On the D-term of the nucleon generalized parton distributions

M. Wakamatsu

Department of Physics, Faculty of Science,
Osaka University,
Toyonaka, Osaka 560-0043, JAPAN

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Abstract

It is known that some of the deeply-virtual-Compton-scattering observables, for instance, the beam-charge asymmetry in the hard electroproduction of real photons on the nucleon, are extremely sensitive to the magnitude of D-term appearing in the parameterization of the generalized parton distributions. We report a theoretical analysis of both the isoscalar and isovector parts of the nucleon D-term within the framework of the chiral quark soliton model, without recourse to the derivative expansion type approximation used in previous works.

Growing attention has recently been paid to the studies of the so-called generalized parton distributions (GPDs) not only because they provide us with the unique means to experimentally access the quark orbital angular momentum in the nucleon but also because they offer the most detailed information on the underlying quark-gluon structure of the nucleon [1] - [7]. However, since the GPDs are functions of three kinematical variables and since they appear as complicated convolution integrals in the cross section formulas of the deeply-virtual Compton scatterings (DVCS), deeply-virtual meson productions (DVMP) etc., a suitable parameterization of them is practically unavoidable. The most popular parameterization of GPDs is to use the double distributions [1],[4] supplemented with the so-called D-term [8]. It turned out that some of the DVCS observables, for instance, the beam-charge asymmetry in the hard electroproductions of real photons on the nucleon, are extremely sensitive to the magnitude of the D-term necessary in the above parameterization [5],[9].

So far, there has been only a limited number of theoretical studies on the nucleon D-term. The first estimate of the D-term is based on the chiral quark soliton model (CQSM), more

Email : wakamatu@phys.sci.osaka-u.ac.jp
precisely, on the CQSM predictions for the GPD $H^{u+d}(x,\xi,t)$ given by Petrov et al. under the derivative-expansion type approximation \cite{10}. Using these predictions, the authors of \cite{5} as well as of \cite{11} estimated several Mellin moments of $H^{u+d}(x,\xi,t)$ at various values of $\xi$ and $t$ and extrapolated them to $t = 0$. The results are then fitted to the Gegenbauer expansion given as

$$\sum_{q=u,d,s,\ldots} D^q(z) = (1 - z^2) \sum_{n=1,odd}^{\infty} d_n C_{3/2}^n(z). \quad (1)$$

This led them to the following estimate for the expansion coefficients:

$$d_1 \simeq -4.0, \quad d_3 \simeq -1.2, \quad d_5 \simeq -0.4, \quad (2)$$

with higher coefficients being small. We recall here that the $D^q(z)$ in the flavor-singlet channel mixes under evolution with the corresponding gluon D-term $D^g(z)$. The numbers quoted in Eq.(2) corresponds to the values at a few GeV scale. At the model energy scale around 600 MeV, their result corresponds to (see the footnote of \cite{12})

$$d_1 \simeq -8.0. \quad (3)$$

Although used in many recent phenomenological analyses of the DVCS and the DVMP processes, such an estimate of the D-term coefficients is of highly qualitative nature. A little more direct estimate of $d_1$, the first coefficient of the Gegenbauer expansion, was made by Schweitzer et al. within the same model \cite{12}. By starting with the model expression for the unpolarized GPD $H(x,\xi,t)$, they derived a closed formula for $d_1$. They also estimate its numerical value by using the derivative-expansion type approximation to find that

$$d_1 \simeq -9.46, \quad (4)$$

at the model energy scale. After a simple estimate of the scale dependence, the author of \cite{12} got the number

$$d_1 \simeq -4.7 \quad \text{at a few GeV}^2, \quad (5)$$

which they claim is qualitatively consistent with the results of \cite{11,5}. Although of preliminary nature, there also exists a lattice QCD study of the generalized form factors $C_2^q(t)$ \cite{13}, forward limit of which are related to $d_1$.

Clearly, in spite of its phenomenological importance, we must say that our knowledge on the precise magnitude of the D-term is still rather poor and uncertain. In view of the circumstance above, we think it useful to evaluate the most important parameter of the D-term, i.e. the first coefficient of the Gegenbauer expansion of the D-term within the framework of the CQSM, without recourse to the derivative-expansion type approximation. In the present paper, we try to estimate not only the isoscalar part of $d_1$ but also its isovector part, which is subleading in
the $1/N_c$ expansion. For the sake of comparison, we also derive the theoretical expression of $d_1^{u+d}$ in the familiar MIT bag model, and see what prediction it gives.

