Transverse-momentum-dependent parton
distributions in a spectator diquark model

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Abstract. Within the framework of a spectator diquark model of the nucleon, involving both scalar and axial-vector diquarks, we calculate all the leading-twist transverse-momentum-dependent parton distribution functions (TMDs). Naive Time-odd densities are generated through a one-gluon-loop rescattering mechanism, simulating the final state interactions required for these functions to exist. Analytic results are obtained for all the TMDs, and a connection with the light-cone wave functions formalism is also established. The model parameters are fixed by reproducing the phenomenological parametrizations of unpolarized and helicity parton distributions at the lowest available scale. Predictions for the other parton densities are given and, whenever possible, compared with available parametrizations.

1. Introduction
Transverse-momentum-dependent parton distributions (TMDs) describe the probability to find, in a hadron, a parton with light-cone longitudinal momentum fraction $x$ and transverse momentum $p_T$ with respect to the direction of the parent hadron momentum, thus encoding low energy non-perturbative information on the QCD partonic structure of the nucleon.

In recent years, great interest has spread around these functions, since intrinsic transverse momentum effects can shed light on various phenomena, such as spin-orbit correlations and single spin asymmetries, and were also shown to be relevant for a correct understanding of multi-jet events at hadron colliders, thus playing a crucial role in LHC physics program \cite{1}. This in turn triggered both theoretical and experimental efforts, aimed at unraveling the TMDs sector. Concerning theory, it was e.g. demonstrated \cite{2} how TMDs which are odd under naive time-reversal transformations (for brevity, T-odd) are not forbidden by QCD time-reversal invariance, and can indeed be non-zero and appear at leading-twist. From the experimental point of view, we still have poor phenomenological information on TMDs, especially concerning their $p_T$ dependence, even though the analysis of azimuthal spin asymmetries, e.g. in semi-inclusive deep inelastic scattering (SIDIS), allowed an extraction of the Sivers function \cite{3}, a T-odd TMD describing how the parton distribution of unpolarized quarks is distorted by the transverse polarization of the parent hadron, and a first attempt to extract the T-odd Boer-Mulders function, i.e. the distribution of transversely polarized partons in an unpolarized hadron, was also presented \cite{4}. All of the above studies assume a flavor-independent Gaussian $|p_T|$-distribution, though there is no compelling physical reason for this choice.

On account of the difficulty to study parton distributions from first principles, it is of great importance to develop simple models able to reproduce the essential qualitative features of
TMDs and to gain insight into their physical content, e.g. concerning partons orbital angular momentum, as well as to perform numerical estimates and simulations on observables of interest, thus guiding further experimental investigations. Among other few TMDs models [5] known in literature, the spectator diquark model [2] here employed owes its appeal in combining relative simplicity and useful versatility, allowing to obtain analytic results for all TMDs. We include both scalar and axial-vector diquarks and we further distinguish between isoscalar (ud-like) and isovector (uu-like) spectators; axial-vector diquarks are essential for a realistic flavor analysis, since the isospin triplet diquark uu flavour combination requires a $S = 1$ spin state. The relative phases necessary to produce T-odd structures are generated by approximating the gauge link operator, ensuring color gauge invariance for the correlators of the theory, with a one gluon-exchange mechanism. We also focus on a possible interpretation of the model results in terms of light-cone wave functions, which in turn gives a more general and physical dress to the results and opens the way to a wider range of applications. The model free parameters are fixed by reproducing the phenomenological parametrizations of unpolarized and longitudinally polarized parton distribution functions (PDFs) at the lowest available scale, and their values employed to obtain numerical results and comparison with available parametrizations for some parton distribution functions of primary interest. As we shall see, an overall good agreement with phenomenology is achieved.

2. Model results for Time-even TMDs

In the framework of parton model, the fundamental theoretical object to study quarks distributions into the proton is the quark-quark correlator, which is the diagonal matrix element, over a bound hadronic state with momentum $P$ and spin $S$, of a bilocal operator built up with light-front quark field operators:

$$
\Phi(x, p_T; S) = \int \frac{d^4 \xi d^4 \xi'}{(2\pi)^8} e^{ip \cdot \xi} \langle P, S | \bar{\psi}(0) \mathcal{U}_{(0, \xi)} \psi(\xi) | P, S \rangle \bigg|_{\xi^+ = 0},
$$

having used light-cone coordinates (so that the quark momentum is given by $p \equiv [p^+, p^-, p_T] = [xP^+, (p^2 + p_T^2)/2xP^+, p_T]$); the gauge link operator, $\mathcal{U}_{(0, \xi)} = P \exp(-ig \int_0^\xi dw \cdot A(w))$ connects the two different space-time points $0, \xi$ by all possible ordered paths followed by the gluon field $A$, coupling to the quark field with strength $g$, and its presence is mandatory in order to ensure color-gauge invariance, while its physical origin lies in the residual interactions between the active quark and the spectator system.

