Three-dimensional imaging of the nucleon in momentum space

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Abstract. Transverse-momentum dependent parton distributions (TMDs) are studied in the
framework of quark models. Results for the six T-even TMDs are obtained from the overlap
of three-quark light-cone wave functions, using both the chiral quark-soliton model and a light-
cone constituent quark model. Furthermore, quark model relations among TMDs are reviewed
and their physical origin is discussed in terms of rotational-symmetry properties of the nucleon
state in its rest frame.

1. Introduction

Transverse-momentum dependent parton distributions (TMDs) have received a great attention
in the last years as they represent key objects to map out the three-dimensional partonic
structure of hadrons in momentum space. The dependence on the transverse momentum of
the quark allows for non-trivial correlations between the orbital angular momentum and the spin
of the quark inside nucleons with different polarization states. TMDs typically give rise to spin and
azimuthal asymmetries in, for instance, semi-inclusive deep inelastic scattering and Drell-Yan
processes, and significant efforts have already been devoted to measure these observables (see
e.g. Ref. [1] for a recent review). However, the extraction of TMDs from experimental data is a
quite difficult task and needs educated Ansätze for fits of TMD parametrizations. To this aim,
model calculations of TMDs play a crucial role and are essential towards an understanding of
the non-perturbative aspects of TMDs.

Studies of the TMDs have been mainly focused on the quark contribution, and predictions
have been obtained within a variety of models [2, 3, 4, 5, 6, 7, 8, 9, 10, 11, 12, 13, 14, 15, 16,
17, 18, 19, 20, 21, 22, 23, 24, 25, 26, 27, 28, 29, 30, 31, 32]. Despite the specific assumptions
for modeling the quark dynamics, most of quark models predict relations among the leading-
twist TMDs. Since in QCD these TMDs are all independent, it is clear that such relations
should be traced back to some common simplifying assumptions in the models. First of all, it
was noticed that they break down in models with gauge-field degrees of freedom. Furthermore,
most quark models are valid at some very low scale and these relations are expected to break
under QCD evolution to higher scales. Despite these limitations, such relations are intriguing
because they can provide guidelines for building parametrizations of TMDs to be tested with
experimental data and can also give useful insights for the understanding of the origin of the different spin-orbit correlations of quarks in the nucleon.

In section 2 we quickly review the formalism for the definition of the leading-twist TMDs, and introduce a convenient representation of the quark-quark correlator in terms of the net-polarization states of the quark and the nucleon. In section 3 we present predictions for the six T-even TMDs using the overlap of three-quark light-cone wave functions (3Q LCWFs) which are modeled within the light-cone version of the chiral quark-soliton model (LCQSM) \(^{21,24,33,34,35,36}\) as well as a light-cone constituent quark model (LCCQM) \(^{22}\). Then we present in section 4 the relations among TMDs observed in most of quark models. In section 5 we identify the assumptions common to all these models and derive the relations by working directly at the amplitude level. An alternative derivation based on the language of wave functions can be found in Ref. \(^{37}\), where we also discuss an additional relation due to SU(6) spin-flavor symmetry. More details about the calculations of the amplitudes can be found in Ref. \(^{24}\).

2. Transverse-Momentum Dependent Parton Distributions

2.1. Definitions

In this section, we review quickly the formalism for the definition of the leading-twist TMDs following the conventions of Refs. \(^{38,39,40}\). Introducing two lightlike four-vectors \(n_\pm\) satisfying \(n_+ \cdot n_- = 1\), we write the light-like components of a generic four-vector \(a\) as \([a^+, a^-, a^\perp]\) with \(a^\pm = a \cdot n_\pm\). The density of quarks can be defined from the following quark-quark correlator

\[
\Phi_{ab}(x, \vec{k}_\perp, \vec{S}) = \int \frac{dz^- d^2 z_\perp}{(2\pi)^3} e^{i(k^+ z^- - \vec{k}_\perp \cdot \vec{z}_\perp)} \langle P, \vec{S}|\psi_b(0)U_{(0, +\infty)}^a(0)U_{(+\infty, z)}^b\psi_a(z)|P, \vec{S}\rangle|_{z^+=0},
\]

where \(k^+ = xP^+\), \(\psi\) is the quark field operator with \(a, b\) indices in the Dirac space, and \(U\) is the Wilson line which ensures color gauge invariance \(^{41}\). The target state is characterized by its four-momentum \(P\) and the direction of its polarization \(\vec{S}\). We choose a frame where the hadron momentum has no transverse components \(P = [P^+, M^2, 0_\perp]\).

