Four-loop QCD analysis of the Bjorken sum rule vs data

V.L. Khandramai\textsuperscript{a}, R.S. Pasechnik\textsuperscript{b}, D.V. Shirkov\textsuperscript{c}, O.P. Solovtsova\textsuperscript{a,c}, O.V. Teryaev\textsuperscript{c}

\textsuperscript{a}Gomel State Technical University, 246746 Gomel, Belarus
\textsuperscript{b}High Energy Physics, Department of Physics and Astronomy, Uppsala University, SE-75121 Uppsala, Sweden
\textsuperscript{c}Bogoliubov Laboratory of Theoretical Physics, Joint Institute for Nuclear Research, 141980 Dubna, Moscow Region, Russia

Abstract

We study the polarized Bjorken sum rule at low momentum transfers in the range $0.22 < Q < 1.73$ GeV with the four-loop N$^3$LO expression for the coefficient function $C_{Bj}(\alpha_s)$ in the framework of the common QCD perturbation theory (PT) and the singularity-free analytic perturbation theory (APT). The analysis of the PT series for $C_{Bj}(\alpha_s)$ gives a hint to its asymptotic nature manifesting itself in the region $Q < 1$ GeV. It relates to the observation that the accuracy of both the three- and four-loop PT predictions happens to be at the same 10% level. On the other hand, the usage of the two-loop APT allows one to describe the precise low energy JLab data down to $Q \approx 300$ MeV and gives a possibility for reliable extraction of the higher twist (HT) corrections. At the same time, above $Q \approx 700$ MeV the APT two-loop order with HT is equivalent to the four-loop PT with HT compatible to zero and is adequate to current accuracy of the data.

PACS: 11.10.Hi, 11.55.Hx, 11.55.Fv, 12.38.Bx, 12.38.Cy

1. Introduction

The higher order perturbative QCD (pQCD) and higher twist corrections become very important, in particular, in observables of the Deep-Inelastic Scattering (DIS) at low momentum transfers $Q \leq 1$ GeV. The most precise low-energy data from the Jefferson Lab \cite{1,2} on one of the main sources of information about the nucleon structure, the Bjorken sum rule (BSR) \cite{3}, are the real challenge to the accuracy of the pQCD expansions. In our previous papers \cite{4,5}, we studied this issue at the three-loop level. In the current paper, we continue this line of investigations and explore the effect of the recent calculation \cite{6} of the four-loop (in $\alpha_s$) contribution to the BSR.

The BSR claims that the difference of the proton and neutron structure functions integrated over all possible values of the Bjorken variable $x$ in the limit of large four-momentum squared of the exchanged virtual photon, $Q^2 \to \infty$, is equal to $g_A/6$, with $g_A = 1.267 \pm 0.004$ \cite{7}, the nucleon axial charge defined from the neutron $\beta$-decay data.

The r.h.s. of Eq. (1) is given by a sum of two series in powers of $1/Q^2$ (OPE higher twists corrections) and in powers of the QCD running coupling $\alpha_s(Q^2)$ (pQCD radiative corrections). Until very recently, the pQCD contribution to BSR was known \cite{8} up to the third order $\sim \alpha_s^3$. So far, the corresponding expression has been used in many studies, in particular, for extraction of the $\alpha_s$ values at low momentum scales \cite{9}.

One of the actual theoretical subjects is the interplay between the higher twists (HT) and higher order pQCD corrections at low $Q$, which has recently been studied in Refs. \cite{4} at the three-loop level. There, it was shown that the satisfactory description of the data down to $Q_{min} \sim 350$ MeV can be achieved within the Analytic Perturbation Theory (APT), the ghost-free modification of pQCD. In the current work we repeat this analysis at the four-loop N$^3$LO level.

The APT approach is based on the causality principle implemented as the analyticity imperative in the complex $Q^2$-plane for the QCD coupling $\alpha_s(Q^2)$ in the form of the Källén-Lehmann spectral representation \cite{10} and on the demand of compatibility with linear integral transformations \cite{11} (for an overview on the APT concept and results, see Ref. \cite{12}). It is well-known that in the APT framework, the theoretical ambiguity associated with pQCD higher-loop corrections is diminished (see Ref. \cite{13}), and results are practically renormalisation scheme independent \cite{14}.

