on the high energy behavior of amplitudes, cross- sections, structure functions, etc.

The *U*-matrix approach [1] allows one to obtain the analytic continuation of a model amplitude from the annihilation to direct channel which satisfies the unitarity. From the physical point of view, the unitarization procedure corresponds to rescattering on various nucleons or quarks depending on the type of a process.

The present data seem to violate the unitarity (power-like behavior of cross-sections or structure functions). It happens because, at present energies, rescattering does not come into play. But, at certain energies, rescattering starts playing an important role and changes the behavior.

In general, there are two ways to take rescattering into account: the multiple scattering theory by Glauber and Sitenko and the *U*-matrix theory. The *U*-matrix approach gives formulas which are more convenient for analytic calculations. Briefly, the *U*-matrix method is as follows: one should construct the *U*-matrix and then solve algebraic equations [1] to find a physical amplitude in the impact parameter representation.

The *U*-matrix method works well in both nonpolarized [2] and polarized cases. This article is devoted to the latter case.

The spin dependence of hadron scattering amplitudes contains an important information on the dynamics of strong interaction. So, the question arising in connection is whether spin effects in hadron-hadron interactions important at high energies? The well-established opinion about the 'death“ of spin effects at high energies is criticized last time. For instance, in [3, 4], the existence of the asymptotic part of the single spin-flip amplitude $f_5$ which could not be neglected comparing with the imaginary part of the nonflip spin amplitude is discussed.

From the QCD (quantum chromodynamics) point of view, hadron-hadron scattering at small transferred momenta is a nonperturbative process. To describe it in a specified kinematic region, hybrid models based on the Regge theory formalism could be applied.

As usually, the pomeron exchange in Regge amplitudes conserves helicity, which in fact causes the 'death“ of spin effects at high energies. However, there are no reasons initially to suppose that helicity is conserved. In other words the single and double helicity flip amplitudes are not small. In this article, we calculate helicity amplitudes for proton-proton scattering using the pomeron model developed in [5], which is modified in such a way that the helicity in a proton-pomeron vertex is not conserved. At quite high energies, the rescattering effect has to be taken into account. For this purpose we use the *U*-matrix method [6]. The transferred momentum dependence of helicity amplitudes is also investigated.

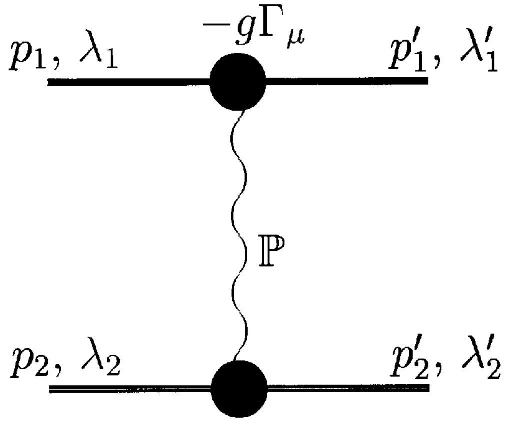
## 1. Kinematics. Helicity Amplitudes

Let’s consider the elastic proton-proton scattering process in the center-of-mass frame (Figure). We denote the momenta of incoming protons as $\mathbf{p}_1$ and $\mathbf{p}_2$, their energies as $\varepsilon_1$ and $\varepsilon_2$, correspondingly.

In the c.m. reference frame, $\mathbf{p}_1 + \mathbf{p}_2 = 0$ and, thus, $\varepsilon_1 = \varepsilon_2 \equiv \varepsilon$ (because the masses of incoming particles are equal).

Analogous relations for energies and momenta of outgoing protons $\varepsilon'_1, \varepsilon'_2$ and $\mathbf{p}'_1, \mathbf{p}'_2$ follow from the energy and momentum conservation laws: $\mathbf{p}'_1 + \mathbf{p}'_2 = 0;\ \varepsilon'_1 = \varepsilon'_2 \equiv \varepsilon$.