TMDs enter the general Lorentz-covariant decomposition of the correlator \(\Phi_{ab}(x, \vec{k}_\perp, \vec{S})\) which, at twist-two level and for a spin-1/2 target, reads

\[
\Phi(x, \vec{k}_\perp, \vec{S}) = \frac{1}{2} \left\{ f_1 \gamma^\perp + \frac{\epsilon^{ij} k_i^* S_j}{M} f_1 T^i \eta^+ + S_z g_{1L} \gamma_5 \eta^+ + \frac{\vec{S}}{M} g_{1T} \gamma_5 \eta^+ \\
+ h_1 \left[ \frac{S^a \cdot d^a}{2} \gamma_5 + S_z h_{1L} \frac{[\gamma_a, d_a]}{2M} \gamma_5 + \frac{\vec{S}}{M} h_{1T} \frac{[\gamma_a, d_a]}{2M} \gamma_5 + i h_{1L} \frac{[\gamma_a, d_a]}{2M} \right] \right\},
\]

where \(\epsilon^{ij} = -\epsilon_{ij} = 1\), and the transverse four-vectors are defined as \(a_\perp = [0, 0, a^\perp]\). The nomenclature of the distribution functions follows closely that of Ref. \(^{38}\), sometimes referred to as “Amsterdam notation”. Among these eight distributions, the so-called Boer-Mulders function \(h_+^T\) \(^{39}\) and Sivers function \(f_1^T\) \(^{12}\) are T-odd, i.e. they change sign under “naive time-reversal”, which is defined as usual time-reversal but without interchange of initial and final states. All the TMDs depend on \(x\) and \(\vec{k}_\perp^2\). These functions can be individually isolated by performing traces of the correlator with suitable Dirac matrices. Using the abbreviation \(\Phi^{[\Gamma]} \equiv \text{Tr} [\Phi \Gamma] / 2\), we have

\[
\Phi^{[\gamma^\perp]}(x, \vec{k}_\perp, \vec{S}) = f_1 - \frac{\epsilon^{ij} k_i^* S_j}{M} f_1 T^i, \\
\Phi^{[\gamma^+ \gamma_5]}(x, \vec{k}_\perp, \vec{S}) = S_z g_{1L} + \frac{\vec{S}}{M} g_{1T}, \\
\Phi^{[\sigma^{ij} \gamma_5]}(x, \vec{k}_\perp, \vec{S}) = S_j^* h_1 + S_z \frac{k_i^*}{M} h_{1L} + S^i \frac{2k_i^* k_{1T} \delta^i_j}{2M^2} h_{1T} + \frac{\epsilon^{ij} k_i^*}{M} h_{1L},
\]

(3a) (3b) (3c)
where \( j = 1, 2 \) is a transverse index, and \( h_1 = h_{1T} + \frac{k_1^2}{2M^2} h_{1T} \).

The correlation function \( \Phi^{(\gamma^\pm)}(x, \vec{k}_\perp, \vec{S}) \) is just the unpolarized quark distribution, which integrated over \( \vec{k}_\perp \) gives the familiar light-cone momentum distribution \( f_1(x) \). All the other TMDs characterize the strength of different spin-spin and spin-orbit correlations. The precise form of this correlation is given by the prefactors of the TMDs in Eqs. \([35, 36]\). In particular, the TMDs \( g_{1L} \) and \( h_1 \) describe the strength of a correlation between a longitudinal/transverse target polarization and a longitudinal/transverse parton polarization. After integration over \( \vec{k}_\perp \), they reduce to the helicity and transversity distributions, respectively. By definition, the spin-orbit correlations described by \( f_{1T}^\perp, g_{1T}, h_{1T}^+, h_{1T}^- \), and \( h_1 \) involve the transverse parton momentum and the polarization of both the parton and the target, and vanish upon integration over \( \vec{k}_\perp \).