The four-loop expression for the pQCD contribution to the Bjorken sum rule became recently available in Ref. \cite{6}. It gives us a reasonable motivation for a new extended QCD analysis of the combined JLab data on $\Gamma_{1-1}(Q^2)$ at low $0.05 < Q^2 < 3.0$ GeV$^2$ accounting for up to $\alpha_s^4$-order in both the (standard) PT and APT approaches.

The paper is organized as follows. In Section 2, we study the higher loop stability of both the PT and APT series and the renormalisation scale dependence of the
higher-order PT expansion for the Bjorken sum rule. Section 3 contains the QCD results on extraction of the higher twist terms from the experimental data at the four-loop level. Summarizing comments are given in the last section.

2. The perturbative QCD contribution

Commonly, one represents the Bjorken integral \( \Gamma_1^{p-n}(Q^2) \) as a sum of the perturbative and the higher twist contributions

\[
\Gamma_1^{p-n}(Q^2) = \frac{gA}{6} \left[ 1 - \Delta_{\text{Bj}}(Q^2) \right] + \sum_{i=2}^{\infty} \frac{\mu_{2i}}{Q^{2i-2}}. \tag{2}
\]

The perturbative QCD correction \( \Delta_{\text{Bj}}(Q^2) \) has a form of the power series in the QCD running coupling \( \alpha_s(Q^2) \). At the up-to-date four-loop (N^3LO) level in the massless case it looks like

\[
\Delta_{\text{Bj}}^{\text{PT}}(Q^2) = \sum_{k \leq 4} c_k \alpha_s^k(Q^2). \tag{3}
\]

Here, the numerical expansion coefficients \( c_k \) in the modified minimal subtraction (MS) scheme, for three active flavors, \( n_f = 3 \), read \( c_1 = 1/\pi = 0.31831 \), \( c_2 = 0.36307 \) \[15\], \( c_3 = 0.65197 \) \[8\] and \( c_4 = 1.8042 \) \[6\]. Besides, the four-loop running coupling \( \alpha_s(Q^2) \) is defined as a solution of the Renormalization Group (RG) equation

\[
\frac{d\alpha_s}{dL} = \beta(\alpha_s); \quad \beta(\alpha_s) = \sum_{0 \leq k \leq 3} \beta_k \alpha_s^{k+2}, \tag{4}
\]

where \( L = \ln(\mu^2/\Lambda^2) \) and \( \beta_k \) are the coefficients of the \( \beta \)-function. For our purposes, it is convenient to represent the \( \beta \)-function in the form

\[
\beta(\alpha_s) = -\beta_0 \alpha_s^2 (1 + b_1 \alpha_s + b_2 \alpha_s^2 + b_3 \alpha_s^3 + \ldots), \tag{5}
\]

with \( b_1 = \beta_1/\beta_0 \), the ratios of the \( \beta \)-function coefficients. For three flavors the coefficients are \( \beta_0 = 9/4\pi = 0.7162 \), \( b_1 = 0.5659 \), \( b_2 = 0.4530 \) \[16\] and \( b_3 = 0.6770 \) \[17\]. In the current analysis we use the exact solutions of the RG equation \[4\] in the \( \text{MS} \)-scheme at the scale \( \mu = Q \).

2.1. Analytic Perturbation Theory

The moments of the structure functions are analytic functions in the complex \( Q^2 \)-plane with a cut along the negative part of the real axis (see, e.g., Ref. \[13\]). The perturbative representation \[4\] violates these analytic properties due to the unphysical singularities of \( \alpha_s(Q^2) \) for \( Q^2 \to 0 \). To resolve the issue, we apply the APT method \[10, 16\], which allows one to combine the RG invariance with proper analytical properties of the RG-invariant coupling and observables. In particular, the four-loop APT expansion for the perturbative part \( \Delta_{\text{Bj}}(Q^2) \) is given by

\[
\Delta_{\text{Bj}}^{\text{APT}}(Q^2) = \sum_{k \leq 4} c_k A_k(Q^2). \tag{6}
\]

Here the coefficients \( c_k \) are the same as in Eq. \( 3 \), and the functions \( A_k(Q^2) \) are defined through the spectral functions \( g_k(\sigma) \equiv \text{Im} \left[ \alpha_s^k(-\sigma - i\epsilon) \right] \) by the spectral integral

\[
A_k(Q^2) = \frac{1}{\pi} \int_0^\infty d\sigma \frac{g_k(\sigma)}{\sigma + Q^2}. \tag{7}
\]

Note, the first function, \( A_1(Q^2) \), is the analytic coupling, \( \alpha_{\text{APT}}(Q^2) = A_1(Q^2) \). At large momentum transfers, all the functions \( A_k(Q^2) \) become proportional to the \( k \)-th power of the usual perturbative coupling \( \alpha_s(Q^2) \) and the expansion \( 3 \) reduces to the power series \( 5 \). However, at small enough \( Q \leq 1 - 2 \text{ GeV} \) the properties of the non-power expansion \( 6 \) become considerably different from the PT power series \( 3 \) (see, e.g., Ref. \[14\] for details).