Let’s choose the $z$-axis of the reference frame along the momentum $\mathbf{p}_1$ and parametrize the 4-momenta of incoming and outgoing particles as: $p_1 = (\varepsilon, 0, 0, |\mathbf{p}|)$, $p_2 = (\varepsilon, 0, 0, -|\mathbf{p}|)$, $p'_1 =$




Pomeron exchange contribution into elastic pp-scattering amplitude

$= (\varepsilon, 0, 0, |\mathbf{p}| \cos\theta)$, $p'_2 = (\varepsilon, 0, 0 - |\mathbf{p}| \cos\theta)$, where $\theta$ is the scattering angle, $|\mathbf{p}| \equiv |\mathbf{p}_1| = |\mathbf{p}_2| = |\mathbf{p}'_1| = |\mathbf{p}'_2|$. We denote proton helicities as $\lambda_1, \lambda_2, \lambda'_1, \lambda'_2$, $|\lambda_i| = 1/2$. In the chosen reference frame, the bispinors of incoming and outgoing particles could be written in the following form:

$$u_1 = \begin{pmatrix} \sqrt{\varepsilon + m}\, w_1 \\ 2\lambda_1 \sqrt{\varepsilon - m}\, w_1 \end{pmatrix},$$

$$\bar{u}'_1 = (\sqrt{\varepsilon + m}\, w'^*_1;\ -2\lambda'_1 \sqrt{\varepsilon - m}\, w'^*_1), \qquad (1)$$

$$u_2 = \begin{pmatrix} \sqrt{\varepsilon + m}\, w_2 \\ 2\lambda_2 \sqrt{\varepsilon - m}\, w_2 \end{pmatrix},$$

$$\bar{u}'_2 = (\sqrt{\varepsilon + m}\, w'^*_2;\ -2\lambda'_2 \sqrt{\varepsilon - m}\, w'^*_2),$$

where the spinors $w$ are

$$w_1 = \begin{pmatrix} \sqrt{1/2 + \lambda_1}\, e^{-i\pi(1/2 - \lambda_1)} \\ \sqrt{1/2 - \lambda_1} \end{pmatrix},$$

$$w_2 = \begin{pmatrix} \sqrt{1/2 - \lambda_2}\, e^{-i\pi(1/2 - \lambda_2)} \\ \sqrt{1/2 + \lambda_2} \end{pmatrix}, \qquad (3)$$

$$w'_1 = \begin{pmatrix} \sqrt{1/2 + \lambda'_1 \cos\theta}\, e^{-i\pi(1/2 - \lambda'_1)} \\ \sqrt{1/2 - \lambda'_1 \cos\theta} \end{pmatrix},$$

$$w'_2 = \begin{pmatrix} \sqrt{1/2 - \lambda'_2 \cos\theta}\, e^{-i\pi(1/2 - \lambda'_2)} \\ \sqrt{1/2 + \lambda'_2 \cos\theta} \end{pmatrix}. \qquad (4)$$

The elastic pp-scattering process is described by 16 helicity amplitudes $T^{\lambda'_1 \lambda'_2}_{\lambda_1 \lambda_2}$. However, only 5 are independent due to the invariance under space reflection, time reversal, and particles indentity. We follow [6] and choose these amplitudes as

$$T_1 \equiv T^{++}_{++}, \qquad (5)$$

$$T_2 \equiv T^{++}_{--}, \qquad (6)$$

$$T_3 \equiv T^{+-}_{+-}, \qquad (7)$$

$$T_4 \equiv T^{+-}_{-+}, \qquad (8)$$

$$T_5 \equiv T^{++}_{+-}. \qquad (9)$$

As one could see from definitions (5) – (9), $T_1$ and $T_3$ are helicity nonflip amplitudes, $T_2$ and $T_4$ double helicity-flip amplitudes, and $T_5$ single helicity-flip amplitude.

We study the behavior of these amplitudes in the domain of high energies $\varepsilon$ and small transferred momenta $t = -4|\mathbf{p}|^2 \sin^2\theta/2$. In this domain, the contribution of nondiffractive processes into helicity amplitudes is negligible. The behaviour is controlled by a pomeron exchange.

## 2. The Model of Pomeron

Let's consider quite high energies to neglect the contributions of reggeons different from pomeron to the process amplitude. This simplification restricts the applicability of the obtained formula to analysis of the experimental data related mainly to intermediate energies. However, it is not a principal point and was done to simplify calculations.

From the Regge pole theory point of view, pomeron is an exchange of trajectory with vacuum quantum numbers and positive signature, which determines the asymptotic behavior of cross-sections at high energies. Experimental data show the pomeron trajectory intercept to be greater than unity.

In the framework of QCD, any consistent description of pomeron is absent, which leads to the need of a model approach.

