Azimuthal asymmetry in SIDIS off nuclei as probe for $\hat{q}$

Jian-Hua Gao

School of Space Science and Physics,
Shandong University at Weihai, Weihai 264209, China and
Department of Modern Physics, University of Science
and Technology of China, Hefei, Anhui 230026, China

Andreas Schäfer

Institut für Theoretische Physik, Universität Regensburg, Regensburg, Germany

Jian Zhou

Institut für Theoretische Physik, Universität Regensburg, Regensburg, Germany and
Department of Physics, Barton Hall,
Temple University, Philadelphia, PA 19122, USA

Abstract

The energy loss parameter $\hat{q}$ is one of the fundamental transport parameters of hadronic matter. Using the twist-4 collinear approach, we show that the $\cos \phi$ azimuthal asymmetry in unpolarized semi-inclusive deeply inelastic scattering (SIDIS) off a large nucleus at intermediate transverse momentum is a sensitive observable for its determination. The effect is due to the suppression of the azimuthal asymmetry by final-state multiple scattering.
I. INTRODUCTION

The detailed understanding of the properties of hot and cold nuclear matter, as described by its transport parameters is one of the topical problems of QCD, in particular in connection with high energy heavy-ion experiments at RHIC and LHC. As deeply inelastic scattering (DIS) is one of the theoretically cleanest processes in QCD it is tempting to use this probe to determine some of them. We argue that this is indeed possible for the transport parameter $\hat{q}$ if one focusses on transverse momentum dependent observables which are especially sensitive to transverse momentum broadening while partons travel through hadronic matter. The transverse momentum broadening effect arises from final state multiple parton interactions, which are enhanced in nuclear matter. Quite a number of different theoretical approaches have been formulated to describe this phenomenon [1–8] and a variety of precise experimental observations will be needed to decide which of these formulations are incorrect and which are equivalent. One common parameter appearing in all of these approaches is the parton transport parameter $\hat{q}$ which controls parton energy loss or transverse momentum broadening squared per unit of propagation length [1]. Therefore, the calculation and measurement of this transport parameter is an important step toward understanding the intrinsic properties of nuclear matter, both cold and hot.

The multiple parton re-scattering in a large nucleus not only leads to energy loss and transverse momentum broadening but also to other nuclear effects, for example, the $\cos \phi$ and $\cos 2\phi$ azimuthal asymmetries of unpolarized semi-inclusive DIS cross sections. These $\cos \phi$ and $\cos 2\phi$ -modulations are sometimes referred to as Cahn effect. In Ref. [9], it was shown that the existence of intrinsic transverse parton momenta in the unpolarized distribution and fragmentation functions can generate such modulations at low transverse momentum. Later, Cahn effect based descriptions of azimuthal asymmetries were formulated in terms of twist-2 and twist-3 transverse momentum dependent parton distributions (TMDs) [10–13]. At large transverse momentum, such asymmetries result primarily from hard gluon radiation [14] which can be calculated in the framework of collinear factorization. Similar processes are also responsible for the azimuthal angle dependence of Drell-Yan dilepton production [15, 16]. In the intermediate transverse momentum region, $\Lambda_{QCD} \ll P_\perp \ll Q$, both, collinear factorization and transverse momentum dependent factorization are supposed to apply [17]. Indeed, in this special kinematic region, the match between leading power TMD factorization...
and collinear factorization has been made explicit for the cos2\(\phi\) asymmetry of the Drell-Yan lepton pair angular distribution in [18, 19]. In contrast however, it is known already since a while that this equivalence does not extend to sub-leading power TMD factorization which suffers from severe problems [12, 20] and yields results different from that calculated in collinear factorization at intermediate transverse momentum [21].