To fix the normalization convention of the relevant quantities, it would be convenient to start our investigation with the nucleon matrix elements of the quark and gluon parts of the (symmetric) QCD energy-momentum tensor parameterized by four form factors as [14], [2]

$$
\langle p' | \hat{T}_{\mu\nu}^q(0) | p \rangle = \bar{N}(p') \left[ M_2^q(t) \frac{P_\mu P_\nu}{M_N} + J_{q,g}(t) \frac{i P_{\mu (\sigma \nu)} \Delta^o}{M_N} \right. \\
+ \left. \frac{d_{q,g}(t)}{2 M_N} \left( \Delta_\mu \Delta_\nu - g_{\mu\nu} \Delta^2 \right) + \bar{c}_{q,g}(t) g_{\mu\nu} \right] N(p),
$$

(6)

where $P = (p' + p)/2$, $\Delta = p' - p$, and $t = \Delta^2$. Here $\hat{T}_{\mu\nu}^q = \bar{q} \gamma_{(\mu} \nabla_{\nu)} q$ is the QCD energy-momentum tensor of the quark with flavor $q$, while $\hat{T}_{\mu\nu}^g = G_{a\mu}^a G_{a\nu}^a + \frac{1}{4} G^2$ is the corresponding gluon part. (The form factor $\bar{c}(t)$ accounts for nonconservation of the separate quark and gluon parts of the energy-momentum tensor. They must satisfy the constraint $\sum_q \bar{c}_{q}(t) + c_{g}(t) = 0$ due to the conservation of the total (quark plus gluon) energy-momentum tensor.) The form factors in Eq. (6) are related to the 2nd Mellin moments of the familiar unpolarized GPDs $H^q(x, \xi, t)$ and $E(x, \xi, t)$ as

$$
\sum_q \int_{-1}^1 x H^q(x, \xi, t) \, dx = M_2^Q(t) + \frac{4}{5} d^Q(t) \xi^2,
$$

(7)

$$
\sum_q \int_{-1}^1 x E^q(x, \xi, t) \, dx = 2 J^Q(t) - M_2^Q(t) - \frac{4}{5} d^Q(t) \xi^2.
$$

(8)

Here, the suffix $Q$ denotes the summation over all quark flavors, for example, $J^Q(t) \equiv \sum_q = u, d, s, \ldots J^q(t)$. (Practically, we confine here to the light-quark components of two flavors, which means that $Q = u + d$.) The sum of the above two equations with $t = 0$ gives the famous Ji’s angular momentum sum rule [2, 3] :

$$
\sum_q \int_{-1}^1 x \left( H^q(x, \xi, 0) + E^q(x, \xi, 0) \right) \, dx = 2 J^Q,
$$

(9)

with $J^Q$ being the total angular momentum carried by the quark fields in the nucleon.

The interest of our present study is the forward limit of $d^Q(t)$, which just corresponds to the first coefficients in the Gegenbauer expansion of the so-called D-term, i.e.

$$
d_1^{u+d} \equiv d^Q(0).
$$

(10)

According to Polyakov [14], the constants $d_1^{u+d}$ can be expressed as

$$
d_1^{u+d} = - \frac{M_n}{2} \int d^3 r \, T^Q_{ij}(r) \left( r^i r^j - \frac{1}{3} \delta^{ij} r^2 \right),
$$

(11)
Here $T^Q_{\mu\nu}(\mathbf{r})$ is the net quark contribution to the static energy-momentum tensor density of the nucleon defined by

$$T^Q_{\mu\nu}(\mathbf{r}) = \frac{1}{2M_N} \int \frac{d^3\Delta}{(2\pi)^3} e^{-i\mathbf{\Delta}\cdot\mathbf{r}} \langle p', S | \hat{T}^Q_{\mu\nu}(0) | p, S \rangle,$$

with $\hat{T}^Q_{\mu\nu}(0)$ the quark part of the QCD energy-momentum tensor operator. As expected, the D-term is seen to contain valuable information on the distribution of energy-momentum tensor inside the nucleon.

