The Next Yukawa Phase of QCD at RHIC
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QCD predicts the existence of a new partonic Yukawa phase above a critical temperature $T_c \sim 150$ MeV. I review some of the key observables which will be used in heavy ion reactions to search for this new phase at RHIC. Systematics of collective observables ($N_{ch}, E_\perp$), flow patterns, meson interferometry, jet quenching, and $J/\psi$ suppression are discussed. This talk is updated with first data from RHIC.

1. Yukawa in QCD

In 1935 H. Yukawa[1] proposed a theory of nuclear forces based on the exchange of a massive boson

$$V_Y(r, m) = \alpha_{\text{eff}} e^{-mr}/r$$

He estimated that $m \sim 100$ MeV to account for the short range $\sim 2$ fm of the nuclear force, and in 1947 Powell confirmed his theory. Yukawa’s meson theory forms the basis for the current effective theory of nuclear forces[2].

In the field of electrolytes, Debye and Hückle[3] had already discovered the Yukawa potential in a completely different context. The polarizability of the medium in the presence of an external charge density leads to a non-linear self consistent equation with the same solution as eq. (1), but in that case the effective mass is the Debye electric screening mass

$$\mu^2 = 4\pi \sum_q q^2 |q|/T.$$  

A conductive medium transforms Coulomb power law forces into the Yukawa form.

In nuclear theory, the Yukawa meson mass results from the finite gap of the elementary excitations (pions, ...) of the physical QCD vacuum. In this talk, I discuss current efforts to try to drill a perturbative hole into the nonperturbative vacuum using RHIC to see the breakdown of Yukawa’s meson theory. As we review below, in the deconfined, chirally symmetric phase of QCD at high temperatures, the Debye-Huckle mechanism transforms the mesonic Yukawa potentials into a color-electric screened gluonic Yukawa potential in a quark-gluon plasma (QGP).

In QCD, the color potential between partons is approximately Coulombic at small distances due to the asymptotic freedom property of non-Abelian gauge theories. However, below a critical temperature, $T_c \sim 150$ MeV, the effective potential between the colored partons has a long range linear confining term $\kappa r$ with a huge ”string” tension, $\kappa \sim 1$
GeV/fm. In this confining phase of QCD, the heavy $q \bar{q}$ potential is well parameterized by the Lüscher form

$$V_L(r, 0) = -\frac{\alpha_L}{r} + \kappa r,$$  \hspace{1cm} (3)

(as long as dynamical quark pair production is ignored). The Coulombic part, with strength $\alpha_L = \pi/12$, arises from the zero point quantum fluctuations of the string. As the temperature increases, but remains below the deconfinement transition, $T < T_c$, the enhanced fluctuations due to thermal agitation of the string modifies the effective potential into the approximate Gao form \[5\]

$$V_G(r, T) = -\frac{\alpha_L}{r} \left[ 1 - \frac{2}{\pi} \tan^{-1}(2rT) \right] + \left[ \kappa - \frac{\pi}{3} T^2 \left( 1 - \frac{2}{\pi} \tan^{-1}\left( \frac{1}{2rT} \right) \right) \right] r + \cdots$$  \hspace{1cm} (4)

The decrease of the effective string tension, $\kappa(T)$, predicted above has been measured via lattice QCD calculations\[6\] as shown in Fig.(1). However, the “measured” string tension is found to decrease faster than predicted in eq.(4) near the critical temperature. Note that in Fig.(1) $T \approx T_c \exp(11/6(\beta - \beta_c))$ with $\beta_c = 4.0729$ for this lattice.

