The Lattice Calculation of Moments of Structure Functions

R. Horsley

School of Physics
The University of Edinburgh

1 Introduction

Much of our knowledge about QCD and the structure of hadrons (mainly nucleons) has been gained from Deep Inelastic Scattering (DIS) experiments such as $eN \rightarrow eX$ or $\nu N \rightarrow \mu^- X$. The (inclusive) cross sections are determined by structure functions $F_1$ and $F_2$ when summing over beam and target polarisations (and an additional $F_3$ when using neutrino beams), and $g_1$, $g_2$ when both the beam and target are suitably polarised. Structure functions are functions of the Bjorken variable $x$ and $Q^2$, the large space-like momentum transfer from the lepton. (Another class of structure functions – the transversity $h_1$ – can be measured in Drell-Yan processes or certain types of semi-inclusive processes.)

As a direct theoretical computation of structure functions does not seem to be possible, we must turn to the Wilson Operator Product Expansion (OPE) which relates moments of structure functions to (nucleon) matrix elements in a twist (ie operator [dimension - spin]) or Taylor expansion in $1/Q^2$. So first defining bilinear quark operators

$$\mathcal{O}_q^{\Gamma; \mu_1 \cdots \mu_n} = \frac{1}{\Gamma} \Gamma^{\mu_1 \cdots \mu_i} \hat{D}^{\mu_{i+1}} \cdots \hat{D}^{\mu_n} q,$$

(1)

where $q$ is taken to be either a $u$ or $d$ quark and $\Gamma$ is an arbitrary Dirac gamma matrix we have for the nucleon matrix elements, the Lorentz decompositions ($s^2 = -m_N^2$)

$$\frac{1}{2} \sum_s \langle \vec{p}, \vec{s} | \mathcal{O}_q^{\Gamma; \mu_1 \cdots \mu_n} | \vec{p}', \vec{s}' \rangle = 2 \nu_n^{(q)} \left[ p^{\mu_1} \cdots p^{\mu_n} - tr \right],$$

$$\langle \vec{p}, \vec{s} | \mathcal{O}_q^{\gamma_5; \sigma \mu_1 \cdots \mu_n} | \vec{p}', \vec{s}' \rangle = \alpha_n^{(q)} \left[ s^{\sigma} p^{\mu_1} \cdots p^{\mu_n} - tr \right],$$

$$\langle \vec{p}, \vec{s} | \mathcal{O}_q^{\gamma_5; \sigma \mu_1 \cdots \mu_n} | \vec{p}', \vec{s}' \rangle = \frac{nd_n^{(q)}}{n+1} \left[ (s^{\sigma} p^{\mu_1} - p^{\sigma} s^{(\mu_1)}) p^{\mu_2} \cdots p^{\mu_n} - tr \right],$$

where

$$\nu_n^{(q)} = \frac{1}{2} \sum_s \langle \vec{p}, \vec{s} | \mathcal{O}_q^{\Gamma; \mu_1 \cdots \mu_n} | \vec{p}', \vec{s}' \rangle$$

$$\alpha_n^{(q)} = \frac{1}{2} \sum_s \langle \vec{p}, \vec{s} | \mathcal{O}_q^{\gamma_5; \sigma \mu_1 \cdots \mu_n} | \vec{p}', \vec{s}' \rangle$$

and $d_n^{(q)}$ are the co-efficients for the Lorentz decompositions.
\[
\langle \bar{p}, s | \bar{O}_q^\gamma \gamma_5 \bar{\sigma}^{\mu_1 \ldots \mu_n} | \bar{p}, s \rangle = \frac{\epsilon^{(q)}}{m_N} \left[ (s^\sigma p^{\mu_1} + p^\sigma s^{\mu_1}) p^{\mu_2} \cdots p^{\mu_n} - \text{tr} \right],
\]

where the symmetrisation/anti-symmetrisation operations on the operator indices also indicates that they are traceless (which gives them a definite spin). \(v_n, a_n, d_n\) and \(t_n\) can be related to moments of the structure functions. For example, we have for \(v_n\) and \(F_2\)