Azimuthal asymmetries in SIDIS off nuclei are affected by final state interactions and thus provide an alternative way to study properties of the nuclear medium. In fact, the nuclear dependence of the angular distribution of Drell-Yan lepton pairs has been calculated both in the small x and intermediate x region, using the collinear twist-4 formalism [22] and the color glass condensate (CGC) model [23]. Nuclear dependent azimuthal asymmetries in SIDIS have also been investigated recently based on TMD factorization [24, 25]. The central ingredient of the treatment in Ref. [24] is the relation between the nucleon twist-3 TMDs and the nuclear ones. In this paper, we extend that earlier work to a kinematic region of relatively large transverse momenta where the process can be treated within perturbative QCD. To be more specific, we use the collinear twist-4 approach to calculate the nuclear dependence of the cos\(\phi\) azimuthal asymmetry in SIDIS at intermediate transverse momentum \(\Lambda_{QCD} \ll P_{J\perp} \ll Q\), where \(P_{J\perp}\) is the final state jet transverse momentum, and \(Q\) is the virtual photon momentum. We restrict ourselves to intermediate transverse momentum and the current fragmentation region because here the calculation can be significantly simplified with the help of general power counting rules valid in light cone gauge [26, 27] as we will demonstrate in the subsequent section. As result of our explicit calculation, we will show that the nuclear dependent part of the asymmetry in this specific kinematic region is directly proportional to the parton transverse momentum broadening in a nucleus.

II. AZIMUTHAL ASYMMETRY IN SIDIS OFF NUCLEUS

The parton model cross section for the unpolarized semi-inclusive DIS process \(e(l) + p/A(P) \rightarrow e(l') + J(P_J) + X\) takes the general form [12, 29],

\[
\frac{d\sigma}{dx_B dz dy d^2P_{J\perp}} = \frac{4\pi\alpha^2_{em}s}{Q^4} \left\{ (1 - y + \frac{y^2}{2})F_T + (1 - y)F_L \\
+ (2 - y)\sqrt{1 - y \cos \phi_J} F_{\cos \phi_J} + (1 - y) \cos(2\phi_J) F_{\cos 2\phi_J} \right\}
\] (1)
where we use the conventions of Ref. [12]. We define \( q = l - l' \) as the virtual photon momentum and its virtuality as \( Q^2 = -q^2 \), while \( x_B = Q^2 / 2P \cdot q \), \( z = P \cdot P_J / P \cdot q \) and \( y = P \cdot q / P \cdot l \) are the common DIS variables. The azimuthal angle between the transverse momentum of the outgoing parton \( P_{J\perp} \) and the leptonic plane is denoted by \( \phi_J \). Four structure functions \( F \), depending on \( x_B, Q^2, z \) (the fraction of the photon energy carried by the jet) and \( P_{J\perp} \), encode the QCD structure of the target and the dynamics of the partonic subprocess. It is convenient to use light-cone coordinates for which \( P^\mu = P^+ p^\mu \), \( q^\mu = -x_B p^\mu + n^\mu Q^2 / (2x_B P^+) \) with \( p = (1, 0, 0, 0) \) and \( n = (0, 1, 0, 0) \). At large \( P_{J\perp} \), the four functions \( F \) can be calculated in collinear twist-2 factorization. As stated above, we restrict ourself to the asymmetry at intermediate transverse momentum \( \Lambda_{QCD} \ll P_{J\perp} \ll Q \) in the current fragmentation region where power counting rules can be applied. In this kinematic region, the power behavior of \( F_T, F_L, F_{\cos \phi_J} \) and \( F_{\cos 2\phi_J} \) is \( 1/P_{J\perp}^2, 1/Q^2, 1/QP_{J\perp}, 1/Q^2 \) respectively. The \( \cos \phi \) asymmetry is determined by the ratio between the functions \( F_{UU}^{\cos \phi_J} \) and \( F_{UU,T} \), which read [14, 21],

\[
F_{T}^{\text{twist}-2} = \frac{1}{P_{J\perp}^2} \frac{\alpha_s C_F}{2\pi^2} \sum_\alpha x_B e_\alpha^2 \left[ 2 \ln \frac{Q^2}{P_{J\perp}^2} \delta(1-z) + \frac{1+z^2}{(1-z)+} \right] f_1^\alpha(x_B)
\]

\[
F_{\cos \phi_J}^{\text{twist}-2} = -\frac{1}{zP_{J\perp} Q} \frac{\alpha_s C_F}{2\pi^2} \sum_\alpha x_B e_\alpha^2 \left[ 2 \ln \frac{Q^2}{P_{J\perp}^2} \delta(1-z) + \frac{2z^2}{(1-z)+} \right] f_1^\alpha(x_B)
\]

where \( f_1(x) \) is the normal leading power collinear parton distribution. The index \( \alpha \) runs over flavors of quarks and antiquarks with fractional charge \( e_\alpha \). To extract the nuclear effect we are interested in, one has to go beyond the leading twist treatment and take into account twist-4 contributions.

