Transverse Expansion and High $p_T$ Azimuthal Asymmetry at RHIC

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Rapid 3+1D transverse plus Bjorken collective expansion in $A + A$ collisions at ultra-relativistic energies is shown to reduce substantially the azimuthal asymmetry resulting from jet quenching. While the azimuthal asymmetry in non-central collisions, $v_2(p_T > 2 \text{ GeV/c}) \sim 0.15$ reported by STAR at RHIC, can be accounted for by spatially anisotropic jet energy loss through a 1+1D expanding gluon plasma with $dN^0/dy \sim 1000$, we show that if rapid transverse collective expansion of the plasma is assumed, then the asymmetry due to jet quenching may be reduced below the observed level. Possible implications of this effect are discussed.

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I. INTRODUCTION

Preliminary STAR data from RHIC [1, 2] suggest a large transverse azimuthal asymmetry $v_2 \sim 0.10 - 0.15$ at moderate high $2 < p_T < 5 \text{ GeV/c}$. In Refs. [3, 4] some of us proposed that the saturation of $v_2(p_T)$ could be interpreted as a manifestation of asymmetric jet energy loss that arises due to the different jet propagation path lengths in non-central ($b > 0 \text{ fm}$) collisions of heavy ions. We applied the GLV [7–10] energy loss formalism to compute the asymmetry and found that un-

II. ENERGY LOSS OF JETS IN TRANSVERSELY EXPANDING MEDIUM

The explicit solution to the problem of energy loss of a hard ultra-relativistic jet ($c = 1, z = \tau$) produced inside a hot and dense medium at position $\vec{x}_0$ and propagating in direction $\vec{v}$ was obtained in [8] to all orders in opacity

$$\chi = \int_0^\infty d\tau \sigma_0(\tau)\rho(\vec{x}_0 + \vec{v}(\tau - \tau_0), \tau).$$

It was shown that for gluon plasma of transverse thickness on the order of the nuclear size and density $\rho \leq 50 \text{ fm}^{-3}$ (relevant at the currently available heavy ion collider energies) the dominant contribution comes from the first order in the opacity expansion [11, 12, 24]. In the case of 1+1D Bjorken longitudinal expansion with initial plasma density $\rho_0 = \rho(0 \equiv \tau_0)$ and formation time $\tau_0 \equiv \tau_0$, i.e.

$$\rho(z = \tau) = \rho_0\left(\frac{\tau_0}{\tau}\right)^\alpha,$$

it is possible to obtain a closed form analytic formula [8] (under the strong asymptotic no kinematic bounds assumption) for the energy loss due to the dominant first order term [15]. For a hard jet penetrating the quark-gluon plasma

$$\frac{d\Delta E(1)}{dx} = \frac{2C_R\alpha_s}{\pi} E \int_{\tau_0}^\infty \frac{d\tau}{\lambda(\tau)} f(Z(x, \tau)),$$

where $x \simeq \omega/E$ is the momentum fraction of the radiated gluon and the formation physics function $f(Z(x, \tau))$ is defined in [8] to be

$$f(x, \tau) = \int_0^\infty \frac{du}{u(1 + u)} [1 - \cos(u Z(x, \tau))].$$

With $Z(x, \tau) = (\tau - \tau_0)\mu^2(\tau)/2xE$ being the local formation physics parameter, two simple analytic limits of Eq. (4) can be obtained. For $x \gg x_c = \ldots$
\[ \mu(\tau_0^2)^2 L^{1-\frac{2\alpha}{\lambda(\tau_0)}} / 2E = L \mu^2(L)/2E, \]

in which case the formation length is large compared to the size of the medium, the small \( Z(x, \tau) \) limit applies leading to \( f(Z) \approx \pi Z/2 \). The interference pattern along the gluon path becomes important and accounts for the non-trivial dependence of the energy loss on \( L \). When \( x \ll x_c \), i.e. the formation length is small compared to the plasma thickness, one gets \( f(Z) \approx \log Z \). In the \( x \gg x_c \) limit \( \bar{\Theta} \) the radiative spectrum Eq. (3) becomes

\[ \frac{d\Delta E^{(1)}_{x \gg x_c}}{dx} \approx \frac{C_R \alpha_s}{2(2-\alpha)} \frac{\mu(\tau_0^2)^2 \lambda(\tau_0)}{L^2-\alpha} \frac{1}{x}. \] (5)