\[
\int_0^1 dx x^{n-2} F_2(x, Q^2) = \frac{1}{3} \sum_{f=u,d,g,\ldots} E^{(\overline{f}\gamma_5)}_{F_2:n}(\mu^2/Q^2, g^{\overline{f}\gamma_5}) v_n^{(\overline{f}\gamma_5)}(\mu) + O(1/Q^2),
\]

and similar relations hold between \(g_1\) and \(a_n; g_2\) and a linear combination of \(a_n\) and \(d_n; h_1\) and \(t_n\). Although the OPE gives \(v_n\) from \(F_1\) (or \(F_2\)) for \(n = 2, 4, \ldots\); \(v_n\) from \(F_3\) for \(n = 3, 5, \ldots\); \(a_n\) from \(g_1\) for \(n = 0, 2, \ldots\); \(a_n, d_n\) from \(g_2\) for \(n = 2, 4, \ldots\), other matrix elements can be determined form semi-exclusive experiments, for example \(a_1\) by measuring \(\pi^\pm\) in the final state.

While the Wilson coefficients, \(E^{\overline{f}\gamma_5}(1, g^{\overline{f}\gamma_5}(Q))\) are known perturbatively (typically two to three loops) and determine how the moments change with scale, the ‘initial condition’ ie the matrix element is non-perturbative in nature. The only known way of determining them from QCD in a model independent way is via Lattice Gauge Theory (LGT). In this talk we review our status (QCDSF and UKQCD Collaborations) of some aspects of these determinations including some higher twist results. A more general review may be found in (Göckeler et al 2002a). We shall also restrict ourselves here to forward matrix elements (and so not consider the form factors and the more embracing Generalised Parton Densities or GPDs). Determining moments of structure functions is an active field of research at present, see for example (Dolgov et al 2002, Guagnelli et al 2004, Ohta et al 2004).

2 The Lattice Approach

The lattice approach involves first Euclideanising the QCD action and then discretising space-time with lattice spacing \(a\). The path integral then becomes a very high dimensional partition function, which is amenable to Monte Carlo methods of statistical physics. This allows ratios of three-point to two-point correlation functions to be defined,

\[
R_{\alpha\beta}(t; \tau; \bar{p}) = \frac{(N_{\alpha}(t; \bar{p})\bar{O}_q(\tau)\bar{N}_\beta(0; \bar{p}))}{(N(t; \bar{p})\bar{N}(0; \bar{p}))} \propto (N_{\alpha}(\bar{p})|\bar{O}_q|N_{\beta}(\bar{p})) ,
\]

where \(N_{\alpha}\) is some suitable nucleon wavefunction (with Dirac index \(\alpha\)) such as

\[
N_{\alpha}(t; \bar{p}) = \sum_{\vec{x}} e^{-i\vec{p}\cdot\vec{x}} \epsilon^{ijk} u_{\alpha}(\vec{x}, t)|u_{\beta}(\vec{x}, t)(C\gamma_5)_{\beta\alpha} d_{\gamma}(\vec{x}, t) .
\]

The proportionality holds for \(0 \ll \tau \ll t \ll \frac{1}{2}N_T\) for a lattice of size \(N_T^3 \times N_T\). There are two basic types of diagrams to compute in eq. (4): the first is a quark insertion in one of the nucleon quark lines (‘quark line connected’), while in the second type the operator interacts only via gluon exchange with the nucleon (‘quark line disconnected’). Due to gluon UV fluctuations these latter diagrams are numerically difficult to compute. However
The Lattice Calculation of Moments of Structure Functions

