Hadron correlations from recombination

Rainer J Fries
School of Physics and Astronomy, University of Minnesota,
Minneapolis, MN 55455
E-mail: fries@physics.umn.edu

Abstract. Quark recombination is a successful model to describe the hadronization of a
decorrelated quark gluon plasma. Jet-like dihadron correlations measured at RHIC provide
a challenge for this picture. We discuss how correlations between hadrons can arise from
correlations between partons before hadronization. An enhancement of correlations through
the recombination process, similar to the enhancement of elliptic flow is found. Hot spots from
completely or partially quenched jets are a likely source of such parton correlations.

We start with a review of the quark recombination model. We discuss the scaling and
amplification of elliptic flow and the role of higher Fock states in hadrons. In the second part
a possible scenario of parton correlations is introduced and the resulting dihadron correlations
are calculated.

1. Review of Quark Recombination
In high energy collisions of hadrons or nuclei, partons are scattered out of the bound states.
Quantum chromodynamics (QCD) does not permit the existence of these colored quarks or
gluons as freely propagating states. Instead they have to be grouped into new color singlets that
eventually become hadrons. This hadronization is a non-perturbative process in QCD and its
dynamics is not fully understood.

For a certain class of processes involving a large momentum transfer, e.g. single inclusive
hadron production at high transverse momentum \( P_T \), it has been shown that hadronization can
be parameterized by universal, process-independent fragmentation functions. They describe the
probability that a hadron is created from a parton with large momentum in the vacuum [1].
Fragmentation functions have been measured and are well suited to describe hadron production
at transverse momenta \( P_T > 1 \text{ GeV}/c \) in \( p + p \) collisions at RHIC energies.

However, several puzzling observations lead to the conclusion that the fragmentation process
is not sufficient to explain hadron production in \( \text{Au} + \text{Au} \) collisions at RHIC at transverse
momenta of several GeV/c. First, baryons are too abundant. E.g. the proton/pion ratio
measured by PHENIX is about 1 between 2 and 4 GeV/c [2], while the value predicted by
leading twist fragmentation is \( \approx 0.2 \) [3]. Second, the elliptic flow measured for baryons and
mesons in the same region of \( P_T \) exhibits a pattern incompatible with fragmentation [4, 5]. In
addition, new measurements of dihadron correlations reveal clear deviations from their vacuum
values. Jet cones are visible, but significantly altered in \( \text{Au} + \text{Au} \) [6].

It is not a surprise that a hadronization picture that assumes a single parton fragmenting
in the vacuum has to fail in heavy ion collisions, where a hot and and very dense fireball of
decorrelated quarks and gluons is created. In the vacuum fragmentation starts by production of
additional partons through radiation. In a medium hadronization can directly proceed to the formation of bound states from the medium partons themselves. In the simplest approximation this recombination or coalescence involves the valence quarks of the hadrons: three quarks recombine into baryons, quark-antiquark pairs into mesons.

The yield of hadrons from a given parton system can be calculated starting from a convolution of Wigner functions [3]. For a meson with valence (anti)quarks \( a \) and \( b \) we have

\[
\frac{dN_M}{d^3P} = C_A \int \sum d\sigma \Phi_M \otimes W_{ab}.
\]

(1)

Here \( W_{ab} \) is the 2-particle Wigner function for partons \( a \), \( b \), and \( \Phi_M \) is the Wigner function of the meson. For practical purposes the parton Wigner function is usually approximated by a product of single particle phase space distributions \( W_{ab} = w_a w_b \). Several implementations of recombination applied to a quark gluon plasma have been discussed in the literature [7, 3, 8, 9], see [10] for a review.

Thermal, or at least exponential, parton spectra \( w \sim e^{-P/T} \) play a special role. In their case recombination is always more effective than fragmentation. The resulting hadron distribution is also exponential with the same slope since

\[
N_M \sim w_a w_b \sim e^{-xP/T} e^{-(1-x)P/T} = e^{-P/T}.
\]

(2)

where \( x \) gives the momentum fraction of parton \( a \). Here we assume that the hadron is very fast, so that the momentum \( P \) is much larger than the mass \( M \). In addition, for three-parton systems \( w_a w_b w_c \sim e^{-P/T} \) as well, hence baryon yields from thermal partons are naturally of the same order of magnitude as meson yields. This holds for arbitrary numbers of partons as was pointed out recently. The probability for the ejection of a cluster on \( n \) partons with large momentum \( P \) from a thermal ensemble of partons is independent of \( n \) [11]. It should be added that fragmentation eventually has to dominate over recombination at very high momentum.