The spectator approximation consists in inserting a completeness relation and truncating, at tree-level, the sum over final states to a single on-shell spectator state with mass $M_X$ and diquark quantum numbers; Eq. (1) can then be explicitly written as

$$
\Phi(x, p_T; S) \sim \frac{1}{(2\pi)^3} \frac{1}{2(1-x)P^+} \tilde{\mathcal{M}}^{(0)}(S) \mathcal{M}^{(0)}(S),
$$

where $m$ is the active quark mass, and the on-shell condition $(P - p)^2 = M_X^2$ for the spectator diquark implies an off-shell condition for the quark ($M$ being the hadron mass):

$$
p^2 = \tau(x, p_T) = -\frac{p_T^2 + L_X^2(m^2)}{1 - x} + m^2,
$$

with $L_X^2(m^2) = xM_X^2 + (1-x)m^2 - x(1-x)M^2$.

In order to match proton spin, the diquark can have either spin 0 (scalar, $X = s$) or spin 1 (axial-vector, $X = a$), as well as isospin 0 (isoscalar $ud$-like system) or isospin 1 (isovector $uu$-like system). The tree-level “scattering amplitude” $\mathcal{M}^{(0)} = \langle P - p | \bar{\psi}(0) | P, S \rangle$, which is indeed a Dirac spinor, is constructed from the Feynman diagram in figure 1 (left), supposing
a nucleon-quark-scalar diquark vertex $\mathcal{Y}_s = ig_s(p^2)1$ and a nucleon-quark-axial vector diquark vertex $\mathcal{Y}_a^\nu = ig_a(p^2)/\sqrt{2}\gamma^\nu\gamma_5$, where $g_X(p^2)$ is a suitable form factor. Moreover, dealing with axial-vector diquarks, there is an intrinsic model arbitrariness in the expression of the sum over polarization states, $d^{\mu\nu} \equiv \sum_{\lambda_\alpha} z^{\nu}_{(+)}(\lambda_\alpha) z^{\mu}_{(-)}(\lambda_\alpha)$, and consequently of the diquark propagator (intervening in T-odd TMDs calculation) as well. In Ref. [6], a complete systematics has been presented: analytic results are derived for all leading-twist TMDs with pointlike, dipolar and exponential form factors, and three different choices of the axial-vector diquark polarization sum, with different physical contents, are probed. Here we will limit ourselves to dipolar and exponential form factors, and three different choices of the axial-vector diquark polarization sum, with different physical contents, are probed. We will limit ourselves to the expressions obtained using a dipolar form factor, $g_X^{dip}(p^2) = (p^2 - m^2)/|p^2 - \Lambda_X^2|^2$ (with $\Lambda_X$ a cut-off parameter), which smoothly suppress the influence of high $|p_T|$, and choosing the so-called light-cone polarization sum (see Ref. [6] for details), physically corresponding to the propagation of only diquark transverse polarization degrees of freedom (the other choices also involve longitudinal as well as scalar polarizations). The theoretical motivation behind this last preference is that it guarantees, unlike the other two, that in DIS processes the axial-vector diquark (which, being charged, can interact with the probing virtual photon) only contributes to the longitudinal structure function $F_L$ (the same situation which naturally occurs for the scalar case): this surely implies a violation of the Callan-Gross relation, but leaves unchanged the leading-order interpretation of the transverse structure function $F_T$ as a charge-weighted sum of quark distributions, thus limiting the phenomenological impact represented by the presence of the diquarks.

![Figure 1](image-url)

**Figure 1.** Left: tree-level spectator model diagram for the calculation of T-even leading-twist parton densities for an active quark (solid line), with a scalar or axial-vector diquark (dashed line). Right: interference between the one-gluon exchange diagram in eikonal approximation and the tree level diagram.

