Hadronization in heavy ion collisions: recombination or fragmentation?

R J Fries†, B Müller†, C Nonaka† and S A Bass†‡

† Physics Department, Duke University, P.O.Box 90305, Durham, NC 27708, USA
‡ RIKEN BNL Research Center, Brookhaven National Laboratory, Upton, NY 11973, USA
E-mail: rjfries@phy.duke.edu

Abstract. We show that hadron production in relativistic heavy ion collisions at transverse momenta larger than 2 GeV/c can be explained by the competition of two different hadronization mechanisms. Above 5 GeV/c hadron production can be described by fragmentation of partons that are created perturbatively. Below 5 GeV/c recombination of partons from the dense and hot fireball dominates. This can explain some of the surprising features of RHIC data like the constant baryon-to-meson ratio of about one and the small nuclear suppression for baryons between 2 to 4 GeV/c.

Submitted to: JPG

PACS numbers: 25.75.Dw,24.85.+p

The Relativistic Heavy Ion collider (RHIC) has provided exciting data about hadron production at transverse momenta of a few GeV/c in central Au+Au collisions. The production of pions at high \( P_T \) was found to be suppressed compared to the scaled yield from \( p + p \) collisions [1]. This jet quenching effect can be understood by final state interaction of fast partons with the dense and hot medium produced in central heavy ion collisions. Fast partons lose energy via induced bremsstrahlung before they can fragment into high \( P_T \) hadrons [2]. The suppression effect is dramatic and can be as large as a factor of 6 above \( P_T = 5 \) GeV/c.

On the other hand the suppression of protons and antiprotons seems to be much less [3]. Experimental data from PHENIX show a proton/pion ratio of about 1 between 1.5 GeV/c and 4 GeV/c [4]. This is surprising since the production of protons and antiprotons is usually suppressed compared to the production of pions because of the much larger mass. At high transverse momentum this can be understood in terms of perturbative quantum chromodynamics (pQCD) [5]. The fragmentation functions \( D_{a \rightarrow h}(z) \) describe the probability for a parton \( a \) with momentum \( p \) to turn into a hadron with momentum \( zp, 0 < z < 1 \). These fragmentation functions were measured for pions and protons, mainly in \( e^+e^- \) collisions, and give a ratio \( p/\pi^0 < 0.2 \) for \( P_T > 2 \) GeV/c when used in \( p + p \) and \( N + N \) collisions. The energy loss of partons in a medium can be taken into account by a rescaling of the parton momentum [6]. However this should affect baryons and mesons in the same way.
The lack of nuclear suppression for baryons is a challenge for our understanding of hadron production. The currently accepted picture assumes that a parton with large transverse momentum is produced in a hard scattering reaction between initial partons, propagates through the surrounding hot matter and loses energy by interactions, and finally hadronizes. Apparently the creation and interaction of a parton will happen independently of its later fate during hadronization. Therefore any unusual behavior that is different between baryons and mesons can only be attributed to hadronization itself. We propose to use recombination of quarks from the surface of the hot fireball as an alternative hadronization mechanism

In the fragmentation process a parton with transverse momentum $p_T$ is leaving the interaction zone while still being connected with other partons by a color string. The breaking of the string creates quark antiquark pairs which finally turn into hadrons. The distribution of one of these hadrons, which is bound to have less transverse momentum $P_T = z p_T$, is then described by a fragmentation function. The average value $\langle z \rangle$ is about 0.5 for pions in $p+p$ collisions. In other words the production of a, say, 5 GeV/$c$ pion has to start with a 10 GeV/$c$ parton in average, which are rare to find due to the steeply falling parton spectrum. Jet quenching even enhances the lack of high $p_T$ partons. On the other hand, the 5 GeV/$c$ pion could be produced by the recombination of a quark and an antiquark with about 2.5 GeV/$c$ each in average. 2.5 GeV/$c$ and 10 GeV/$c$ are separated by orders of magnitude in the parton spectrum. The price to pay is of course that two of these partons have to be found close to each other in phase space. However we do have a densely populated phase space in central heavy ion collisions at RHIC where we even expect the existence of a thermalized quark gluon plasma.

Recombination of quarks has been considered before in hadron collisions \cite{8} and was also applied to heavy ion collisions \cite{9}. In QCD the leading particle effect in the forward region of a hadron collision is well known. This is the phenomenon that the production of hadrons that share valence partons with the beam hadrons are favored in forward direction. It has been realized that this can only be explained by recombination. This has fueled a series of theoretical work, see e.g. \cite{10} and references therein. Recently the recombination idea for heavy ion collisions, stimulated by the RHIC results, has been revived for elliptic flow \cite{11} and hadron spectra and ratios \cite{7,12}.