The mean energy loss (to first order in \( \chi \)) integrates to

\[ \Delta E^{(1)}_x = \frac{C_R \alpha_s}{2(2-\alpha)} \frac{\mu(\tau_0^2)^2 \lambda(\tau_0)}{L^{2-\alpha} \pi C} \times \]

\[ \times \left( \log \frac{2E}{\mu(\tau_0)^2 x_0^1/2} + \cdots \right). \] (6)

The logarithmic enhancement with energy comes from the \( x_c < x < 1 \) region \( \bar{\Theta} \). In the case of sufficiently large jet energies \( (E \to \infty) \) this term dominates. For parton energies \( < 20 \text{ GeV} \), however, corrections to this leading log \( 1/x_c \) expression that can be exactly evaluated numerically from the GLV expression and found to be comparable in size. The effective \( \Delta E/E \) in this energy range is approximately constant \( \bar{\Theta} \). It was also recently shown that the effect of gluon absorption is negligible for jet energies \( E \geq 5 \mu \sim 3 \text{ GeV} \) \( \bar{\Theta} \).

In the above expression, transverse expansion can only be very crudely modeled taking \( \alpha > 1 \). To derive an improved analytic expression taking transverse flow into account, we consider next an asymmetric expanding sharp elliptic density profile the surface of which is defined by

\[ \frac{x^2}{(R_x + v_x \tau)^2} + \frac{y^2}{(R_y + v_y \tau)^2} = 1. \] (7)

The area of this elliptic transverse profile increases with time \( \tau \), as

\[ A_\perp(\tau) = \pi (R_x + v_x \tau)(R_y + v_y \tau). \] (8)

Consider the plasma density seen by a jet in direction \( \phi_0 \) starting from \( x_0 = (x_0, y_0) \) inside the ellipse with a specified initial density \( \tau_0 \rho_0 = 1/(\pi R_x R_y) \) \( dN^y/dy \). The density along its path is

\[ \rho(\tau, \phi_0; x_0, y_0) = \frac{1}{\pi} \frac{dN^y}{dy} \left( \frac{1}{\tau} \right) \left( \frac{1}{R_x + v_x \tau} \right) \left( \frac{1}{R_y + v_y \tau} \right) \theta \left( 1 - \frac{(x_0 + \tau \cos \phi_0)^2}{(R_x + v_x \tau)^2} - \frac{(y_0 + \tau \sin \phi_0)^2}{(R_y + v_y \tau)^2} \right). \] (9)

Analytic expressions can be obtained only for asymptotic jet energies when the kinematic boundaries can be ignored. In that case Eq. (3) reduces to

\[ \frac{d\Delta E^{(1)}}{dx} \approx \frac{C_R \alpha_s}{2x} \int_0^\infty d\tau \frac{\mu^2(\tau)}{\lambda(\tau)} \theta \left( \frac{2x E}{\tau \mu^2(\tau)} - 1 \right), \] (10)

where we kept only the dominant \( \log 1/x_c, Z < 1 \) regime. By integrating over \( x \) the total energy loss is for this density profile azimuthally dependent.

\[ \Delta E^{(1)}(\phi_0) \approx \frac{9\pi C_R \alpha_s^3}{4} \int_0^\infty d\tau \rho(\tau, \phi_0) \tau \log \frac{2E}{\tau \mu^2(\tau)}, \] (11)

which is a linear \( (\log \mu) \) weighted line integral over the symmetric density. This integral is still convergent in spite of the singularity of the density at \( \tau = 0 \).

We consider in more detail the analytically tractable case of a sharp expanding elliptic cylinder. We approximate the assumed \( \phi_0 \) independent screening \( \mu(\tau) \approx gT(\tau) = 2(\rho(\tau)/2)^{1/3} \) since \( g \approx 2 \) and the \( (16C(3)/\pi^2)T^3 \approx 2T^3 \) for gluon plasma. We define \( \tau(\phi_0) \) as the escape time to reach the expanding elliptic surface from an initial point \( x_0 \) in the azimuthal direction \( \phi_0 \):

\[ \frac{(x_0 + \tau \cos \phi_0)^2}{(R_x + v_x \tau)^2} + \frac{(y_0 + \tau \sin \phi_0)^2}{(R_y + v_y \tau)^2} = 1. \] (12)