2.2. The \( Q^2 \)-dependence

Now we analyze the \( Q^2 \)-dependence of the BSR in the framework of both the PT and APT approaches in different orders (NLO, N^2LO and N^3LO) of the perturbative expansions \[3 \] and \[4 \] respectively. As a normalization point, we use the most accurate \( \alpha_s \)-value at \( Q = M_Z \), \( \alpha_s(M_Z) = 0.1184 \pm 0.0007 \) \[4, 10\]. In order to take into account flavor thresholds, we apply the matching conditions for the values of \( \alpha_s(Q^2) \) which are rather nontrivial in higher PT orders (see Refs. \[24, 21, 22\]). Following to analysis in Ref. \[23\], our matched calculation for the four-loop \( \text{MS} \)-coupling gives \( \Lambda^{(n_f=3)} = 336 \pm 10 \text{ MeV} \). Note, we obtain practically the same results, but with larger errors, if we choose the pseudo-observable value \( R(M_Z^2) = 1.03904 \pm 0.00087 \) as a normalization point \[24\] which leads to the four-loop running coupling equal to \( \alpha_s(M_Z) = 0.1190 \pm 0.0026 \).

In Fig. \[1\] we illustrate the behavior of the perturbative part of the BSR in different orders in \( \alpha_s \) in both PT and APT approaches. The PT curves in different orders (NLO, N^2LO and N^3LO) practically (at about 1 % accuracy) coincide with each other, so we represent the APT result by a single dash-dotted line in Fig. \[1\]. For completeness, we also show here the combined SLAC and JLab data on \( \Gamma_1^{p-n}(Q^2) \) used in our analysis. The SLAC data points \[25\] are denoted by squares, the JLab CLAS Hall A 2002 data – by downward pointing triangles, the JLab CLAS Hall B 2003 data – by diamonds \[2\], and the most recent JLab data \[1\] – by circles. The horizontal dotted line represents the limiting value \( \Gamma_1^{p-n}(Q^2 \to \infty) = g_A/6 \).

One can see that at \( Q^2 \geq 0.7 \text{ GeV}^2 \) the four-loop approximation describes the data quite well. Moreover, the corresponding curve passes close to the central values of several data points, although the experimental accuracy

\[\text{This correction is defined by the coefficient function, } \Delta_{\text{Bj}} = 1 - C_{\text{Bj}}(\alpha_s).\]
(which is of the same order as both the three- and four-loop contributions) does not allow one to make a definite choice between four- and three-loop approximations.

At the same time, at \( Q^2 \lesssim 0.7 \text{ GeV}^2 \) the four-loop approximation describes the data equally bad as the three- and two-loop ones. This is a signal of the necessity to account for HT contributions, and it will be strongly dependent on the order of PT used for its extraction \( \delta \).

This changes when APT is applied and the higher-loop stability is achieved. This is a well-known feature of APT free from unphysical singularities. At the same time, the deviation of APT curve from the data shows for necessity of the HT contribution which in this case is quite stable \( \delta \).

This situation may be considered as an indication of the transition of PT series to the asymptotic regime (while APT series remains convergent) for \( Q^2 \sim 0.7 \text{ GeV}^2 \). Let us explore this possibility in more detail.

2.3. Convergence of the PT and APT expansions

Clearly, at low \( Q^2 \) a value of the strong coupling is quite large, questioning the convergence of perturbative QCD series. The PT power series truncated after four-loop order (c.f. Eq. (8)) reads

\[
\Delta_{\text{PT}}^{\text{Bj}}(\alpha_s) = 0.3183 \alpha_s + 0.3631 \alpha_s^2 \\
+ 0.6520 \alpha_s^3 + 1.804 \alpha_s^4 = \sum_{i \leq 4} \delta_i(\alpha_s),
\]

where \( \delta_i \) is the \( i \)-th term. The qualitative resemblance of the coefficients pattern to the factorial growth did not escape our attention although the more definite statements, if possible, would require much more efforts. This observation allows one to estimate the value of \( \alpha_s \sim 1/3 \) providing a similar magnitude of three- and four-loop contributions to the BSR.

