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THE QUARK DISTRIBUTION AMPLITUDE OF THE PROTON:-

A LATTICE COMPUTATION OF THE LOWEST TWO MOMENTS

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ABSTRACT

We compute the zeroth and first moments of the quark distribution amplitude of the proton on a $20 \times 10^2 \times 40$ lattice, using Wilson fermions in the quenched approximation and at $\beta = 6.0$. We find that the average fraction of momentum carried by each of the three valence quarks is approximately the same, in contradiction to the results obtained using QCD sum rules. For the "decay constant" of the proton we obtain the result:

$$f_p = (2.9 \pm 0.6 \text{ (statistical)}) \times 10^{-3} \text{ GeV}^2.$$

In this letter we present the results of a lattice computation of the zeroth and first moments of the quark distribution amplitude of the proton. The quark distribution amplitude, $\phi(x_i, Q^2)$, contains the non-perturbative information required to predict hard exclusive processes in QCD, such as hadronic form factors or elastic scattering at large momentum transfers. Q^2 is the square of the momentum transfer and the label i takes values from 1 to n_v , where n_v

is the number of valence partons. $\sum_{i=1}^{n_v} x_i = 1$ and in a physical gauge x_i is interpreted as the fraction of momentum carried by the valence parton i . In perturbation theory, for sufficiently large values of Q^2 , we can compute the behaviour of ϕ with Q^2 , as well as its asymptotic form, $\phi_{as} = \phi(x_i, Q^2 = \infty)^{(1,2)}$. However as the evolution with Q^2 is only logarithmic, the knowledge of the asymptotic form is not phenomenologically useful, (in the same way that the fact that deep inelastic structure functions tend to delta functions at $x = 0$ is not useful), and hence it is necessary to evaluate ϕ at a finite momentum transfer.

For the proton there are three distribution amplitudes V, A , and T defined by⁽³⁻⁶⁾

$$|p, \uparrow\rangle = \text{const} \int [dx] \left\{ \frac{1}{2} (V(x_i) - A(x_i)) | u^\uparrow(x_1) u^\downarrow(x_2) d^\uparrow(x_3) \rangle + \frac{1}{2} (V(x_i) + A(x_i)) | u^\downarrow(x_1) u^\uparrow(x_2) d^\uparrow(x_3) \rangle - T(x_i) | u^\uparrow(x_1) u^\uparrow(x_2) d^\downarrow(x_3) \rangle \right\} \quad (1)$$

where $|p, \uparrow\rangle$ is the proton state, the arrows denote chiralities and

$$[dx] \equiv dx_1 dx_2 dx_3 \delta(x_1 + x_2 + x_3 - 1) \quad (2)$$

The three distribution amplitudes are not independent. The symmetry due to the presence of two up quarks and the fact that the three quarks combined to make a state with isospin $\frac{1}{2}$ leads to two relations between the three distribution amplitudes⁽³⁻⁵⁾. Following ref.(1) we take ϕ to be the independent distribution amplitude where

$$\phi \equiv V - A \quad (3)$$

As $Q^2 \rightarrow \infty$, $\Phi(x_i, Q^2) \rightarrow \Phi_{as}(x_i)$ where^(1,2)

$$\Phi_{as}(x_1, x_2, x_3) = 120 x_1 x_2 x_3 \quad (4)$$

which is symmetric in the x_i 's. For phenomenological calculations however, we need to know Φ at a finite value of Q^2 ($Q^2 \approx \text{few GeV}^2$ or so). This problem has been studied using QCD sum rules⁽³⁻⁶⁾, with the conclusion that $\Phi(x_i, Q^2 \approx 1 \text{ GeV})$ is very different from the asymptotic function given in eqn.(4), and in particular is asymmetric in the x_i 's. For example, defining the moments of Φ by

$$\Phi^{(\ell, m, n)} \equiv \int_0^1 [dx] x_1^\ell x_2^m x_3^n \Phi(x_i, Q^2) \quad (5)$$

and choosing a normalisation so that $\int_0^1 [dx] \Phi(x_i, Q^2) = 1$, it was found that $\Phi^{(1,0,0)}$ is about 0.6, whereas $\Phi^{(0,1,0)} \approx \Phi^{(0,0,1)} \approx 0.2$. The distribution amplitude obtained using QCD sum rules can reproduce the data on the magnetic form-factors of the nucleons^(3,7), and in particular their signs, (the asymptotic amplitude of eqn.(4) together with lowest order perturbation theory would yield the wrong sign for the magnetic form factors of both the proton and neutron⁽³⁾). In this letter we report on the results of a lattice computation of the zeroth and first moments of Φ .