In the following we will focus the discussion on the quark contribution to the TMDs, ignoring the gauge-field degrees of freedom and therefore reducing the gauge links in Eq. (1) to the identity. It follows that the T-odd TMDs \( f_{1T}^\perp \) and \( h_1 \) are identically zero in this approach. Our discussions will then cover the six T-even TMDs only.

### 2.2. Net-Polarization Basis

The physical meaning of the correlations encoded in the TMDs becomes especially transparent when expressed in the basis of net polarization for the quark and the nucleon, see Refs. \([24, 37]\). In this basis, the correlator (1) for a given quark flavor \( q \) can be written in a matrix form

\[
\Phi_q^{\mu\nu} = \begin{pmatrix}
  f_1^q & k_M h_{1T}^q & 0 & 0 \\
  k_M^T f_{1T}^q & h_{1T}^q & 0 & 0 \\
  0 & 0 & h_{1T}^+ & k_M g_{1T}^q \\
  0 & 0 & k_M^T h_{1T}^q & g_{1T}^q
\end{pmatrix},
\]

where we introduced the notations \( h_{1T}^\pm = h_1^\pm + \frac{k_1^2}{2M^2} h_{1T} \) and chose for convenience the axes in the transverse plane such that \( \vec{k}_\perp = k_\perp \vec{e}_\perp \). The four-component index \( \mu \) refers to the net polarization of the nucleon: \( \mu = 0 \) stands for an unpolarized \( \frac{1}{2}(\uparrow + \downarrow) \) nucleon, while \( \mu = i = 1, 2, 3 \) stands for a nucleon with net polarization \( \frac{1}{2}(\uparrow - \downarrow) \) in the \( i \)th direction. Likewise, \( \nu = 0 \) stands for an unpolarized quark (\( \Gamma = \gamma^\pm \) ), \( \nu = j = 1, 2 \) stands for a quark with net polarization in the transverse \( j \)th direction (\( \Gamma = i\sigma_j \gamma_5 \)), and \( \nu = 3 \) stands for a quark with net polarization in the \( z \)-direction (\( \Gamma = \gamma^\gamma_5 \)).

## 3. Predictions from the LC\(\chi\)QSM and LCCQM

Using the formalism developed in Ref. \([21]\) for the three-quark Fock sector, we have calculated the six T-even TMDs in the LC\(\chi\)QSM and the LCCQM. In figure we show the transverse moments of these TMDs which represent directly the strength of the corresponding multipole in \( \vec{k}_\perp \)-space as function of \( x \). For a generic TMD \( j(x, \vec{k}_\perp^2) \), these transverse moments are defined as

\[
j^{(0)}(x) = \int d^2k_\perp j(x, \vec{k}_\perp^0), \quad j^{(1/2)}(x) = \int d^2k_\perp \frac{k_\perp^2}{M^2} j(x, \vec{k}_\perp^2), \quad j^{(1)}(x) = \int d^2k_\perp \frac{k_\perp^2}{2M^2} j(x, \vec{k}_\perp^2).
\]

Note that our definition of the transverse \((1/2)\)-moment is twice larger than in \([26]\).

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1 Note this is not a Lorentz index but Einstein’s summation convention still applies.
The shape of the curves within the LC\(\chi\)QSM and the LCCQM are very similar. The size of the helicity and transversity distributions is somewhat smaller in the LCCQM. On the other hand, the peak of the distributions for the other polarized TMDs is larger in the LCCQM than in the LC\(\chi\)QSM, especially for \(h_1^{\perp(1)}\). This pattern for the magnitude of TMDs in the two models can be traced back to the fact that more orbital angular momentum is involved in the LCCQM than in the LC\(\chi\)QSM. Note also that, in the LC\(\chi\)QSM, the structure functions do not vanish at \(x = 0\) while they do vanish in the LCCQM. In the latter, a simple power-law Ansatz for the functional form of the 3Q LCWF is assumed and implies vanishing LCWFs when at least one fraction of quark longitudinal momentum \(x_i\) tends to zero. This is at variance with the wave function of the LC\(\chi\)QSM which comes from the solution of the Dirac equation describing the motion of the quarks in the solitonic pion field. We should also mention that the \(\chi\)QSM was recently applied in Ref. [30] to calculate the unpolarized TMDs, taking into account the whole contribution from quarks and antiquarks without expansion in the different Fock components, and obtaining results for the isoscalar combination \(u + d\) which corresponds to the leading contribution in the \(1/N_c\) expansion.