The formalism of recombination has already been developed in a covariant setup for the process of baryons coalescing into light nuclei and clusters in nuclear collisions \cite{13,14}. We give a brief derivation for the case of mesons. By introducing the density matrix $\hat{\rho}$ for the system of partons, the number of quark-antiquark states that will be measured as mesons is given by

$$ N_M = \sum_{ab} \int \frac{d^3P}{(2\pi)^3} \langle M; P | \hat{\rho}_{ab} | M; P \rangle $$

Here $\langle M; P \rangle$ is a meson state with momentum $P$ and the sum is over all combinations of quantum numbers — flavor, helicity and color — of valence partons that contribute to the given meson $M$. This can be cast in covariant form using a hypersurface $\Sigma$ for hadronization \cite{15}

$$ \frac{dN_M}{d^3P} = C_M \int \frac{d\sigma \cdot u(\sigma)}{(2\pi)^3} \int \frac{d^3q}{(2\pi)^3} w_a \left( \sigma; \frac{P}{2} - q \right) \Phi^W_M(q) w_b \left( \sigma; \frac{P}{2} + q \right). $$
Here $E$ is the energy of the four vector $P$, $d\sigma$ a measure on $\Sigma$ and $u(\sigma)$ is the future oriented unit vector orthogonal to the hypersurface $\Sigma$. $w_a$ and $w_b$ are the phase space densities for the two partons $a$ and $b$, $C_M$ is a degeneracy factor and

$$\Phi_W^M(r, q) = \int d^3r \phi_W^M(r, q)$$

(3)

where $\Phi_W^M(r, q)$ is the Wigner function of the meson [13].

Since the hadron structure is best known in a light cone frame, we write the integral over $q$ in terms of light cone coordinates in a frame where the hadron has no transverse momentum but a large light cone component $P^+$. This can be achieved by a simple rotation from the lab frame. Introducing the momentum $k = P/2 - q$ of parton $a$ in this frame we have $k^+ = xP^+$ with $0 < x < 1$. We make an ansatz for the Wigner function of the meson in terms of light cone wave functions $\phi_M(x)$. The final result can be written as [7]

$$E \frac{N_M}{d^3P} = C_M \int_{\Sigma} d\sigma \frac{P \cdot u(\sigma)}{(2\pi)^3} \int_0^1 dx w_a(\sigma; xP^+) |\phi_M(x)|^2 w_b(\sigma; (1-x)P^+).$$

(4)

For a baryon with valence partons $a$, $b$ and $c$ we obtain

$$E \frac{N_B}{d^3P} = C_B \int_{\Sigma} d\sigma \frac{P \cdot u(\sigma)}{(2\pi)^3} \int_0^1 dx_1 dx_2 dx_3 \delta(x_1 + x_2 + x_3 - 1)$$

$$\times w_a(\sigma; x_1 P^+) w_b(\sigma; x_2 P^+) w_c(\sigma; x_3 P^+) |\phi_B(x_1, x_2, x_3)|^2.$$ (5)

$\phi_B(x_1, x_2, x_3)$ is the effective wave function of the baryon in light cone coordinates.

A priori these wave functions are not equal to the light cone wave functions used in exclusive processes. We are recombining effective quarks in a thermal medium and not perturbative partons in an exclusive process. Nevertheless, as an ansatz for a realistic wave function one can adopt the asymptotic form of the light cone distribution amplitudes

$$\phi_M(x) = \sqrt{30x(1-x)},$$

$$\phi_B(x_1, x_2, x_3) = 12\sqrt{30} x_1 x_2 x_3$$

as a model. However, it turns out that for a purely exponential spectrum the shape of the wave function does not matter. In that case the dependence on $x$ drops out of the product of phase space densities

$$w_a(\sigma; xP^+) w_b(\sigma; (1-x)P^+) \sim e^{-xP^+/T} e^{-(1-x)P^+/T} = e^{-P^+/T}.$$ (7)

One can show that a narrow width approximation, using $\delta$ peaked wave functions that distribute the momentum of the hadron equally among the valence quarks — $1/2$ in the case of a meson and $1/3$ in the case of a baryon — deviates by less than 20% from a calculation using the wave functions in (6). That there is a small deviation can be attributed to the violation of energy conservation. Since recombination is a $2 \to 1$ or $3 \to 1$ process, energy will generally not be conserved if we enforce momentum conservation and a mass shell condition for all particles. However, one can show that violations of energy conservation are of the order of the effective quark masses or $\Lambda_{QCD}$ and can therefore be neglected for transverse hadron momenta larger than 2 GeV/c.
Fragmentation of partons is given by [16]

$$E \frac{d\sigma_h}{d^3P} = \sum_a \int_0^1 \frac{dz}{z^2} D_{a \rightarrow h}(z) E_a \frac{d\sigma_a}{d^3P_a}. \quad (8)$$