The starting point for both the sum-rule and lattice calculations is the expression for the moments of the distribution amplitude in terms of the matrix elements of the following local operators:

$$\begin{aligned} \hat{V}_T^{(\ell, m, n)}(0) \equiv \text{L.T.} \quad & [(iD^{\mu_i})^\ell u(0)]^a C\gamma^\lambda [(iD^{\nu_j})^m u(0)]^b \\ & \times [(iD^{\rho_k})^n \gamma_5 d(0)]_T^c \epsilon^{abc} \quad , \end{aligned} \quad (6a)$$

$$\begin{aligned} \hat{A}_T^{(\ell, m, n)}(0) \equiv \text{L.T.} \quad & [(iD^{\mu_i})^\ell u(0)]^a C\gamma^\lambda \gamma_5 [(iD^{\nu_j})^m u(0)]^b \\ & \times [(iD^{\rho_k})^n d(0)]_T^c \epsilon^{abc} \quad , \end{aligned} \quad (6b)$$

and

$$\begin{aligned} \hat{T}_\tau^{(\ell, m, n)}(o) \equiv \text{L.T.} \left[(iD^{\mu_i})^\ell u(o) \right]^a C(-i\sigma^{\sigma\lambda}) \left[(iD^{\nu_j})^m u(o) \right]^b \\ \times \left[(iD^{\rho_k})^n \gamma_\sigma \gamma_5 d(o) \right]^c \epsilon^{abc}, \end{aligned} \quad (6c)$$

where τ is a spinor label, C is the charge conjugation matrix and a, b and c are colour labels. The carat, $\hat{}$, on the left hand sides is introduced to distinguish the operator from the moment of the distribution amplitude, and

$$(iD^{\mu_i})^n = (iD^{\mu_1}) (iD^{\mu_2}) \dots (iD^{\mu_n}).$$

L.T. stands for "leading twist" and is introduced in eqn.(6) to indicate that the leading twist projection of these operators must be taken. The moments of the distribution amplitudes are given by

$$\langle o | \hat{V}_\tau^{(\ell, m, n)}(o) | p \rangle = -f_p (p^{\mu_i})^\ell (p^{\nu_j})^m (p^{\rho_k})^n p^{\lambda_{N_\tau}} \times V^{(\ell, m, n)} \quad (7a)$$

$$\langle o | \hat{A}_\tau^{(\ell, m, n)}(o) | p \rangle = -f_p (p^{\mu_i})^\ell (p^{\nu_j})^m (p^{\rho_k})^n p^{\lambda_{N_\tau}} \times A^{(\ell, m, n)} \quad (7b)$$

and

$$\langle o | \hat{T}_\tau^{(\ell, m, n)}(o) | p \rangle = 2f_p (p^{\mu_i})^\ell (p^{\nu_j})^m (p^{\rho_k})^n p^{\lambda_{N_\tau}} \times T^{(\ell, m, n)} \quad (7c)$$

where $N_\tau = N_\tau(p)$ is the spinor for the proton at rest.

The symmetry due to the presence of two up quarks implies that $V^{(\ell, m, n)}$ and $T^{(\ell, m, n)}$ are symmetric in ℓ and m , whereas $A^{(\ell, m, n)}$ is anti-symmetric. The fact that the proton has isospin $\frac{1}{2}$ implies that $2T^{(\ell, m, n)} = \phi^{(\ell, n, m)} - \phi^{(m, n, \ell)}$. Equation (7) for $\ell = m = n = 0$, together with the normalisation condition $T(0,0,0) = V(0,0,0) = 1$ provide the definition of the "proton's decay constant", f_p . Thus from the matrix elements (7) the moments of the distribution amplitudes can be obtained.