by considering the Non-Singlet, NS, or $O_u - O_d$ operators, giving matrix elements such as $v_n,NS = v_n^{(u)} - v_n^{(d)}$ in eq. $\mathbb{B}$ then the $f = s$ and $g$ (gluon) terms cancel. (For higher moments however, one might expect that sea effects anyway are less significant as the integral is more weighted to $x \sim 1$.) Although LQCD is in principle an ‘ab initio’ calculation there are, of course, several caveats. First our lattice ‘box’ must be large enough to fit our correlation functions into. A continuum limit $a \to 0$ must be taken. A chiral extrapolation must be made from simulations often at the strange quark mass or larger down to the almost massless $u/d$ quarks, or until we can match to Chiral Perturbation Theory, $\chi$PT (the problem there being that the radius of convergence of $\chi$PT is not known). Also to save CPU time, the fermion determinant in the action (representing $n_f$ quark flavours) is often discarded - the ‘quenched’ approximation. Finally in addition to all the above problems the matrix element must be renormalised, in order to be able to compare with the phenomenological $\overline{MS}$ results.

To attempt to address some of these issues we have generated data sets (Bakeyev et al 2004)

1. $O(a)$-improved Wilson fermions (‘clover fermions’) in the quenched approximation at three couplings $\beta \equiv 6/g^2 = 6.0, 6.2$ and 6.4 (Göckeler et al 2004) corresponding to lattice spacings $a^{-1} \sim 2.12, 2.91$ and $3.85$ GeV. (This checks lattice discretisation errors, which should be $O(a^2)$.) The pseudoscalar mass, $m_{ps}$, lies between 580 MeV and 1200 MeV.

2. Unquenched clover fermions at $m_{ps}$ down to $\sim 560$ MeV in order to see if there are any discernable quenching effects. Various couplings are used, $\beta = 5.20, 5.25, 5.29$ and 5.40 with lattice spacings ranging from $a^{-1} \sim 1.61$ GeV to 2.4 GeV.

3. Wilson fermions at one fixed lattice spacing, $a^{-1} \sim 2.12$ GeV in the quenched approximation at pseudoscalar masses, $m_{ps}$, down to $\sim 310$ MeV, ($m_{ps}/m_V \sim 0.4$) in order to try to match to chiral perturbation theory. This lattice fermion formulation has discretisation errors of $O(a)$.

4. Overlap fermions, in the quenched approximation at one lattice spacing $a^{-1} \sim 2.09$ GeV down to $m_{ps}$ of about 440 MeV. These have a chiral symmetry even with finite lattice spacing and hence have better chiral properties than either Wilson or clover fermions (and also have discretisation errors of $O(a^2)$).

Note that the physical pion mass is about $m_\pi \sim 140$ MeV and we use the force scale $r_0 = 0.5$ fm $\equiv (394.6$ MeV $)^{-1}$ to set the scale. These results cover various patches of ($m_{ps}, a, n_f$) space. This is however not completely satisfactory. Overlap fermions, although the best formulation of lattice fermions known, are very expensive in CPU time, and are only just beginning to be investigated, eg (Galletly et al 2003, Ohta et al 2004), to which we refer the reader to for more details.

The results obtained from eq. $\mathbb{E}$ are, of course, bare results and must be renormalised. We shall not discuss this further here, just noting that many one-loop perturbative results are known; but are generally not very satisfactory as lattice perturbative series do not appear to converge very fast. (The convergence can be helped using tadpole-improvement.) A preferred non-perturbative method is also available, (Martinelli et al 1995) and the results presented here will have the $\tilde{Z}s$ determined by this method.
All our results are for hadrons containing light (i.e., $u/d$) quarks. Reaching this limit is extremely costly in CPU time (and except for the overlap formulation, other problems connected with the non-chiral nature of the fermion formulation may arise). Much work has been done recently on chiral perturbation theory and it would be highly desirable to be in a region where these results can be matched to lattice results and then the limit $m_{ps} \to m_\pi$ can be taken. Although one should take the continuum and chiral limits separately (and preferably in that order) we shall try here a variant procedure of using a simultaneous "plane" fit containing both limits. This is because at present the unquenched, data set 2 is less complete than the quenched data set 1 and this procedure at least allows for a direct comparison of results. (For set 1 these different fit procedures lead to similar results.) Practically we might thus expect that for a quantity $Q$ of interest

$$Q = F^Q(r_0 m_{ps}) + d^Q_s (a/r_0)^s. \quad (6)$$

$F^Q(r_0 m_{ps})$ describes the (chiral) physics and the discretisation errors are $O(a^s)$ where $s = 1$ for Wilson fermions and $s = 2$ for clover fermions. Naively one might expect a Taylor series expansion for $F^Q_N$ to be sufficient, i.e.

