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STUDY OF FACTORIZATION IN QCD WITH POLARIZED BEAMS AND Λ PRODUCTION AT LARGE p_T *

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ABSTRACT

We discuss what aspects of the leading orders in QCD could be tested in the large p_T reactions $\vec{p}\, p \to \vec{\Lambda} + X$, $\bar{p}\, \vec{p} \to \vec{\Lambda} + X$ and $\vec{\bar{p}}\, \vec{p} \to \Lambda + X$.

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Recently there has been a considerable interest in polarized proton(anti proton) physics extended to large p_T, since this provides an additional means of testing QCD. In particular it will provide valuable information about factorization beyond the leading logarithmic order. Already much work (1) has been done on what one might call reflected asymmetries A_{LL}^{ii}, which refers to both beam and target polarized longitudinally. This is depicted in Fig. 1 (a).

Recently, experiments (2) have demonstrated that the polarization of Λ's can be detected at transverse momentum beyond 2 GeV/c. Together with a polarized $\bar{p}$ beam or polarized proton target, this will provide a valuable mean of measuring transmitted polarization A_{LL}^{if} (see Fig. 1 (b)), which we shall argue can also be calculated in the leading order of QCD.

The parton model has been used with some striking success in a large number of processes including hadron production at large transverse momentum $(AB \to C + X)$. The main feature of the naive parton model is the factorization of parton densities (distribution inside hadrons and fragmentation into hadrons). In the last two years or so it has been demonstrated (3,4) that this factorization holds in the leading order in QCD and in some sense even beyond (3) the leading order.

The naive parton model factorization corresponds to being able to split the discontinuity associated with the hard process in question, blockwise into the Green's functions connected with the various parton densities and elementary cross-sections, the whole process being linked together through the parton momenta. This fact allows us to simply extend the parton analysis to spin asymmetries at short distances, showing the way spin information is transmitted between the different parts of the process. Further, at least in the leading order in QCD, which maintains the block-wise factorization, the same formulae hold, with the appropriate replacements of parton densities by the scale dependent counterparts. The latter satisfy evolution equations governed by anomalous dimensions. The way this works is reviewed in reference (5) and we only briefly summarize it here. Consider an arbitrary short distance process involving at least two visible hadrons A and B, with B in the initial or final state. Denoting the associated parton densities by $D_A^a(x_a, s_A, s_a)$ and $D_B^b(x_b, s_B, s_b)$ respectively, one can write the cross-section in the form

$$\sigma^{AB\ldots}(p_A, s_A, p_B, s_B; \ldots) = \sum_{a, b, \ldots} \sum_{s_a, s_b, \ldots} \int dx_a \int dx_b \ldots D_A^a(x_a, s_A, s_a) \tag{1}$$

$$D_B^b(x_b, s_B, s_b) \ldots \sigma^{ab\ldots}(x_a p_A, s_a; x_b p_B, s_b, \ldots) .$$

where the corresponding elementary parton cross-section $\sigma^{ab\ldots}$ is evaluated to lowest order in QCD and one integrates over the available parton phase space. By defining the helicity differential or asymmetry $\Delta_a D_A^a(s_a) \equiv \frac{1}{2}\left[D_A^a(\rightarrow) - D_A^a(\leftarrow)\right]$ this cross section can be written in the form

$$\frac{1}{4}\sigma^{AB}(s_A, s_B, \ldots) = \int D_A^a(s_A)\, D_B^b(s_B) \ldots \sigma^{ab\ldots} \tag{2}$$

$$+ \int \Delta_a D_A^a(s_A, s_a)\, D_B^b(s_B) \ldots \Delta_a \sigma^{ab\ldots}(s_a, \ldots)$$

$$+ \int D_A^a(s_A)\, \Delta_b D_B^b(s_B, s_b) \ldots \Delta_b \sigma^{ab\ldots}(s_b, \ldots)$$

$$+ \int \Delta_a D_A^a(s_A, s_a)\, \Delta_b D_B^b(s_B, s_b) \ldots \Delta_a \Delta_b \sigma^{ab\ldots}(s_a, s_b, \ldots) .$$