As a warm-up, let us first evaluate $d_1$ in a simple model of baryons, i.e. the MIT bag model. We must calculate

$$d_1^{u+d}(\text{MIT}) = -\frac{M_N}{2} \langle \Psi_{gs} | r^2 [\alpha \cdot \hat{\mathbf{r}} \cdot \hat{\mathbf{p}} \cdot \hat{\mathbf{r}} - \frac{1}{3} \alpha \cdot \hat{\mathbf{p}} ] | \Psi_{gs} \rangle,$$

with $\alpha^i = \gamma^0\gamma^i$ the standard Dirac matrices. Here

$$\Psi_{gs}(\mathbf{r}) = N \left( \begin{array}{c} j_0(kr) | (l = 0) j = 1/2, m \rangle \\ -i j_1(kr) | (l = 1) j = 1/2, m \rangle \end{array} \right),$$

with $N^{-2} = 2R^2(kR-1)j_0(kR)$, $\omega_0 \equiv kR \approx 2.043$

is the ground state wave function with $R$ being the bag radius. After some manipulation, we easily find that

$$d_1^{u+d}(\text{MIT}) = -M_N N^2 k \left\{ \int_0^R j_1(kr) r^4 j_1(kr) \, dr - \int_0^R j_0(kr) r^4 j_2(kr) \, dr \right\}.$$

For an order of magnitude estimate for $d_1^{u+d}(\text{MIT})$, here we use the bag model parameter adopted by Jaffe and Ji [15], i.e. $M_N R \approx 4.0 \omega_0$. Eq.(16) then gives

$$d_1^{u+d}(\text{MIT}) \approx -0.716,$$

which turns out to be many times smaller in magnitude than the previous estimates based on the CQSM, $d_1^{u+d} \approx -(8.0 \sim 9.5)$.

Next, we derive the theoretical expression for $d_1$ within the CQSM. We first note that the energy-momentum tensor in the CQSM formally takes the same form as that of QCD, since the effective pion degrees of freedom contained in it is not an independent fields of quarks. The leading order contribution to the isoscalar $d_1$ comes from the zeroth order term in the collective rotational velocity of the soliton [16],[17], so that it is given in the following form :

$$d_1^{u+d} = -\frac{M_N}{2} N_c \sum_{n \leq 0} \langle n | r^2 [\alpha \cdot \hat{\mathbf{r}} \cdot \hat{\mathbf{p}} \cdot \hat{\mathbf{r}} - \frac{1}{3} \alpha \cdot \hat{\mathbf{p}} ] | n \rangle.$$
The symbol $\sum_{n\leq 0}$ denotes the summation over the occupied states (i.e., the discrete valence level ($n = 0$) plus the negative-energy Dirac-sea orbitals ($n < 0$)) in the hedgehog mean field.

In the above equation, $\vec{p} = \frac{1}{2}(\vec{p} + \vec{p})$ with $\vec{p}$ and $\vec{p}$ being the momentum operators respectively acting on the initial and final state wave functions. After some manipulation, $d_1^{u+d}$ can be transformed into a form, which is convenient for numerical calculation:

$$d_1^{u+d} = -\frac{\sqrt{4\pi}}{\sqrt{6}} M_N N_c \sum_{n\leq 0} \langle n | r^2 \left[ [Y_2(\hat{r}) \times \frac{\hat{p}}{p}]^{(1)} \times \alpha \right]^{(0)} | n \rangle. \quad (19)$$

It is an easy exercise that this precisely coincides with the following expression

$$d_1^{u+d} = -\frac{5}{4} N_c M_N \sum_{n\leq 0} \langle n | \gamma^0 \gamma^3 \left\{ \frac{\hat{x}}{x} P_2(\cos \theta) \right\} | n \rangle, \quad (20)$$

derived by Schweitzer et al. in the same model by starting with the expression for the nucleon unpolarized GPD $H^{u+d}(x, \xi, t)$. First pointing out that the valence quark contribution to $d_1^{u+d}$ vanishes identically, i.e. $d_1^{u+d}(\text{val}) = 0$, they estimated the contribution of the polarized Dirac sea by means of a kind of derivative-expansion type approximation. They thus find that $d_1^{u+d} = d_1^{u+d}(\text{sea}) \simeq -9.46$ at the model energy scale around 600 MeV.