A simple model of pomeron was proposed in [5]. In this model, pomeron is treated as an $SU(3)$-scalar photon. Pomeron-proton interaction vertices are described by the expression $(-g F_1(t) \gamma_\mu)$ which contains the coupling constant $g$, dipole formfactor of proton $F_1(t)$ and Dirac $\gamma$-matrices. Such a look of vertices conserves helicity.

But, by analogy with electrodynamics, the contribution of the isoscalar formfactor $F_2(t)$ was neglected in [5].

It was pointed out in [7] that, in a case of strong interactions, such an asssumption is not valid. A vertex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function has to be chosen in general form: $(-g\Gamma_\mu)$, where $\Gamma_\mu = F_1(t)\gamma_\mu - F_2(t)\sigma_{\mu\nu}q^\nu/2m$.

It will be shown below that the term containing the isoscalar formfactor $F_2(t)$ is responsible for the helicity nonconservation and very important for the description of spin effects.

## 3. Single Scattering Amplitudes

According to [7], we write the helicity amplitude $T^{\lambda'_1\lambda'_2}_{\lambda_1\lambda_2}$ in the following form:

$$T^{\lambda'_1\lambda'_2}_{\lambda_1\lambda_2} = i\frac{g}{2s_0}(\bar{u}'_2\Gamma^\mu u_2)(\bar{u}'_1\Gamma_\mu u_1)\left(\frac{s}{s_0}e^{-i\pi/2}\right)^{\alpha_p(t)-1}, \tag{10}$$

where $u$ are bispinors defined in (1) and (2), $s=(p_1+p_2)^2$, $t=(p_1-p'_1)^2=-q^2$, $\alpha_p(t)$ is the pomeron trajectory, which was chosen in the form $\alpha_p(t) = 1+\Delta+\alpha' t$. The exponent in (10) is a signature multiplier (here, the pomeron signature $\xi = +1$ is taken into account).

Utilizing the expressions for bispinors (1) in (10), the following expressions for the amplitudes $T_1$, $T_2$, $T_3$, $T_4$, $T_5$ could be obtained:

$$T_1(s,t) = i\frac{g}{2s_0}\left(\frac{s}{s_0}e^{-i\pi/2}\right)^{\Delta+\alpha' t}\{2sF_1^2(t) + tF_2^2(t)\}, \tag{11}$$

$$T_2(s,t) = -T_4(s,t) = i\frac{g}{2s_0}t\left(\frac{s}{s_0}e^{-i\pi/2}\right)^{\Delta+\alpha' t}\times$$

$$\times\left\{\frac{4m^2}{s}F_1^2(t) - \frac{s}{2m^2}F_2^2(t)\right\}, \tag{12}$$

$$T_3(s,t) = i\frac{g}{2s_0}\left(\frac{s}{s_0}e^{-i\pi/2}\right)^{\Delta+\alpha' t}\{2sF_1^2(t) - tF_2^2(t)\}, \tag{13}$$

$$T_5(s,t) = i\frac{g}{2s_0}2m\sqrt{-t}\left(\frac{s}{s_0}e^{-i\pi/2}\right)^{\Delta+\alpha' t}\times$$

$$\times\left\{F_1^2(t) + \frac{s}{2m^2}F_1(t)F_2(t) - \frac{t}{4m^2}F_2^2(t)\right\}. \tag{14}$$

To simplify calculations, we take an exponential parametrization of formfactors [7]:

$$F_1(t) = e^{bt}, \quad F_2(t) = \mu e^{bt}, \tag{15}$$

where $\mu$ is the isoscalar anomalous magnetic moment of nucleon.

Then we substitute formfactors (15) into (11) – (14) and neglect terms of order $t/s$ and $m^2/s$:

$$T_1(s,t) = T_3(s,t) = iG\frac{s}{s_0}e^{-Bq^2}; \tag{16}$$

$$T_2(s,t) = -T_4(s,t) = iG\frac{s}{s_0}\frac{\mu^2q^2}{4m^2}e^{-Bq^2}; \tag{17}$$

$$T_5(s,t) = iG\frac{s}{s_0}\frac{\mu q}{2m}e^{-Bq^2}, \tag{18}$$

where $B = 2b + \alpha'(\ln s/s_0 - i\pi/2)$, $G = g\left(\frac{s}{s_0}e^{-i\pi/2}\right)^\Delta$.

Note that, at $\mu = 0$, the spin-flip amplitudes $T_2$, $T_4$, $T_5$ appear to be of a higher order of smallness in energy than spin-nonflip amplitudes.