We have computed f_p , $\phi^{(1,0,0)}$, $\phi^{(0,1,0)}$ and $\phi^{(0,0,1)}$ on a $10^2 \times 20 \times 40$ lattice using Wilson fermions in the quenched approximation, and at a value of the coupling constant corresponding

to $\beta = 6/g_0^2 = 6.0$. We use the same 28 gauge field configurations¹, and the corresponding quark propagators, as those which were used in our studies of weak matrix elements and of hadronic structure, (for a discussion of the properties of these configurations see refs. (8) and (9)). The quark propagators have been generated at three values of the Wilson hopping parameter, $K = 0.1515$, $K = 0.1530$ and $K = 0.1545$. All our results have then to be extrapolated to the chiral limit ($K = 0.1563(2)$)(8,9). We now present our results together with a brief description of the calculation.

(i) Evaluation of f_p

There is clearly some freedom as to which operator to choose with which to evaluate f_p . Following ref.(10) we use the operator $\hat{f}$ defined by,

$$\hat{f}_\alpha = (\gamma_0 \hat{f}^0 - \frac{1}{4} \gamma_\nu \hat{f}^\nu)_\alpha \quad (8)$$

where

$$\hat{f}_\alpha^\nu = \epsilon^{abc} (u^a C \gamma^\nu \gamma_5 d^b) u_\alpha^c \quad (9a)$$

$$= \frac{1}{2} [\hat{V}_\alpha^{(0,0,0)} + \hat{A}_\alpha^{(0,0,0)} - \hat{T}_\alpha^{(0,0,0)}] \quad (9b)$$

α is a spinor label and (9b) can be derived from (9a) using the Fierz identity. The subtraction of the trace in equation (8) projects out the leading twist component. Thus for a proton at rest

$$\langle 0 | \hat{f}_\alpha | p \rangle = -\frac{9}{8} f_p m_p (\gamma_N^0)_\alpha \quad (10)$$

where m_p is the mass of the proton.

We evaluate f_p by computing the correlation function,

$$C_{(0)}(t_x) \equiv \sum_{y=y'=1}^2 \sum_{\underline{x}} \langle 0 | T(\hat{f}_y(x) \bar{J}_{p_{y'}}(0)) | 0 \rangle \quad (11)$$

where y and y' are spinor indices, t_x is the time coordinate of x , and J_p is an interpolating operator for the proton,

¹With the exception of $\phi(0,0,1)$ which was computed using only the subset of 15 configurations of ref.(9).

$$J_{p_y} \equiv \epsilon^{abc} (u^a C \gamma^5 d^b) u^c_y \quad (12)$$

Inserting a complete set of states between the operators in (11) and taking t_x to be sufficiently large so that only the single proton state gives a significant contribution, we have

$$C_{(0)}(t_x) = \frac{9}{4} \sqrt{Z_p} m_p f_p e^{-m_p t_x} \quad (13)$$

where $\sqrt{Z_p}$ is given by

$$\langle p | \bar{J}_{p,y}(0) | 0 \rangle \equiv \sqrt{Z_p} \bar{N}_y. \quad (14)$$

Z_p and m_p are determined by computing the correlation function (11) with $\hat{f}$ replaced by J_p ,

$$\begin{aligned} P(t_x) &\equiv \sum_{y=y'=1}^2 \sum_{\underline{x}} \langle 0 | T (J_{p_y}(x) \bar{J}_{p_{y'}}(0)) | 0 \rangle \\ &= -2 Z_p e^{-m_p t_x}. \end{aligned} \quad (15)$$

Thus f_p can be determined using equations (13) and (15).

In table 1 we present the results for the dimensionless quantity f_p/m_p^2 obtained by calculating the correlation functions (13) and (15) with $t_x = 7-11$. The results correspond to the operator $\hat{f}$ (see eqn (8)) renormalised in the lattice renormalisation scheme at a scale of about 2 GeV (i.e. $a^{-1} \approx 2 \text{ GeV}^{(8,9)}$). The results are very pleasing. They have small statistical errors (which increase slowly as t_x increases), are only very mildly dependent (if at all) on K , and the results obtained by performing the calculations at different values of t_x are consistent. In the last column of table 1 we present the values obtained by extrapolating the results at $K=0.1515, 0.1530$ and 0.1545 linearly in $1/K$ to the chiral limit ($K_c \approx 0.1563(2)$)^(8,9). The results obtained at different times are highly correlated and cannot be combined in quadrature. We take the result for $t_x = 8$ as our best estimate