However, we find no reason why $d_1^{u+d}(\text{val})$ vanishes. In fact, the relevant operator appearing in Eq. (19) is a positive parity operator with the total grand spin $K$ being zero (here $K = J + \frac{1}{2} \tau$), and there is no selection rule which enforces its matrix element between the valence quark state with the quantum numbers $K^P = 0^+$ to vanish. (Although not so obvious, the operator appearing in Eq. (20) also contains the $K = 0$ component.) Here, we shall calculate this discrete valence level contribution explicitly. We also try to evaluate the contribution of the deformed Dirac sea, without recourse to the derivative expansion type approximation. This is possible with use of the discretized momentum basis of Kahana, Ripka and Soni [18], [19]. In the present analysis, we use the self-consistent soliton solutions obtained in [21] within the double-subtraction Pauli-Villars regularization scheme [20]. The model in the chiral limit contains only one parameter $M$, i.e. the dynamical quark mass. Here, we use the value $M = 400$ MeV. Since the energy-momentum-tensor distribution carried by the quark fields is expected to be sensitive to the value of the pion mass, we shall also investigate the pion mass dependence of $d_1^{u+d}$. The effective model lagrangian, which incorporates the finite pion mass effects, is given in [20]. Self-consistent soliton solutions are prepared in [21] for several values of pion mass. As pointed out in that paper, favorable physical predictions of the model are obtained by using the value of $M = 400$ MeV and $m_\pi = 100$ MeV, since this set gives a self-consistent solution close to the phenomenologically successful one obtained with $M = 375$ MeV and $m_\pi = 0$ MeV in the single-subtraction Pauli-Villars regularization scheme in the previous studies of nucleon parton distribution functions [22] - [28]. We first show the theoretical prediction for $d_1^{u+d}$ corresponding to this favorable parameter set. It gives

$$d_1^{u+d} = d_1^{u+d}(\text{val}) + d_1^{u+d}(\text{sea}), \quad (21)$$
with
\[ d_{1}^{u+d}(\text{val}) \simeq 0.66, \quad d_{1}^{u+d}(\text{sea}) \simeq -5.51. \tag{22} \]

We confirm that the dominant contribution to \( d_{1}^{u+d} \) comes from the quarks in the negative-energy Dirac-sea orbitals. This deformed Dirac-sea contribution is large and negative. However, we also find that the quarks in the discrete valence level gives nonzero and positive contribution, which partially cancels the Dirac-sea contribution. As pointed out in [5], the large and negative prediction for the D-term coefficient is a characteristic feature of the CQSM, which maximally incorporates the spontaneous breaking of the chiral symmetry. In fact, the dominant contribution from the polarized Dirac sea is also viewed as simulating the \( t \)-channel exchange of two pions with the quantum numbers \( J^{PG} = 0^{++}, 2^{++}, \cdots \) [6]. The final model predictions for \( d_{1}^{u+d} \) obtained as the sum of the valence and the Dirac-sea contribution is
\[ d_{1}^{u+d} \simeq -4.85, \tag{23} \]
which is a little smaller than the prediction \( d_{1}^{u+d} \simeq -8.0 \) obtained from the numerically evaluated \( H_{u+d}(x, \xi, t) \) [11,10] and the prediction \( d_{1}^{u+d} \simeq -9.46 \) obtained based on the derivative-expansion type approximation with neglect of the valence level contribution in the same model [12]. In view of the difference of the soliton profile functions used in all these analyses (note that our result corresponds to \( d_{1}^{u+d} \simeq -6.2 \) in the chiral limit), the qualitative agreement is encouraging, and it confirms the unique feature of the CQSM, which takes account not only of three valence quarks but also infinitely many Dirac-sea quarks in the mean potential. This is clear, if one compares the above predictions of the CQSM with that of the MIT bag model, \( d_{1}^{u+d}(\text{MIT}) \simeq -0.716 \). We also recall that the lattice QCD simulation performed by the QCDSF collaboration in the heavy pion region around \( m_{\pi} \sim 800 \text{MeV} \) [13] gives fairly small number: \( d_{1}^{u+d}(\text{QCDSF}) = \frac{5}{4} C_{2}^{u+d}(0) \simeq -0.25 \pm 0.13 \). Although the lattice QCD prediction quoted here corresponds to the energy scale around \( Q^{2} \simeq 4 \text{(GeV)}^{2} \), there seems to be more difference in magnitude than explained by its scale dependence.