By properly projecting the tree level model correlator of Eq. (2) on various Dirac structures, all Time-even leading twist TMDs can be calculated explicitly. Focusing on the unpolarized distribution, for instance, we can write (with $\gamma^+ = (\gamma^0 + \gamma^3)/\sqrt{2}$)

$$f_1(x, p_T^2) = \frac{1}{4} \text{Tr} \left[ (\Phi(x, p_T; S) + \Phi(x, p_T; -S)) \gamma^+ \right] + \text{h.c.}$$

$$= \frac{1}{(2\pi)^3} \frac{1}{8(1-x)p^+} \text{Tr} \left[ (M(0)(S)M(0)(S) + M(0)(-S)M(0)(-S)) \gamma^+ \right] + \text{h.c.}$$

(4)

thus obtaining, with the choices just discussed, the following results:

$$f_1^{(s)}(x, p_T^2) = \frac{g_s^2}{(2\pi)^3} \frac{[(m + xM)^2 + p_T^2]}{2[p_T^2 + L_s^2(\Lambda_s^2)]^4} (1-x)^3,$$

(5)

$$f_1^{(a)}(x, p_T^2) = \frac{g_a^2}{(2\pi)^3} \frac{[p_T^2(1 + x^2) + (m + xM)^2 (1-x)^2]}{2[p_T^2 + L_s^2(\Lambda_s^2)]^4} (1-x),$$

(6)

for the scalar and axial-vector contributions, respectively.
Besides this diagrammatic approach, T-even TMDs can also be derived by exploiting the so-called overlap representation in terms of light-cone wave functions (LCWFs) [7]; these are the frame independent interpolating functions between hadrons and quark/gluon degrees of freedom, and within our model can be defined as follows:

\[
\psi^{\lambda N}_{\lambda_q}(x, p_T) = \sqrt{\frac{p^+}{(P-p)^+}} \frac{\bar{u}(p, \lambda_q)}{p^2 - m^2} \gamma_\mu U(P, \lambda_N),
\]

(7)

\[
\psi^{\lambda N}_{\lambda_q \lambda_u}(x, p_T) = \sqrt{\frac{p^+}{(P-p)^+}} \frac{\bar{u}(p, \lambda_q)}{p^2 - m^2} \varepsilon^\mu (P-p, \lambda_u) \gamma_\mu U(P, \lambda_N),
\]

(8)

for scalar and axial-vector diquarks respectively, in terms of the quark and nucleon free spinors, \(u\) and \(U\), and of the axial-vector diquark polarization vector \(\varepsilon^\mu\) (see Ref. [7, 8]). The indices \(\lambda_N\), \(\lambda_q\) and \(\lambda_u\) refer to the helicity of the nucleon, the quark and the axial-vector diquark respectively, and are constrained by angular momentum conservation to the spin sum rule \(\lambda_N = \lambda_q (+\lambda_u) + L_z\), where \(L_z\) is the projection of the relative orbital angular momentum between the quark and the diquark. The explicit spectator model expressions for LCWFs can then be used to recover the results in eqs. (5) and (6) in terms of LCWFs overlaps:

\[
f_1^{q(s)}(x, p_T^2) = \frac{1}{16\pi^3} \frac{1}{2} \sum_{\lambda_N=\pm} \sum_{\lambda_q=\pm} |\psi^{\lambda N}_{\lambda_q}|^2,
\]

(9)

\[
f_1^{q(a)}(x, p_T^2) = \frac{1}{16\pi^3} \frac{1}{2} \sum_{\lambda_N=\pm} \sum_{\lambda_q=\pm} \sum_{\lambda_u=\pm} |\psi^{\lambda N}_{\lambda_q \lambda_u}|^2,
\]

(10)

and similar formulae hold for every T-even TMD: the approach is indeed completely general. We refer to [6] for details; for the sake of completeness and in view of its importance, we report here also the model results for the chiral-odd TMD transversity distribution, i.e. the transverse spin distribution in a tranversely polarized hadron:

\[
h_1^{q(s)}(x, p_T^2) = g_s^2 \frac{(m + xM)^2 (1-x)^3}{2 (2\pi)^5 [p_T^2 + L_z^2(\Lambda_s^2)]^4},
\]

(11)

\[
h_1^{q(a)}(x, p_T^2) = -g_s^2 \frac{p_T^2 x(1-x)}{(2\pi)^5 [p_T^2 + L_z^2(\Lambda_s^2)]^4}.
\]

(12)

By integration over \(p_T\), finally, analytic model expressions can be derived also for the usual integrated PDFs, \(f_1(x)\), \(g_1(x)\) and the transversity \(h_1(x)\).