$$\frac{f_p^{\text{LATT}}}{m_p^2} = (4.2 \pm 0.8) 10^{-3} \quad (16)$$

t_x	$\frac{f_p}{m_p^2}$			
	$K = 0.1515$	$K = 0.1530$	$K = 0.1545$	K_c
7	$(3.8 \pm 0.3)10^{-3}$	$(3.9 \pm 0.4)10^{-3}$	$(4.1 \pm 0.6)10^{-3}$	$(4.2 \pm 0.7)10^{-3}$
8	$(3.8 \pm 0.3)10^{-3}$	$(3.9 \pm 0.5)10^{-3}$	$(4.1 \pm 0.6)10^{-3}$	$(4.2 \pm 0.8)10^{-3}$
9	$(3.9 \pm 0.3)10^{-3}$	$(4.0 \pm 0.5)10^{-3}$	$(4.4 \pm 0.8)10^{-3}$	$(4.5 \pm 1.0)10^{-3}$
10	$(4.1 \pm 0.4)10^{-3}$	$(4.3 \pm 0.6)10^{-3}$	$(4.8 \pm 1.0)10^{-3}$	$(5.0 \pm 1.2)10^{-3}$
11	$(4.1 \pm 0.4)10^{-3}$	$(4.4 \pm 0.7)10^{-3}$	$(5.0 \pm 1.4)10^{-3}$	$(5.2 \pm 1.7)10^{-3}$

TABLE 1: Values of f_p/m_p^2 computed on the lattice at three values of the Wilson hopping parameter $K = 0.1515, 0.1530$ and 0.1545 and extrapolated to the chiral limit ($K_c = 0.1563(2)$).

The superscript LATT is included to remind us that the result is obtained in the lattice renormalisation scheme. Now f_p in the $\overline{MS}$ renormalisation scheme is about 21% smaller than in the lattice scheme at the same renormalisation scale for $\beta = 6.0^{(10)}$ so that we find (at a renormalisation scale of about 2 GeV)

$$f_p^{\overline{MS}} = (2.9 \pm 0.6) \times 10^{-3} \text{ GeV}^2 \quad (17)$$

The error quoted in (17) is the statistical error only. The main systematic error is the uncertainty in the inverse lattice spacing. We have used the proton's mass to set the scale in (17); had we used the rho mass we would have obtained an answer for $f_p^{\overline{MS}}$ which is about a factor of 2 larger. Our result (17) is smaller (by a factor of about 3) than the previous lattice result ⁽¹¹⁾, which was obtained on a $8^3 \times 16$ lattice at $\beta = 5.7$, (which was found in ref.(11) to correspond to an inverse lattice spacing of about 0.85 GeV), also with Wilson fermions and using the proton's mass to set the scale. The anomalous dimension of the operator $\hat{f}_\alpha$ (eqn (8)) is very small⁽¹⁰⁾ so that the one-loop perturbative evolution from 0.85 GeV to about

2 GeV is negligible. Much of the discrepancy between our result for f_p and that in ref (11) can however be accounted for by the different method of extrapolation to the chiral limit (we chose to work with the dimensionless quantity f_p/m_p^2 which is almost independent of K , whereas in ref (11) the results for $a^2 f_p$ are extrapolated linearly in the quark mass).

f_p has also been computed using QCD sum rules⁽³⁻⁶⁾ with the result $f_p = (5.0 \pm 0.3) \times 10^{-3} \text{ GeV}^2$. Given the inherent systematic uncertainties in both the lattice and sum rule calculations, we conclude that the results are consistent.

(ii) Evaluation of $\phi^{(1,0,0)}$, $\phi^{(0,1,0)}$ and $\phi^{(0,0,1)}$

We now compute the matrix elements $\langle 0 | \hat{\phi}^{(\ell,m,n)}(0) | p \rangle$ where $\ell + m + n = 1$ and $\hat{\phi}^{(\ell,m,n)} \equiv \gamma^{(\ell,m,n)} - \hat{A}^{(\ell,m,n)}$ (see equation (6)). These operators have two Lorentz indices and by choosing these suitably we can ensure that the leading twist part is projected out. To this end we choose to work with the operators (10)

