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ON GENERATION MECHANISMS OF SPIN EFFECTS IN QCD AT LARGE DISTANCES

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At present, great interest in high energy physics stems from the spin phenomena. The modern theory of strong interaction has many difficulties there [1]. They are very crucial for QCD too. Really, the perturbative QCD leads to small polarization at large momenta transfer [2]. This result is in contradiction with the experiment. This contradiction may be due to the fact that the experimentally investigated region is far from asymptotics, and the perturbative theory can not be used here.

The large-distance effects play an important role in investigations of high energy hadron scattering at small angles. As a result, different dynamical models [3-5] are used usually for studying of these processes. Some of them [4-5] lead to the spin effects which do not disappear as $s \to \infty$. In this region the contribution comes from the t-channel exchange with the vacuum quantum numbers (pomeron). So investigation of the pomeron spin structure is very important.

The amplitude with the vacuum quantum numbers in the t-channel is due to the two-gluon exchange [6]. The nonperturbative properties of the theory, which are important in this case [7], were taken into account in the model [8] for the case of qq scattering. It was shown that taking into account the full matrix structure of the two-gluon amplitude leads to the spin-flip amplitude growing as s

$$\frac{|T_{flip}|}{|T_{non-flip}|} \simeq \frac{m\ |t|}{\ln s/s_0\ \ a(m,t)}. \qquad (1)$$

Here m is a constituent quark mass and a is a linear function of $|t|$ at large $|t|$.

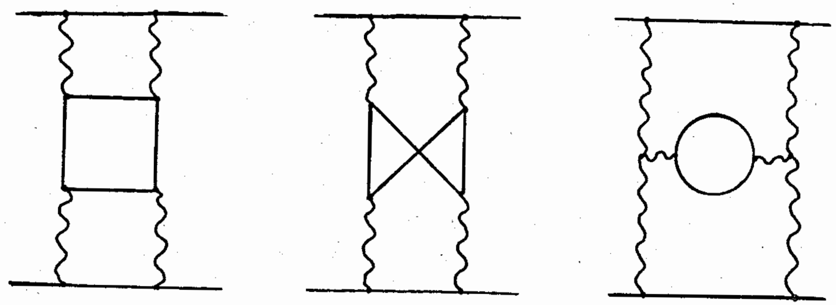
In this paper on the basis of [8,9] it is shown that the quark loops in the t-channel gluon exchange and $q\bar{q}$ sea contribution lead to the spin-flip amplitude growing as s. Physical mechanisms lead to such a behaviour of the spin-flip amplitude will be discussed too.

We shall calculate the qq scattering at high energies and fixed momenta transfer. In our model we suppose, following /7,8/, that nonperturbative effects are important in this region and the following representations for the quark and gluon propagators:

$$G^{q}(p)=i(\hat{p}+m)D(-p^2), \qquad G^{g}_{\alpha\beta}(q)=-ig_{\alpha\beta}F(-q^2) \qquad (2)$$

are used. The problems of normalization of F and gauge invariance in this case are discussed in /7,10/.

We shall investigate the quark-loop contribution to the imaginary part of two-gluon ladder diagrams(fig 1).



(a) (b) (c)

Fig.1 The quark-loop contribution to the ladder gg-amplitude.

For the diagram, fig.1a, for a definite flavor in the quark loop we have

$$\mathrm{Im}T^{a} \simeq c^{a}\frac{g^{8}}{2(2\pi)^{8}}\int d^4k d^4q d^4l \;\delta[(p-k)^2-m^2]\delta[(p'+l)^2-m^2]$$
$$\delta[(k+q)^2-\sigma^2]\delta[(q+l)^2-\sigma^2]\; H^{a}\; S^{a}_{\lambda\mu,\nu\sigma}\; N^{\lambda\mu,\nu\sigma}, \qquad (3)$$

were m and σ are the masses of a constituent quark and a quark in the loop, respectively, c^{a} is a color factor,