Calculations for Au+Au collisions at RHIC energies assume a thermalized system of constituent quarks with a temperature \( T \) around the phase transition temperature of the quark gluon plasma and strong radial flow. Recombination is applied to this phase. To describe the high-\( P_T \) spectrum of hadrons this has to be supplemented by a pQCD calculation using fragmentation and taking into account energy loss [7, 3]. Alternatively one can use another model to describe the power-law tail of the \( P_T \) spectrum, e.g. recombination of hard partons or soft-hard recombination [8, 12].

Fig. 1 shows the \( P_T \) spectrum of \( \pi^0 \), \( p \), \( K_0^0 \) and \( \Lambda + \bar{\Lambda} \) in central Au+Au collisions obtained in [3] with a model including recombination and fragmentation. The agreement with available data is excellent for \( P_T > 2 \) GeV/c. We note that the hadron spectra exhibit an exponential shape up to about 4 GeV/c for mesons and up to about 6 GeV/c for baryons, where recombination of thermal quarks dominates. Above those values, the spectra follow a power-law and production is dominated by fragmentation. Recombination naturally leads to a ratio \( p/\pi \approx 1 \) and to values of \( R_{AA} \) that are larger for baryons than for mesons.

2. Elliptic Flow and Higher Fock States

The most striking evidence for quark recombination is provided by measurements of elliptic flow at RHIC. Let us assume the parton phase exhibits an elliptic flow \( v_2(p_T) \) just before hadronization. Recombination makes a prediction for elliptic flow of any hadron species after recombination [15, 3]:

\[
v_2(P_T) = n v_2^n(P_T/n).
\]

(3)

Here \( n \) is the number of valence quarks for the hadron. Hence recombination predicts very different behavior for mesons \( (n = 2) \) and baryons \( (n = 3) \). Different mesons should follow the same scaling law, even if the masses are very different.
Figure 1. Spectra of $\pi^0$, $p$, $K^0_s$ and $\Lambda + \bar{\Lambda}$ as a function of $P_T$ at midrapidity in central Au+Au collisions at $\sqrt{s} = 200$ GeV [3]. Dashed lines are hadrons from recombination of the thermal phase, dotted line is pQCD with energy loss, solid line is the sum of both contributions. Data are from PHENIX ($\pi^0$, $p$) [13, 2] and STAR ($K^0_s$, $\Lambda + \bar{\Lambda}$) [14].

Fig. 2 shows the measured elliptic flow $v_2$ for several hadron species in a plot with scaled axes $v_2/n$ vs $P_T/n$. All data points (with exception of the pions) fall on one universal curve. This is an impressive confirmation of the quark scaling rule and the entire recombination picture. The pions are shifted to lower $P_T$, because most pions in the detectors, even at intermediate $P_T$, are not from hadronization, but from secondary decays of hadrons. It was shown that $\rho$ recombination with subsequent decay of the $\rho$ resonance into 2 pions can largely account for the shift [16].

The scaling law for elliptic flow seems to put a sharp constraint on the number of partons recombining. But indeed this is not true as was pointed out recently [11]. Since there is no perturbative scale in a hadronizing quark gluon plasma with temperature $T_c \sim \Lambda_{QCD}$, the quarks are effective degrees of freedom similar to constituent quarks. Therefore valence quark configurations are dominating, but higher Fock states in the hadron, i.e. additional gluons or quark-antiquark pairs could still play a role in the recombination process. However, it was already pointed out above that the number of partons does not show up in hadron spectra because of the egalitarian nature of thermal recombination.