$$F^Q_N(x) = Q(0) + c^Q x^2 + \ldots, \quad (7)$$

where $x = r_0 m_{ps}$. Over the last few years expressions for $F^Q_N$ have been found

$$F^Q_N(x) = Q(0) \left( 1 - c^Q x^2 \ln(x/r_0 \Lambda^2) \right) + \ldots, \quad (8)$$

showing the existence of a chiral logarithm $\sim m_q \ln m_q$ (including the quenched case). For $v_{n:NS}$, $a_{n:NS}$, $t_{n:NS}$ the constant $c^Q_N$ is known (and positive), see eg (Chen et al 1997). One expects most effect of the chiral logarithm for $t_{1;NS}$ and least for $a_{0;NS}$. The chiral scale, $\Lambda$, is usually taken to be $\sim 1$ GeV. The range of validity of the expansion, eq. (8) is not known; one might expect that for $m_{ps} > \Lambda$, pion loops are suppressed, leading to a smooth variation in $m_q$ (i.e., constituent quarks, while for $m_{ps} < \Lambda$ non-linear behaviour would be seen. Thus building in some of the constituent or heavy quark mass expectations, an equation of the form

$$F^Q_N(x) = Q(0) \left( 1 - c^Q x^2 \ln \left( \frac{x^2}{(x^2 + (r_0 \Lambda)^2)} \right) \right) + c^Q x^2, \quad (9)$$

has been proposed (Detmold et al 2001).

Present (numerically) investigated matrix elements include $v_2 \equiv \langle x \rangle$ (also part of the momentum sum rule: $\sum_q \langle x \rangle^{(q)} + \langle x \rangle^{(q)}) = 1$, $v_3 \equiv \langle x^2 \rangle$, $v_4 \equiv \langle x^3 \rangle$, $a_0 = 2 \Delta q$ (also occurring in neutron decay, as well as the Bjorken sum rule, as $\Delta u - \Delta d = 2 A$ and connected with quark spin), $a_1 = 2 \Delta q^{(2)}$, $a_2 = 2 \Delta q^{(3)}$, $t_0 = 2 \delta q$, $t_1 = 2 \delta q^{(2)}$ and $d_2$. We shall only discuss $v_n$, $n = 1, 2, 3, a_0/2, t_0/2$ and $d_2$ here.

3 Results of Continuum/Chiral Extrapolations for some twist two operators

We now show some results. We start by considering $v_{2;NS}(2 \text{GeV})$ for $n = 2, 3, 4$. In Fig. 1 we show $v_{2;NS}$ from data set 1, together with a fit using eqs. (6) and (7). We see
The Lattice Calculation of Moments of Structure Functions

Figure 1. $v_{2,NS}^{MS} (2\text{ GeV})$ versus $(r_0 m_{ps})^2$ (upper plot) and versus $(a/r_0)^2$ (lower plot) using data set 1. Filled circles, squares and diamonds represent the three lattice spacings corresponding to $\beta = 6.0, 6.2, 6.4$. The chiral limit $(r_0 m_{ps})^2 = 0$ is shown as a short-dashed line, while the physical pion mass is denoted by the long-dashed line. Also shown as a dot-dashed line is the mass of a hypothetical $\Xi_s$ meson calculated as $\sim \sqrt{2m_K}$. The MRST phenomenological value is denoted by a star.

that $O(a^2)$ discretisation errors are small and seem to be relatively benign. By this we mean that the only limiting factor with the extrapolation is the amount of data available. We shall (thus) in future assume that this limit is not a problem. This does not seem to be the case with the chiral extrapolation where the data seems to strongly favour a linear extrapolation rather than the $\chi$PT result in eq. (8). The value found in the chiral limit is about 50% larger than the MRST phenomenological value, (Martin et al 2002). (Note however that there are exciting hints that overlap fermions may be closer to the phenomenological value, (Galletly et al 2003, Gürtler et al 2004).)

