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ON THE COMPTON TWIST-3 ASYMMETRIES

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1 Introduction

It is known that large asymmetries are observed in processes with a single polarized particle. For instance, asymmetries observed in pion production with a high transverse momentum are of an order of ten or more percent[1]. For this reason single asymmetries can be very convenient for the study of polarization effects in QCD. However, for their generation a mass parameter and an additional imaginary phase are necessary. As a result, these asymmetries are absent in the massless QCD twist-2 approach and twist-3 contributions appear in the leading approximation.

A decade ago [2, 3] a self-consistent approach to the single asymmetries in QCD was proposed. As a result, the parton-like expression was obtained. A short-distance part is calculable in perturbative QCD with slightly modified Feynman rules. A long-distance contribution is described by a new two-argument parton matrix elements, the so-called quark-gluon correlators. The latter should, in principle, be determined experimentally from a "partonometer" process, just like the ordinary parton distributions is determined from the deep inelastic scattering.

The experimental study of quark-gluon correlators was recently approved by the RHIC Spin Collaboration [4]. The comparison with the QCD predictions involves a number of the short-distance subprocesses. Up to now, the Compton subprocesses $\gamma N\uparrow\rightarrow \gamma X$ and $gN\uparrow\rightarrow \gamma X$ are calculated [3, 5, 7]. Apart of the quark-gluon correlators, mentioned above, there are also pure gluonic correlators, giving rise to the single asymmetry in the subprocess $qN\uparrow\rightarrow qX$ [8]. In the present article we try to add a new line to this short list. We confine ourselves to the quark-gluon correlators and calculate the asymmetry in the essentially non-Abelian subprocess $gN\uparrow\rightarrow gX$, contributing significantly to asymmetry of pions and jets.

There is also an important "pure theoretical" problem. The approach [2, 3] is based on the Ellis, Furmansky, Petronsio (EFP) factorization scheme by using an axial gauge[9]. More recently similar results were obtained by J. Qiu and G. Sterman [6, 7] who used the so-called special propagators technique and covariant gauge. Although these approaches coincide in the most essential points, there is a significant discrepancy. Sterman and Qiu proposed as the main source of large asymmetries the so-called gluonic poles. Their manifestation in the framework of the EFP approach requires the singular behavior of one of the correlators, whose origin is not yet clear enough.

The article is organized as follows. In section 2 the basic features of the twist-3 approach to single asymmetries are discussed and the results of some previous calculation are presented. Section 3 is devoted to the calculation of the non-Abelian subprocess and to the careful analysis of gauge invariance. The contributions of the gluonic and fermionic poles are compared in section 4. The asymmetry of dilepton production is calculated, manifesting a smooth interpolation between these poles. As a result, the correlator singularities, mentioned above, seem to be abandoned by the hadron density matrix positivity. The possible resolution of these controversies and the qualitative predictions for the observable asymmetries are presented in the conclusions.

2 Imaginary phase and single asymmetries in the twist–3 subprocesses

Consider a hard inclusive process with a transverse polarized nucleon. The term in the cross section proportional to twist-3 correlators can be expressed in the form [3]

$$d\sigma_s = \int dx_1 dx_2 \frac{1}{4} Sp[S_\mu(x_1,x_2)T_\mu(x_1,x_2)], \qquad (1)$$

where $S_\mu(x_1,x_2)$ is the coefficient function of parton subprocesses with two quark and one gluon legs (Fig. 1); $T_\mu(x_1,x_2)$ depends on

parton correlators:

$$T_\mu(x_1,x_2) = \frac{M}{2\pi}(\hat p_1 \gamma^5 s_\mu b_A(x_1,x_2) - i\gamma_\rho \epsilon^{\rho\mu s p_1} b_V(x_1,x_2)), \qquad (2)$$

where $\epsilon^{\rho\mu s p_1} = \epsilon^{\rho\mu\alpha\beta} s_\alpha p_{1\beta}$, s_μ is the covariant hadron polarization vector and M is the hadron mass. Two-argument distributions

$$b_A(x_1,x_2) = \frac{1}{M}\int \frac{d\lambda_1 d\lambda_2}{2\pi} e^{i\lambda_1(x_1-x_2)+i\lambda_2 x_2} \times$$
$$\times \langle p_1,s|\bar\psi(0)\hat n \gamma^5 (D(\lambda_1)s)\psi(\lambda_2)|p_1,s\rangle, \qquad (3)$$