$$\hat{\phi} \equiv \frac{i}{2} \left[\gamma^4 (\hat{\phi}^{44} - \hat{\phi}^{11}) - \gamma^1 (\hat{\phi}^{41} + \hat{\phi}^{14}) \right] \quad (18)$$

where we have now exhibited the Lorentz indices but suppressed the labels (ℓ, m, n) . (18) represents 3 equations, one for each of the three cases $(\ell, m, n) \equiv (1, 0, 0)$, $(0, 1, 0)$ and $(0, 0, 1)$. The moments of the distribution amplitude are determined by computing the coefficient functions

$$C_{(1)}^{(\ell,m,n)}(t_x) \equiv \sum_{\gamma=\gamma'=1}^2 \sum_{\underline{x}} \langle 0 | T (\hat{\phi}_{\gamma}^{(\ell,m,n)}(x) \bar{J}_{p_{\gamma'}}(0)) | 0 \rangle \quad (19)$$

where γ and γ' are spinor indices and t_x is the time component of x . The interpolating operator J_p is defined in equation (12). For sufficiently large values of t_x we have

$$\left| \frac{C_{(1)}^{(\ell,m,n)}(t_x)}{C_{(0)}^{(\ell,m,n)}(t_x)} \right| = \frac{4}{9} m_p \phi^{(\ell,m,n)} \quad (20)$$

Thus by computing these simple ratios the $\phi^{(\ell,m,n)}$'s can be determined.

In tables 2-4 we present the results for $\phi^{(1,0,0)}$, $\phi^{(0,1,0)}$ and $\phi^{(0,0,1)}$ respectively, for t_x in the range 7 to 11. From these tables it can readily be seen that the three moments $\phi^{(\ell,m,n)}$ with $\ell + m + n = 1$ are equal to very good precision (at least in the lattice renormalisation scheme). For example at $t_x=8$ we find

$$\begin{aligned} \frac{\phi^{(1,0,0)}}{\phi^{(0,1,0)}} &= 0.99 \pm 0.03 \quad \text{at} \quad K = 0.1515 \\ &= 0.97 \pm 0.06 \quad \text{at} \quad K = 0.1530 \\ &= 0.87 \pm 0.07 \quad \text{at} \quad K = 0.1545 \end{aligned} \quad (21)$$

and¹

$$\begin{aligned} \frac{\phi^{(1,0,0)}}{\phi^{(0,0,1)}} &= 1.04 \pm 0.06 \quad \text{at} \quad K = 0.1515 \\ &= 1.03 \pm 0.07 \quad \text{at} \quad K = 0.1530 \\ &= 0.96 \pm 0.10 \quad \text{at} \quad K = 0.1545 \end{aligned} \quad (22)$$

It is only for the leading twist projections that the ratios are close to 1. This is not the case for each of the four terms on the right hand side of equation (18) separately. Using QCD sum rules⁽³⁻⁶⁾ one finds values for the ratios (21) and (22) which are in the range 2.5 - 3.

The results presented in tables 2-4 are in the lattice renormalisation scheme at a renormalisation scale of about 2 GeV. The renormalisation of the operators $\hat{\phi}^{(\ell,m,n)}$ with $\ell + m + n = 1$ has been studied in detail in ref.(10). These three operators mix with a large number of other operators with the same quantum numbers, however all the non-diagonal mixing coefficients are very small. We have neglected these non-diagonal mixing terms. Using the results of ref.(10) we estimate that the moments $\phi^{(1,0,0)}$, $\phi^{(0,1,0)}$ and $\phi^{(0,0,1)}$ computed in the $\overline{MS}$ scheme at a renormalisation scale of about 2 GeV

¹In equation (22) both the numerator and denominator are computed using the subset of 15 configurations for which results for $\phi^{(0,0,1)}$ are available.

t_x	$\phi(1,0,0)$		
	$K = 0.1515$	$K = 0.1530$	$K = 0.1545$
7	0.38 ± 0.02	0.38 ± 0.03	0.34 ± 0.05
8	0.34 ± 0.03	$0.33 \begin{smallmatrix} + 0.06 \\ - 0.03 \end{smallmatrix}$	0.27 ± 0.05
9	0.34 ± 0.03	$0.33 \begin{smallmatrix} + 0.06 \\ - 0.03 \end{smallmatrix}$	0.28 ± 0.05
10	$0.31 \begin{smallmatrix} + 0.03 \\ - 0.01 \end{smallmatrix}$	$0.29 \begin{smallmatrix} + 0.04 \\ - 0.02 \end{smallmatrix}$	0.23 ± 0.07
11	$0.29 \begin{smallmatrix} + 0.04 \\ - 0.02 \end{smallmatrix}$	$0.27 \begin{smallmatrix} + 0.04 \\ - 0.03 \end{smallmatrix}$	$0.24 \begin{smallmatrix} + 0.08 \\ - 0.04 \end{smallmatrix}$