3. Model results for Time-odd TMDs

It is a well known fact that, by means of tree level amplitudes only, no Time-odd effect can be simulated, since there is no residual interaction between the active quark and the spectators, i.e. we lack the interference between two competing channels, coupled to the same final state, required to produce the complex amplitude whose imaginary part would give the T-odd contribution. We could anyway go further and consider the single-gluon-exchange scattering amplitude (in eikonal approximation) \(\mathcal{M}^{(1)}\), as the leading-twist \(\mathcal{O}(\alpha_s)\) approximation of the gauge link operator [9]; its interference with the tree-level term \(\mathcal{M}^{(0)}\) thus allows to compute the Sivers and Boer-Mulders T-odd TMDs in the model. For the Sivers case, e.g., we have:

\[
\frac{\varepsilon^{ij}_{T \mu T} S_{T ij}}{M} f_1^{T} (x, p_T^2) = -\frac{1}{(2\pi)^3} \frac{1}{8(1-x)p^+} \text{Tr} \left[ (\mathcal{M}^{(1)}(S) \mathcal{\bar{N}}^{(0)}(S) - \mathcal{M}^{(1)}(-S) \mathcal{\bar{N}}^{(0)}(-S) ) \gamma^+ \right] + \text{h.c.},
\]

(13)
where the transverse tensor $\varepsilon^{ij} = \varepsilon^{\mu
u j} n_{+\mu} n_{-\nu}$ only non-zero components are $\varepsilon^{12} = -\varepsilon^{21} = 1$.

After introducing proper Feynman rules for the loop integral involved in $M^{(1)}$, its imaginary part can be computed applying Cutkosky rules for the corresponding box diagram (see figure 1, right), by cutting, i.e. putting on shell, the eikonal and the axial-vector diquark propagators; the final model results read (considering, as before, dipolar form factor and light-cone polarization sum):

$$f_{1T}^{q(s)}(x, p_T^2) = -\frac{g_2^2}{4} \frac{M e_{2A}}{(2\pi)^3} \frac{(1-x)^3}{L_2^2(L_2^2 + L_9^2)^3} (m + xM),$$
$$f_{1T}^{q(a)}(x, p_T^2) = -\frac{g_2^2}{4} \frac{M e_{2A}}{(2\pi)^3} \frac{(1-x)^2}{L_2^2(L_2^2 + L_9^2)^3} (m + xM),$$

for the Sivers function (with $e_c^2 = 4\pi C_F a_s$, while $1_3 C_F = \sum a_t^a t^a = 1_3 (4/3)$, $t^a$ being the SU(3) Gell-Mann color matrices), while for the Boer-Mulders function we find:

$$h_{2T}^{q(s)}(x, p_T^2) = f_{1T}^{q(s)}(x, p_T^2),$$
$$h_{2T}^{q(a)}(x, p_T^2) = f_{1T}^{q(a)}(x, p_T^2)/x.$$ 

As already seen for T-even TMDs, T-odd densities as well do admit an overlap representation in terms of tree level LCWFs, but while for T-even functions this simply appeared as a sum of products of LCWFs with the same arguments, we are now forced to consider convolution integrals of LCWFs with different arguments, including a suitable operator in order to describe the needed final-state-interactions (FSI) produced by the gluon rescattering. By comparing the overlap expressions with those obtained within the Feynman diagram approach, we found that this FSI kernel has a universal form (i.e., the same for scalar and axial-vector diquarks, and for Sivers and Boer-Mulders functions), reminiscent of the virtual gluon propagator:

$$G(x, x', p_T, p_T') = -\frac{i}{2\pi (p_T - p_T')^2} \frac{C_F a_s}{\delta(x - x')}.$$ 

### 4. Numerical results

In order to give predictive power to the model, we must fix its free parameters, namely the quark mass $m$, the nucleon-quark-diquark coupling $g_X$, the diquark mass $M_X$, and the cutoff $\Lambda_X$ in the nucleon-quark-diquark dipolar form factor, for $X = s, a$ scalar and axial-vector diquarks. A good fit to phenomenological parametrizations also requires to distinguish between the two isospin states of the axial-vector diquark, i.e. the axial-vector isoscalar state with $I_3 = 0$ (corresponding to the $ud$ system, label $a$) and axial-vector isovector diquark with $I_3 = 1$ (uu system, label $a'$).