The pion cloud interpretation of the Dirac-sea contribution to \( d_{1}^{u+d} \) in the CQSM may also be confirmed by investigating its pion mass dependence. We show in Fig.1 the CQSM prediction for \( d_{1}^{u+d} \) in dependence of the pion mass. One sees that the magnitude of \( d_{1}^{u+d} \) rapidly decreases as \( m_{\pi} \) increases, showing a tendency to match the very small lattice QCD prediction obtained in the heavy pion region at least qualitatively. Since the valence quark contribution to \( d_{1}^{u+d} \) is less sensitive to the variation of the pion mass, it can be interpreted as the reduction of the pion cloud effects as the model parameter \( m_{\pi} \) increases.

Now, we turn to the discussion of the isovector \( d_{1} \). It was emphasized in [5] that the isovector D-term is suppressed relative to the isoscalar one by a factor of \( 1/N_{c} \), so that it is negligible in the large \( N_{c} \) limit. Since \( N_{c} = 3 \) in reality, however, it is not self-evident whether it is in fact numerically small or not. Here, we try to evaluate this subleading term in the
1/$N_c$ expansion in an explicit manner. In conformity with the observation above, $d_{1}^{u-d}$ in the CQSM survives only at the 1st order in the collective rotational velocity of the soliton, so that the answer is given in the following form, i.e. as a double sum over the single quark orbitals in the hedgehog mean field:

$$
d_{1}^{u-d} = -\frac{\sqrt{4\pi}}{\sqrt{6}} M_N \frac{N_c}{6} \sum_{m>0,n\leq 0} \frac{1}{E_m - E_n} \langle m \parallel \tau \parallel n \rangle \\
\times \langle m \parallel r^2 \left[ [Y_2(\hat{r}) \times \hat{P}^{(1)}] \times \alpha \right]^{(0)} \tau \parallel n \rangle.
$$

(24)

Let us first show the prediction obtained with the favorable set of the model parameters, i.e. $M = 400$ MeV and $m_\pi = 100$ MeV. This gives

$$
d_{1}^{u-d} = d_{1}^{u-d}(\text{val}) + d_{1}^{u-d}(\text{sea}) \simeq 0.33 + 0.01 \simeq 0.34.
$$

(25)

We find that, in contrast to the isoscalar case, the contribution of the discrete valence level is dominant, while that of the deformed Dirac-sea is almost negligible. We also confirm that net $d_{1}^{u-d}$ is much smaller in magnitude than $d_{1}^{u+d}$, in conformity with the expectation based on the large $N_c$ counting. Still, our explicit calculation shows that the difference of $d_{1}^{u}$ and $d_{1}^{d}$ takes nonzero positive value owing to the presence of the discrete valence level contribution. Also interesting is how the isovector $d_{1}$ depends on the pion mass. We show in Fig.2 the CQSM prediction for $d_{1}^{u-d}$ in dependence of the pion mass. Contrary to the isoscalar case, the
magnitude of $d_{1}^{u-d}$ is an increasing function of $m_{\pi}$. This peculiar behavior of $d_{1}^{u-d}$ resembles the pion mass dependence of the generalized form factor $B_{20}^{u-d}(t)$ at $t = 0$, or equivalently the isovector gravito anomalous magnetic moment of the nucleon, so that it may have a similar origin. (See Fig.5 of [21], and the explanation around there.)

In summary, we have carried out a theoretical analysis of the most important constants that characterize the nucleon D-term, i.e. the first coefficients $d_{1}$ of its Gegenbauer expansion, within the framework of the CQSM, without recourse to the derivative-expansion type approximation used in the previous studies. We gave predictions not only for the leading isoscalar part of $d_{1}$ but also for the subleading isovector part of $d_{1}$ in the $1/N_{c}$ expansion. Our treatment makes it possible to estimate the contribution of the discrete valence level and that of the negative-energy Dirac-sea levels separately. We found that, as for the isoscalar $d_{1}^{a+d}$, the contribution of the deformed Dirac-sea is large and negative and dominates over the small positive contribution from the discrete valence level. This reconfirms the unique feature of the CQSM, which takes good account of the effects of the pion cloud generated by the spontaneous chiral symmetry breaking of the QCD vacuum. The predicted value of $d_{1}^{a+d} \simeq - (4.9 \sim 6.2)$ at the model energy scale around 600 MeV qualitatively supports the previous estimate $d_{1}^{a+d} \simeq - (8.0 \sim 9.6)$ obtained in the same model based on the derivative-expansion type approximation. We have also found that $d_{1}^{u-d} \simeq 0.34$, which we confirm is much smaller than $d_{1}^{a+d}$ but cannot be completely neglected. We hope that the theoretical
analysis carried out here will give useful constraints on the future analyses of high-energy DVCS and DVMP processes using the double distribution parameterization of the nucleon GPDs supplemented with the D-term.