$$b_V(x_1,x_2) = \frac{i}{M}\int \frac{d\lambda_1 d\lambda_2}{2\pi} e^{i\lambda_1(x_1-x_2)+i\lambda_2 x_2} \epsilon^{\mu s p_1 n} \times$$
$$\times \langle p_1,s|\bar\psi(0)\hat n D_\mu(\lambda_1)\psi(\lambda_2)|p_1,s\rangle \qquad (4)$$

are real and dimensionless. They possess symmetry properties which follow from T-invariance

$$b_A(x_1,x_2) = b_A(x_2,x_1),\ b_V(x_1,x_2) = -b_V(x_2,x_1). \qquad (5)$$

In the case of the standard parton picture the imaginary phase can appear only in the one-loop approximation. Consequently, any single asymmetry is proportional to α_s. The situation is different in the case of twist-3 subprocess. Besides the "usual" imaginary part of the propagator (Fig. 1a)

$$\frac{1}{x_1(s+u)+t+i\epsilon} = P\frac{1}{x_1(s+u)+t} - i\pi\delta(x_1(s+u)+t)$$

($(p_1x_2+p_2)^2 = x_2 s$, $\hat u = (p_1x_2-p_3)^2 = x_2 u$, $t = (p_2-p_3)^2$) which corresponds to cut on the missing mass $M_X^2 = (p_1+p_2-p_3)^2$, one should take into account another one. It arises, for instance, in the diagram of Fig. 1a where the gluon is attached to a Born subprocess. The imaginary part of the propagator

$$\frac{1}{x_2 s + i\epsilon} = P\frac{1}{x_2 s} - i\pi\delta(x_2 s) \qquad (6)$$

leads to the required imaginary phase in the hadron-quark amplitude. Consider, for definiteness, only the diagrams which have cuts in s or u to the left of the cut in M_X^2. The effect of mirror diagrams is just to cancel the real part of the amplitude and to double the imaginary one. The $u-$ cut corresponds to an antiquark contribution, as usual [3].

The single asymmetry generated by the twist-3 correlators does not depend on α_s, since α_s is included in the distribution function, that contains $< \psi g A \psi >$. This results in large (proportional not to α_s but to the hadron mass) polarization effects.

The twist-3 asymmetries for the direct photon production in the Compton subprocess were calculated a few years ago. The simplest one is the Abelian asymmetry for the process $\gamma N\uparrow\rightarrow \gamma X$[3]

$$A_{\gamma\gamma} = \frac{b_A(0,x) - b_V(0,x)}{f(x)} \frac{x_F(1-x_F)}{(1+x_F^2)} \frac{2Mp_T}{m_T^2} = \\ = \frac{b_A(0,x) - b_V(0,x)}{x f(x)} (1+x_F^2)^{-1} \frac{2Mp_T}{s}. \qquad (7)$$

Although the diagrams calculation results in the factor $C_F = (N^2 - 1)/2N$, it is included in the definition of the correlators [3]. Consequently, all the asymmetries below should be divided by the same factor. The simplest non-Abelian asymmetry for $gN\uparrow\rightarrow \gamma X$ is [5]

$$A_{g\gamma} = \frac{b_A(0,x) - b_V(0,x)}{f(x)} \times \\ \times \frac{x_F(1-x_F)(C_F - (x_F+1)C_A/2)}{C_F(1+x_F^2)} \frac{2Mp_T}{m_T^2} = \\ = \frac{b_A(0,x) - b_V(0,x)}{x f(x)} \frac{[C_F - (x_F+1)C_A/2]}{C_F(1+x_F^2)} \frac{2Mp_T}{s}. \qquad (8)$$

$x_F = -u/s, x = -t(s+u), m_T^2 = ut/s, C_A = N$. (In the articles [3, 5] the sign of x_F is wrong in some places; the correct result was first obtained in [6].).The quantity $f(x)$ is an "ordinary" spin-averaged quark distribution. Here and below the expressions for "raw" asymmetries are presented: all parts related to unpolarized hadrons are omitted. To pass to the hadron case, one should in (8) change $t, s \rightarrow ty, sy$ (y being the gluon momentum fraction) and integrate over y (separately!) the properly normalized numerator and denominator. Note that asymmetry (7) is the natural "partonometer" for the correlators (the additional y integration is absent). The coefficients of 1 and x_F in the expression (8) are proportional to the color factors of s– and u– channel diagrams $C_F - C_A/2$ and C_F, respectively. This is a consequence of the symmetry of the Abelian result under interchange of u and s.