4. Model Relations
In QCD, the eight TMDs are a priori all independent. The T-even TMDs obtained in the LC\(\chi\)QSM and LCCQM are however not all independent. They are found to satisfy three flavor-
independent relations. Two of them are linear, while one is quadratic in the TMDs

\[ g_{1L}^q - \left[ h_{1L}^q + \frac{\vec{k}_{\perp}^2}{2M_T} h_{1T}^q \right] = 0, \]  
\[ g_{1T}^q + h_{1L}^q = 0, \]  
\[ (g_{1T}^q)^2 + 2h_{1L}^q h_{1T}^q = 0. \]  

A further flavor-dependent relation involves both polarized and unpolarized TMDs

\[ \mathcal{D}^q f_1^q + g_{1L}^q = 2h_{1L}^q, \]  

where, for a proton target, the flavor factors with \( q = u, d \) are given by \( \mathcal{D}^u = \frac{2}{3} \) and \( \mathcal{D}^d = -\frac{1}{3} \).

For example, from the relation (6), we see that the difference between helicity and transversity distributions \( g_{1L}^q - h_{1T}^q \) is directly connected to the pretzelosity distribution \( h_{1T}^q \). The difference being smaller in the L\( ^\chi \)QSM, it follows immediately that the pretzelosity distribution is also smaller compared to the LCCQM.

These relations are not specific to the L\( ^\chi \)QSM and LCCQM as part or all of them have also been observed in a large majority of other low-energy quark models\(^2\) like the bag model \(^{26,27}\), some quark-diquark models \(^\text{[2,12,13]}\), and the covariant parton model \(^\text{[11]}\). Note however that there also exist models where the relations are not satisfied, like in some versions of the spectator model \(^\text{[25]}\) and the quark-target model \(^\text{[29]}\). The relations should then be a consequence of some underlying assumptions not shared by the models of Refs. \(^\text{[25]}\) and \(^\text{[29]}\). Here we focus only on the flavor-independent relations \(^\text{[6]-[8]}\), since it is already clear that the flavor dependence in the relation \(^\text{[9]}\) requires specific assumptions for the spin-isospin structure of the nucleon state, like \( SU(6) \) spin-flavor symmetry. A discussion of the flavor-dependent relation can be found in Ref. \(^\text{[37]}\).

The interest in these relations is purely phenomenological. In order to interpret the experimental data sensitive to TMDs, one needs inputs from educated models and parametrizations for the extraction of these distributions. It is therefore particularly interesting to see to what extent the relations \(^\text{[6]-[8]}\) can be useful as approximate relations, which provide simplified and intuitive notions for the interpretation of the data. Note that some preliminary calculations in lattice QCD give indications that the relation \(^\text{[7]}\) may indeed be approximately satisfied \(^\text{[43,44]}\).

5. Derivation of the Flavor-Independent Relations

We show in this section that the flavor-independent relations \(^\text{[6]-[8]}\) can easily be derived, once the following assumptions are made:

(i) the probed quark behaves as if it does not interact directly with the other partons (i.e. one works within the standard impulse approximation) and there are no explicit gluons;

(ii) the quark light-cone and canonical polarizations are related by a rotation with axis orthogonal to both \( \vec{k}_{\perp} \) and the light-cone direction;

(iii) the target has spherical symmetry in the canonical-spin basis.

From these assumptions, one realizes that the flavor-independent relations have essentially a geometrical origin, as was already guessed in the context of the bag model almost a decade ago \(^\text{[45]}\). We note however that the spherical symmetry is a sufficient but not necessary condition for the validity of the individual flavor-independent relations. The assumptions (i)-(iii) are satisfied by all models where the relations have been observed and can be considered as necessary

\(^2\) Other expressions can be found in the literature, but are just combinations of the relations \(^\text{[6]-[8]}\).
conditions [37]. Consistently, the few known models where the relations are absent [25, 29] fail with at least one of the above three conditions.