We take \( \omega(\phi_0) = 2 \tau(\phi_0)(\rho(\tau(\phi_0))/2)^{2/3} \) to estimate an upper bound on the logarithmic enhancement factor. Performing the remaining integrals one gets:

\[ \Delta E^{(1)}(\phi_0) \approx \frac{9C_R \alpha_s^3}{4} \frac{dN^y}{dy} \log \frac{E}{\omega(\phi_0)} \times \]

\[ \times \int_0^\infty d\tau \left( \frac{1}{R_x + v_x \tau} \right) \left( \frac{1}{R_y + v_y \tau} \right) \theta \left( 1 - \frac{(x_0 + \tau \cos \phi_0)^2}{(R_x + v_x \tau)^2} - \frac{(y_0 + \tau \sin \phi_0)^2}{(R_y + v_y \tau)^2} \right) \]

\[ = \frac{9}{4} \frac{C_R \alpha_s^3}{R_x R_y} \frac{dN^y}{dy} \left( \frac{1}{\omega(\phi_0)} \right) \frac{E}{a_x - a_y} \log \frac{E}{\omega(\phi_0)}. \] (13)

where \( a_x = v_x/R_x, a_y = v_y/R_y \). This expression is the central result of this paper and provides a simple analytic generalization that interpolates between pure Bjorken \( 1+1 \) expansion for small \( a_{x,y} \), and \( 3+1 \) expansion at large \( a_{x,y} \). In the special case of pure Bjorken (longitudinal) expansion with \( v_x = v_y = 0 \)

\[ \Delta E^{(1)}_{Bj}(\phi_0) = \frac{9C_R \alpha_s^3}{4} \frac{dN^y}{dy} \tau(\phi) \log \frac{E}{\omega(\phi_0)}. \] (14)

In this case, the energy loss depends \( \text{linearly} \) on \( \tau(\phi) \).

Another special case is azimuthally \( \text{isotropic} \) expansion with \( R_x = R_y = R \) and \( v_x = v_y = v_r \). Taking also the longitudinal Bjorken expansion into account leads in this case to

\[ \Delta E^{(1)}_{3D}(\phi_0) = \frac{9C_R \alpha_s^3}{4} \frac{dN^y}{dy} \frac{\tau(\phi_0)}{1 + v_r \tau(\phi_0)/R} \log \frac{E}{\omega(\phi_0)}. \] (15)
We note that for a jet originating near the center of the medium and fully penetrating the plasma the enhanced escape time due to expansion \( \tau = R/(1 - v_T) \) compensates for the \( 1/(1 + v_T\tau(\phi_0)/R) \) dilution factor. Therefore, in this isotropic case, the extra dilution due to transverse expansion has in fact no effect of the total energy loss

\[
\Delta E_{3D}^{(1)}(b = 0 \text{ fm}) \approx \Delta E_{3D}^{(1)}(b = 0 \text{ fm})
\]  
(16)

modulo logarithmic factors. We note that in performing the line integral Eq. (13), the logarithm dependence on the cut-off, \( \omega(\phi_0) \) was neglected. For our numerical estimates we approximate \( \omega(\phi_0) \) by its average. Our first important conclusion is that the inclusive azimuthally averaged jet quenching pattern in central collisions is approximately independent of transverse expansion.

It is important to verify Eq. (12) for more realistic density profiles since it implies that the overall suppression of the high \( p_T \) particle spectra is not affected by transverse flow. We have checked numerically that this approximate line integral Eq. (11) independence of the transverse expansion velocity \( v_T \) holds for transverse density profiles with diffuse edges, namely a Gaussian profile

\[
\rho(\tau, \phi_0; x_0, y_0) = \frac{dN^g}{dy} \left(\frac{1}{\tau}\right) \frac{1}{\pi(R_x^{eff} + v_x\tau)(R_y^{eff} + v_y\tau)} \exp\left[\frac{-(x_0 + \tau \cos\phi_0)^2}{(R_x^{eff} + v_x\tau)^2} - \frac{(y_0 + \tau \sin\phi_0)^2}{(R_y^{eff} + v_y\tau)^2}\right]
\]  
(17)

and a curious (\( \phi_0 \)) asymmetric exponential profile

\[
\rho(\tau, \phi_0; x_0, y_0) = \frac{dN^g}{dy} \left(\frac{1}{\tau}\right) \frac{1}{4(R_x^{eff} + v_x\tau)(R_y^{eff} + v_y\tau)} \exp\left[\frac{-(x_0 + \tau \cos\phi_0)^2}{(R_x^{eff} + v_x\tau)^2} - \frac{(y_0 + \tau \sin\phi_0)^2}{(R_y^{eff} + v_y\tau)^2}\right].
\]  
(18)