Only the first term in Eq. (2) survives in the leading order of QCD since the remaining terms involve single asymmetries of collinear parton-hadron systems and/or single asymmetries of the Born parton cross sections. The double spin differential of this cross section is in leading order given by

$$\frac{1}{4}\Delta_A \Delta_B \sigma^{AB\ldots}(s_A, s_B, \ldots) = \int \Delta_A \Delta_a D_A^a(s_A, s_a)\, \Delta_B \Delta_b D_B^b(s_B, s_b) \ldots \Delta_a \Delta_b \sigma^{ab\ldots}(s_a, s_b, \ldots) \tag{3}$$

where

$$\Delta_a \Delta_b \sigma(s_a, s_b) = \frac{1}{4}\left[\sigma(\rightarrow, \rightarrow) - \sigma(\rightarrow, \leftarrow) - \sigma(\leftarrow, \rightarrow) + \sigma(\leftarrow, \leftarrow)\right] \equiv \Delta_2 \sigma . \tag{4}$$

Equation (3) corresponds to last term in Eq. (2), the other giving non-leading contributions. The double spin differentials of the density functions satisfy the evolution equations

$$\Delta_A \Delta_a D_A^a(x, Q^2, s_A, s_a) = \Delta_A \Delta_a D_A^a(x, Q_0^2, s_A, s_a) + \int_{Q_0^2}^{Q^2} \frac{dk^2}{k^2} \frac{\alpha_s(k^2)}{2\pi} \int_x^1 \frac{dz}{z}\, \Delta_a \Delta_b P_{ab}(z, s_a, s_b)\, \Delta_A \Delta_b D_A^b\left(\frac{x}{z}, k^2, s_A, s_b\right) \tag{5}$$

where the branching kernels are those of Ref. 6 and will be given below for the cases we shall be interested in. $D_A^a(x, Q_0^2, s_A, s_a)$ is the primordial distribution defined at some chosen scale Q_0^2.

Turning to the reaction $\bar{p}p \rightarrow \Lambda + \bar{S} + X$, where $\bar{S}$ is an optional away-side strange particle $(\bar{\Lambda}, \bar{K}, \ldots)$ trigger, we have at large p_T the mechanisms shown in Fig. 2. At least as we go to large values of $x_T = 2p_T/\sqrt{s}$, the $q\bar{q}$ annihilation mechanism in Fig. 2 (a) will be the primary source of strange particles, because of the dominance of valence quarks as the parton momenta x go to unity. At lower x_T however, the gluon-gluon mechanism (Fig. 2(b)) will be a non-negligible background. Another source of strange quark jets is the hard scattering off an s quark in the sea distribution of the $\bar{p}$ or p (see Fig. 2(c)). However this again is only likely to be important at very small x_T . In any case the accompanying strange particle $\bar{S}$ will be in the beam-target jet system, so such a background could be eliminated by using the away side $\bar{S}$ trigger option . Finally gluons can fragment into a pair of strange quarks, so the mechanism in Fig. 2(d), could also produce a Λ

or a $\bar{\Lambda}$. This kind of process involves pair creation and subsequent fragmentation into a pair of heavy strange particles; which will also have the effect of pushing the Λ into the small x_T region. Further this background could also be eliminated by the away side trigger option.