At high energies, the corrections to amplitudes due to rescattering must be taken into account. Here, we describe multiple scattering in the framework of the *U*-matrix theory [1].

## 4. Multiple Scattering Amplitudes. *U*-matrix

Relations for unitarized helicity amplitudes $f^{\lambda'_1\lambda'_2}_{\lambda_1\lambda_2}$ have the most simple form in the impact parameter representation.

In this representation, the following set of equations could be written (see [6]):

$$f^{\lambda'_1\lambda'_2}_{\lambda_1\lambda_2}(s,\rho) = T^{\lambda'_1\lambda'_2}_{\lambda_1\lambda_2}(s,\rho) +$$

$$+ i\rho(s)\sum_{\nu_1\nu_2}T^{\nu_1\nu_2}_{\lambda_1\lambda_2}(s,\rho)f^{\lambda'_1\lambda'_2}_{\nu_1\nu_2}(s,\rho). \tag{19}$$

In the case of pp-scattering, this set is

$$f_1 = T_1 + i\rho\{T_1f_1 + T_2f_2 - 2T_5f_5\},$$

$$f_2 = T_2 + i\rho\{T_2f_1 + T_1f_2 - 2T_5f_5\},$$

$$f_3 = T_3 + i\rho\{T_3f_3 + T_4f_4 - 2T_5f_5\},$$

$$f_4 = T_4 + i\rho\{T_4f_3 + T_3f_4 + 2T_5f_5\},$$

$$f_5 = T_5 + i\rho\{f_5(T_1+T_2) + T_5(f_3 - f_4)\}, \tag{20}$$

where $\rho(s)\to 1$ as $s\to\infty$.

It's easy to solve the set (20). The leading order terms of a solution (at $\rho \to 1$) are

$$f_1(s,\rho) = \frac{T_1}{1 - iT_1}, \tag{21}$$

$$f_2(s,\rho) = \frac{T_2(1 - iT_3) - 2iT_5^2}{(1 - iT_1)^2(1 - iT_3)}, \tag{22}$$

$$f_3(s,\rho) = \frac{T_3}{1 - iT_3}, \tag{23}$$

$$f_4(s,\rho) = \frac{T_4(1 - iT_1) + 2iT_5^2}{(1 - iT_3)^2(1 - iT_1)}, \tag{24}$$

$$f_5(s,\rho) = \frac{T_5}{(1 - iT_1)(1 - iT_3)}. \tag{25}$$

Equations (20) contain the helicity amplitudes $T_1(s,\rho),\ldots,T_5(s,\rho)$ in the impact parameter representation which are related with amplitudes (11) – (14) by the transformation:

$$T^{\lambda'_1\lambda'_2}_{\lambda_1\lambda_2}(s,\rho) = \frac{2}{s}\int_0^\infty dq\, q\, J_{|\lambda-\lambda'|}(\rho q)\, T^{\lambda'_1\lambda'_2}_{\lambda_1\lambda_2}(s,t), \tag{26}$$

where $\lambda = \lambda_1 - \lambda_2$, $\lambda' = \lambda'_1 - \lambda'_2$. It's easy to see that the spin-nonflip $T_1$ and double spin-flip amplitude $T_2$ are transformed with the Bessel function $J_0$, the single spin-flip amplitude $T_5$ is transformed with $J_1$ and the double spin-flip amplitude $T_4$ is transformed with $J_2$.

By carring out simple calculations, one gets:

$$T_1(s,\rho) = T_3(s,\rho) = i\frac{G}{s_0 B}e^{-\rho^2/4B}, \tag{27}$$

$$T_2(s,\rho) = i\frac{G\mu^2}{4s_0 m^2 B^2}\left(1 - \frac{\rho^2}{4B}\right)e^{-\rho^2/4B}, \tag{28}$$

$$T_4(s,\rho) = -i\frac{G\mu^2}{4s_0 m^2 B^2}\frac{\rho^2}{4B}e^{-\rho^2/4B}, \tag{29}$$

$$T_5(s,\rho) = i\frac{G\mu}{4s_0 m B^2}\rho\, e^{-\rho^2/4B}. \tag{30}$$

So, expressions (21) – (25) together with (27) – (30) define the unitarized helicity amplitudes for pp-scattering in the impact parameter representation.