$$H^{a}= F[-(k+r)^2]F[-(k-r)^2]F[-(l+r)^2]F[-(l-r)^2]$$
$$D[-(q+r)^2+\sigma^2]D[-(q-r)^2+\sigma^2],$$
$$S^{a}_{\lambda\mu,\nu\sigma}= \mathrm{Sp}\{\gamma_{\lambda}[\hat{q}+\hat{k}+\sigma]\gamma_{\mu}[\hat{q}-\hat{r}+\sigma]\gamma_{\sigma}[\hat{q}+\hat{l}+\sigma]\gamma_{\nu}[\hat{q}+\hat{r}+\sigma]\}, \qquad (4)$$
$$N^{\lambda\mu,\nu\sigma}=\bar{u}(p-r)\gamma^{\lambda}(\hat{p}-\hat{k}+m)\gamma^{\mu}u(p+r)\;\bar{u}(p'+r)\gamma^{\nu}(\hat{p}'+\hat{l}+m)\gamma^{\sigma}u(p'-r).$$

In the light-cone variables: $k=(xp_+,k_-,k_\perp)$, $l=(yp_+,l_-,l_\perp)$, $q=(-zp_+,q_-,q_\perp)$, $p_\pm=p_0\pm p_z$ integration over k_-,l_-,q_-,y can be performed with the help of δ-functions. This leads to the cut-off over transverse momenta in the upper limit: $q_\perp^2,k_\perp^2,l_\perp^2\simeq s$. The main contribution to the amplitude with the non-spin-flip in the down quark line comes from the following term of the matrix element/8/:

$$N_1^{\nu\sigma}= \bar{u}^{+}(p'+r)\gamma^{\nu}(\hat{p}'+\hat{l}+m)\gamma^{\sigma}u^{+}(p'-r) \simeq 4p'^{\nu}p'^{\sigma}.$$

The spin-flip matrix element in the upper quark line has a more complicated structure. It can be decomposed into the sum of symmetric and antisymmetric parts

$$N^{\lambda\mu}\simeq \bar{u}\,[\hat{S}^{\lambda\mu}+\hat{A}^{\lambda\mu}]\,u, \qquad (5)$$
$$\hat{S}^{\lambda\mu}\simeq s^{\lambda}\gamma^{\mu}+s^{\mu}\gamma^{\lambda}+g^{\lambda\mu}\hat{k} \qquad \hat{A}^{\lambda\mu}\simeq i\varepsilon^{\lambda\mu\delta\rho}s_{\delta}\gamma_{\rho}\gamma_{5}-im\,\sigma^{\lambda\mu}.$$

Here and in what follows $s=p-k$ and Δ is a momentum transfer. Using the results from /8/ we can calculate the spin-flip matrix elements of $\hat{S}^{\lambda\mu}$ in terms of the light-cone variables

$$<\hat{S}^{\lambda\mu}a_{\lambda}v_{\mu}>_{\mathrm{flip}}= m\Delta\;[(sa)v_+/p_+ +(sv)a_+/p_+ +(av)k_+/p_+],$$
$$<\hat{S}^{\lambda\mu}g_{\lambda\mu}>_{\mathrm{flip}}= 2m\Delta\;[1+k_+/p_+]. \qquad (6)$$

For the matrix element of the antisymmetric part of (5) we have

$$<\hat{A}^{\lambda\mu}q_{\lambda}b_{\mu}>_{\mathrm{flip}}=m\;\{k_x[(q_+b_-)-(b_+q_-)]+k_+[q_-b_x-b_-q_x]+$$
$$+[b_xq_+/p_+-q_xb_+/p_+](\Delta^2/2-p_+k_-)\}. \qquad (7)$$

The quantities a,v,b,q in (6,7) are some vectors. As a result, for the spin-flip matrix element of the diagram (fig.1a) we have

$$\langle S^a_{\lambda\mu,\nu\sigma} N^{\lambda\mu,\nu\sigma}\rangle_{flip} = -16m\Delta(z-y)s^2\Big\{-q_\perp^2\frac{x^3-2x^2z-x^2+2xz^2+2xz-2z^2}{(x-z)} +$$
$$+k_\perp^2\frac{z(x^2z-x^2-2xz^2+2xz+2z^3-3z^2+z)}{(1-x)(x-z)} - 2(q_\perp k_\perp)\frac{x^2z-x^2-2xz^2+xz+2z^3-z^2}{(x-z)}$$
$$-2k_xq_x(x+z)+\frac{\Delta^2}{4}(x-z)^2-m^2\frac{z(x^2-3xz+2z^2)}{(1-x)} -\sigma^2\frac{x^3-2x^2+3xz-2z^2}{(x-z)}\Big\}. \quad (8)$$