The principal mechanism in Fig. 2(a) leads to a very large transmitted asymmetry A_{LL}^{if} for $\vec{\bar{p}}\vec{p} \to \vec{\Lambda} + X$ or $\vec{\bar{p}}p \to \vec{\Lambda} + X$ in the leading order in QCD, given by the formula (written for the first case, the other case being simply related)*

$$A_{LL}^{if} = \frac{\sum_{q=u,d} \int \Delta_z D_p^q(x_a, Q^2)\; D_{\bar{p}}^{\bar{q}}(x_b, Q^2)\; \Delta_z \bar{D}_s^{\Lambda}(x_c, Q^2)\; a_{LL}^{if} \left(\frac{d\hat{\sigma}}{d\hat{t}}\right)_{q\bar{q}\to s\bar{s}}}{\sum_{q=u,d} \int D_p^q(x_a, Q^2)\; D_{\bar{p}}^{\bar{q}}(x_b, Q^2)\; \bar{D}_s^{\Lambda}(x_c, Q^2) \left(\frac{d\hat{\sigma}}{d\hat{t}}\right)_{q\bar{q}\to s\bar{s}}} \qquad (6)$$

* $A_{LL}^{if}(\vec{\bar{p}}p \to \vec{\Lambda} + X) = - A_{LL}^{if}(\vec{\bar{p}}\vec{p} \to \vec{\Lambda} + X)$.

where the integration over the parton phase space is given by $\int dx_a\, dx_b\, dx_c^{-1}\, \hat{s}/\pi \;\; \delta(\hat{s} + \hat{t} + \hat{u})$, with $\hat{s} = x_a x_b s$, $\hat{t} = x_a/x_c\, t$ and $\hat{u} = x_b/x_c\, u$ $\left(t = (p_a - p_\Lambda)^2 \right.$ and $\left. u = (p_b - p_\Lambda)^2 \right)$.

The value Q^2 is taken to be of order p_T^2 in the estimate we make . However its choice is clearly dependent on the role of the next to leading order logarithmic corrections to which we shall turn in a moment. The basic Born $q\bar{q} \to s\bar{s}$ cross section is given by

$$\left(\frac{d\hat{\sigma}}{d\hat{t}}\right)_{q\bar{q}\to s\bar{s}} = \frac{4\pi}{9} \frac{1}{\hat{s}^2} \alpha_s^2(Q^2) \left(\frac{\hat{t}^2 + \hat{u}^2}{\hat{s}^2} \right) \qquad (7)$$

and the transmitted asymmetry is

$$a_{LL}^{if} = \frac{\hat{t}^2 - \hat{u}^2}{\hat{t}^2 + \hat{u}^2} \qquad (8)$$

The corresponding reflected parton asymmetry is $a_{LL}^{ii} = -1$ whereas a_{LL}^{if} varies anti-symmetrically between the values -1 and $+1$ as a function of $-\hat{t}/\hat{s}$ (see Fig. 3).

The helicity differentials of the parton densities will be principally determined by the primordial value at some chosen scale Q_0^2 . As yet, we know nothing about these functions except for models based on phenomenological constraints. However they are directly measurable in experiments like $\vec{e}\vec{p} \to e' X$ and in the case of the fragmentation functions in $\vec{e}^+ e^- \to \vec{\Lambda} + X$ or $e\vec{p} \to e'\vec{\Lambda} + X$. For a detailed discussion we refer to Refs. 5 and 7. At this point it is worth mentioning that from the very crude approximation in which we write the asymmetry at fixed x_T in the form $\langle A_{LL}^{if} \rangle \simeq \left(\langle \Delta_z D_p^q \rangle / \langle D_p^q \rangle\right)\left(\langle \Delta_z \bar{D}_s^{\Lambda} \rangle / \langle \bar{D}_s^{\Lambda} \rangle\right) \langle a_{LL}^{if} \rangle$ we see that if the measured asymmetry turns out to be large then the spin differentials of both the valence quarks distributions inside the proton (i.e. antiproton) and the fragmentation function

$s \to \Lambda$ must necessarily also be large.

The models we have used to make estimates are the so called conservative model of Sivers et al.[1] and the Carlitz-Kaur[8] model. The former is based in the idea that the leading valence quark carries all the helicity (which is presumably true only as $x \to 1$) and is normalized according to the Bjorken sum rule $2\int_0^1 dx\ (\Delta_2 D_p^u - \Delta_2 D_p^d + \Delta_2 D_p^{\bar u} - \Delta_2 D_p^{\bar d}) = G_A/G_V$ (i.e. the ratio of axial vector and vector coupling constants).