TABLE 2: Values of $\phi(1,0,0)$ computed at three values of the Wilson hopping parameter $K = 0.1515, 0.1530$ and 0.1545 .

t_x	$\phi(0,1,0)$		
	$K = 0.1515$	$K = 0.1530$	$K = 0.1545$
7	0.36 ± 0.02	0.36 ± 0.03	0.33 ± 0.04
8	0.35 ± 0.02	0.35 ± 0.03	0.31 ± 0.05
9	0.32 ± 0.02	$0.31 \begin{smallmatrix} + 0.04 \\ - 0.02 \end{smallmatrix}$	0.27 ± 0.03
10	0.31 ± 0.03	$0.30 \begin{smallmatrix} + 0.05 \\ - 0.03 \end{smallmatrix}$	$0.26 \begin{smallmatrix} + 0.05 \\ - 0.03 \end{smallmatrix}$
11	$0.30 \begin{smallmatrix} + 0.04 \\ - 0.02 \end{smallmatrix}$	$0.28 \begin{smallmatrix} + 0.06 \\ - 0.04 \end{smallmatrix}$	$0.23 \begin{smallmatrix} + 0.08 \\ - 0.04 \end{smallmatrix}$

TABLE 3: Values of $\phi(0,1,0)$ computed at three values of the Wilson parameter K .

t_x	$\phi(0,0,1)$		
	$K = 0.1515$	$K = 0.1530$	$K = 0.1545$
7	0.38 ± 0.02	0.39 ± 0.02	$0.43 \begin{smallmatrix} + 0.01 \\ - 0.03 \end{smallmatrix}$
8	0.36 ± 0.02	0.38 ± 0.02	$0.41 \begin{smallmatrix} + 0.01 \\ - 0.03 \end{smallmatrix}$
9	0.33 ± 0.01	0.34 ± 0.01	0.36 ± 0.02
10	0.34 ± 0.01	0.36 ± 0.02	$0.37 \begin{smallmatrix} + 0.02 \\ - 0.04 \end{smallmatrix}$
11	0.36 ± 0.01	0.37 ± 0.01	0.35 ± 0.05

TABLE 4: Values of $\phi(0,0,1)$ computed at three values of the Wilson parameter K .

are each about 9% larger than those in the lattice scheme which were presented in tables 2-4. There is virtually no change in the ratios (21) and (22) between the lattice and $\overline{MS}$ schemes.

We end this letter with a brief summary. Our main results are presented in equation (17) and in tables (2-4) (increased by 9% if the operators are renormalised in the $\overline{MS}$ renormalisation scheme). The statistical errors are very small particularly for the ratios (21) and (22). It is disturbing that the results of our lattice calculations disagree with those obtained using QCD sum rules⁽³⁻⁶⁾. The systematic uncertainties corresponding to these two methods of estimating the matrix elements (7) are very different. Our results indicate that these systematic errors must be large in at least one of the two cases¹ and cast serious doubt upon the validity of the

¹ For the pion's quark distribution amplitude the lattice and sum rule results are compatible⁽¹²⁾, however in that case the subtraction of power divergences induces a large ($\approx 50\%$) error in the lattice computation.

wavefunctions proposed in refs. (3-6). The phenomenological situation is also not completely clear. Using lowest order perturbation theory for the hard scattering amplitude it seems that asymmetric distribution amplitudes of the form obtained by using QCD sum rules are required to explain the data on hard exclusive processes (in particular to reproduce correctly the signs of the magnetic form factors of the nucleons⁽³⁾). On the other hand, the effective momentum scale at the parton level in exclusive processes at currently available values of momentum transfer is small^(13,14) so that higher order corrections, and perhaps non-leading twist contributions, are likely to be significant.

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