The sum runs over all parton species $a$ and $\sigma_a$ is the cross section for the production of parton $a$ with momentum $P_a = P/z$. We use a leading order (LO) pQCD calculation of $\sigma_a$ [17] together with LO KKP fragmentation functions [18]. Energy loss of the partons is taken into account by a shift of the parton spectrum by

$$\Delta p_T = \sqrt{\lambda} p_T. \quad (9)$$

We summarize that the transverse momentum dependent yield of mesons from recombination can be written as $\sim C_M w^2(P_T/2)$ in the simple narrow width approximation, whereas from fragmentation we expect $\sim D(z) \otimes w(P_T/z)$. For an exponential parton spectrum $w = e^{-P_T/T}$ the ratio of recombination to fragmentation is

$$R/F = \frac{C_M}{\langle D \rangle} e^{-P_T/T} \left(1 - \frac{1}{\langle z \rangle}\right) \quad (10)$$

where $\langle D \rangle < 1$ and $\langle z \rangle < 1$ are average values of the fragmentation function and the scaling variable. Therefore $R/F > 1$. In fact, recombination always wins over fragmentation from an exponential spectrum (as long as the exponential is not suppressed by small fugacity factors). The same is true in the case of baryons.

Now let us consider a power law spectrum $w = A(P_T/\mu)^{-\alpha}$ with a scale $\mu$ and $\alpha > 0$. Then the ratio of recombination over fragmentation is

$$R/F = \frac{C_M A}{\langle D \rangle} \left(\frac{4}{\langle z \rangle}\right)^\alpha \left(\frac{P_T}{\mu}\right)^{-\alpha} \quad (11)$$

and fragmentation ultimately has to win at high $P_T$. We also note that we can expect a constant baryon/meson ratio from recombination, when the parton spectrum is exponential. In this case the ratio is only determined by the degeneracy factors

$$\frac{dN_B}{dN_M} = \frac{C_B}{C_M} \quad (12)$$

For our numerical studies we consider fragmentation of perturbative partons and recombination from a thermal phase

$$w_{th} = e^{-P_T \cosh(\eta-y)/T} e^{-v^2/2\Delta^2} \quad (13)$$

with an effective temperature $T$. $\eta$ is the space-time rapidity and $\Delta \approx 2$ the width of the rapidity distribution. We fix the hadronization hypersurface $\Sigma$ by the condition $\tau_f = \sqrt{t^2 - z^2} = \text{const.}$. We set $\tau_f = 5$ fm/c. A two phase parton spectrum with a perturbative power law tail and an exponential part at low transverse momentum is also predicted by parton cascades like VNI/BMS [19].

The parameters, $\lambda$ for the average energy loss and $T$ for the slope of the exponential parton spectrum, are determined by a fit to the inclusive charged hadron spectrum measured by PHENIX [20]. We obtain $\lambda \approx 1$ GeV and $T \approx 350$ MeV. The temperature contains the effect of a blue-shift because of the strong radial flow. An additional normalization factor of about $1/30$ for the recombination part is necessary to describe the data. This is due to the use of an effective temperature which gives a too large particle number compared to the physical temperature that we expect to be around 175 MeV.
Hadronization in heavy ion collisions: recombination or fragmentation?

Figure 1. Left: the charged hadron spectrum in Au+Au collisions at $\sqrt{s} = 200$ GeV as a function of $P_T$. Contributions from fragmentation (dashed), recombination (dotted) and the sum of both (solid line) are shown. Data are from the PHENIX collaboration [20]. Right: the spectrum of up quarks.

In Figure 1 we show the inclusive charged hadron spectrum using recombination and fragmentation for pions, kaons, protons and antiprotons. We note that the crossover between the recombination dominated and fragmentation dominated regimes is around 5 GeV/$c$. It is in fact earlier for pions ($\sim 4$ GeV/$c$) than for protons ($\sim 6$ GeV/$c$). In Figure 1 we also give the spectrum of up quarks as an example for the partonic input of our calculation. We note that the crossover between the perturbative and the thermal domain is here around 3 to 3.5 GeV/$c$. It is characteristic for recombination that features of the parton spectrum are pushed to higher transverse momentum in the hadron spectrum.

In Figure 2 we give the proton/pion ratio and the nuclear modification factor $R_{AA}$ for pions, protons and charged hadrons. The proton/pion ratio shows a plateau between 2 and 4 GeV/$c$ in accordance with experiment and a steep decrease beyond that. This decrease has not yet been seen and is a prediction of our work.