One of the most interesting cases arises as a solution of Vlasov equation for a free-streaming massive partons. We consider the ideal Bjorken case with perfect correlation between kinetic and space-time rapidities: \( y = \eta \) with \( y = \tanh^{-1}(p_z/p_0) \), \( \eta = \tanh^{-1}(z/t) \) and transverse propagation at mid-rapidity (\( \tau = t \)). Given initial transverse density \( \rho_0(\vec{x}) \) and with azimuthally isotropic angular distribution of plasma particles with a fixed magnitude of velocity, at time \( \tau \) one gets

\[
\rho(\vec{x}, \tau) = \frac{\tau_0}{\tau} \int \frac{d\Omega_{v_T}}{2\pi} \rho_0(\vec{x} - \vec{v}_T\tau).
\]  
(19)

For Gaussian initial density the evolution is actually given by

\[
\rho(r, \phi, \tau) = \frac{dN^g}{dy} \left(\frac{1}{\tau}\right) \frac{1}{\pi R_x^{eff} R_y^{eff}} \int \frac{d\Omega_{v_T}}{2\pi} \exp \left[ -\frac{(r \cos\phi_0 - v_T\tau \cos\phi)^2}{(R_x^{eff})^2} - \frac{(r \sin\phi_0 - v_T\tau \sin\phi)^2}{(R_y^{eff})^2} \right].
\]  
(20)

This solution is rather different from Eq. (17) because the density expands as a thick shell rather than homogeneously. The same observation applies for the sharp elliptic profile Eq. (3). The general line integral that is needed assuming free Vlasov expansion but averaging over initial coordinates and the angle of the internal particle motion for fixed azimuthally isotropic \( v_T \) is

\[
L(\hat{n}) = \int d^2x_0 \tau_0 \rho_0(\vec{x}_0, \tau_0) \int_{\tau_0}^{\infty} d\tau \rho(\vec{x}_0 + \hat{n}\tau, \tau) = \tau_0 \int d^2x_0 \frac{\tau_0 \rho_0(\vec{x}_0, \tau_0)}{2\pi} \int_{\tau_0}^{\infty} d\tau \rho_0(\vec{x}_0 + (\hat{n} - \vec{v}_T)\tau, \tau_0).
\]  
(21)

The explicit averaging over the jet production point \( \vec{x}_0 \) in this equation can only be performed numerically. However for a jet originating near the center of the medium one can derive expression similar to Eq. (3)

\[
\Delta E_{\text{Vlas.}3+1D}^{(1)}(\phi_0) \approx \frac{9}{4} \frac{C_R \alpha_3^2}{(R^{eff})^2} \frac{dN^g}{dy} \log \frac{E}{\omega(\phi_0)} \int_0^{\infty} d\tau \int d\phi \frac{1}{2\pi} \exp \left( -\tau \left( -1 - 2v_T \cos\phi + v_T^2 \right) \right) \]  
(22)

The equivalent escape factor (enhanced propagation length modulated by the dilution of the medium) is given in this Vlasov scenario by

\[
\tau(\phi_0)_{\text{vlas.}} = \frac{R^{eff}}{\sqrt{\pi}(1 + v_T)} K \left( \frac{4v_T}{(1 + v_T)^2} \right),
\]

where \( K(x) \) is the complete elliptic integral of first kind. The escape factor stays within \( \pm 10\% \) of \( R^{eff} \) for flow velocities up to \( v_T = 0.8 \). As \( v_T \rightarrow 1 \), \( \tau(\phi_0)_{\text{vlas.}} \) diverges only logarithmically. In case of non-central collisions, \( b \neq 0 \), and averaging over the jet production points, \( \langle \Delta E_{3D}^{(1)} \rangle_{\vec{x}_0} \), the validity of Eqs. (14, 22) can be confirmed numerically.

In non-central collisions, the azimuthal asymmetry of the energy loss can be expanded in Fourier series and characterized as

\[
\Delta E_{3D}^{(1)}(\phi) = \Delta E(1 + 2\delta_2(E) \cos 2\phi + \cdots).
\]  
(23)

From Eq. (13), we compute numerically the azimuthally averaged mean \( \Delta E \) and the energy loss asymmetry \( \delta_2(E) \) averaging over the uniform elliptic density, Eq. (3), of the jet formation points. We consider only the specific
case of a $E = 10$ GeV gluon jet to illustrate the dependence of these quantities on the mean expansion velocity. The spatial anisotropy is fixed by $R_x = 2.5$ fm and $R_y = 4.7$ fm that approximately corresponds to the same second moment of the spatial anisotropy of the sharp intersecting cylinders considered in Ref. [5] for $b = 7$ fm. The gluon plasma rapidity density that fixes the magnitude of the energy loss was assumed to be $dN^g/dy (b = 7$ fm$) = 325$ with $\alpha_s = 0.2$ for illustration.