\[\text{Figure 1: Perturbative part of the BSR as a function of the momentum transfer squared } Q^2 \text{ in different orders in both the APT and standard PT approaches against the combined set of the Jefferson Lab [1, 2] and SLAC [25] data.}\]

\[\text{Figure 2: The } Q^2\text{-dependence of the relative contributions at the four-loop level in the PT approach. Four-loop PT order overshoots the three-loop one at } Q^2 \lesssim 2 \text{ GeV}^2 \text{, so it does not improve the accuracy of the PT series compared to the three-loop one.}\]

\[\text{Figure 3: The } Q^2\text{-dependence of the relative contributions of separate terms in the four-loop expansion (8) } N_i(Q^2) = \delta_i(Q^2)/\Delta_{\text{Bj}}(Q^2).\]

\[\text{To test that, we present in Fig. 2 the relative contributions of separate terms in the four-loop expansion (8) } \]

\[\text{As it is seen from Fig. 2 in the region } Q^2 < 1 \text{ GeV}^2 \text{ the dominant contribution to the pQCD correction } \Delta_{\text{Bj}}(Q^2) \text{ comes from the four-loop term } \sim \alpha_s^4 \text{. Moreover, its relative contribution increases with decreasing } Q^2. \text{ In the region } Q^2 > 2 \text{ GeV}^2 \text{ the situation changes – the major contribution comes from one- and two-loop orders there.}\]
Analogous curves for the APT series given by Eq. (6) are presented in Fig. 3.

Figures 2 and 3 demonstrate the essential difference between the PT and APT cases, namely, the APT expansion obeys much better convergence than the PT one. In the PT case, the higher order contributions are stable at all $Q^2$ values, and the one-loop contribution gives about 70%, two-loop – 20%, three-loop – not exceeds 5%, and four-loop – up to 1%.

One can see that the four-loop PT correction becomes equal to the three-loop one at $Q^2 = 2$ GeV$^2$ and noticeably overshooting it (note that the slopes of these contributions are quite close in the relatively wide $Q^2$ region) for $Q^2 \sim 1$ GeV$^2$ which may be considered as an extra argument supporting an asymptotic character of the PT series in this region.

In the PT case, the contribution of the higher loop corrections is not so large as in the PT one. The four-loop order in APT can be important, in principle, if the theoretical accuracy to better than 1% will be required.

### 2.4. The $\mu$-scale dependence

As it is known, any observable obtained to all orders in pQCD expansion should be independent of the renormalisation scale $\mu$, but in any truncated-order perturbative series the cancelation is not perfect, such that the pQCD predictions depend on the choice of the $\mu$-scale (for a review see, e.g., Ref. [19]).

In order to estimate this dependence of $\Gamma^{n-\mu}_I$ on the unphysical renormalization-scale parameter $\mu$, we use the four-loop expression for the coefficient function $C_{B_0}(\mu^2/Q^2)$ recently published in Ref. [6]. One commonly introduce the dimensionless parameter $x_\mu (\mu^2 = x_\mu Q^2)$, which we have chosen to change within the interval $x_\mu = 0.5 \div 2$ (see, for example, the analysis in Ref. [24]), and compare the $\mu$-scale ambiguities between the three- and four-loop PT series.

In Fig. 4 the perturbative part of the BSR is plotted as a function of $Q^2$ in three- and four-loop orders of PT series corresponding to $x_\mu$ in the interval $0.5 \div 2$. The width of the arising strip for the four-loop approximation is similar to the one for the three-loop approximation in the highest JLab region $Q^2 \sim 3$ GeV$^2$ [4], so these approximations provide the description of the data with comparable accuracy, as discussed above. Thus, the four-loop result does not improve the data description noticeably in the low-energy domain.

At the same time, for $Q^2 \leq 1$ GeV$^2$, where PT does not allow the description of the data, the inclusion of the four-loop contribution leads to a stronger $\mu$-dependence.

These observations provide yet other arguments supporting the mentioned transition to asymptotic PT series at $Q^2 \sim 1$ GeV$^2$.