As we have seen, the TMDs can be expressed in simple terms using light-cone polarization. On the other hand, rotational symmetry is easier to handle in terms of canonical polarization, which is the natural one in the instant form. We therefore write the TMDs in the canonical-spin basis, and then impose spherical symmetry. But before that, we need to know how to connect light-cone helicity to canonical spin.

5.1. Connection between Light-Cone Helicity and Canonical Spin
Relating in general light-cone helicity with canonical spin is quite complicated, as the dynamics is involved. However, the common approach in quark models is to assume that the mutual interactions among quarks can be neglected during the interaction process with an external probe. In this case the connection simply reduces to a rotation in polarization space with axis orthogonal to both \( \vec{k}_\perp \) and \( \vec{e}_z \). The quark creation operator with canonical spin \( \sigma \) can then be written in terms of quark creation operators with light-cone helicity \( \lambda \) as follows

\[
q^\dagger_\sigma = \sum_\lambda D^{(1/2)\ast}_{\sigma \lambda} q^\dagger_\lambda \quad \text{with} \quad D^{(1/2)\ast}_{\sigma \lambda} = \left( \begin{array}{cc}
\cos \frac{\theta}{2} & -\hat{k}_R \sin \frac{\theta}{2} \\
\hat{k}_L \sin \frac{\theta}{2} & \cos \frac{\theta}{2}
\end{array} \right),
\]

(10)

where \( \hat{k}_{RL} = (k_x \pm ik_y)/k_\perp \). Note that the rotation does not depend on the quark flavor. The angle \( \theta \) between light-cone and canonical polarizations is usually a complicated function of the quark momentum \( k_\perp \) and is specific to each model. It contains part of the model dynamics. The only general property is that \( \theta \to 0 \) as \( k_\perp \to 0 \). Due to our choice of a reference frame where the target has no transverse momentum, the light-cone helicity and canonical spin of the target can be identified, at variance with the quark polarizations.

5.2. TMDs in Canonical-Spin Basis
The four-component notation introduced in section [22] is very convenient for discussing the rotation between canonical spin and light-cone helicity at the amplitude level. One can easily see that the canonical tensor correlator \( \Phi^\mu_\nu_{Cq} \) is related to the light-cone one in Eq. (4) as follows

\[
\Phi^\mu_\nu_{Cq} = \Phi^\mu_\nu q \circ O_{\rho} \circ \nu,
\]

(11)

with the orthogonal matrix \( O_{\rho} \), representing the rotation at the amplitude level, given by (remember that we chose \( \hat{k}_\perp = k_\perp \vec{e}_y \))

\[
O_{\rho} = \begin{pmatrix}
1 & 0 & 0 & 0 \\
0 & 1 & 0 & 0 \\
0 & 0 & \cos \theta & -\sin \theta \\
0 & 0 & \sin \theta & \cos \theta
\end{pmatrix}.
\]

(12)

The canonical tensor correlator then takes the form

\[
\Phi^\mu_\nu_{Cq} = \begin{pmatrix}
\frac{\vec{f}_q^T}{M} & \frac{k_M h_{1L}^q}{M} & 0 & 0 \\
\frac{k_M h_{1L}^q}{M} & 0 & 0 & 0 \\
0 & 0 & \frac{h_{1T}^q}{M} & \frac{g_{1L}^q}{M} \\
0 & 0 & \frac{g_{1T}^q}{M} & \frac{h_{1L}^q}{M}
\end{pmatrix},
\]

(13)

where we introduced the notations

\[
\begin{pmatrix}
\frac{\vec{g}_{1L}^q}{M} \\
\frac{h_{1L}^q}{M}
\end{pmatrix} = \begin{pmatrix}
\cos \theta & -\sin \theta \\
\sin \theta & \cos \theta
\end{pmatrix} \begin{pmatrix}
\frac{g_{1L}^q}{M} \\
\frac{h_{1L}^q}{M}
\end{pmatrix},
\]

(14)

\[
\begin{pmatrix}
\frac{\vec{g}_{1T}^q}{M} \\
\frac{h_{1T}^q}{M}
\end{pmatrix} = \begin{pmatrix}
\cos \theta & -\sin \theta \\
\sin \theta & \cos \theta
\end{pmatrix} \begin{pmatrix}
\frac{g_{1T}^q}{M} \\
\frac{h_{1T}^q}{M}
\end{pmatrix}.
\]

(15)
Comparing Eq. (13) with Eq. (4), we observe that the multipole structure is conserved under the rotation (11). However, the rotation from light-cone to canonical polarizations affects the strength of some multipoles, see Eqs. (14) and (15).