We shall take the ansatz in Ref. 1, namely $\Delta_2 D_p^u = .44(D_p^u)_{valence}$ and $\Delta_2 D_p^d = -.35\ (D_p^d)_{valence}$. In the Carlitz-Kaur model, on the other hand, the valence quarks loose their polarizations as $x \to 0$ through interactions with the sea. This has the effect of polarizing the gluons. The detailed form we used is given in Ref. 9, i.e. $\Delta_2 D_p^u = \left(D_p^u - \frac{2}{3} D_p^d\right)_{valence}$ and $\Delta_2 D_p^d = -\frac{1}{3}\left(D_p^d\right)_{valence}$. We have no models for the spin differential of the fragmentation $s \to \Lambda$. However again one expects the leading s quark to give all its helicity to the Λ as $x \to 1$. We shall make the simple ansatz $\Delta_2 \bar D_s^\Lambda = \bar D_s^\Lambda$ *. The curve we plot can be scaled down according to one's choice for $\Delta_2 \bar D_s^\Lambda$ relative to $\bar D_s^\Lambda$. The above primordial distributions evolve away from Q_0^2 with the following kernels in Eq. (5)

$$\Delta_2 P_{qq}(z) = \left(\frac{4}{3}\right)\left[\frac{1+z^2}{(1-z)_+} + \frac{3}{2}\ \delta(z-1)\right] ,$$
$$\Delta_2 P_{qG}(z) = \left(\frac{1}{2}\right)\left[z^2 - (1-z)^2\right] \qquad (9)$$

By virtue of the fact that the driving asymmetry of the subprocess is antisymmetric about $\theta_{trigg.} = 90°$ in the CM , the full asymmetry also vanishes at $\theta_{trigg.} = 90°$ and is a non trivial function of $\theta_{trigg.}$. In Fig. 4(a) and (b) we give the asymmetry as a function of p_T for respectively $\theta_{trigg.} = 60°$ and $\theta_{trigg.} = 30°$.

* For the fragmentation function $\bar D_s^\Lambda$ we assume a form $x\bar D_s^\Lambda(x) \sim (1-x)^3$.

In each case we show the predictions of the two extreme models we discussed above. In Figs. 4(a) and (b) we have not shown the effect of the gluon contribution (Fig. 2(b)) which depletes the asymmetry. This effect is shown for the conservative model in Fig. 4(c) as a function of x_T , where as expected the asymmetry vanishes as $x_T \to 0$.

If in addition to the transmitted asymmetry , the reflected asymmetry is measured with the polarized proton i.e. $\vec p \vec p \to \Lambda + X$, then we have some additional information. In particular $a_{LL}^{ii} = -1$ for the $q\bar q \to s\bar s$ mechanism so the overall asymmetry is proportional to $\langle(\Delta_2 D)^2\rangle / \langle D^2\rangle$ and

$$\frac{A_{LL}^{if}}{A_{LL}^{ii}} \sim -\frac{\langle D_p\rangle}{\langle\Delta_2 D_p\rangle}\frac{\langle\Delta_2 \bar D_s^\Lambda\rangle}{\langle\bar D_s^\Lambda\rangle} \qquad \langle a_{LL}^{if}\rangle \qquad (10)$$

Hence, in an average sense this is proportional to the transmitted asymmetry a_{LL}^{if}. For completeness we show the expected reflected asymmetry of this reaction in Fig. 5.