We calculate the energy loss and its asymmetry also by using a full hydrodynamic calculation of Ref. [14]. In this case we use the parametrization eBC of Ref. [14] to initialize the system and treat gluon number as conserved current to calculate the density evolution needed in the line integral Eq. (11). We again average over the jet formation points but in this case their density is not constant but given by the number of binary collisions per unit area as in the Woods-Saxon geometry of Ref [5].

We plot in Figs. 1 and 2 the dependence of $\Delta E$ and $\delta_2(E = 10$ GeV$)$ on the mean transverse expansion velocities, fixing the velocity asymmetry from hydrodynamic results to be $v_x = v_T (1 + \nu_a)$ and $v_y = v_T (1 - \nu_a)$. Boosted thermal sources simulations and hydrodynamic calculations indicate that $v_T \approx 0.5 - 0.6$ and $\nu_a \approx 0.05 - 0.10$. The analytic computation and the hydro simulation in Fig. 1 were normalized to the same mean energy loss for the $v_T = 0$ point. We see that $\Delta E$ is weakly dependent on transverse flow for the schematic but analytically tractable uniform expanding elliptic density and we have obtained similar results for a wide variety of transverse profiles Eqs. (14 and 21). Full hydrodynamic calculation of the $\phi$ and $x_0$ averaged line integral Eq. (11) is also consistent with approximately constant mean energy loss within $\pm 5\%$ for both central and semi-central collisions.

However, the azimuthal asymmetry of the energy loss is strongly reduced for realistic hydrodynamic flow velocities. This implies a much smaller $v_2$ at high $p_T$ than obtained in Ref. [3] where transverse expansion was not considered. Hydrodynamic calculations with transverse expansion predict similar reduction of the averaged line integral Eq. (11). We note that factor of 1.5 - 2 difference in $\delta_2(E)$ without transverse flow in Fig. 2 reflects the difference between the sharp cylinder geometry and the Wood-Saxon geometry as shown in Ref. [5].

![Gluon Jet, $E = 10$ GeV](image1)

**FIG. 1.** The azimuthally averaged energy loss for a 10 GeV gluon jet propagating through a 3+1D elliptic expanding plasma is plotted as a function of the mean transverse flow velocity, $v_T$. The initial profile is a homogeneous ellipse approximating $Au + Au$ at $b = 7$ fm impact parameter. Transverse velocity asymmetry is fixed to be $\nu_a = 0.1$. 1+1D and 3+1D hydro calculations are shown.

![Gluon Jet, $E = 10$ GeV](image2)

**FIG. 2.** The second harmonic of the jet energy loss for a 10 GeV gluon jet propagating through a 3+1D elliptic expanding plasma is plotted as a function of the mean transverse flow velocity, $v_T$, is shown. The initial profile is chosen as in Fig. 1. We study 3 different transverse velocity asymmetries, i.e. $\nu_a = 0.0, 0.1, 0.2$. Hydrodynamic 1+1D and 3+1D calculations using diffuse Wood-Saxon geometry are also included.

### III. HIGH $p_T$ AZIMUTHAL ASYMMETRY SCENARIOS

The results of the previous section suggest that if strong transverse expansion of the plasma can be confirmed (and other scenarios including soft string fragmentation, baryon junctions [13,20,27] and classical Yang-Mills [28] evolution rejected) then the high transverse
momentum asymmetry predicted by energy loss calculations may be reduced (by a factor of 2 to 4). This would leave open the understanding of the present STAR data in terms of the eikonal picture. Interesting new possibilities arise to explain the moderate to high \( p_T \) dependence of the transverse azimuthal asymmetry \( v_2(p_T > 2 \text{ GeV/c}) \). We suggest that the key to converging to the correct dynamical picture is the flavor dependence of both the asymmetry and the \( p_T \) differential particle multiplicities. We discuss three possible scenarios for high \( p_T \) hadron production with expected \( v_2(p_T) \) behavior shown in Fig. 3.

![Graph showing three possible scenarios of the \( v_2(p_T) \) behavior](image)

**FIG. 3.** Three possible scenarios of the \( v_2(p_T) \) behavior are shown. Error bands are suggestive of the theoretical uncertainties within each framework (e.g. the density of the medium or the transverse expansion velocity \( v_T \)).