On the other hand, in the presence of more partons there seems to be no simple scaling of the meson and baryon elliptic flow. A Fock state with $k$ partons would contribute a term $C_k k v_2(P_T/k)$ to the hadron elliptic flow $v_2(P_T)$, where $C_k$ is the probability for this state determined by the hadron wave function. However, a numerical evaluation shows that even a 30% admixture of Fock states with an additional gluon ($|q\bar{q}g\rangle$ for mesons, $|qqqg\rangle$ for baryons) leads to results that are compatible with experimental data after a slight adjustment of the parton elliptic flow. Note that the elliptic flow of partons before hadronization is not a prediction.
Figure 2. Elliptic flow $v_2$ for $\pi^+, p$, $K_0^*$ and $\Lambda$ as a function of $P_T$ scaled by the number of valence quarks $n$ vs $P_T/n$. The data follows a universal curve, impressively confirming the quark scaling law predicted by recombination. Deviations for the pions are discussed in the text. Data are taken from PHENIX ($\pi^+, p$) [4] and STAR ($K_0^*$, $\Lambda$) [5].
that the strong jet quenching observed at RHIC is due to energy loss of high-\(p_T\) partons in the medium. The energy loss is estimated to be up to 14 GeV/fm for a 10 GeV parton [20]. This means that most jets except those close to the surface are completely stopped, dumping their energy and momentum into a small cell of the medium. This results in a dramatic local heating, creating a hot spot in the fireball. Moreover, the directional information of the jet is preserved. Subsequent evolution of the fireball might lead to partial diffusion of the hot spot, but the remaining correlations should be jet-like.

It has been pointed out that correlations can also emerge from soft-hard recombination, where a parton in a jet picks up soft partons from the surrounding medium. The emerging hadron is hence correlated with other hadrons in the jet [21]. For the case of dimeson production, let us consider all possible scenarios: (1) Both mesons come from fragmentation, either of the same or of two different jets. We will denote this process by F-F. (2) Both mesons emerge from recombination of soft partons. We call this SS-SS. (3) One meson recombines, the other emerges from the coalescence of a soft and a hard parton (from a jet), thus named SH-SS. (4) One meson comes from recombination while the other is from fragmentation (F-SS). (5) Both mesons can come from soft-hard recombination (SH-SH). (6) One meson fragments while the other is from soft-hard recombination (F-SH).

Here we are mostly interested in correlations coming from pure recombination (SS-SS). However, conventional jet correlations (from F-F) have to be taken into account when comparing to data. Three of the processes are schematically depicted in Fig. 3.

### 4. A Model for Correlations

Let us now discuss a simple model for the SS-SS process discussed in [19, 22]. We start with a generalization of the factorization formula for \(n\)-parton distributions given above. We want to assume that correlations are a small effect and restrict ourselves to 2-particle correlations \(C_{ij}\), so that a 4-parton Wigner function can be written as

\[
W_{1234} \approx w_1 w_2 w_3 w_4 (1 + \sum_{i<j} C_{ij}).
\]

The correlation functions \(C_{ij}\) between parton \(i\) and parton \(j\) are arbitrary, but we want to assume that they vary slowly with momentum and that they are only non-vanishing in a subvolume \(V_c\) of the fireball.

Our picture of hot spots induced by quenched jets motivates the following Gaussian ansatz in rapidity and azimuthal angle:

\[
C_{ij} = c_0 S_0 f_0 e^{-(\phi_i - \phi_j)^2/(2\phi_0^2)} e^{-(y_i - y_j)^2/(2y_0^2)}
+ c_\pi S_\pi f_\pi e^{-(\phi_i - \phi_j + \pi)^2/(2\phi_0^2)} e^{-(y_i - y_j)^2/(2y_0^2)}.
\]
Here $\phi_{0,\pi}$ and $y_{0,\pi}$ are the widths of the Gaussians in azimuth and rapidity, respectively. The two terms of the sum correspond to correlations initiated by an energetic parton ($\phi = 0$) and its recoil partner ($\phi = \pi$). $c_0$ and $c_\pi$ give the strength of the near side and far side correlations, while the functions $f_{0,\pi}(P_{T,i}, P_{T,j})$ describe the transverse momentum dependence of the correlations. The functions $S_{0,\pi}(\sigma_i, \sigma_j)$ parameterize the spatial localization of the parton correlations on the hypersurface $\Sigma$. For simplicity we assume that $S_{0,\pi} = 1$, if $\sigma_i, \sigma_j \in V_c$ and $S_{0,\pi} = 0$ otherwise.

The 2-meson yield is given by a convolution of the partonic Wigner function $W_{1234}$ with the Wigner functions $\Phi_A, \Phi_B$ of the mesons with an additional integration over the hadronization hypersurface $\Sigma$ [19]

$$E_A E_B \frac{dN_{AB}}{d^3 P_A d^3 P_B} = C_{AB} \int_\Sigma d\sigma \Phi_A \otimes W_{1234} \otimes \Phi_B.$$  \hspace{1cm} (6)

We restrict our discussion to near side correlations. We also work in a fixed $P_T$ window and will therefore neglect the transverse momentum dependence by setting $f_0 = 1$, absorbing the normalization into the constant $c_0$. We assume $c_0 \ll 1$ which permits omitting quadratic terms like $c_0^2$ or $c_0 v_2$.