The estimates we made were in the leading order in QCD. However it is expected from the analysis of Drell-Yan by Altarelli, Ellis and Martinelli[10] and a recent analysis[11] of non leading effects in large p_T processes, that these may play a non negligible role. On the other hand there are statements in the literature[12], that the all orders factorization theorem in reference (3) implies that one can always write a factorized formula of the type

$$\sigma^{AB\ldots} = \int D_A^a\ D_B^b \ldots\ \sigma^{ab\ldots} \qquad (11)$$

where each object can be systematically expanded in powers of α_s beyond the leading order . This statement has however many ambiguities, not least the one exhibited in Eq. (1), where the additional spin correlation terms (involving single spin asymmetries) should enter. It is therefore important to examine polarization asymmetries at large p_T, to check how spin information is trasmitted between the pieces, which one supposes factorize. This will improve our understanding of what precisely is meant by factorization beyond the leading order. Concerning the latter, some interesting suggestions have been made recently (13,14,15), that the most important next to leading order correction to the parton densities as $x \to 1$ can be incorporated in the evolution equations by the simple substitution $\alpha_s(k^2) \to \alpha_s(k^2(1-x))$ (corresponding to $k^2 \to k_T^2$), so the running coupling constant appears also under the integration over the parton momentum x (see Eq. (5) for the structure of the evolution equations). In addition by modifying Q^2 with a function $f(x_a, x_b, ...)$ (i.e. $Q^2 \to Q^2 f(x_a, x_b, ...)$) in Eqs. (7) and (11) one can absorb the effects of the next to leading logarithms in the large p_T process, a point stressed in reference (11). The double asymmetry measurement is an ideal testing ground for such a proposal, because it has to work in both the numerator and denominator in formulae like Eq. (6). Since, if it does not, the asymmetry will clearly be sensitive to the choice of Q^2, because of the claim in reference (11) that the next to leading logarithms can have a large effect on the large p_T cross-section (i.e. the denominator in Eq. (6)) depending on how we choose the variable Q^2. This means we can try to minimize the effects in the denominator. However, if the procedure is arbitrary, it is likely to affect the numerator differently. The same applies to the suggestion in references (13,14,15) for the evolution equations.

For the above reasons we believe the studies of asymmetries will be a very fruitful means of understanding the parton factorization property of QCD, which underwrites much of the present day phenomenology.

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Figure Captions

Fig. 1 a) Reflected asymmetry

b) Transmitted asymmetry

Fig. 2 Mechanisms for $\vec{p}p \to \Lambda + X$

Fig. 3 Transmitted asymmetry $a_{LL}^{if}(z)$ for $q\vec{q} \to s\bar{s}$

Fig. 4 Plots of transmitted asymmetry for $\vec{p}\vec{p} \to \vec{\Lambda} + X$ (a) $\theta_{trig.} = 60°$, (b) $\theta_{trig.} = 30°$ and (c) $\theta_{trig.} = 60°$ with (solid) and without (dashed) gluon contribution

Fig. 5 Plot of reflected asymmetry for $\vec{\bar{p}}\vec{p} \to \Lambda + X$ for $\theta_{trig.} = 60°$