---

2One can find that an account for four-loop contribution leads to a decrease of the $\mu$-dependence if $Q^2 \geq 5$ GeV$^2$ which is currently outside the JLab kinematical range, but will be accessible by JLab after the scheduled upgrade.
In the above analysis, we normalized $\alpha_s$ at the Z-boson mass scale and then fixed the value of the $\Lambda$ parameter separately in each order in $\alpha_s$ approximation (it was sufficient for understanding the role of the fourth order in the PT/APT perturbative series). However, the corresponding values of the $\Lambda$ parameter extracted in this way may be different from ones obtained in the direct QCD analysis of the experimental data on the moments of the structure functions (see, e.g., Ref. [26]). Having this in mind, we investigate additionally the sensitivity of the extracted values of the higher twist term $\mu_4$ to the QCD scale parameter $\Lambda$ in various orders of PT. In the framework of APT, the sensitivity of $\mu_4$ to the $\Lambda$ parameter is weak, and it does not depend on the order of the loop expansion. Correspondingly, the values of the higher twist coefficients turn out to be considerably more precise than those extracted in the PT approach (see also Table 1).

Table 1: Results of higher twist extraction from the JLab data on BSR in various (NLO, $N^2$LO, $N^3$LO) orders of PT and all orders of APT.

| Method | $Q^2_{min}$ | $\mu_4/M^2$ | $\mu_6/M^4$ | $\mu_8/M^6$ |
|--------|-------------|-------------|-------------|-------------|
| PT NLO | 0.5         | -0.028(5)   | -           | -           |
| PT $N^2$LO | 0.66      | -0.014(7)   | -           | -           |
| PT $N^3$LO | 0.71      | 0.006(9)    | -           | -           |
| APT     | 0.47        | -0.050(4)   | -           | -           |

In Fig. 7 we show values of the coefficient $\mu_4$ extracted from the JLab data using the PT at different orders at $Q^2_{min} = 0.66$ GeV$^2$ with error bands. Vertical lines denote the corresponding uncertainty ranges in $\Lambda$-parameter. The ranges corresponding to $N^2$LO and $N^3$LO approximations have similar sizes and overlap with each other, so the four-loop result does not improve the stability w.r.t. $\Lambda$ variations compared to the three-loop one.

In Fig. 7 we show values of the coefficient $\mu_4$ extracted from the JLab data using two-, three- and four-loop PT at $Q^2_{min} = 0.66$ GeV$^2$ vs the parameter $\Lambda$. One can see that the PT does not lead to a stable result for extracted $\mu_4$ value with respect to $\Lambda$ variations. The extracted higher twist coefficient $\mu_4$ changes quite strongly between different orders of the PT expansion. And it happens in both in absolute value and sign, namely, at $\Lambda > 320$ MeV the higher twist coefficient becomes positive in the four-loop PT order. This sensitivity of the higher twist term $\mu_4$ to variations of the $\Lambda$ becomes stronger at higher PT orders.

On the other hand, these data tell us that the absolute value of $\mu_4$ decreases with the order of PT and just
at four-loop order becomes compatible to zero. This may be considered as a manifestation of duality between higher orders of PT and HT (see Ref. [4] and references therein). Moreover, when PT series manifests the asymptotic behavior (i.e. becomes most close to exact result), the HT (which may be considered as a contribution completing the PT series) can be reduced to zero.

4. Summary and Conclusion

In this work, we performed the QCD analysis of the precise low energy JLab data on the NLO PT order and extracted the OPE higher twist terms using the four-loop expression for the QCD correction to the Bjorken integral \( \Delta_{\text{Bj}} \) published recently in Ref. [6].

Our main observations are:

i) The four-loop approximation provides good description of the data for the highest JLab \( Q^2 \sim 3 \text{ GeV}^2 \). For several data points there is an impression that the four-loop approximation is better than the three-loop one. At the same time, the order of magnitude of both these contributions is the same as an experimental error, so a more precise statement can hardly be made.

ii) For lower \( Q^2 \leq 0.7 \text{ GeV}^2 \) the four-loop PT contribution does not help to describe the data. Meanwhile, as it was shown earlier [4], the APT application leads to higher order stability of the HT extraction. In turn, this results in accurate data description down to \( Q^2 \sim 0.1 \text{ GeV}^2 \) always at the two-loop APT level (see Fig. [5]).

iii) The magnitude of HT decreases with an order of PT and becomes compatible to zero at the four-loop level.

Our concluding impression is that all these features may indicate that the asymptotic nature of the QCD PT series is the same as an experimental error, so a more precise statement can hardly be made.