Using Boltzmann distributions for the single parton distributions $w_i$ with temperature $T$, radial flow rapidity $\eta_T$, and a boost-invariant hypersurface $\Sigma$ at proper time $\tau$, Eqs. (6), (4) and (5) lead to

$$\frac{dN_{AB}}{\prod_{i=A,B} [P_{Ti} dP_{Ti} d\phi_i d\eta_i]} = \left(1 + 2 \hat{c}_0 + 4 \hat{c}_0 e^{-(\Delta \phi)^2 / (2 \phi_0)^2}\right) \times \prod_{i=A,B} \left[h_i(P_{Ti}) (1 + 2 v_2 (P_{Ti}) \cos(2 \phi_i))\right].$$  \hspace{1cm} (7)

$v_{2i}$ is the elliptic flow coefficient for hadron $i$ and $\Delta \phi = |\phi_A - \phi_B|$ is the relative azimuthal angle. The factor $Q = 4$ in front of the Gaussian term is called the amplification factor. We have introduced the short notation

$$h_i(P_T) = C_i \frac{T \tau A_T}{(2\pi)^3 M_T I_0 \left(P_T \sinh \eta_T \right)} K_1 \left(\frac{\sum_{j=1,2} \sqrt{m_j^2 + P_T^2 / 4} \cosh \eta_T}{T}\right) \hspace{1cm} (8)$$

where $m_j$ are the masses of the recombining (anti)quarks, $M_T$ is the transverse mass of the meson, $A_T$ is the transverse area of the fireball and $C_i$ is the degeneracy factor for meson $i$. The integration over the correlation volume has been absorbed into the correlation strength leading to the new parameter $\hat{c}_0 \approx c_0 \nu_c / (\tau A_T)$.

The single meson spectra take the simple form

$$\frac{dN_i}{P_{Ti} dP_{Ti} d\phi_i d\eta_i} = h_i(P_{Ti}) (1 + \hat{c}_0) (1 + 2 v_2 (P_{Ti}) \cos(2 \phi_i)).$$  \hspace{1cm} (9)

We note that the only modification introduced by correlations consists of a moderate rescaling $1 \to 1 + \hat{c}_0$ [3]. For small numerical values of $\hat{c}_0$ this can be easily absorbed in the overall normalization, so that a good description of single particle spectra can still be achieved with a consistent set of parameters for the parton phase.

The amplification factor $Q = 4$ implies an enhancement of the correlations for hadrons compared to partons. The effect is essentially the same as in the amplification of elliptic flow by the number $n$ of valence quarks. In the case of 2-parton correlations, $Q$ counts the number of possible correlated parton pairs between the $n_A$ (anti)quarks of hadron $A$ and the $n_B$ (anti)quarks of hadron $B$. Apparently

$$Q = n_A n_B,$$  \hspace{1cm} (10)
thus $Q = 6$ for a meson-baryon pair and $Q = 9$ for a baryon-baryon pair.

We define the background subtracted associated yield per trigger particle as a function of relative azimuthal angle

$$Y_{AB}(\Delta \phi) = \frac{1}{N_A} \left( \frac{dN_{AB}}{d(\Delta \phi)} - \frac{d(N_A N_B)}{d(\Delta \phi)} \right).$$

(11)

The uncorrelated background is given by

$$2\pi \frac{d(N_A N_B)}{d(\Delta \Phi)} = N_A N_B (1 + 2\hat{c}_0 + 2\bar{v}_2 A \bar{v}_2 B \cos(2\Delta \phi)).$$

(12)

We note that the background is not constant as a function of $\Delta \phi$ because of the elliptic flow. The terms proportional to $1 + 2\hat{c}_0$ cancel in (11) and hence

$$2\pi N_A Y_{AB}(\Delta \Phi) = Q\hat{c}_0 e^{-\frac{(\Delta \phi)^2}{2\phi_0^2}} N_A N_B.$$  

(13)

The $N_i$ are single particle yields in the kinematic window of the trigger meson or associated meson

$$N_i = 2\pi \int dy_i \int dP_T T^i h_i(P_T).$$

(14)

Correlations in rapidity have not been studied here. Since we integrate over the rapidities, any residual effects can be absorbed into the constant $\hat{c}_0$. We note that the result in (13) also holds for meson-baryon and baryon-baryon correlations with the appropriate amplification factor $Q$.