To investigate the slope of differential cross-sections, it's necessary to pass to the transferred momentum representation in (21) – (25) using the transformation inversed to (26):

$$f^{\lambda'_1\lambda'_2}_{\lambda_1\lambda_2}(s,t) = \frac{s}{2}\int_0^\infty d\rho\,\rho\, J_{|\lambda-\lambda'|}(\rho q)\, f^{\lambda'_1\lambda'_2}_{\lambda_1\lambda_2}(s,\rho). \tag{31}$$

The integral in (31) is not analytically calculable for each helicity amplitude, but several simplifications could be applied to our specific problem.

Numerical values of the quantities from the expressions for helicity amplitudes are as follows:

$$\alpha' = 0.25\,\text{GeV}^{-2};\ \Delta = 0.08;\ m \approx 1\,\text{GeV};\ s_0 \approx 1\,\text{GeV}^2; \tag{32}$$

$$\mu \sim 0.1;\ b \approx 2.087; \tag{33}$$

$$G = 0.489 \text{ at } \sqrt{s} = 53\,\text{GeV};\ G = 1.108 \text{ at } \sqrt{s} = 3\cdot 10^4\,\text{GeV}. \tag{34}$$

$\mu^2$ appears to be a small parameter and the integrand in (31) could be developed into series in it. Here, we've left only terms of the lowest order:

$$f_1(s,t) = isR_B^2\int_0^\infty dz\, J_0(2R_B q\sqrt{z})\frac{c_1}{e^z + c_1} = -isR_B^2\sum_{r=0}^\infty\frac{(R_B^2 t)^r}{r!}\mathrm{Li}_{r+1}(-c_1), \tag{35}$$

$$f_2(s,t) = isR_B^2\int_0^\infty dz\, J_0(2R_B q\sqrt{z})\times\frac{c_1c_2e^z((1-z)e^z + (1+z)c_1)}{(e^z+c_1)^3} = isR_B^4 c_2 t\sum_{r=0}^\infty\frac{(R_B^2 t)^r}{r!}\mathrm{Li}_{r+1}(-c_1), \tag{36}$$

$$f_4(s,t) = -isR_B^2\int_0^\infty dz\, J_2(2R_B q\sqrt{z})\times\frac{c_1c_2ze^z(e^z + 3c_1)}{(e^z+c_1)^3} = isR_B^4 c_2 t\sum_{r=0}^\infty\frac{(R_B^2 t)^r}{r!}(\mathrm{Li}_{r+1}(-c_1) - 2\mathrm{Li}_{r+2}(-c_1)). \tag{37}$$

$$f_5(s,t) = isR_B^2 \int_0^\infty dz\, J_1(2R_B q\sqrt{z}) \frac{c_1\sqrt{c_2}\sqrt{z}\, e^z}{(e^z+c_1)^2} =$$

$$= -isR_B^3\sqrt{c_2}\sqrt{-t}\sum_{r=0}^{\infty}\frac{(R_B^2 t)^r}{r!}\mathrm{Li}_{r+1}(-c_1), \qquad (38)$$

where $R_B^2 = |B|$ is the square of the interaction radius, $c_1 = |G|/(s_0 R_B^2)$, $c_2 = \mu^2/(4m^2 R_B^2)$, $\mathrm{Li}_{r+1}(-c_1) = -c_1/r!\int_0^\infty dz\, z^r/(e^z+c_1)$ is a polylogarithm of order $r+1$. The imaginary parts of $G$ and $B$ are neglected in comparison with their moduli.

Expressions (35) – (38) allow us to analyze the dependence of the imaginary parts of helicity amplitudes on transferred momentum $t$.

Taking into account that $c_1 = |G|/(s_0 R_B^2) \approx 0.1$, we substitute the polylogarithm by the first term of its expansion into a series: $\mathrm{Li}_{r+1}(-c_1) \sim -c_1$. Then we get:

$$f_1(s,t) = i|G|\frac{s}{s_0}e^{R_B^2 t}; \qquad (39)$$

$$f_2(s,t) = -f_4(s,t) = -i|G|\frac{s}{s_0}\frac{\mu^2 t}{4m^2}e^{R_B^2 t}; \qquad (40)$$

$$f_5(s,t) = i|G|\frac{s}{s_0}\frac{\mu\sqrt{-t}}{2m}e^{R_B^2 t}. \qquad (41)$$

These expressions correspond to the absence of rescattering and coincide with (16) – (18).