The same situation persists for the higher moments $v_{3,NS}^{MS}$ and $v_{4,NS}^{MS}$. In Figs. 2 and 3 we show these moments and compare the results with the MRS phenomenological values. In all cases we find that the moments are too large in comparison with phenomenological result. It is not clear why this is so, again the quarks seem to be acting more like constituent quarks rather than current quarks. Possible causes are quenching and/or a chiral extrapolation from too heavy a quark mass. We first consider possible quenching effects. In Fig. 4 we consider $v_{2,NS}^{MS}$ again, but this time for $n_f = 2$ flavours using data set 2. No real difference is seen in comparison to the quenched case. Indeed for other matrix elements considered a similar situation prevails.
To try to examine the situation at smaller quark mass, we now turn to the data set 3. Most of the above results have a quark mass at the strange quark mass (or heavier). In this data set we have generated quenched Wilson data at one lattice spacing, at light pion masses down to $310$ MeV. In Fig. 5 we show $\nu_{3;NS}^{\text{MS}}(2 \text{ GeV})$ versus $(r_0 m_{ps})^2$. In comparison with the previous pictures note that the $x$-scale only runs to $(r_0 m_{ps})^2 \sim 3.0$. Again (except possibly for the lightest pion mass) the data seems rather linear (and constant). Also shown is a forced fit from eq. (9), leaving $\Lambda_\chi$ and $c_{\chi}^{\nu 2}$ free but constrained to go through the MRST phenomenological value at $m_{ps} = m_\pi$. Ignoring the lightest quark mass point, this is just possible; however it is very unnatural giving, for example, $\Lambda_\chi \sim 500$ GeV which is a very low value.

A similar situation holds for the axial $a_0;NS/2 = g_A$ and tensor charge $t_0;NS/2$. In Figs. 6 and 7. Again the results in both cases seem very linear. Indeed from eq. (9) due to the negative sign, we must have the lattice data decreasing to the phenomenological value. This is certainly not the case here (although experimentally $t_0;NS$ is not known, we expect a similar situation as for $g_A$; indeed in the non-relativistic limit $t_0;NS/2 \rightarrow g_A$). Later work including the $\Delta$ as well as the $N$ in chiral perturbation theory (Detmold et al 2002, Hemmert et al 2003) reduce the $c_{\chi}^{g_A}$ coefficient, but this still is a problem. We show only linear fits, giving for $g_A$ a value somewhat lower than the experimental one.
The Lattice Calculation of Moments of Structure Functions

Figure 4. $v_{2rNS}^2(2\text{ GeV})$ versus $(r_0m_{ps})^2$ for unquenched fermions using data set 2. \(\beta = 5.20\) results are (filled) circles; 5.25 squares; 5.29, diamonds; 5.40 down triangle. Otherwise the same notation as in Fig. 7.

Figure 5. $v_{2rNS}^2(2\text{ GeV})$ versus $(r_0m_{ps})^2$ for Wilson fermions from data set 3. The ‘fit’ uses eq. 9. Otherwise the same notation as in Fig. 7.

(There are two possible caveats: for clover fermions the continuum extrapolation may have significant $O(a^2)$ effects, in distinction to $v_{n;NS}$ and also there may or may not be larger finite volume effects present both in the data and theoretically, see for example the discussion in (Cohen 2001).) But at present the same general picture emerges as for the unpolarised moments.