In the quantity $R_{AA}$ the effect of jet quenching is manifest for all hadrons at large transverse momentum. For pions $R_{AA}$ grows only moderately at low $P_T$ in accordance with PHENIX data [21]. On the other hand recombination is much more important for protons, because they are suppressed in the fragmentation process. Therefore $R_{AA}$ nearly reaches a value of 1 below 4 GeV/$c$ for protons.

In summary, we have discussed recombination as a possible hadronization mechanism in heavy ion collisions. We have shown that recombination can dominate over fragmentation up to transverse momenta of 5 GeV/$c$. Recombination provides a natural explanation for the baryon/meson ratio and the nuclear suppression factors observed at RHIC.

Acknowledgments. This work was supported by RIKEN, Brookhaven National Laboratory, DOE grants DE-FG02-96ER40945 and DE-AC02-98CH10886, and by the Alexander von Humboldt Foundation.

[1] K. Adcox et al. [PHENIX], Phys. Rev. Lett. 88, 022301 (2002); C. Adler et al. [STAR], nucl-ex/0210033
Hadronization in heavy ion collisions: recombination or fragmentation?

Figure 2. Left: the $p/\pi^+$ ratio exhibits a transition between 4 and 6 GeV from the recombination dominated regime ($p/\pi^+ \sim 1$) to the fragmentation dominated regime ($p/\pi^+ \sim 0.1$). Right: the nuclear modification factor $R_{AA}$ for pions (solid), charged hadrons (dashed) and protons (dotted line). Data points are $R_{AA}$ for $\pi^0$ from PHENIX [21].

[2] M. Gyulassy and X. Wang, Nucl. Phys. B 420, 583 (1994); R. Baier et al., Nucl. Phys. B 484, 265 (1997).
[3] K. Adcox et al. [PHENIX], Phys. Rev. Lett. 88, 242301 (2002); C. Adler et al. [STAR], Phys. Rev. Lett. 86, 4778 (2001); C. Adler et al. [STAR], Phys. Rev. Lett. 89, 092301 (2002); K. Adcox et al. [PHENIX], Phys. Rev. Lett. 89, 092302 (2002);
[4] T. Chujo [PHENIX Collaboration], arXiv:nucl-ex/0209027.
[5] J. C. Collins and D. E. Soper, Nucl. Phys. B 194, 445 (1982).
[6] X. F. Guo and X. N. Wang, Phys. Rev. Lett. 85, 3591 (2000); Nucl. Phys. A 696, 788 (2001).
[7] R. J. Fries, B. Müller, C. Nonaka and S. A. Bass, Phys. Rev. Lett. 90, 202303 (2003).
[8] K. P. Das and R. C. Hwa, Phys. Lett. B 68, 459 (1977); Erratum ibid. 73, 504 (1978); R. G. Roberts, R. C. Hwa and S. Matsuda, J. Phys. G 5, 1043 (1979).
[9] C. Gupta et al., Nuovo Cim. A 75, 408 (1983); T. S. Biro, P. Levai and J. Zimanyi, Phys. Lett. B 347, 6 (1995); T. S. Biro, P. Levai and J. Zimanyi, J. Phys. G 5, 1561 (2002).
[10] J. C. Anjos, J. Magnin and G. Herrera, Phys. Lett. B 523, 29 (2001); E. Braaten, Y. Jia and T. Mehen, Phys. Rev. Lett. 89, 122002 (2002).
[11] S. A. Voloshin, nucl-ex/0210014; Z. W. Lin and C. M. Ko, Phys. Rev. Lett. 89, 202302 (2002); D. Molnar and S. A. Voloshin, nucl-th/0302014.
[12] V. Greco, C. M. Ko and P. Levai, Phys. Rev. Lett. 90, 202302 (2003).
[13] C. B. Dover et al., Phys. Rev. C 44, 1636 (1991).
[14] R. Scheibl and U. W. Heinz, Phys. Rev. C 59, 1585 (1999).
[15] F. Cooper and G. Frye, Phys. Rev. D 10, 186 (1974).
[16] J. F. Owens, Rev. Mod. Phys. 59, 465 (1987).
[17] R. J. Fries, B. Müller, and D. K. Srivastava, Phys. Rev. Lett. 90, 132301 (2003); D. K. Srivastava, C. Gale and R. J. Fries, Phys. Rev. C 67, 034903 (2003).
[18] B. A. Kniehl, G. Kramer and B. Pötter, Nucl. Phys. B 582, 514 (2000).
[19] S. A. Bass, B. Müller, and D. K. Srivastava, Phys. Lett. B 551, 277 (2003).
[20] J. Jia [PHENIX], nucl-ex/0209029.
[21] S. Mioduszewski [PHENIX], nucl-ex/0210021.