3 Non-Abelian Compton subprocess and Gauge Invariance

Calculations of $A_{\gamma g}$ do not differ from calculations for $A_{g\gamma}$ and after all transformations we have

$$A_{\gamma g} = \frac{b_A(0,x) - b_V(0,x)}{f(x)} \times \\ \times \frac{(1-x_F)(C_F x_F - (x_F+1)C_A/2)}{C_F(1+x_F^2)} \frac{2Mp_T}{m_T^2} = \\ = \frac{b_A(0,x) - b_V(0,x)}{x f(x)} \frac{[C_F - (x_F+1)/x_F C_A/2]}{C_F(1+x_F^2)} \frac{2Mp_T}{s}. \qquad (9)$$

Note that the expressions in square brackets in the second equalities in (8) and (9) differ by interchange of x_F and $1/x_F$. It is natural since processes $gN\uparrow\rightarrow \gamma X$ and $\gamma N\uparrow\rightarrow gX$ differ by interchange of s and u. As we have now hadrons in the final state it is necessary to take into account fragmentation processes. However, it is possible to measure the asymmetry of the gluon jet. This allows one to avoid complications connected with hadronization of the final gluon.

The calculations of A_{gg} are more complicated because of the essentially non-Abelian nature of the gluon Compton subprocess. New diagrams of the type in Fig. 1b appear with the three-gluon vertex to the right of the cut in M_X^2. Calculating 18 diagrams (in two of which the colour factors are equal to zero.), instead of 10 for $A_{\gamma g}$ and $A_{g\gamma}$, we get

$$\begin{aligned} A_{gg} = & \frac{b_A(0,x) - b_V(0,x)}{f(x)} \frac{(C_F - C_A/2)(1-x_F)}{C_F(1+x_F^2)} \times \\ & \times \frac{((x_F^4+1)C_A/2 - C_F x_F(1-x_F)^2)}{(x_F(C_F - C_A/2) - (1+x_F^2)C_F/2)} \frac{M p_T}{m_T^2} = \\ = & \frac{b_A(0,x) - b_V(0,x)}{x f(x)} \frac{(C_F - C_A/2)}{C_F(1+x_F^2)} \times \\ & \times \frac{((x_F^3 + 1/x_F)C_A/2 - C_F(1-x_F)^2)}{(x_F(C_F - C_A/2) - (1+x_F^2)C_F/2)} \frac{M p_T}{s}. \end{aligned} \quad (10)$$

Note that the expressions (8), (9) and (10) reproduce the inclusive Compton asymmetry (7) in the "Abelian" limit $C_A = 0$, $C_F = 1$.

The calculations are performed using an axial-type gauge. In this gauge the gluon density matrix and the numerator of gluon propagator are

$$\rho_{\mu\nu}(k) = -g_{\mu\nu} + a\frac{k_\mu n_\nu + k_\nu n_\mu}{(kn)}, \quad (11)$$

where n is the gauge vector $((nA) = 0,\ n^2 = 0)$. The parameter a is introduced to control the result. In fact, we performed the calculations keeping it free. The b_A term is gauge invariant, because the Feynman rules for it generate on-shell amplitudes with almost standard external fermion lines. Therefore, it should not depend on a. This is not the case for the coefficient of b_V that is gauge-dependent and contains a polynomial a-dependence. Note that it does not mean the gauge dependence of observable quantities because the EFP factorization scheme is valid just for the axial gauge. One gets this gauge setting $a = 1$. The coefficient of b_V appears then to be equal, up to a sign, to that of b_A. As a result, the asymmetries are proportional to $b_A - b_V$. This combination in different processes (7)-(10) appears due to the positive t'–channel parity [5, 10] (t' corresponds to the forward scattering amplitude and, of course, is equal to zero) of the four-photon state. The change of two photons to two gluons in a color singlet state does not violate this property.