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Note added

After submission of the present paper, two papers [29] and [30] appeared where the constant $d_{1}^{u+d}$ (in addition to some other nucleon form factors of the energy momentum tensor) was calculated in a different manner within the same model with different regularization scheme. In the footnote 4 of [29], it was pointed out that the statement $d_{1}^{u+d}(val) = 0$ made in [12] is incorrect, which is consistent with our observation in the present paper.

References

[1] D. Mueller, D. Robaschik, B. Geyer, F.M. Dittes and J. Horejsi, Fortsch. Phys. 42 (1994) 101.
[2] X. Ji, Phys. Rev. Lett. 78 (1997) 610 ; X. Ji, Phys. Rev. D55 (1997) 7114.
[3] X. Ji, J. Phys. G24 (1998) 1181.
[4] A.V. Radyushkin, hep-ph/0101225.
[5] K. Goeke, M.V. Polyakov and M. Vanderhaeghen, Prog. Part. Nucl. Phys. 47 (2001) 401.
[6] M. Diehl, Phys. Rep. 388 (2003) 41.
[7] A.V. Belitsky and A.V. Radyushkin, Phys. Rep. 418 (2005) 1.
[8] M.V. Polyakov and C. Weiss, Phys. Rev. D60 (1999) 114017.
[9] HERMES Collaboration : F. Ellinghaus, Nucl. Phys. A711 (2002) 171c.
[10] V.Yu. Petrov, P. Pobylitsa, M.V. Polyakov, I. B"ornig, K. Goeke and C. Weiss, Phys. Rev. D57 (1998) 4325.
[11] N. Kivel, M.V. Polyakov and M. Vanderhaeghen, Phys. Rev. D63 (2001) 114014.
[12] P. Schweitzer, S. Boffi and M. Radici, Phys. Rev. D66 (2002) 114004.

[13] QCDSF Collaboration: M. Göckeler, R. Horsley, D. Pleiter, P.E.L. Rakow, A. Schäfer, G. Schierholz, and W. Schroers, Phys. Rev. Lett. 92 (2004) 042002.

[14] M.V. Polyakov, Phys. Lett. B555 (2003) 57.

[15] R.L. Jaffe and X. Ji, Nucl. Phys. B375 (1992) 527.

[16] D.I. Diakonov, V.Yu. Petrov, and P.V. Pobylitsa, Nucl. Phys. B306 (1988) 809.

[17] M. Wakamatsu and H. Yoshiki, Nucl. Phys. A524 (1991) 561.

[18] S. Kahana and G. Ripka, Nucl. Phys. A429 (1984) 462.

[19] S. Kahana, G. Ripka and V. Soni, Nucl. Phys. A415 (1984) 351.

[20] T. Kubota, M. Wakamatsu and T. Watabe, Phys. Rev. D60 (1999) 014016.

[21] M. Wakamatsu and Y. Nakakoji, Phys. Rev. D74 (2006) 054006.

[22] D.I. Diakonov, V.Yu. Petrov, P.V. Pobylitsa, M.V. Polyakov and C. Weiss, Nucl. Phys. B480 (1996) 341.

[23] D.I. Diakonov, V.Yu. Petrov, P.V. Pobylitsa, M.V. Polyakov and C. Weiss, Phys. Rev. D56 (1997) 4069.

[24] H. Weigel, L. Gamberg and H. Reinhardt, Mod. Phys. Lett. A11 (1996) 3021.

[25] L. Gamberg, H. Reinhardt and H. Reinhardt, Phys. Rev. D58 (1998) 054014.

[26] M. Wakamatsu and T. Kubota, Phys. Rev. D60 (1999) 034020.

[27] M. Wakamatsu, Phys. Rev. D67 (2003) 034005 ; Phys. Rev. D67 (2003) 034006.

[28] M. Wakamatsu, Phys. Lett. B646 (2007) 24.

[29] K. Goeke, J. Grabis, J. Ossmann, M.V. Polyakov, P. Schweitzer, A. Silva, and D. Urbano, hep-ph/0702030.

[30] K. Goeke, J. Grabis, J. Ossmann, M.V. Polyakov, P. Schweitzer, A. Silva, and D. Urbano, hep-ph/0702031.