For temperatures above $T_c$, the lattice data in Fig.1 reveal the predicted Yukawa phase of QCD is predicted. The heavy quark potential mutates into a short range generalized Yukawa form, which on the lattice is fit with the form

$$V_L(r, T, d) = -\frac{\alpha(T)T}{(rT)^{d/2}} e^{-\mu_L(T)r/2}.$$  \hspace{1cm} (5)

Note that $d_L = 2, \mu_L(T) = 2m_E(T)$ correspond to the ideal Yukawa form. The perturbative thermal QCD chromo-electric Debye mass $m_E = \mu(T)/2$ is \[7\]

$$m_E(T) = g(T)T \left( \frac{N_c}{3} + \frac{N_F}{6} \right)^{1/2}$$  \hspace{1cm} (6)

for $N_c$ colors and $N_F$ flavors. For $N_c = 2, N_f = 0$ in Fig. (1), we expect $\mu_L(T) = 2m_E = 1.6g(T)T$ as shown by the solid line in Fig.(1). The fits to the lattice QCD data in Fig.(1b) from \[8\] show that for $T > 2T_c$ $\mu_L \approx 2.5T$ is not far from the pQCD estimate. However, the exponent $d_L \approx 1.5$ is significantly below the value 2 expected from pQCD. An even more striking nonperturbative deviation is seen near $T_c$, at which point $d < 1$ and $\mu \sim T/2$. This suggests a rather long range interaction that may be the precursor of the confinement transition.

The predicted thermodynamic properties of the QCD Yukawa phase are shown in Fig.3 from ref.\[9\] for 2 flavor $12^3 \times 6$ IQCD. A present caveat about all lattice results is that the pion is still too massive to allow contact with the “known” thermodynamic properties of ordinary hadronic/nuclear matter below $T_c$. Two striking features of Fig.(3) suggest two key observable signatures of this phase transition in nuclear collisions. First, the entropy density $\sigma(T) = (\epsilon + p)/T$ increase very rapidly with $T$ in a narrow interval $\Delta T/T_c < 0.1$. Second, the plasma becomes extremely soft $p/\epsilon \ll 1$ and $c_s^2 \ll 1$ near $T_c$. As we review below, the first feature may lead to time delay (the QGP stall) measurable via hadronic interferometry. The second feature may lead to interesting non-linear collective flow observables in nuclear collisions. The experimental verification of these fundamental
Figure 1. The heavy quark potential in the confined phase of SU(2) quenched QCD (top) compared to that in the deconfined Yukawa phase above $T_c$ (bottom). Results are for a $32^3 \times 4$ lattice from Karsch et al[6]. Eq.(4) fits the confined lattice ”data” well, but the QCD string tension decreases more rapidly near $T_c$. For $T > T_c$ the potential is screened by the deconfined gluons (in this quenched calculation) and acquires the generalized Yukawa form (5). Here $T \approx T_c \exp(11/6(\beta - \beta_c))$ increases with $\beta$ and $\beta_c = 4.0729$. 
Figure 2. The chromo-electric Debye screening mass, $\mu(T) = 2m_E(T)$ and the effective exponent, $d(T)$, of the effective Yukawa potential in the deconfined phase versus $T/T_c$ is shown from Karsch et al [6] for SU(2) quenched QCD. The lines show expected dependence based on thermal pQCD. Note that near $T_c$, the range $2/\mu$ is much larger than predicted by pQCD and that $d \approx 1$ implies an especially long range interaction there.

Figure 3. Thermodynamic energy density ($\epsilon/T^4$ top curves left), pressure ($3p/T^4$ lower curves left) and speed of sound squared (right) from lattice QCD (2 flavor $12^3 \times 6$) from the MILC collab [8]. The curves are zero quark mass extrapolations. Note the rapid reduction of the pressure and speed of sound as $T_c$ is approached from above.
predictions of QCD is one of the primary goals of the heavy ion experimental program at Brookhaven and CERN. In the following sections, I review first how deep into the QGP phase RHIC may be able to reach, and then discuss several signatures that will be used to test the QCD predictions in such experiments.

2. Initial conditions in A+A

In order to see the partonic Yukawa phase, we must first heat the vacuum to about 100 times the energy density of nuclear matter. At SPS energies the low $p_\perp$ physics makes it impossible to predict the initial conditions. However, with increasing energy pQCD begins to provide increasingly more reliable theoretical basis for predicting those initial conditions. For highly boosted nuclei with $E_{cm} > 100m_N$, time dilation effectively freezes out the quantum chromo fluctuations inside the nuclei while the two pass through each other. Au beams at collider energies can be thought of as well collimated, ultra dense beams of partons. This (chromo Weizsacker-Williams\cite{9}) gluon cloud contains a very large number,