One may calculate the b_V term in an alternative way using a kinematical identity:

$$\epsilon^{\rho\mu s p_1} = -p_1^\rho \epsilon^{\mu s p_1 n} + p_1^\mu \epsilon^{\rho s p_1 n}. \quad (12)$$

The first term in the r.h.s. generates the Feynman rules with $\hat{p}_1$ on the external line [6, 7], just like for the on-shell particle. It is therefore natural that the corresponding term is a–independent and may be calculated in the Feynman gauge. The result of this calculation exactly coincides (for all the above mentioned processes) with the whole answer in the axial gauge. The remaining n–dependence through $\epsilon^{\mu s p_1 n}$ appears to be fictitious. To verify this, we performed the calculation for the most general n obeying the restrictions $n^2 = 0$, $(np_1) = 1$,

$$n^\mu = (2\frac{t}{su}\gamma(\gamma-1) + \beta^2/2)p_1^\mu + \frac{2}{s}\gamma p_2^\mu + \frac{2}{u}(\gamma-1)p_3^\mu + \beta s^\mu. \quad (13)$$

The dependence on free parameters γ, β really cancels out. In fact, one can expand E^μ as follows:

$$tr[E^\mu \hat{p}_1] = p_1^\mu E_1 + p_2^\mu E_2 + p_3^\mu E_3 + s^\mu E_s. \quad (14)$$

The combination of interest is then just

$$tr[E^\mu \hat{p}_1]\epsilon^{\mu s p_1 n} = \epsilon^{p_1 p_2 p_3 s}(E_3 \frac{2}{s}\gamma - E_2 \frac{2}{u}(\gamma - 1)) \quad (15)$$

A straightforward calculation shows that the $n-$ dependence cancels due to the gauge-invariant relation $E_3 = E_2 s/u$. It is equivalent to the following one:

$$tr[E^\mu \hat{p}_1]p_1^\mu = 0. \qquad (16)$$

It is just the collinear on-shell Ward identity, and it should be valid in the general case for a rather simple reason. Generally speaking,

$$E^\mu(x_1, x_2)p_1^\mu = \frac{E(x_1) - E(x_2)}{x_1 - x_2}. \qquad (17)$$

Consider the double cut in x_1, x_2 required by the single asymmetry (Sect. 2). The r.h.s. is zero , at least for $x1 \neq x2$. The case of equal x's will be considered in the next Section.

To clarify the $n-$independence, consider the following ansatz for E^μ compatible with (16)

$$tr[E^\mu \hat{p}_1] = \epsilon^{\mu p_1 \alpha\beta} E_{\alpha\beta} \qquad (18)$$

Substituting the latter into (15) and expanding the product of two ϵ's, one should note that n enters only into the scalar product $(pn) = 1$, clearly manifesting the $n-$independence.

All the gauge dependence should therefore be attributed to the contribution of the second term in (12). The direct calculation really leads to the $n-$dependent expression. However, it is equal to zero (for any n) in the axial gauge used to derive the factorized formula (1). This seems quite natural because the collinear Ward identity is again applicable. Note the important difference between the two terms in the r.h.s. of (12): the application of the Ward identity leads to the $n-$independence of the first one and to the vanishing of the second one.

4 Fermionic and gluonic poles

All the presented results are related to contributions of "fermionic poles" contributions (diagrams of the type in Fig.1). They come

from the phase space region in which the hadron momentum fraction carried by a quark tends to zero ($x_2 \to 0$). These results coincide, up to definition of correlators and kinematical variables, with similar contributions calculated by Qiu and Sterman [6, 7]. In their work "gluonic poles" were also introduced (diagrams of Fig. 1c type). They are related to the phase space region in which the hadron momentum fraction carried by a gluon tends to zero ($(x_1 - x_2) \to 0$). These "gluonic poles" are considered as the main contribution to asymmetry. The contribution of "gluonic poles" appears to be equal to zero in our approach, however.

In the paper [7] it is proportional to the correlator $T_V^F(x, x)$:

$$T_V^F(x_1, x_2) = \int \frac{d\lambda_1 d\lambda_2}{4\pi} e^{i\lambda_1(x_1 - x_2) + i\lambda_2 x_2} \epsilon^{\mu s p_1 n} \times$$
$$\times \langle p_1, s|\bar\psi(0)\hat{n} F_{\mu n}(\lambda_1)\psi(\lambda_2)|p_1, s\rangle \qquad (19)$$

It does not appear in our approach because the use of axial gauge immediately leads to:

$$g T_V^F(x_1, x_2) = \frac{x_1 - x_2}{2} M b_V(x_1, x_2). \qquad (20)$$

Although the Feynman gauge is adopted in [7], all these correlators are gauge invariant, because the change of the gluon strength tensor F under the gauge transformation does not influence the color–singlet matrix element related to T_F.