Before comparing to data we have to find a way to describe the presence of correlations from the F-F process. This requires a study in the framework of dihadron fragmentation functions. These functions are under investigation [23]. Here we choose a simple model that factorizes dihadron fragmentation functions into a product of single hadron fragmentation functions. Details of this model are discussed in [19].

Soft-hard recombination is not part of the original recombination model in [7, 3], but it is taken into account by other groups [8, 12]. The relative importance of soft-hard recombination is not clear, since both approaches fit the experimental data well. It was argued that soft-hard processes are a good source of hadron correlations. We will not discuss correlations from soft-hard recombination further here, but note that they are seizable in [21], while they are found to be a negligible contribution in [19].

For the numerical evaluation of the model we use the set of parameters found in [3] to fit the single hadron spectra and elliptic flow measured at RHIC. We use the minijet calculation in [24] and KKP fragmentation functions [25]. We choose the windows 2.5 GeV/$c \leq P_{T_A} \leq 4.0$ GeV/$c$ for trigger particles, 1.7 GeV/$c \leq P_{T_B} \leq 2.5$ GeV/$c$ for associated particles and $|y_A|, |y_B| < 0.35$ as in the recent analysis of the PHENIX experiment [18]. We fix the azimuthal correlation width to be $\phi_0 = 0.2$, in rough agreement with the experimental data. In order to have a measure of the absolute strength of the correlation we follow [18] and integrate $Y_{AB}$ over the near-side peak

$$Y_{AB}^{\text{cone}} = \int_0^{0.94} d(\Delta \Phi) Y_{AB}(\Delta \Phi).$$

(15)

Fig. 4 shows the integrated associated yield of hadrons for the case that the trigger is a baryon (proton or antiproton) and a meson (pion or kaon) for different centralities. We consider the following cases: (i) F-F, i.e. fragmentation only. (ii) F-F and SS-SS with $\hat{c}_0 = 0$. Keeping fragmentation and turning on recombination without correlations dilutes the signal as expected. The effect is particularly strong for baryon triggers. (iii) F-F and SS-SS with $\hat{c}_0 = 0.08 \times 100/N_{\text{part}}$. We remember that $\hat{c}_0$ contains a space-time integral which can depend
Figure 4. The associated yield $Y_{cone}^{AB}$ for baryon triggers (right panel) and meson triggers (left panel) as a function of $N_{\text{part}}$ compared with PHENIX data. Squares are F-F only, diamonds are F-F and SS-SS with $\hat{c}_0 = 0$, circles are F-F and SS-SS with $\hat{c}_0 = 0.08 \times 100/N_{\text{part}}$.

on centrality. The scaling of $\hat{c}_0$ corresponds to a correlation volume $V_c$ that is independent of centrality. This leads to a nearly constant $Y_{cone}^{AB}$ which is in qualitative agreement with PHENIX measurements [18].

We note that this simple model for correlations can reproduce associated yields in qualitative agreement with the data. The parameters of the parton phase can be chosen consistently so that the excellent fits of single hadron spectra and elliptic flow are preserved. It should be emphasized that the recombination model does not make predictions about the partonic phase. Quantities in the parton phase like radial and elliptic flow, etc. are fit to describe the data. A stringent test for the recombination model is provided by measuring different hadron species. Recombination should be able to describe all hadron species using one universal parameterization of the parton phase. This was demonstrated earlier for observables like hadron spectra, ratios, nuclear modification factors and elliptic flow. It was shown here that it is not difficult to introduce correlations between partons that lead to hadron correlations consistent with experiment. The next step would be to do a more precise study of correlations between identified hadrons and search for a universal parameterization of parton correlations that can fit these results. Fig. 5 shows a possible results obtained with the model above including F-F and SS-SS with the same parameters but now for different combinations of identified trigger and associated hadrons.

5. Conclusions
I have reviewed the quark recombination model applied to heavy ion collisions. There is ample evidence that hadronization in heavy ion collisions is dominated by recombination from a thermalized parton phase. The quark counting rule for elliptic flow, impressively confirmed by experiment, might prove to be an important cornerstone to make the case for the quark gluon plasma.