1. If experiments confirm an approximate saturation of \( v_2 \approx 0.10 - 0.15 \) at high \( p_T \) for both pions and protons together with a baryon to meson ratio dropping back below unity this may suggest that the transverse flow might be significantly smaller than the values used in the thermal model description of relativistic heavy ion collisions. The assumption of very large transverse expansion velocities may be a caveat since it is not supported by pion interferometry as pointed out in Ref. [29]. In this case the energy loss asymmetry plays significant role, producing comparable asymmetries for \( p \) and \( \pi \). Alternatively, appropriate description of particle production could be provided by dissipative relativistic hydrodynamics [32] or inelastic parton cascades [33] with large partonic cross sections \( \sigma \sim 10 - 40 \text{ mb} \).

2. Another possibility arises if the pion multiplicity at moderate high transverse momenta is dominated by the quenched pQCD spectra but the baryon production at moderate transverse momenta \( 2 < p_T < 5 \text{ GeV/c} \) is controlled by a non-perturbative mechanism such as baryon junctions (which exhibit azimuthal asymmetry and flow) or hydro (with baryon-chemical potential which can be inferred from the baryon transport picture). A distinctive characteristic of this picture is that one expects qualitatively different \( v_2(p_T) \) behavior for \( p \) and \( \pi \) and a baryon/meson anomaly limited to the intermediate transverse momentum window [15-22]. Above \( p_T = 2 \text{ GeV/c} \) the proton and anti-proton \( v_2 \) may keep growing due to flow effects, whereas the pion \( v_2 \), being driven by asymmetrically quenched pQCD, may start decreasing because strong transverse expansion (if confirmed) reduces the energy loss azimuthal anisotropy. The comparable contribution of baryons and mesons in this region [12,14,21,34] leads to cancellation of those effects in the inclusive charged particle elliptic flow, thus producing a plateau in the \( 2 < p_T < 5 \text{ GeV/c} \) window. At high transverse momenta of \( p_T \approx 6 - 10 \text{ GeV/c} \) the pQCD pions dominate the charged particle multiplicity, thus leading to a gradual decrease of \( v_2 \). The rate at which the high \( p_T \) azimuthal asymmetry vanishes is determined by the mean transverse expansion velocity \( v_T \), the velocity asymmetry (see Fig. 2), and the pQCD fractional contribution to the baryon differential \( p_T \) multiplicities. The schematic \( v_2 \) behavior in this scenario is illustrated in Fig. 4. It predicts a detectable difference in the moderate high transverse momentum behavior of the elliptic flow for pions and protons.

3. Last, but not least, there are predictions based on boosted thermal sources and ideal hydrodynamics. The predicted \( v_2(p_T) \) grows continuously as a function of the transverse momentum and saturates at a maximum value. This was parametrized from hydro simulations [33] in Ref. [1] by

\[
v_2(p_T) = \tanh[|p_T/(10 \pm 2 \text{ GeV/c})|].
\]

Similarly the baryon/meson ratio (assuming equal freeze-out temperatures) is a monotone function of \( p_T \) and saturates at \((p/\pi)_{\text{max}} = 2\) from simple spin counting.

Precision data on identified hadron spectra at high \( p_T \) at RHIC is needed to shed light on which if any of these dynamical scenarios predicts the correct flavor dependence of the azimuthal asymmetry and the proton to pion ratio.
We have analyzed the effect of possible strong transverse flow \( v_T \sim 0.6 \) on the transverse azimuthal asymmetry generated by jet energy loss. We have shown that while the overall magnitude of the quenching is only weakly dependent on \( v_T \), the azimuthal asymmetry \( v_2 \) of the quenched pQCD component is significantly reduced. In order to account for the observed saturation of \( v_2 \) of charged hadrons, scenario 2 based on \([5,19]\) is forced into predicting an anomalous asymmetry of the baryon component alone. This scenario interprets the anomalous baryon to pion ratio in terms of baryon junction dynamics \([19,21]\). It differs from scenarios 1 and 3 in predicting that the pion asymmetry is expected to vanish more quickly with increasing \( p_T \), while the baryon asymmetry is expected to be near maximal up to \( \sim 4 \ldots 7 \) GeV/c. Hydrodynamic and dissipative transport models expect a much less dramatic dependence on \( p_T \).

**IV. CONCLUSIONS**

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