It is easy to see that the integrals over $d^2k_\perp$, $d^2l_\perp$ are convergent in the upper limit. The main logarithmic asymptotics of (3) is connected with the integration over $d^2q_\perp$ near $q_\perp^2 \approx s$. In this region the momenta squared in the quark loop are large and we can use the asymptotically free quark propagators D. It is convenient to write (8) in the form

$$\langle S^a_{\lambda\mu,\nu\sigma} N^{\lambda\mu,\nu\sigma}\rangle_{flip} = m\Delta s^2[a^a(x,z)\ q_\perp^2 + \Psi^a(x,z,q_\perp,k_\perp,\Delta)]. \quad (9)$$

As a resalt of integration over $q_\perp^2$ we find

$$T_{flip}(s,t) = im\Delta s[\ln s/s_0\ T_1(t)+T_0(t)+\dots], \quad (10)$$

where

$$T_1(t) = 2N^f\frac{\alpha_s^4\pi}{(2\pi)^4}\int d^2l_\perp\ F[-(l+r)^2]F[-(l-r)^2]\times$$
$$\times\int_y^1\frac{dx}{x^2(1-x)}\int d^2k_\perp\ F[-(k+r)^2]F[-(k-r)^2]I^a(x). \quad (11)$$

Here

$$I^a(x) = c^a\int_y^x\frac{dz\ (x-z)}{(z-y)}a^a(x,z). \quad (12)$$

The gluon propagators in (11) depend on the variables

$$-(k\pm r)^2=\{x^2m^2+[\vec{k}_\perp\pm(1-x)\vec{r}_\perp]^2\}/(1-x); \quad -(l\pm r)^2=(\vec{l}_\perp\pm\vec{r}_\perp)^2.$$

In calculations of T_0 one must take into account the nonperturbative properties of the theory.

Similar calculations can be performed for the nonplanar graph (fig.1b). In this case, the spin-flip matrix element has a term proportional to $q_\perp^2$ as in eq. (9)

$$\langle S^b_{\lambda\mu,\nu\sigma} N^{\lambda\mu,\nu\sigma}\rangle_{flip} = m\Delta s^2[a^b(x,z)\ q_\perp^2 + \Psi^b(x,z,q_\perp,k_\perp,\Delta)]. \quad (13)$$

Expression (11) is fulfilled for the leading logarithmic term of the nonplanar graph but instead of (12) we have

$$I^b(x) = c^b\int_y^x dz\ a^b(x,z),$$

where

$$b(x,z) = 16[2z^2(1-x)-2zx(1-x)-x^3+x^2].$$

As a result of integration over x, we obtain

$$I^a(x) = -I^b(x) = -\frac{32}{3}x^3(1-x)c^{(a,b)}. \quad (14)$$

It is easy to see that in the case of quantum electrodynamics the leading logarithmic terms compensate each other because of the absence of colour factors in this case.

This compensation does not occur in the sum of diagrams drawn in fig 1a,b for the colour singlet exchange in QCD where we have:

$$I(x) = I^a(x)+I^b(x) = -\frac{16}{9}x^3(1-x)$$

because in QCD there exists the third diagram with a quark loop in the s-channel gluon propagator (fig 1c) which contributes to the imaginary part of the spin-flip amplitude.

The calculation of this contribution shows that the spin-flip matrix element in the diagram (fig 1c) numerator has a term proportional to $q_\perp^2$ as before. The integration of the δ-functions leads to the appearance of $q_\perp^2$ in the s-channel gluon propagators. As a result, the asymptotical equation (10,11) is fulfilled for this contribution. For the integral I^c in the case of the colour singlet exchange in the t-channel we have:

$$I^c(x) = c^c\int_y^x dz\ \frac{(x-z)(z-y)}{x^2}a^c(x,z) = \frac{16}{9}x^3(1-x).$$

Thus, in QCD we obtain the compensation of leading logarithms for the sum of diagrams (fig. 1a-c). All these *lns* terms are determined by short distances in t-channel quark propagators in the diagrams (fig.1a,b) and in s-channel gluon propagators in the diagram (fig.1c). As a result, the diagrams fig.1a-c, have the same topological structure in the $q_\perp^2\to\infty$ limit.