4 Some higher twist operator results

4.1 Twist three

The prime example is given by $d_2$, which can be determined from $g_2$, the first moment of which is a linear combination of $a_2$ and $d_2$. The operators for the $a_n$ moments have twist two, but $d_n$ corresponds to twist three and is thus of particular interest. A 'straightforward' lattice computation, (Göckeler et al 1996), gave rather large values for $d_2$ (where
\[ d_2^{(p)} = Q_u^2 d_2^{(u)} + Q_d^2 d_2^{(d)} \]. A recent experiment, (E155 Collaboration 1999), however indicated that this term was very small, which would mean that \( g_2 \) is almost completely determined by \( g_1 \) (the Wandzura–Wilczek relation). This problem was traced in (Göckeler et al 2000a), to a mixing of the original operator with a lower-dimensional operator. This additional operator mixes \( \propto 1/a \) and so its renormalisation constant must be determined non-perturbatively. In (Göckeler et al 2000a) this procedure was attempted, and led to results qualitatively consistent with the experimental values. Note that this is only a problem when using Wilson or clover like fermions, as we would expect the additional operator to appear like \( \sim m_q \pi \sigma \bar{D}q \) and hence vanish in the chiral limit. Thus there should be no mixing if one uses overlap fermions.

### 4.2 Twist four

Potential higher twist effects are present in the moment of a structure function, see eq. 3. These \( O(1/Q^2) \) terms are composed of dimension 6, four quark matrix elements. A general problem is the non-perturbative mixing of these operators with the previous dimen-
sion 4 operators. At present results are restricted to finding combinations of these higher twist operators which do not mix from flavour symmetry. For the nucleon the $SU_F(3)$ flavour symmetry group must be considered, i.e., taking mass degenerate $u$, $d$ and $s$ quarks, (Göckeler et al 2002b) giving

$$\int_0^1 dx F_2(x, Q^2)|_{Nachtmann}^{I=1} = -0.0005(5) \frac{m_2^2 \alpha_s(Q^2)}{Q^2} + O(\alpha_s^2),$$

for quenched Wilson fermions (i.e., part of data set 3). To access this moment experimentally needs very exotic combinations of moments from the measurement of the $p$, $n$, $\Lambda$, $\Sigma$ and $\Xi$ baryons and is not possible. Nevertheless this term is very small in comparison with the leading twist result, and might hint that higher twist contributions are small.

5 Miscellaneous Pion and Lambda results

Moments for the pion and rho structure functions were computed in (Best et al 1997), for unimproved Wilson fermions. Using the Schrödinger Functional method, $v_2$ was recently calculated for the pion, (Guagnelli et al 2004) for both unimproved and $O(\alpha)$-improved fermions. A higher twist matrix element for the pion has also been computed for quenched Wilson fermions (i.e., using part of data set 3)

$$\int_0^1 dx F_2(x, Q^2)|_{Nachtmann}^{I=2} = 1.67(64) \frac{f_2^2 \alpha_s(Q^2)}{Q^2} + O(\alpha_s^2),$$

where the $SU_F(2)$ flavour symmetry group gives the combination $F_2^{I=2} = F_2^{\pi^+} + F_2^{\pi^-} - 2F_2^{\pi^0}$. This is again a rather small number, so although a rather exotic combination of pion matrix elements it might indicate that higher twist terms are small.

Finally there have been results for moments of $\Lambda$ structure functions, (Göckeler et al 2002b) again using Wilson fermions (i.e., part of data set 3). These are potentially useful results as one can compare with nucleon spin structure and check violation of $SU_F(3)$ symmetry. First indications are that there is no evidence of flavour symmetry breaking in the matrix elements i.e., that $\Lambda$ and $p$ are related by an $SU(3)_F$ flavour transformation.