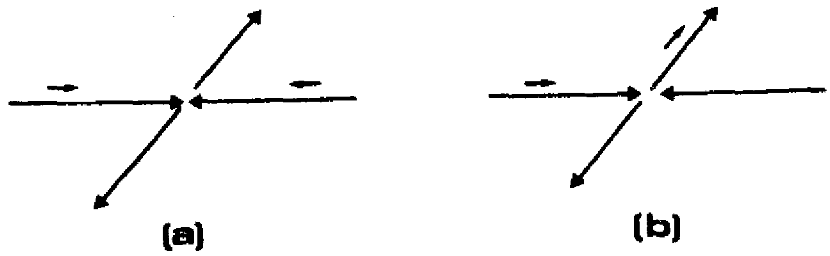


Fig 1

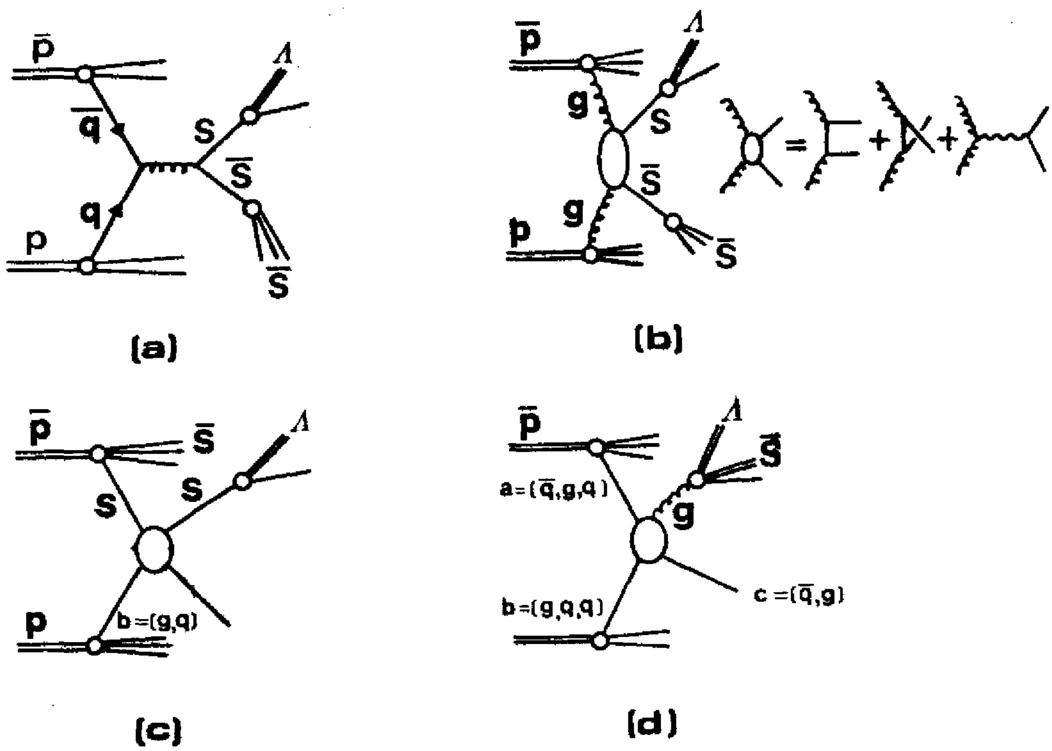