There are many terms in spin-flip matrix elements which do not contain $q_\perp^2$(see Eq.(8) for the diagram in fig.1a e.g.). They contribute to the nonlogarithmic term T_0 in (10). It is difficult to believe in full compensation of all these terms because of their different structure. Now, the accurate calculation of these contributions are impossible in QCD because the diagram propagators are in the nonperturbative region.

To show this let us calculate the nonlogarithmic terms for the sum of diagrams (fig.1a,b) in the case of quantum electrodynamics. Let us rewrite T_0 in the form

$$T_0(t)= \frac{2}{(2\pi)^4 m^4}\, \phi(t). \tag{15}$$

For simplicity we shall calculate the function $\phi(t)$ in (15) for zero momenta transfer. We shall conclude that a term proportional to s exists in the spin-flip amplitude if $\phi(0)$ is not equal to zero. Let us introduce the photon mass λ to avoid singularities at t=0. The calculations show for the diagram (fig.1a)

$$\phi^a(0)=\phi^a_\lambda(0)/\lambda^2+\phi^a_0(0). \tag{16}$$

The behavior different from (16) takes place for the diagram of fig. 1b:

$$\phi^b(0)=\phi^b_\lambda(0)/\lambda^2+\phi^b_{ln}(\ln\lambda^2/m^2)+\phi^b_0(0). \tag{17}$$

The existence of logarithmic terms in (17) is connected with the presence of $l_\perp^2$ terms in the diagram numerator. These terms are absent in the planar graph (see Eq.(8)). As a result, we obtain integrals of the form:

$$\Psi(r_\perp,\lambda)=\int\frac{d^2 l_\perp\, l_\perp^2\, f(l_\perp^2)}{[(l_\perp+r_\perp)^2+\lambda^2]\,[(l_\perp-r_\perp)^2+\lambda^2]}\,, \tag{18}$$

where $f(l_\perp^2)$ decreases as $l_\perp^2\to\infty$. The integrals (18) in the $r_\perp\to 0$ limit lead to the following behaviour:

$$\Psi(0,\lambda)=\ln(\lambda^2/m^2). \tag{19}$$

The REDUCE program allows us to calculate the terms $\phi^{a,b}_\lambda$ and ϕ^b_{ln} in (16,17) up to the end. As a result of these

calculations the $1/\lambda^2$ terms compensate each other in the sum of contributions (fig.1a,b) but logarithmic terms are not equal to zero

$$\phi^b_{ln}(\ln(\lambda^2/m^2))= \frac{\pi^3}{54}\,\alpha_s^4\,[537-112\pi^2-180\ln(\lambda^2/m^2)]\ln(\lambda^2/m^2). \tag{20}$$

Thus, the planar and nonplanar contributions have different dependences on momenta transfer. In the case of nonzero t, the logarithmic form $\ln(\lambda^2/m^2)$ turn to $\ln(\Delta^2/m^2)$. As a result there is no singularity at λ=0 in the spin-flip amplitude at all $|t|$.

So we see that in quantum electrodynamics the contributions proportional to s are not compensated in the sum of diagrams (fig.1a,b) and the asymptotic form

$$T_{flip}(t)\simeq i\,\frac{\sqrt{|t|}\; s}{m^3}\,\phi(t) \tag{21}$$

is true for the spin- flip amplitude as $s\to\infty$, t-fixed. Thus this amplitude is suppressed only logarithmically with respect to the spin-non- flip amplitude. In order to conclude this, we use the results from /11/ for the spin-non-flip amplitude. The behaviour like (21) must be correct in QCD too.