6 Conclusions

Clearly the computation of many matrix elements giving low moments of structure functions is possible. We would like to emphasise that a successful computation is a fundamental test of QCD – this is not a model computation. There are however many problems to overcome: finite volume effects, renormalisation and mixing, continuum and chiral extrapolations and unquenching. At present although overall impressions are encouraging, still it is difficult to re-produce experimental/phenomenological results of (relatively) simple matrix elements, e.g., $v_2$. But progress is being made by the various groups working in the field. For example in comparison to our previous results, (Göckeler et al 1996) there are now non-perturbative $Z$s and considerations of both chiral and continuum extrapolations and some unquenched results are now available. While the
continuum extrapolation seems to be ‘just’ a matter of more data at smaller lattice spacing, the chiral extrapolation does seem to present a problem, with no sign of any chiral logarithms being seen as predicted by \( \chi \)PT. Clearly everything depends on the data and the quest for better results should continue. To leave the region where constituent quark masses give a reasonable description of the data (ie linearity) unfortunately requires pion masses rather close to the physical pion mass. In this region fermions with better chiral properties will probably be needed, such as overlap, which in turn will need much faster machines.

Acknowledgements

I wish to thank my co-workers in the QCDSF and UKQCD Collaborations: A. Ali Khan, T. Bakeyev, D. Galletly, M. Göckeler, M. Gürtler, P. Hägler, T. R. Hemmert, A. C. Irving, B. Joó, A. D. Kennedy, B. Pendleton, H. Perlt, D. Pleiter, P. E. L. Rakow, A. Schäfer, G. Schierholz, A. Schiller, W. Schroers, T. Streuer, H. Stüben, V. Weinberg and J. M. Zanotti for a pleasant and profitable collaboration.

The numerical calculations have been performed on the Hitachi SR8000 at LRZ (Munich), on the Cray T3E at EPCC (Edinburgh) (Allton et al 2002) on the Cray T3E at NIC (Jülich) and ZIB (Berlin), as well as on the APE1000 and Quadrics (QH2b) at DESY (Zeuthen). We thank all institutions. This work has been supported in part by the EU Integrated Infrastructure Initiative Hadron Physics, contract number RII3-CT-2004-506078 and by the DFG (Forscherguppe Gitter-Hadronen-Phänomenologie).

References

Allton C R et al, 2002, Phys Rev D65 054502, [hep-lat/0107021]
Bakeyev T et al, 2004, Nucl Phys Proc Suppl 128 82, [hep-lat/0311017]
Best C et al, 1997, Phys Rev D56 2743, [hep-lat/970314]
Chen J-W et al, 1997, Phys Lett B523 107, [hep-ph/0105197]
Cohen T D, 2001, Phys Lett B529 50, [hep-lat/0112014]
Detmold W et al, 2001, Phys Rev Lett 87 172001, [hep-lat/0103006]
Detmold W et al, 2002, Phys Rev D66 054501, [hep-lat/0208001]
Dolgov D et al, 2002, Phys Rev D66 034506, [hep-lat/0201021]
E155 Collaboration, 1999, Phys Lett B458 529, [hep-ex/9901006]
Galletly D et al, 2003, Nucl Phys Proc Suppl 129 453, [hep-lat/0310028]
Göckeler M et al, 1996, Phys Rev D53 2317, [hep-lat/9508004]
Göckeler M et al, 2000a, Phys Rev D63 074506, [hep-lat/0001109]
Göckeler M et al, 2000b, Nucl Phys B623 287, [hep-lat/0010308]
Göckeler M et al, 2002a, Nucl Phys Proc Suppl B119 398, [hep-lat/0209111]
Göckeler M et al, 2002b, Phys Lett B545 112, [hep-lat/0208017]
Göckeler M et al, 2004, [hep-ph/0410187]
Guagnelli M et al, 2004, [hep-lat/0405027]
Gürtler M et al, 2004, [hep-lat/0409164]
Hemmert T R et al, 2003, Phys Rev D68 075009 [hep-lat/0303002]
Martin A D et al, 2002, Eur Phys J C23 73, [hep-ph/0110215]
Martinelli G et al, 1995, Nucl Phys B445 81, [hep-lat/9411010]
Ohta S et al, 2004, [hep-lat/0411008]