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ON THE COMPTON NON-ABELIAN TWIST-3 ASYMMETRY

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It is known that large asymmetries appear in processes with a single polarized particle. For instance, asymmetries observed in pion production with high transverse momentum are of the order of ten or more percent[1]. Because of this fact single asymmetries can be very convenient for the study of polarization effects in QCD. However, for their appearance a mass parameter and an additional imaginary phase are necessary. As a result, these asymmetries are absent in the leading approach of QCD and for their calculation one has to take into account high twist contributions. In this letter we calculate single asymmetries for processes $\gamma N\uparrow\rightarrow gX$ and essentially non-Abelian $gN\uparrow\rightarrow gX$ and give some qualitative predictions for processes which can be studied in existing and future accelerators.

Some years ago transverse polarization in quantum chromodynamics was investigated in the framework of the Ellis, Furmansky, Petronsio factorization scheme [2]. The hadron density matrix in QCD was analized to establish that the mass of a polarized hadron is the true mass parameter of transverse polarization [3]. It was shown that contributions to single asymmetries proportional to two-argument distribution functions (parton correlation densities) can receive imaginary parts even from Born subprocess [4]. Recently similar results were obtained by G. Sterman and J. Qiu [6] who used the so-called special propagators technique. The correspondence and discrepancy between these approaches will be discussed below.

The term in the cross section proportional to correlators can be expressed in the form [4]

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

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

$$T_\mu(x_1,x_2) = \frac{M}{2\pi}(\hat{p}_1\gamma^5 s_\mu b_A(x_1,x_2) + i\gamma_\rho\epsilon^{\rho\mu s p_1} b_V(x_1,x_2)), \tag{2}$$

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

$$b_A(x_1,x_2) = b_A(x_2,x_1), \qquad b_V(x_1,x_2) = -b_V(x_2,x_1). \tag{3}$$

In the case of the standard parton picture the imaginary part can appear only in one-loop approximation, and consequently asymmetry is proportional to α_s. The situation is different in the case of correlators. Besides the usual imaginary part of the propagator (Fig. 1a)

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

($(p_1x_2+p_2)^2 = x_2s$, $\hat{u}=(p_1x_2-p_3)^2 = x_2u$, $t=(p2-p3)^2$) which corresponds to cut on M_X^2, it is necessary to take into account the additional imaginary part which appears for instance in the diagram of the Fig. 1a , where the gluon is attached to a Born subprocess. The imaginary part of the propagator

$$\frac{1}{x_2s+i\epsilon} = P\frac{1}{x_2s} - i\pi\delta(x_2s) \quad (4)$$

leads to the appearance of an additional imaginary part in the hadron-quark amplitude. We consider diagrams which have cuts in s or u to the left of the cut in M_X^2. The $u-$ cut corresponds to an antiquark contribution, as usual [4].

The asymmetry in this case does not depend on α_s, since α_s is included in the distribution function, which contains $< \bar\psi g A\psi >$. This results in large (not proportional to α_s but proportional to the hadron mass) single asymmetries. The twist-3 asymmetries of the direct photons produced in the Compton subprocess were calculated a few years ago. The Abelian asymmetry for the process $\gamma N\uparrow\to\gamma X$ [4] is

$$A_{\gamma\gamma} = \frac{b_A(0,x)-b_V(0,x)}{f(x)}\frac{x_F(1-x_F)}{(1+x_F^2)}\frac{2Mp_T}{m_T^2} = \frac{b_A(0,x)-b_V(0,x)}{xf(x)}(1+x_F^2)^{-1}\frac{2Mp_T}{s}. \quad (5)$$

The non-Abelian asymmetry for $gN\uparrow\to\gamma X$ [5] is

$$A_{g\gamma} = \frac{b_A(0,x)-b_V(0,x)}{f(x)}\frac{x_F(1-x_F)(C_F-(x_F+1)C_A/2)}{(1+x_F^2)}\frac{2Mp_T}{m_T^2} = \frac{b_A(0,x)-b_V(0,x)}{xf(x)}\frac{[C_F-(x_F+1)C_A/2]}{(1+x_F^2)}\frac{2Mp_T}{s}. \quad (6)$$