5.3. Spherical Symmetry
We are now ready to discuss the implications of spherical symmetry in the canonical-spin basis. Spherical symmetry means that the canonical tensor correlator has to be invariant, i.e. $O_R^\ast \Phi_{Cq}O_R = \Phi_{Cq}$ under any spatial rotation $O_R = \left( \begin{array}{ccc} 1 & 0 \\ 0 & R \end{array} \right)$ with $R$ the ordinary $3 \times 3$ rotation matrix. It is equivalent to the statement that the tensor correlator has to commute with all the elements of the rotation group $\Phi_{Cq}O_R = O_R\Phi_{Cq}$. As a result of Schur’s lemma, the canonical tensor correlator must have the following structure

$$\Phi_{Cq}^{\mu\nu} = \begin{pmatrix} A^q & 0 & 0 & 0 \\ 0 & B^q & 0 & 0 \\ 0 & 0 & B^q & 0 \\ 0 & 0 & 0 & B^q \end{pmatrix}. \quad (16)$$

Comparing this with Eq. (13), we conclude that spherical symmetry implies

$$f_1^q = A^q,$$

$$g_{1L}^q = h_{1T}^+ = h_{-1T}^- = B^q, \quad (17a)$$

$$g_{1T}^q = h_{1L}^\perp = f_{1T}^\perp = h_{1L}^\perp = 0. \quad (17b)$$

Spherical symmetry in the canonical-spin basis implies that only the monopole structures have non-vanishing amplitude in the canonical tensor correlator $\Phi_{Cq}$. Note however that because of the rotation connecting canonical and light-cone polarizations, see Eq. (11), higher multipole amplitudes are non-vanishing in the tensor correlator $\Phi_q$. It follows that spherical symmetry imposes some relations among the multipole structures in the light-cone helicity basis, and therefore among the TMDs. Inserting the constraints (17b) and (17c) into Eqs. (14) and (15), we automatically obtain the flavor-independent relations (6)-(8).

6. Conclusions
In this work we presented a study of the transverse-momentum dependent parton distributions in the framework of quark models. We discussed the predictions in the LC\(\chi\)QSM for the six T-even TMDs using the overlap representation in terms of 3Q LCWFs, and compared them with the ones in the LCCQM. We then focused the discussion on model relations among TMDs, showing that they have essentially a geometrical origin, and can be traced back to properties of rotational invariance of the system. In particular, we identified the conditions which are sufficient for the existence of the flavor-independent relations.

We presented a derivation of the relations based on the representation of the quark correlator entering the definition of TMDs in terms of the polarization amplitudes of the quarks and nucleon. Such amplitudes are usually expressed in the basis of light-cone helicity. However, in order to discuss in a simple way the rotational properties of the system, we introduced the representation in the basis of canonical spin and showed how the two bases are related.

Finally, we remark that the model relations are not expected to hold identically in QCD where TMDs are all independent. However, they provide simplified and intuitive notions for the interpretation of the spin and orbital angular momentum structure of the nucleon. As such, they can be useful for phenomenological studies to build up simplified parametrizations of TMDs to be fitted to data. Furthermore, the comparison with the experimental data will tell us the degree of accuracy of such relations, giving insights for further studies towards more refined quark models.
Acknowledgments
This work was supported in part by the Research Infrastructure Integrating Activity “Study of Strongly Interacting Matter” (acronym HadronPhysics2, Grant Agreement n. 227431) under the Seventh Framework Programme of the European Community, by the Italian MIUR through the PRIN 2008EKLACK “Structure of the nucleon: transverse momentum, transverse spin and orbital angular momentum”.

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