One can get a nonzero result if $b_V(x_1, x_2)$ has a pole at $x_1 = x_2$. To avoid the manipulations with the potentially dangerous quantity $b_V(x, x)$, we calculated the asymmetry for the dilepton pair photoproduction, i.e. $A_{\gamma\gamma}$ with a virtual final photon ($p_3^2 = Q^2 > 0$). It allows a smooth interpolation between fermionic and gluonic poles and may be considered as a "poles partonometer". Averaging over lepton angles we have:

$$A_{\gamma\gamma^*} = \frac{x(1-y)}{x_F(1 - y - x_F)[(x(1-y) - y x_F)^2 + x^2 x_F^2]} \frac{1}{f(x)} \frac{2 M p_T}{s} \times$$

$$\times[(b_A(0,x)-b_V(0,x))(x_F(1-y)^2-\frac{y}{x}x_F^2(2(1-y)+x_F)))+$$
$$+(b_A(0,x)-\frac{y}{x}b_V(0,x))2\frac{y}{x}x_F^3+ \tag{21}$$
$$+(1-y)[(b_A(y,x)+b_V(y,x))\frac{y}{x}(1-y-x_F)^2-$$
$$-(b_A(y,x)-b_V(y,x))x_F^2]],$$

where $y=Q^2/(Q^2-u)$. The u-channel diagrams contribute with the correlators whose arguments $x,y\neq 0$. If $y\to x$, the pole in $b_V(x,y)$ results in the infinite rise of asymmetry. This contradicts the positive definiteness of the density matrix. Note that the coefficients of b_A and b_V are no more equal because the virtual photon does not have a definite parity.

This "proof" of the nonexistence of gluonic poles is qualitatively close to the renormalization group approach. Some restrictions on the large distance contribution are imposed by the short-distance contribution. To understand this result better, we calculated the contribution of gluonic poles for all the mentioned processes. The result appears to be $n-$ dependent.

The application of the identity (12) leads to decomposition of the gluonic pole contribution. We pay the most attention to the analysis of the first term. It is gauge invariant ($a-$dependence cancels out because E^μ is projected onto $\hat{p}_1$) but $n-$dependent. This means that the on-shell coefficient function is not orthogonal to p contrary to the case of fermion poles (16). It may be caused only by the fact that the r.h.s. of the equation

$$tr[E^\mu(x_1,x_2)\hat{p}_1]p_1^\mu=\frac{tr[E(x_1)\hat{p}_1]-tr[E(x_2)\hat{p}_1]}{x_1-x_2}. \tag{22}$$

obtains a double cut; one of the cuts is produced by x_1-x_2 in the denominator. However, it is possible to treat poles like that by using the principal value prescription [11], resulting in the zero double cut. The n-dependence may be considered as another argument in favor of the nonexistence of gluonic poles.

Note that a pole of eq. (22) type appears directly in the expression for the gluonic pole hard scattering part, calculated by Qiu and Sterman [7]. The coefficient of the $\delta-$function derivative (see eq. (5.11a) of Ref. [7]) is proportional to the l.h.s. of (22). Changing it by the r.h.s., one may reproduce the result [7] if and only if one of the cuts is produced by the denominator. Note that it should then be proportional to the Born spin-averaged cross-section of the Compton effect because the hadron polarization completely decouples. Really, formula (5.11b) of Ref. [7] manifests the familiar expression.

If, however, the pole must really be treated by using the principal value prescription as mentioned above, the corresponding contribution of gluonic poles turns to zero! The remaining term in eq. (5.11a) [7], proportional to the $\delta-$function itself, may also be transformed by using the Ward identities. They are, however, more complicated and require further investigation.

Consider now the second term in (12). It also results in the $n-$dependent expression. Moreover, it appears to be divergent when $\beta\to 0,\gamma\to 1$, i.e. n becomes directed along the incoming gluon p_2. This divergence may be cancelled by the additional pole coming from the gauge-dependent piece of the gluon propagator [12].

A possible qualitative reason for the absence of the gluonic poles is as follows: the dominance of the gluonic poles was related [7] to the fact that "soft gluons are emitted much more readily than soft quarks". However, this emission is known to result in the infrared divergence and may somehow cancel for the observable quantities.