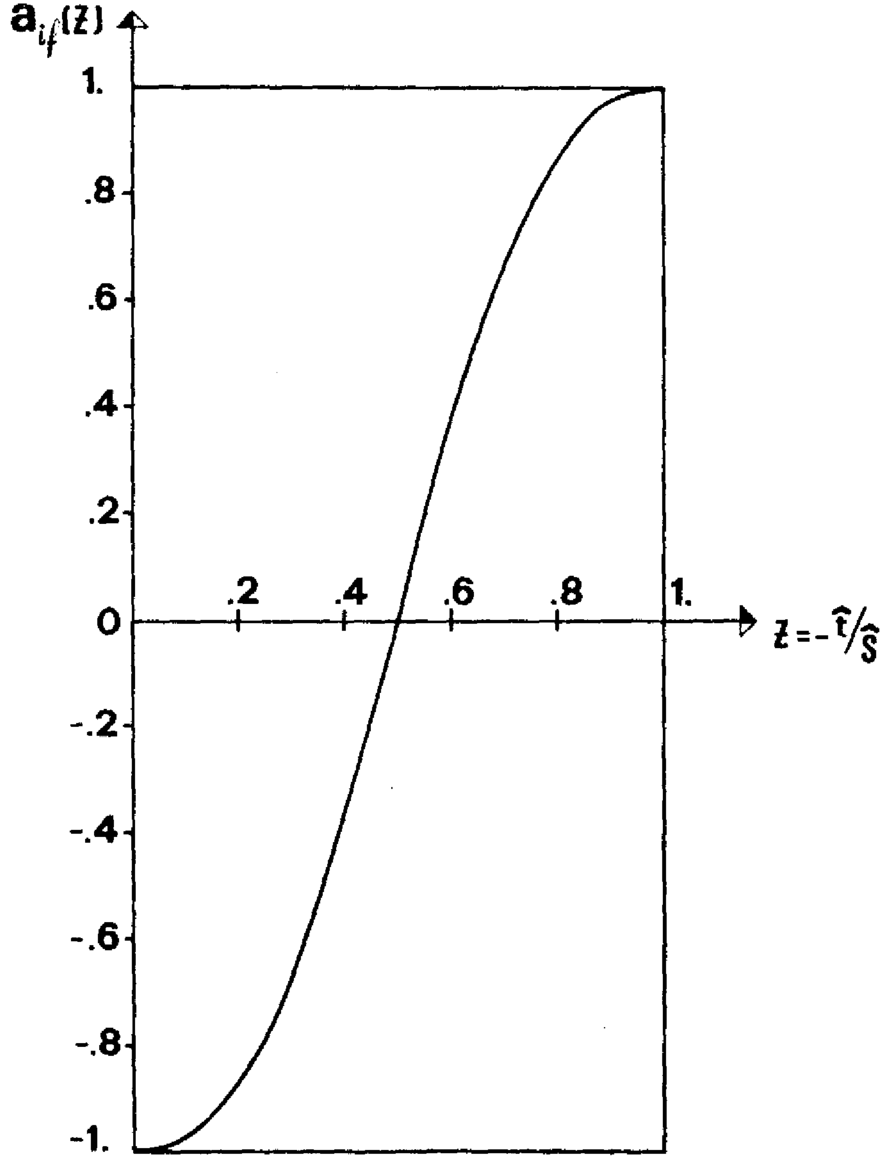
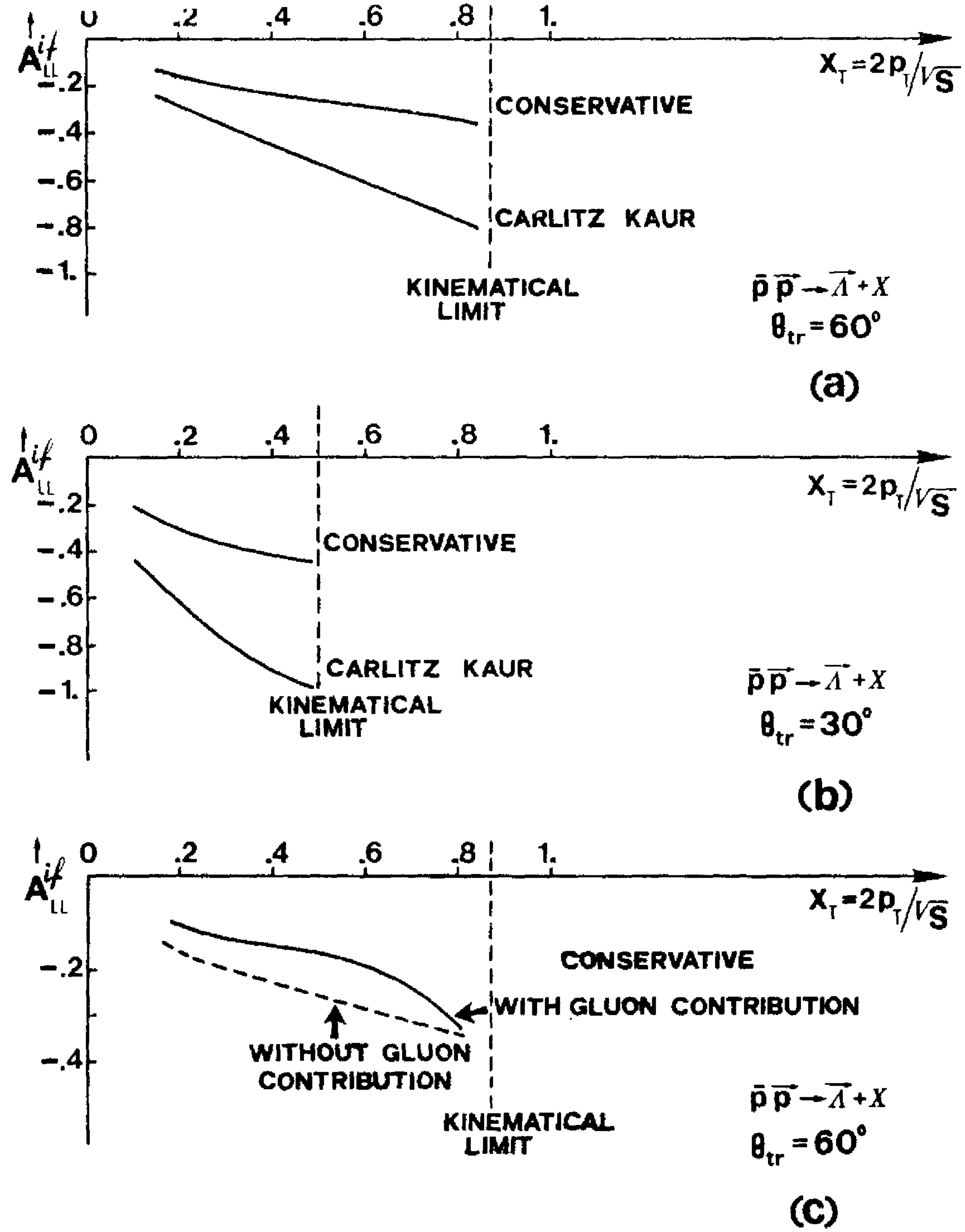