The machinery of collinear higher twist factorization was pioneered already in the early 1980s [30], and later frequently applied in hadron spin physics [31] and nuclear physics [3, 22, 32]. The higher twist collinear approach has been well established in both the covariant and the light cone gauge. In order to better classify the contributions according to power counting rules, we carry our calculation out in the light cone gauge with retarded boundary conditions. For retarded boundary conditions, certain collinear twist-4 correlators can be directly related to the moment of corresponding TMD distributions.

Following the standard procedure, the higher twist contributions can be systematically isolated by expanding the hard part in the parton intrinsic transverse momentum and including the diagrams with transverse polarized gluon exchange between the struck parton
and the target remnant. At twist-4 level, the correlators associated with these two types of contributions are of the form \( \langle \bar{\psi} \partial_{\perp} \partial_{\perp} \psi \rangle \), \( \langle \bar{\psi} \partial_{\perp} A_{\perp} \psi \rangle \) and \( \langle \bar{\psi} A_{\perp} A_{\perp} \psi \rangle \). The general power counting rule \([26, 27]\) states that the diagrams with one additional transversely polarized gluon exchange are suppressed by one additional power of \( \Lambda_{QCD}/Q \) in the current fragmentation region where \( P_{J\perp} \ll Q \), as long as final state interactions at \( y^- = +\infty \) have been removed by imposing retarded boundary conditions \([28]\). Therefore, the possible leading powers of the corresponding hard parts convoluted with the correlators \( \langle \bar{\psi} \partial_{\perp} \partial_{\perp} \psi \rangle \), \( \langle \bar{\psi} \partial_{\perp} A_{\perp} \psi \rangle \) and \( \langle \bar{\psi} A_{\perp} A_{\perp} \psi \rangle \) are suppressed by the powers of \( \Lambda_{QCD}^2/P_{J\perp}^2 \), \( \Lambda_{QCD}^2/Q P_{J\perp} \) and \( \Lambda_{QCD}^2/Q^2 \) respectively. The explicit calculation shows that the leading power contribution from \( \langle \bar{\psi} \partial_{\perp} \partial_{\perp} \psi \rangle \) drops out in the azimuthal angle dependent cross section such that its sub-leading part \( \Lambda_{QCD}^2/Q P_{J\perp} \) generates the non-vanishing \( \cos \phi \) asymmetry. The hard part associated with the correlator \( \langle \bar{\psi} \partial_{\perp} A_{\perp} \psi \rangle \) contributes to the \( \cos \phi \) asymmetry with the same power \( \Lambda_{QCD}^2/Q P_{J\perp} \). As a result, we can eventually neglect all diagrams with two transversely polarized gluon exchanges and are left with the contributions from the correlators of the first two types.

To be more specific, the following three twist-4 collinear correlators enter the calculation,

\[
ff_1(x) = \int \frac{dy^-}{4\pi} e^{ixP^+ y^-} \langle P | \bar{\psi}(0) \gamma^+ (\gamma^5 + i\partial_{\perp})^2 (\gamma^5) \psi(y) | P \rangle d\rho \sigma \\
\varphi_1(x) = \int \frac{dy^-}{4\pi} e^{ixP^+ y^-} \langle P | \bar{\psi}(0) \gamma^+ (\gamma^5 + i\partial_{\perp})^2 (\gamma^5) \psi(y) | P \rangle d\rho \sigma \\
\tilde{\varphi}_1(x) = \int \frac{dy^-}{4\pi} e^{ixP^+ y^-} \langle P | \bar{\psi}(0) \gamma^5 (\gamma^5 + i\partial_{\perp})^2 (\gamma^5) \psi(y) | P \rangle d\rho \sigma
\]

where \( D_{\rho} = -i \partial_{\rho} + A_{\rho} \). Since the gauge is completely fixed by choosing retarded boundary conditions in the light cone gauge, all three correlators can be uniquely brought into a gauge invariant form. The nuclear dependence has been encoded in the above twist-4 matrix elements. It will become evident when we relate them to the relevant moment of the corresponding nuclear TMDs. The perturbative calculation of the hard coefficients associated with these twist-4 correlators is straightforward. The functions \( F_T \) and \( F_{\cos \phi_J} \) can thus be expressed as convolutions of the hard coefficients and the twist-4 correlators given above.