5 Discussion and Conclusions

The calculated twist-3 asymmetries are not the corrections, like the twist-4 ones in the spin averaged case, but the leading order contributions.

Our results provide some qualitative predictions for the observable asymmetries. The single asymmetries for pion (jets) production predominate over the direct photon asymmetries in the region of target fragmentation ($x_F \sim 0$). It can be seen from (7), (8), (9) and (10) that $A_{\gamma\gamma}$ and A_{gg} differ in sign from $A_{\gamma g}$ and $A_{g\gamma}$. The dependences on x_F of A_{gg} and $A_{g\gamma}'$ (the kinematical factor $2Mp_T/m_T^2$ is omitted) are shown in Fig. 2.

The asymmetry in the lepton pair photoproduction is of especial interest as a direct probe of quark gluon correlators in a wide range of its arguments. It is possible to perform such experiments at HERA (HERMES) and CERN (HELP) in the nearest future. This asymmetry seems to manifest that the existence of gluon poles contradicts the nucleon density matrix positivity.

Another argument against gluon poles comes from the fact that their contribution appears to depend on the vector n fixing the axial gauge and transverse direction. This property, in its turn, is related to the nonzero double cut of the leading twist coefficient function. The latter seems to be fictitious because the relevant pole resulting from the collinear Ward identities may be treated by using the principal value prescription. This treatment nullifies at least one of the gluonic poles contributions calculated by Qiu and Sterman.

Here one meets some non-commutativity in calculating the imaginary part and applying the Ward identity. If one first calculate the double cut of E^μ, the contraction with p^μ results in a nonzero expression. The latter is related to the imaginary part of the denominator produced by the collinear Ward identity. It is possible, however, to apply the Ward identity to the amplitude in the Euclidian region and to perform the analytic continuation afterwards. It seems that it is natural to treat the mentioned denominator by the principal value prescription leading to a zero gluon pole contribution. In our opinion, the second recipe is supported by the fundamental role of Ward identities [13]. From a more formal point of view, there is just the non-commutativity of the double cut and limit $x_1 - x_2 \to 0$. As this limit is just the result of double cut, the