The model scale is another free parameter, but we choose to identify it with the lowest possible $Q^2$ at which parametrizations are available: hence, we fitted model results to the $f_1^u$ and $f_1^d$ ZEUS parametrizations $[10]$ (ZEUS2002) at $Q_0^2 = 0.3 \text{ GeV}^2$ and to the $g_1^u$ and $g_1^d$ leading-order (LO) data from Ref. $[11]$ (GRSV2000) at $Q^2 = 0.26 \text{ GeV}^2$. The global normalizations are fixed by imposing that $\int dx \int d^2p_T f_1^{q(a)}(x, p_T^2) = 1$. The relation between quark flavors, i.e. the functions entering the fit, and diquark types, i.e. our model results, is written as

$$f_1^u = c_s^2 f_1^{u(s)} + c_a^2 f_1^{u(a)}, \quad f_1^d = c_a'^2 f_1^{d(a')}.$$ 

where the weight factors $c_s$, $c_a$ and $c_a'$ are also considered free parameters, and their values fixed by the global fit. The reason for this approach is twofold: first, this model involves non-zero (and relativistically enhanced) LCWFs with $L_z = 1$, hence the quark-diquark system can have a non-vanishing relative orbital angular momentum, and thus the proton wave function no longer displays a simple $SU(4)$ symmetry; secondly, a fundamental limitation of the diquark model
consists in the intrinsic impossibility to simultaneously fulfill the quark number and momentum sum rules: as a matter of fact, on the one hand the total number of quarks “seen” in this model is only one, since the other two are always confined into the diquark state, and on the other hand the charged diquarks can be probed in DIS, thus contributing to the overall momentum. On the basis of such considerations, we decided to avoid imposing the momentum sum rule and to let the fit choose the values of the parameters $c_X$ (see figure 4 and table 1 of [6] for the results of the fit and the corresponding values of the model parameters).

We now present a brief selection of plots, in order to highlight some crucial features of our results. In figure 2 we show the $p_T^2$ dependence of the unpolarized TMD: we observe that $f_1^u$ displays a non-monotonic behavior for $x \leq 0.02$, which is due to the presence of LCWFs with non-zero orbital angular momentum; indeed, the contribution of LCWFs with $L_z = 1$, dominant at high $x$, falls linearly with $p_T^2$ as $p_T^2 \to 0$, unlike the one from $L_z = 0$ LCWFs. This simple example shows how the study of the transverse-momentum-dependence of TMDs can already expose some effects due to orbital angular momentum. We also note that in our model the average quark transverse momentum decreases as $x$ increases, and that $d$ quarks on average carry less transverse momentum than $u$ ones. Though dealing with a relatively simple model, we thus already acknowledge how the widely accepted assumption of a factorized, flavor-independent quark transverse-momentum distribution is inadequate. A significant violation of the $x \cdot p_T$ factorization has also been recently obtained in a chiral quark soliton model [17].

![Figure 2](image-url)  
**Figure 2.** The $p_T^2$ dependence of $f_1(x, p_T^2)$ for $u$ (left) and $d$ quark (right), for different values of $x$.

In figure 3 we show the results for up and down quarks transversity distribution $h_1(x) \equiv \int d^2p_T h_1(x, p_T^2)$. Model functions are represented by the dashed lines and assumed to refer to the model scale $Q_0^2 = 0.3$ GeV$^2$, while the solid lines are the curves obtained after applying the DGLAP evolution at LO up to the scale $Q^2 = 2.5$ GeV$^2$ [12]; the latter scale pertains to the parametrization of Ref. [13], whose uncertainty band is depicted as a shaded area. We find reasonable agreement with the phenomenological data, with the maxima in the correct positions and a somewhat too small result for the up quark at small $x$. An interesting feature, typical of this model, is the change of sign at $x \sim 0.5$ in $h_1^u$ (it cannot be reproduced by parametrizations, whose functional form does not allow for a sign change): indeed, the axial-vector diquark contribution is negative (see Eq. (12)) and becomes dominant over the scalar one at sufficiently high $x$. In figure 4 the $|p_T|$-shape (at $x = 0.1$) of $h_1(x, p_T^2)$ is shown, together with the parametrization in Ref. [13]: we have again a satisfying agreement between model results and phenomenology, at least below $|p_T| \sim 0.3$ GeV/c (moreover, we stress that in this case no TMD-evolution was applied to model functions). Also $h_1^u(x = 0.1, p_T^2)$ changes sign, at $|p_T| \sim 0.5$ GeV/c, due to the dominance of the negative $S=1$ diquark contribution at high $|p_T|$. 