Experimental findings of jet-like hadron correlations at intermediate $P_T$ are not in contradiction with the existence of recombination. Correlations among partons can be introduced into the existing formalism and naturally lead to correlations among hadrons. Jet-like correlations in the medium can be created by jets that are quenched by the medium and leave a hot spot. Another possible origin for hadron correlations are soft-hard recombination.
**Figure 5.** The same as Fig. 4 with $\hat{c}_0 = 0.08 \times 100/N_{\text{part}}$ but for trigger and associated hadrons identified as mesons or baryons.

**Acknowledgments**

I would like to thank the organizers and participants for an inspiring workshop. This work was supported by DOE grant DE-FG02-87ER40328.

**References**

[1] Collins J C and Soper D E 1982 *Nucl. Phys.* B **194** 445
[2] Adcox K et al PHENIX Coll. 2002 *Phys. Rev. Lett.* **88** 242301
  Adler S S et al PHENIX Coll. 2003 *Phys. Rev. Lett.* **91** 172301
  Adler S S et al PHENIX Coll. 2004 *Phys. Rev.* C **69** 034909
[3] Fries R J, Müller B, Nonaka C and Bass S A 2003 *Phys. Rev.* C **68** 044902
[4] S. S. Adler et al. PHENIX Coll. 2003 *Phys. Rev. Lett.* **91** 182301
[5] Adams J et al STAR Coll. 2004 *Phys. Rev. Lett.* **92** 052302
[6] Magestro D STAR Coll. 2005 *J. Phys. G: Nucl. Part. Phys.* these proceedings
[7] Fries R J, Müller B, Nonaka C and Bass S A 2003 *Phys. Rev. Lett.* **90** 202303
  Nonaka C, Fries R J and Bass S A 2004 *Phys. Lett.* B **583** 73
  C Nonaka, Müller B, Asakawa M, Bass S A and Fries R J 2004 *Phys. Rev.* C **69** 031902
[8] Greco V, Ko C M and Levai P 2003 *Phys. Rev. Lett.* **90** 202302
  Greco V, Ko C M and Levai P 2003 *Phys. Rev.* C **68** 034904
[9] Hwa R C and Yang C B 2003 *Phys. Rev.* C **67** 064902
[10] Fries R J 2004 *J. Phys. G: Nucl. Part. Phys.* **30** S853
[11] Müller B, Fries R J and Bass S A 2005 *Phys. Lett.* B **618** 77
[12] Hwa R C and Yang C B 2003 *Phys. Rev.* C **70** 024904
  Hwa R C and Yang C B 2004 *Phys. Rev.* C **70** 024905
[13] Adler S S et al PHENIX Coll. 2003 *Phys. Rev. Lett.* **91** 072301
[14] Long H STAR Coll. 2004 *J. Phys. G: Nucl. Part. Phys.* **30** S193
[15] Voloshin S A 2003 *Nucl. Phys.* A **715** 379
  Lin Z W and Ko C M 2002 *Phys. Rev. Lett.* **89** 202302
  Lin Z W and Molnar D 2003 *Phys. Rev.* C **68** 044901
[16] Greco V and Ko C M 2004 *Phys. Rev.* C **70** 024901
[17] Adler C et al STAR Coll. 2003 *Phys. Rev. Lett.* **90** 082302
[18] Adler S S et al PHENIX Coll. 2005 *Phys. Rev.* C **71** 051902
  Sickles A PHENIX Coll. 2004 *J. Phys. G: Nucl. Part. Phys.* **30** S1291
[19] Fries R J, Bass S A and Müller B 2005 *Phys. Rev. Lett.* **94** 122301
[20] Wang X N 2004 *Phys. Rev.* C **70** 031901
[21] Hwa R C and Yang C B 2004 *Phys. Rev.* C **70** 054902
[22] Fries R J 2005 *J. Phys. G: Nucl. Part. Phys.* **31** S379
[23] de Florian D and Vanni L 2004 Phys. Lett. B 578 139
    Majumder A and Wang X N 2004 Phys. Rev. D 70 014007
[24] Fries R J, Müller B, and Srivastava D K 2003 Phys. Rev. Lett. 90 132301
    Srivastava D K, Gale C and Fries R J 2003 Phys. Rev. C 67 034903
[25] Kniehl B A, Kramer G and Pötter B 2000 Nucl. Phys. B 582 514