Let's consider the case of single rescattering, i.e., in formulas (35) – (38), the second term of the polylogarithm expansion is taken into account: $\mathrm{Li}_{r+1}(-c_1) \sim -c_1 + c_1^2/2^{r+1}$. We get

$$\sum_{r=0}^{\infty}\frac{(R_B^2 t)^r}{r!}\mathrm{Li}_{r+1}(-c_1) \sim -c_1 e^{R_B^2 t} + \frac{c_1^2}{2}e^{R_B^2 t/2}. \qquad (42)$$

The appearance of the additional term in (42) due to rescattering leads to a modification of the cross-section slope, defined as the logarithmic derivative of the differential cross-section, $\left.\frac{d}{dt}\ln\frac{d\sigma(s,t)}{dt}\right|_{t=0}$.

Now we calculate the asymptotic behavior of amplitudes at high energies. In the limit $s\to\infty$, the value of the polylogarithm argument in (35) – (38) becomes large and therefore we can use the asymptotics: $\lim_{c_1\to\infty}\mathrm{Li}_{r+1}(-c_1) = -\ln^{r+1}(c_1)/(r+1)!$. In this case:

$$\sum_{r=0}^{\infty}\frac{(R_B^2 t)^r}{r!}\mathrm{Li}_{r+1}(-c_1) =$$

$$= -\ln c_1\sum_{r=0}^{\infty}\frac{(-R_B^2(-t)\ln c_1)^r}{r!(r+1)!} =$$

$$= -\ln c_1\frac{J_1(2R_B\sqrt{(-t)\ln c_1})}{R_B\sqrt{(-t)\ln c_1}}. \qquad (43)$$

Finally, the asymptotic behavior of helicity amplitudes is

$$f_1(s,t) = isR_B^2\Delta\ln s/s_0\frac{J_1(2R_B\sqrt{(-t)\Delta\ln s/s_0})}{R_B\sqrt{(-t)\Delta\ln s/s_0}}, \qquad (44)$$

$$f_2(s,t) = -isR_B^2\frac{\mu^2}{4m^2}t\Delta\ln s/s_0\times$$

$$\times\frac{J_1(2R_B\sqrt{(-t)\Delta\ln s/s_0})}{R_B\sqrt{(-t)\Delta\ln s/s_0}}, \qquad (45)$$

$$f_4(s,t) = -isR_B^2\frac{\mu^2}{4m^2}t\Delta\ln s/s_0\times$$

$$\times\left(\frac{J_1(2R_B\sqrt{(-t)\Delta\ln s/s_0})}{R_B\sqrt{(-t)\Delta\ln s/s_0}} - \right.$$

$$\left. - 2\Delta\ln s/s_0\frac{J_2(2R_B\sqrt{(-t)\Delta\ln s/s_0})}{R_B^2(-t)\Delta\ln s/s_0}\right), \qquad (46)$$

$$f_5(s,t) = isR_B^2\frac{\mu}{2m}\sqrt{-t}\Delta\ln s/s_0\times$$

$$\times\frac{J_1(2R_B\sqrt{(-t)\Delta\ln s/s_0})}{R_B\sqrt{(-t)\Delta\ln s/s_0}}. \qquad (47)$$

As we just have seen, the rescattering leads to a change of the power-like dependence of amplitudes on energy, typical of Regge models, to the logarithmic one.

## Conclusion

In this paper, we have shown that taking the anomalous magnetic moment of proton into account leads to the helicity nonconservation, and both the spin-flip and spin-nonflip amplitudes are of the same order of magnitude in energy $s$.

The *U*-matrix method was utilized to take multiple scattering into account which is important at high energies, especially to obtain the asymptotic behavior of helicity amplitudes. The amplitudes $f_1(s, t)$, $f_2(s, t)$, $f_3(s, t)$, $f_4(s, t)$, $f_5(s, t)$ are calculated in the leading order in the anomalous magnetic moment $\mu^2$.

The obtained asymptotic expressions for helicity amplitudes at high energies evidence that $r_5 = m\,\mathrm{Im}\, f_5/(\sqrt{-t}\,\mathrm{Im}\, f_1) = \mu/2 \approx 5\%$ and is fully defined by the isoscalar nucleon anomalous magnetic moment.

Concluding, we would like to note that the asymptotic behavior of the helicity amplitudes (44) – (47) begins at superhigh energies. To correspond to the energies of modern and projected accelerators, one would consider formulas (35) – (38) rather than (44) – (47).

The experimental verification of the results obtained in this paper demands taking into account of reggeon contributions. It could be easily done in an analogous manner.

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