$x_F=-u/s, x=-t(s+u), m_T^2 = ut/s, C_A = N, C_F = (N^2-1)/2N$. (In the articles [4, 5] the sign of x_F is wrong in some places). $b_A(0,x)$, $b_V(0,x)$ and $f(x)$ are quark-gluon correlators and "ordinary" quark distribution, respectively. 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 change $t,s\to ty, sy$ (y being the gluon momentum fraction) and integrate over y (separately!) the properly normalized numerator and denominator. If we have 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. Note that asymmetry (5) is the natural "partonometer" for the correlators (the additional y integration is absent). The coefficients of 1 and x_F in the expression (6) 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.

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

$$A_{\gamma g} = \frac{b_A(0,x) - b_V(0,x)}{f(x)} \frac{(1-x_F)(C_F x_F - (x_F+1)C_A/2)}{(1+x_F^2)} \frac{2Mp_T}{m_T^2} =$$
$$\frac{b_A(0,x) - b_V(0,x)}{xf(x)} \frac{[C_F - (x_F+1)/x_F C_A/2]}{(1+x_F^2)} \frac{2Mp_T}{s}. \quad (7)$$

Note that the expressions in square brackets in the second equalities in (6) and (7) 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.

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 three-gluon vertex to the right of the cut on 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

$$A_{gg} = \frac{b_A(0,x) - b_V(0,x)}{f(x)} \frac{(C_F - C_A/2)(1-x_F)}{(1+x_F^2)} \times$$
$$\times \frac{((x_F^4+1)C_A/2 - C_F x_F(1-x_F)^2)}{(x_F(C_F - C_A/2) - (1+x_F^2)C_F/2)} \frac{2Mp_T}{m_T^2} =$$
$$\frac{b_A(0,x) - b_V(0,x)}{xf(x)} \frac{(C_F - C_A/2)}{(1+x_F^2)} \times$$
$$\times \frac{((x_F^3+1/x_F)C_A/2 - C_F(1-x_F)^2)}{(x_F(C_F - C_A/2) - (1+x_F^2)C_F/2)} \frac{2Mp_T}{s}. \quad (8)$$

Note that the expressions (6), (7) and (8) reproduce the inclusive Compton asymmetry (5) 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

$$\rho_{\mu\nu}(k) = -g_{\mu\nu} + a\frac{k_\mu n_\nu + k_\nu n_\mu}{(kn)}. \quad (9)$$

where n is the normal vector ($(n \cdot A) = 0$, $n^2 = 0$). 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, which is gauge-dependent and equal, up to a sign, to that of b_A, provided the axial gauge is adopted. Setting $a = 1$, we get the axial gauge and structure $b_A - b_V$. The appearance of the same combination $b_A - b_V$ in different processes is related to the positive t'-channel parity [5, 7] (t' corresponds to the forward scattering amplitude and, of course, is equal to zero). The substitution of two gluons in a color singlet state instead of two photons does not violate this fact. Therefore, neither single $V-$ nor $A-$ correlators, but their definite combination should appear in the result. All the formulas (5)-(8) clearly manifest this property.

All the presented results are related to "fermionic poles" contributions (diagrams of type on Fig.1). It comes from the phase space region in which hadron momentum fraction carried by quark tends to zero ($x_2 \to 0$). These results coincide, up to definition of correlators and kinematical variables, with the "fermionic poles" contributions of Qiu and Sterman [6]. In their work "gluonic poles" were introduced (diagrams of Fig. 1c type). They are related to the phase space region in which hadron momentum fraction carried by gluon tends to zero ($(x_1 - x_2) \to 0$). These authors believe that "gluonic poles" give the main contribution to asymmetry. However the "gluonic poles" contribution appears to be equal to zero in our approach.