At small transverse momentum \( P_{J\perp} \ll Q \), the results take a remarkably simple form,

\[
F_T^{\text{twist-4}} = \frac{1}{P_{J\perp}^2} \frac{\alpha_s C_F}{2\pi^2} \sum_{\alpha} x_B \delta_{\alpha}^2 \left\{ \frac{2}{P_{J\perp}^2} \ln \frac{Q^2}{P_{J\perp}^2} \delta(1 - z) + \frac{1 + z^2}{(1 - z)_+} \right\} f_1^T(x_B) \\
+ \frac{z^2(1 + z^2)}{(1 - z)_+} \frac{1}{P_{J\perp}^2} \int_{-1}^{1} d^2 \varphi_1(x_B)
\]

(4)
Using the QCD equation of motion, one obtains the relation

\[- \Re \varphi_{\perp}(x) + \Im \tilde{\varphi}_{\perp}(x) = x f_{\perp}(x) \] (6)

where

\[ f_{\perp}(x) = \int \frac{dy^-}{4\pi} e^{ixP^+ y^-} \langle P|\bar{\psi}(0) (-i\partial_{\perp}) \psi(y)|P \rangle \] (7)

One should note that the calculations just discussed apply to SIDIS of both nuclear and nucleon targets. All results have the same form and differ only in that the collinear correlators are taken inside of a nucleus or a nucleon. The nuclear dependence of the azimuthal angle asymmetry is thus generated by the nuclear dependence of the collinear twist-4 correlators, which can be best seen from the following relations between the twist-4 correlators and the moments of TMDs with retarded boundary conditions,

\[ f_1(x, k_\perp) = \int d^2 k_\perp e^{ixP^+ y^-} \langle P|\bar{\psi}(0) \gamma^+ L(0, y) \gamma_5 \psi(y)|P \rangle \] (8)

\[ f_{\perp}(x, k_\perp) = \int d^2 k_\perp e^{ixP^+ y^-} \langle P|\bar{\psi}(0) \gamma_5 \gamma_\perp L(0, y) \gamma^+ \psi(y)|P \rangle \] (9)

here, \( f_1(x) \) and \( f_{\perp}(x) \) are the normal leading twist TMD quark distribution and twist-3 TMD distribution, respectively. Their matrix element definitions are given by

\[ f_1(x, k_\perp) = \int dy^- d^2 \vec{y}_\perp e^{ixP^+ y^- + ik_\perp \cdot \vec{y}_\perp} \langle P|\bar{\psi}(0) \gamma^+ L(0, y) \gamma_5 \psi(y)|P \rangle \] (10)

\[ f_{\perp}(x, k_\perp) = \int dy^- d^2 \vec{y}_\perp e^{ixP^+ y^- + ik_\perp \cdot \vec{y}_\perp} \frac{k_\perp}{k_\perp^2} \langle P|\bar{\psi}(0) \gamma_5 \gamma_\perp L(0, y) \gamma^+ \psi(y)|P \rangle \] (11)

For retarded boundary conditions, the transverse gauge links appearing in the above matrix elements become unity. The nuclear dependence of the TMD distributions \( f_1(x, k_\perp) \) and \( f_{\perp}(x, k_\perp) \) have been worked out and given by

\[ f_{1}^{A}(x, k_\perp) \approx \frac{A}{\pi \Delta_{2F}} \int d^2 \ell_\perp e^{-(k_{\perp} - \vec{\ell}_\perp)^2/\Delta_{2F}} f_{1}^{N}(x, \ell_\perp) \] (12)

\[ f_{\perp}^{A}(x, k_\perp) \approx \frac{A}{\pi \Delta_{2F}} \int d^2 \ell_\perp \frac{(k_{\perp} \cdot \vec{\ell}_\perp)}{k_{\perp}^2} e^{-(k_{\perp} - \vec{\ell}_\perp)^2/\Delta_{2F}} f_{\perp}^{N}(x, \ell_\perp) \] (13)

where \( \Delta_{2F} = \int d\xi \hat{q}(\xi) \) with the quark energy loss transport coefficient \( \hat{q} \), which controls parton energy loss in a cold nuclear medium and transverse momentum broadening squared.