Fig 2



$a_{if}(z)$
1.
.8
.6
.4
.2
0
-.2
-.4
-.6
-.8
-1.
.2
.4
.6
.8
1.
$z = -\hat{t}/\hat{s}$

Fig 3



A_{LL}^{if}
$X_T = 2p_T/\sqrt{S}$
CONSERVATIVE
CARLITZ KAUR
KINEMATICAL LIMIT
$\bar{p}\vec{p} \to \vec{\Lambda} + X$
$\theta_{tr} = 60°$
(a)
A_{LL}^{if}
$X_T = 2p_T/\sqrt{S}$
CONSERVATIVE
CARLITZ KAUR
KINEMATICAL LIMIT
$\bar{p}\vec{p} \to \vec{\Lambda} + X$
$\theta_{tr} = 30°$
(b)
A_{LL}^{if}
$X_T = 2p_T/\sqrt{S}$
CONSERVATIVE
WITH GLUON CONTRIBUTION
WITHOUT GLUON CONTRIBUTION
KINEMATICAL LIMIT
$\bar{p}\vec{p} \to \vec{\Lambda} + X$
$\theta_{tr} = 60°$
(c)

Fig 4

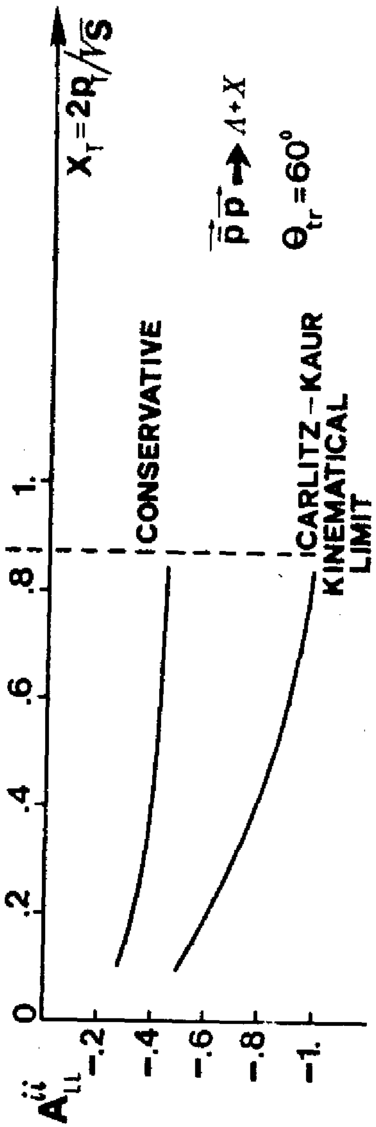


Fig5

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