Let us exemplify by the diagram (Fig.1a) the physical reason which leads to the growing as s contribution to the spin-flip amplitude. This diagram can be decomposed into two quark-quark subgraphs with the two-gluon exchange in t-channel. The corresponding spin-flip amplitude in this case has the following energy dependence:

$$T^{qq}_{flip}(t)\simeq \frac{m\sqrt{|t|}}{s_{qq}}\cdot T^{qq}_{non\text{-}flip},$$

where s_{qq} is the quark subprocess energy. It can be shown that in the integration region which contributes to the T_0 spin-flip amplitude the energies in the up and down subgraph are of the following order of magnitude:

$$s^{up}_{qq}\simeq m^2;\qquad s^{down}_{qq}\simeq s.$$

Thus the up quark subprocess is at low energies and the spin-flip amplitude has not energy suppression in it

$$T^{qq-up}_{flip}(t)\simeq\frac{\sqrt{|t|}}{m}\,T^{qq}_{non\text{-}flip}.$$

The spin-non-flip amplitude growing as s contribute to the down quark subprocess. As a result, we obtain the behaviour (10,11) for the total diagram.

So we can conclude that in QCD the terms $\approx s$ in the spin-flip amplitude are determined by the nonperturbative region in the up subprocess of the diagram. From our point of view the $q\bar{q}$ sea contribution (Fig.2) can be very important here.The method of effective meson lagrangians[13] permits us to replace the $q\bar{q}$ interactions in t-channel by the π-meson exchanges (Fig.3).

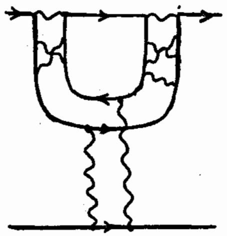


Fig.2. Diagram with the $q\bar{q}$ sea contribution.

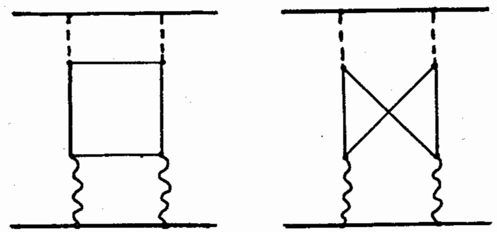


Fig.3. Effective diagrams with the π-meson exchanges.

It is easy to see that in this case the spin-flip and non-flip amplitude in the up quark subprocess are of the same order of magnitude:

$$T^{qq-up}_{non-flip}(t) \simeq \frac{m^2\ x^2}{(1-x)} + \frac{k_\perp^2}{(1-x)} - \frac{\Delta^2}{4}\ .$$

$$T^{qq-up}_{flip}(t) \simeq -m\Delta x.$$

The large magnitude of the quark-meson coupling constant $\alpha_{qqm}\simeq 1$ permits us to conclude that the diagrams (Fig.3) must be important in spin effects in high energy hadron interactions. This can explain the success of the meson-cloud model[5] which take into account similar effects phenomenologically.

Thus, the quark loop effects in gluon t-channel exchange and $q\bar{q}$ sea contributions lead to the spin-flip amplitude growing as s. Similar contribution can be obtained from the nonperturbative diquark state in the wave function for example[14]. In all cases such a behaviour of the spin-flip amplitude is determined by the long-distance effects.

So it is shown that in QCD the spin effects which decrease very slowly (only logarithmically) with energy growth really can be obtained in the $s\to\infty$ limit. They are connected with the nonperturbative contributions. It is necessary to use the theory properties at large distances to obtain some quantitative estimations.

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Голоскоков С.В. E2-90-383

О механизмах возникновения спиновых эффектов в КХД на больших расстояниях

Обсуждены различные физические механизмы, приводящие к амплитуде с переворотом спина, растущей как s, при высоких энергиях и фиксированных переданных импульсах.

Работа выполнена в Лаборатории теоретической физики ОИЯИ.

Препринт Объединенного института ядерных исследований. Дубна 1990

Goloskokov S.V. E2-90-383

On Generation Mechanisms of Spin Effects in QCD at Large Distances

The discussion of the different physical mechanisms which lead to the spin-flip amplitude growing as s at high energies and fixed momenta transfer is done.

The investigation has been performed at the Laboratory of Theoretical Physics, JINR.

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