\[ F_{\cos \phi_j}^{\text{twist-4}} = \frac{-1}{z P_{J} Q} \frac{\alpha_s C_F}{2\pi^2} \sum_{\alpha} x_{B} e_{\alpha}^{2} \left\{ \left[ 2 \ln \frac{Q^2}{P_{J}^2} \delta(1-z) + \frac{2z^2}{(1-z)^{+}} \right] f_{1}^{\alpha}(x_{B}) + \frac{2z^2(1+z^2)}{(1-z)^{+}} P_{J}^2 \right\} \] (5)
per unit of propagation length. The superscripts ‘A’ and ‘N’ denote the nuclear and nucleon TMDs respectively. Using the relations between nucleon TMDs and nuclear TMDs, we find,

$$\int k_{-}^{2} f_{A}^{\perp}(x, k_{-})d^{2}k_{-} - A \int k_{-}^{2} f_{N}^{\perp}(x, k_{-})d^{2}k_{-} = 0$$ \hspace{1cm} (14)

$$\int k_{-}^{2} f_{A}^{\perp}(x, k_{-})d^{2}k_{-} - A \int k_{-}^{2} f_{N}^{\perp}(x, k_{-})d^{2}k_{-} = Af_{N}^{1}(x)\Delta_{2F}$$ \hspace{1cm} (15)

where \( f_{N}^{1}(x) \) is the ordinary integrated parton distribution function of the nucleon. From the above two identities, we can conclude that the difference of the \( \cos \phi_{J} \) azimuthal asymmetries is proportional to the amount of transverse momentum broadening. More precisely, when \( z \neq 1 \), one has,

$$\langle \cos \phi_{J} \rangle_{eA} - \langle \cos \phi_{J} \rangle_{eN} = \frac{(2 - y)\sqrt{1 - y} - y}{1 - y + y^{2}/2} \left( \frac{F^{\text{twist-4}}_{\cos \phi_{J}}|_{eA}}{F^{\text{twist-4}}_{\cos \phi_{J}}|_{eN}} \right) \approx \frac{(2 - y)\sqrt{1 - y}}{1 - y + y^{2}/2} \frac{2z^{3}}{1 + z^{2}} \frac{\Delta_{2F}}{P_{J\perp}Q} .$$ \hspace{1cm} (16)

This is the main result of this paper, which is valid at intermediate transverse momentum \( \Lambda_{QCD} << P_{J\perp} << Q \). Equation (16) provides a direct handle to extract the crucial parameter \( \hat{q} \) from measurements of the azimuthal asymmetry in nuclei and nucleons.

### III. SUMMARY

In summary, we calculated the \( \cos \phi \) azimuthal asymmetry in semi-inclusive DIS off a nuclear target within the collinear twist-4 approach. At intermediate transverse momentum, the nuclear dependence of the azimuthal asymmetry is linked to the \( k_{-}^{2} \)-moment of the two relevant quark TMD distributions. The difference between nucleon TMDs and nuclear TMDs is generated by final state interactions, such that the difference in the azimuthal asymmetries is sensitive to their strength. To be more specific, the difference between the \( \cos \phi \) azimuthal asymmetries in SIDIS off nucleons and nuclei is proportional to the transverse momentum broadening in the latter. Therefore, it provides an alternative way to pin down the transport parameter \( \hat{q} \) by measuring the nuclear dependence of the asymmetry at intermediate transverse momentum.

The approach we developed in this paper can be extended to study the nuclear dependence of azimuthal asymmetries in other processes, such as direct photon production in SIDIS off nuclei and Drell-Yan lepton pair production in high energy \( pA \) scattering.
ACKNOWLEDGMENTS: J. Zhou thanks A. Metz, M. Diehl and H. Avakian for helpful discussion. This work has been supported by BMBF (OR 06RY9191).