Turning to T-odd functions, they are mainly studied in connection with single spin asymmetries (SSA). In figure 5 we report the model results for the first $p_T$-moment of the Sivers function, defined as $f_{T}^{1}(x) \equiv \int d^2 p_T \frac{p_T^2}{2M^2} f_{T}^{1}(x, p_T^2)$, compared with two different available parametrizations [14, 15] extracted from SIDIS data on transverse SSA. The spectator model predicts the correct signs for the $u$ and $d$ quark functions (also in agreement with lattice simulations [16]), and the maxima are reached at about the same $x$ ($\sim 0.3$) as the parametrizations, but the theoretical curve underestimates the experimental one for down quark, while the agreement is more satisfactory for the up quark (anyway, also in this case no scale-evolution was applied). Finally, figure 6 shows the model results for the spin density of unpolarized quarks with flavor $q$ in transversely polarized protons, defined as

$$f_{q/p}^{1}(x, p_T) = f_{q}^{1}(x, p_T^2) - f_{T}^{1}(x, p_T^2) \frac{(\hat{P} \times \hat{p}_T) \cdot S}{M} ,$$

revealing the expected asymmetry in transverse-momentum space (along $\pm \hat{p}_y$, for a proton transverse polarization along $+\hat{x}$, and choosing $\hat{P} = -\hat{z}$), connected with the Sivers function. The distortion, towards positive (negative) values of $p_y$ for $u$ ($d$) quarks, is in agreement with similar lattice results [16] and with the signs of the anomalous magnetic moments $\kappa_q$.

Although no reliable experimental parametrization is yet available for the Boer-Mulders function, also in this case we checked the agreement between our model functions signs (negative for both $u$ and $d$) and the lattice calculations of Ref. [16].
Figure 5. The first $p_T$-moment $x f^{(1)}_{LT}(x)$ of the Sivers function for $u$ (left) and $d$ (right) quark. Solid line for the spectator diquark model results. Darker shaded area for the uncertainty band of parametrizations from Ref. [14], lighter one from Ref. [15].

Figure 6. Model results for the spin density of unpolarized quarks in transversely polarized protons in $p_T$ space at $x = 0.1$, for $u$ (left) and $d$ (right) quarks. Proton transverse polarization along $+\hat{x}$.

References
[1] Hautmann F, Jung H 2008 JHEP 0810:113
[2] Brodsky S, Hwang D S and Schmidt I 2002 Phys. Lett. B 530 99
[3] Anselmino M et al 2005 Preprint hep-ph/0511017
[4] Zhang B, Lu Z, Ma B-Q and Schmidt I 2008 Phys. Rev. D 77 054011
[5] Pasquini B, Cazzaniga S and Boffi S 2008 Phys. Rev. D 78 034025
[6] Bacchetta A, Conti F, Radici M 2008 Phys. Rev. D 78 074010
[7] Brodsky S J, Hwang D S, Ma B-Q and Schmidt I 2001 Nucl. Phys. B 593 311
[8] Lepage G P, Brodsky S J, 1980 Phys. Rev. D 22 2157
[9] Boer D, Mulders P J and Pijlman F 2003 Nucl. Phys. B 667 201
[10] ZEUS, Chekanov S et al 2003 Phys. Rev. D 67 012007
[11] Gluck M, Reya E, Stratmann M and Vogelsang W 2001 Phys. Rev. D 63 094005
[12] Hirai M, Kumano S and Miyama M 1998 Comput. Phys. Commun. 111 150
[13] Anselmino M et al 2007 Phys. Rev. D 75 054032
[14] Anselmino M et al 2009 Eur. Phys. J. A 39 89
[15] Collins J C et al 2005 Preprint hep-ph/0510342
[16] QCDSF, Gockeler M et al 2007 Phys. Rev. Lett. 98 222001
[17] Wakamatsu M 2009 Preprint arXiv:0903.1886[hep-ph]