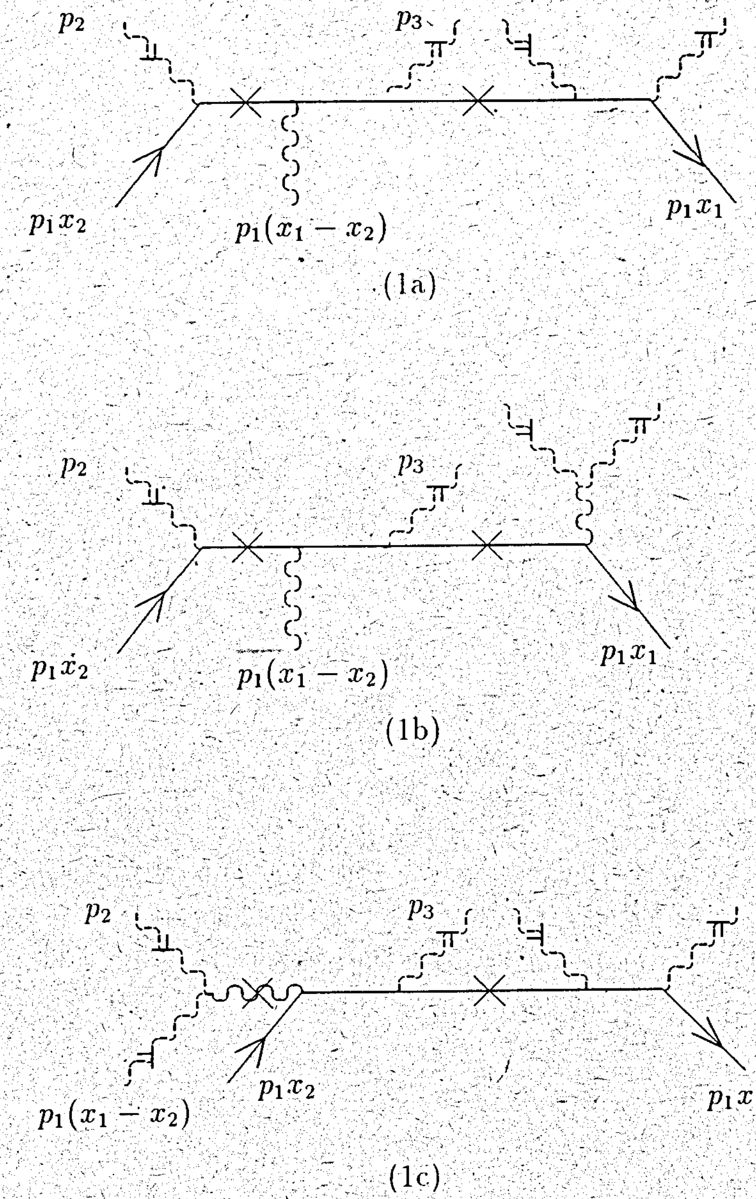

Figure 1: The typical diagrams with "fermionic" (a,b) and "gluonic" (c) poles.

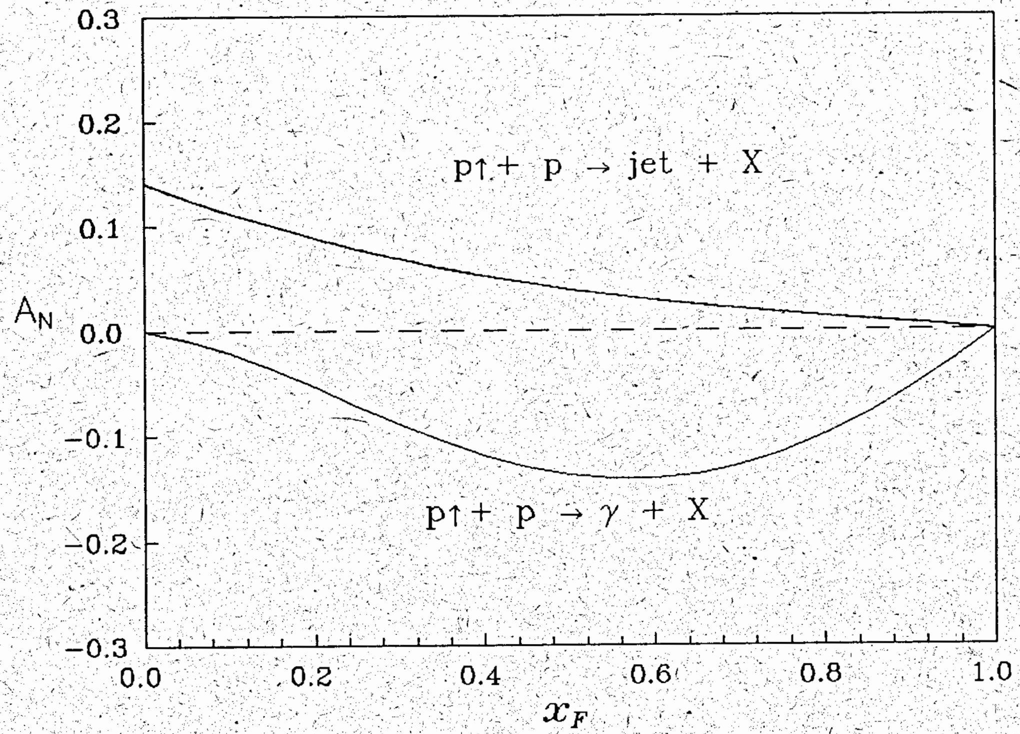

Figure 2: The parton asymmetries for the processes of direct γ and jets inclusive production.

latter may be performed first, leading again to a zero result.

However, the final resolution of the puzzle of gluon poles requires further investigation. The planned and proposed experiment can also help because the results depend strongly on the gluonic poles existence. It is very interesting that such fine and intimate properties of QCD can, in principle, be tested experimentally.

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Комптоновские асимметрии твиста-3

Вычислен вклад «фермионных полюсов» в одиночную асимметрию твиста-3 глюонного комптоновского процесса. Существование «глюонных полюсов» противоречит положительной определенности матрицы плотности. Представлены качественные предсказания для асимметрий прямых фотонов и струй.

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On the Compton Twist-3 Asymmetries

The «fermionic poles» contribution to the twist-3 single asymmetry in the gluon Compton process is calculated. The «gluonic poles» existence seems to contradict the density matrix positivity. Qualitative predictions for the direct photon and jets asymmetries are presented.

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

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