### Acknowledgments

We are thankful to S.V. Mikhailov and K.G. Chetyrkin for valuable discussions as well as to A.V. Sidorov and D.B. Stamenov for useful comments and to V.V. Skalabuz for interest in the work and stimulating discussions.

This work was partly supported by the Russian presidential grant Scien. School–3810.2010.2, the RFBR grants 09-02-00732, 09-02-01149, 11-01-00182, the BelRFFR-JINR grant F10D-001, and by the Carl Trygger Foundation.

### References

[1] K.V. Dharmawandane et al., Phys. Lett. B 641 (2006) 11, arXiv:nucl-ex/0605028.
[2] R. Fatemi et al., Phys. Rev. Lett. 91 (2003) 222002, arXiv:nucl-ex/0306019.
[3] R.S. Pasechnik, D.V. Shirkov, O.V. Teryaev, Phys. Rev. D 78 (2008) 071902, arXiv:hep-ph/0808.0066.
[4] R.S. Pasechnik, D.V. Shirkov, O.V. Teryaev, O.P. Solovtsova, V.L. Khandramai, Phys. Rev. D 81 (2010) 016010, arXiv:hep-ph/0911.3297.
[5] R.S. Pasechnik, J. Soffer, and O.V. Teryaev, Phys. Rev. D 82 (2010) 076007, arXiv:hep-ph/1009.3355.
[6] P.A. Baikov, K.G. Chetyrkin, J.H. Kuhn, Phys. Rev. Lett. 104 (2010) 132004, arXiv:hep-ph/1001.3606.
[7] K. Nakamura et al. (Particle Data Group), J. Phys. G 37 (2010) 075021.
[8] S.A. Larin, J.A.M. Vermaseren, Phys. Lett. B 259 (1991) 345.
[9] J. Ellis, M. Karliner, Phys. Lett. B 341 (1995) 397, arXiv:hep-ph/9407287.
[10] A. Deur et al., Phys. Lett. B 650 (2007) 244; Phys. Lett. B 665 (2008) 349, arXiv:nucl-ex/0803.4119; Phys. Rev. D 78 (2008) 032001, arXiv:nucl-ex/0802.3198.
[11] D.V. Shirkov, LL. Solovtsov, JINR Rapid Comm. No.2 76-96 (1996) 5; Phys. Rev. Lett. 79 (1997) 1209, arXiv:hep-th/9704353.
[12] K.A. Milton, LL. Solovtsov, Phys. Rev. D 55 (1997) 5295, arXiv:hep-ph/9611438.
[13] D.V. Shirkov, TMP TH, 1999, 438, arXiv:hep-th/9810246.
[14] D.V. Shirkov, LL. Solovtsov, Theor. Math. Phys. 150 (2007) 132, arXiv:hep-ph/0611229.
[15] D.V. Shirkov, Eur. Phys. J. C 22 (2001) 331, arXiv:hep-ph/0107282.
[16] I.L. Solovtsov, D.V. Shirkov, Phys. Lett. B 442 (1998) 344, arXiv:hep-ph/9711251.
[17] K.A. Milton, LL. Solovtsov, O.P. Solovtsova, Phys. Rev. B 415 (1997) 104, arXiv:hep-ph/9706409; Phys. Rev. D 60 (1999) 016001, arXiv:hep-ph/9809513.
[18] S.G. Gorishny, S.A. Larin, Phys. Lett. B 172 (1986) 109.
[19] O.V. Tarasov, A.A. Vladimirov, A.Yu. Zharkov, Phys. Lett. B 93 (1980) 429.
[20] S.A. Larin and J.A.M. Vermaseren, Phys. Lett. B 303 (1993) 334, arXiv:hep-ph/9302028.
[21] T. van Ritbergen, J.A.M. Vermaseren, S.A. Larin, Phys. Lett. B 400 (1997) 379, arXiv:hep-ph/9701390.
[22] I.L. Solovtsov, Phys. Lett. B 422 (1998) 344, arXiv:hep-ph/9711251.
[23] K.A. Milton, LL. Solovtsov, Phys. Rev. D 60 (1999) 016001, arXiv:hep-ph/9809513.
[24] A. Deur et al., Phys. Lett. B 650 (2007) 244.
[25] A. Deur et al., Phys. Rev. Lett. 93 (2004) 212001, arXiv:nucl-ex/0306019.
[26] E. Leader, A. V. Sidorov, D. B. Stamenov, Phys. Rev. D 82 (2010) 114018, arXiv:hep-ph/1010.0574.