Our calculations give for the "gluonic poles" contributions results (for all processes) which are proportional to $b_V(x,x)$. In the paper [6] they are proportional to the correlators $T_V^F(x,x)$ which have no counterpart in our axial-gauge based approach. However, although the Feynman gauge is adopted in [6], all these correlators are gauge invariant, because the change of the gluon strength tensor F under the gauge transformation does not contribute to the color singlet matrix element related to T_F. The use of the equations of motion lead to $T_V^F(x_1,x_2) \sim (x_1 - x_2)T_V^D(x_1,x_2) \sim (x_1 - x_2)b_V(x_1,x_2)$. Consequently, the coefficients attached to b_V must be equal to zero for coincidence with results of [6] (and we can get nonzero result if $b_V(x_1,x_2)$ has a pole at $x_1 = x_2$). However, these coefficients are nonzero. Therefore, we believe that b_V has no pole and $b_V(x,x) = 0$ (otherwise, we get a meaningless infinite result). We conclude that either the Ellis, Furmanski, Petronzio approach is inadequate for the gluonic poles contribution, or the latter is absent altogether.

Our results allow us to give some qualitative predictions for the obtained 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 (5), (6), (7) and (8) that $A_{\gamma\gamma}$ and A_{gg} differ by a sign from $A_{\gamma g}$ and $A_{g\gamma}$. The dependencies on x_F of A_{gg} and $A_{g\gamma}$ (the kinematical factor Mp_T/m_T^2 is omitted) are shown in Fig. 2. These predictions can probably be seen experimentally.

The planned polarization program of the STAR collaboration at RHIC is most adapted for measuring the direct photon and jets asymmetries [8]. As a result one can check new important predictions of QCD and resolve its controversies.

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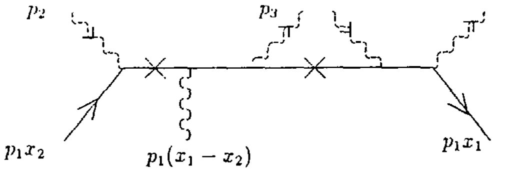
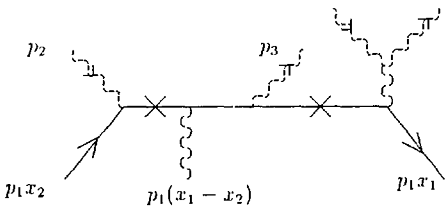


(1a)



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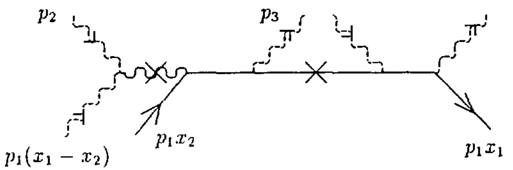

(1c)

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

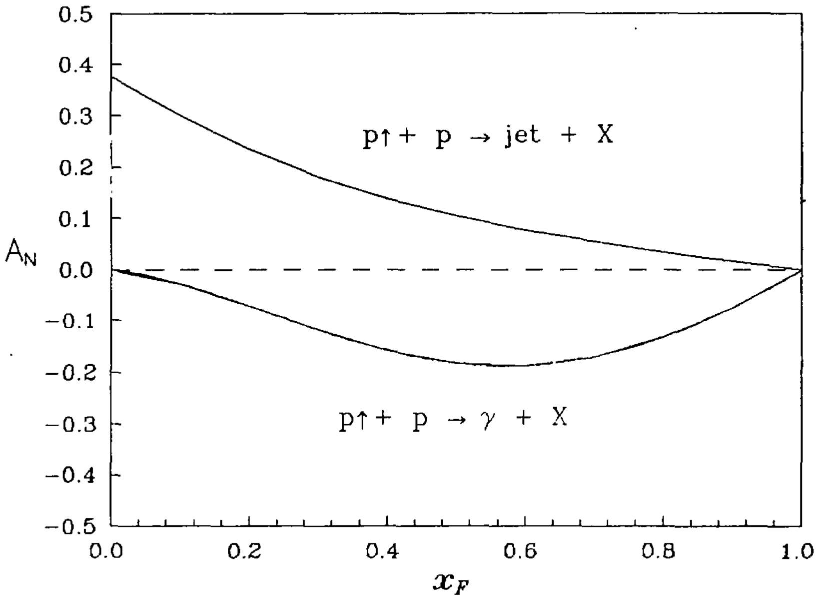

Figure 2. The parton asymmetries for the processes of direct γ and jets inclusive production.

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