[1] R. Baier, Y. L. Dokshitzer, A. H. Mueller, S. Peigne and D. Schiff, Nucl. Phys. B 484, 265 (1997).
[2] G. T. Bodwin, S. J. Brodsky and G. P. Lepage, Phys. Rev. D 39, 3287 (1989).
[3] M. Luo, J. -w. Qiu, G. F. Sterman, Phys. Lett. B279, 377-383 (1992); Phys. Rev. D49, 4493-4502 (1994); Phys. Rev. D50, 1951-1971 (1994).
[4] X. F. Guo, Phys. Rev. D 58, 114033 (1998).
[5] U. A. Wiedemann, Nucl. Phys. B 588, 303 (2000)
[6] R. J. Fries, Phys. Rev. D 68, 074013 (2003).
[7] A. Majumder, B. Muller, Phys. Rev. C77, 054903 (2008).
[8] F. D’Eramo, H. Liu, K. Rajagopal, Nucl. Phys. A855, 457-460 (2011).
[9] R. N. Cahn, Phys. Lett. B78, 269 (1978).
[10] P. J. Mulders, R. D. Tangerman, Nucl. Phys. B461, 197-237 (1996).
[11] D. Boer, P. J. Mulders, Phys. Rev. D57, 5780-5786 (1998).
[12] A. Bacchetta, M. Diehl, K. Goeke, A. Metz, P. J. Mulders, M. Schlegel, JHEP 0702, 093 (2007).
[13] Z. T. Liang and X. N. Wang, Phys. Rev. D 75, 094002 (2007)
[14] H. Georgi, H. D. Politzer, Phys. Rev. Lett. 40, 3 (1978).
[15] C. S. Lam, W. -K. Tung, Phys. Rev. D18, 2447 (1978). Phys. Rev. D21, 2712 (1980).
[16] J. C. Collins, Phys. Rev. Lett. 42, 291 (1979).
[17] X. Ji, J. -W. Qiu, W. Vogelsang et al., Phys. Rev. Lett. 97, 082002 (2006).
[18] D. Boer, Phys. Rev. D60, 014012 (1999).
[19] J. Zhou, F. Yuan, Z. -T. Liang, Phys. Lett. B678, 264-268 (2009).
[20] L. P. Gamberg, D. S. Hwang, A. Metz, M. Schlegel, Phys. Lett. B639, 508-512 (2006).
[21] A. Bacchetta, D. Boer, M. Diehl et al., JHEP 0808, 023 (2008).
[22] R. J. Fries, B. Müller, A. Schäfer, E. Stein, Phys. Rev. Lett. 83, 4261-4264 (1999). R. J. Fries, A. Schäfer, E. Stein, B. Müller, Nucl. Phys. B582, 537-570 (2000).
[23] F. Gelis, J. Jalilian-Marian, Phys. Rev. D76, 074015 (2007).
[24] J. -H. Gao, Z. -t. Liang, X. -N. Wang, Phys. Rev. C81, 065211 (2010).
[25] Y. -k. Song, J. -h. Gao, Z. -t. Liang, X. -N. Wang, Phys. Rev. D83, 054010 (2011).
[26] J. C. Collins and D. E. Soper, Nucl. Phys. B 193, 381 (1981) [Erratum-ibid. B 213, 545 (1983)].
[27] X. D. Ji, J. P. Ma and F. Yuan, Phys. Rev. D 71, 034005 (2005).
[28] X. D. Ji and F. Yuan, Phys. Lett. B 543, 66 (2002). A. V. Belitsky, X. Ji and F. Yuan, Nucl. Phys. B 656, 165 (2003).
[29] K. Hagiwara, K. -i. Hikasa, N. Kai, Phys. Rev. D27, 84 (1983).
[30] R. K. Ellis, W. Furmanski and R. Petronzio, Nucl. Phys. B 207, 1 (1982); Nucl. Phys. B 212, 29 (1983).
[31] A. V. Efremov and O. V. Teryaev, Phys. Lett. B 150, 383 (1985). R. L. Jaffe, X. -D. Ji, Phys. Rev. Lett. 67, 552-555 (1991); Nucl. Phys. B375, 527-560 (1992). J. w. Qiu and G. Sterman, Phys. Rev. Lett. 67, 2264 (1991); Nucl. Phys. B 378, 52 (1992); Phys. Rev. D 59, 014004 (1999). H. Eguchi, Y. Koike and K. Tanaka, Nucl. Phys. B 763, 198 (2007). J. Zhou, F. Yuan and Z. T. Liang, Phys. Rev. D 81, 054008 (2010).
[32] X. -N. Wang, X. -f. Guo, Nucl. Phys. A696, 788-832 (2001); Phys. Rev. Lett. 85, 3591-3594 (2000).
[33] Z. T. Liang, X. N. Wang and J. Zhou, Phys. Rev. D 77, 125010 (2008).