$$G_A(x,p_\perp^2) \sim A/x^{1+\delta},$$

of almost on-shell collinear gluons with longitudinal momentum fraction $x = p_0/E_{cm} \ll 1$. As the clouds pass through each other, many of the (virtual) partons scatter and form a dominantly gluonic plasma on a very fast time scale $1/p_\perp \ll 1$ fm/c. The number of gluons pairs (mini-jets) extracted from the nuclei by this mechanism at rapidities $y_i$ and transverse momentum $\pm k_\perp$ can be calculated in pQCD from the well known expression \cite{10,13}

$$\frac{dN_{AB \rightarrow ggX}}{dy_1dy_2dk_\perp^2} = K x_1 G_A(x_1,k_\perp^2) x_2 G_B(x_2,k_\perp^2) \frac{d\sigma_{gg \rightarrow gg}}{dk_\perp^2} T_{AB}(\vec{b}),$$

(7)

where $x_1 = x_+ (\exp(y_1) + \exp(y_2))$ and $x_2 = x_+ (\exp(-y_1) + \exp(-y_2))$, with $x_+ = k_\perp/\sqrt{s}$, and where the pQCD $gg \rightarrow gg$ cross section for scattering with $t = -k_\perp^2(1 + \exp(y_2 - y_1))$ and $y_2 - y_1 = y$ is given by

$$\frac{d\sigma_{gg}}{dt} = \frac{9}{8} \frac{4\pi\alpha^2}{k_\perp^4} (1 + e^y + e^{-y})^3 \left( e^{y/2} + e^{-y/2} \right)^6.$$  

(8)

For atomic numbers, $A \gg 1$, the geometrical amplification, $T_{AB}(\vec{b}) \sim 30/\text{mb}$, enhances by orders of magnitude the gluon density relative to $pp$. The factor $K \sim 2$ approximates next to leading order contributions.

For symmetric systems, $A + A$, with $G_A \approx AG$, the inclusive gluon jet production cross section is obtained by integrating over $y_2$ with $y_1 = y$ and $k_\perp$ fixed. To about 50% accuracy, the single inclusive gluon rapidity density in central collisions is approximately \cite{13}

$$\frac{dN}{dy dt} \approx \frac{A^2}{\pi R^2} 2N_g(x_\perp,t) x_\perp G(x_\perp,t) \frac{d\sigma_{gg}}{dt} \propto A^{4/3}$$

(9)

where $N_g(x_\perp,t) = \int^{x_\perp}_0 dx G(x,t)$. This copious mini-jet production mechanism is believed to the dominant source of gluon plasma production at RHIC and higher energies.

Recent upper bound estimates of the total gluon rapidity density as a function of the CM energy from EKRT\cite{14} are shown in Fig.\(\text{(4)}. The differential yields are integrated down to a transverse momentum scale $p_0 \sim 1 - 2$ GeV. This scale separates the “soft” nonperturbative beam jet fragmentation domain from the calculable perturbative one.
Figure 4. Mini-jet initial rapidity density of gluons, \( N_i \), produced in central A+A collisions as a function of the pQCD cutoff, \( p_0 \), from \[10\]. The dashed “saturation” curve corresponds to \( (p_0^2 R^2) \). \( (N_i \) is dimensionless; \( [mb] \) is a typo). The magnitude of possible hydrodynamic transverse energy loss due to longitudinal work in an ideal \( p = \epsilon/3 \) quark gluon plasma is shown by the difference between \( E_{T_i} \) and \( E_{T_f} \) in the right panel.

above. The curve marked saturation\[12\] is an upper bound marking the point where the transverse gluon density of mini-jets becomes so high that the newly liberated gluons completely fill the nuclear area, i.e., \( dN/dy \approx \frac{p_0^2 R^2}{\tau_0} \). A current hypothesis is that at that point, higher order gluon absorption may limit the further increase of the gluon number. At RHIC energies EKRT estimates yield up 1500 gluons per unit rapidity. Our conservative estimates together with X.N. Wang\[14,15\] only gives \(< 500 \) gluons per unit rapidity when initial and final state radiation is also taken into account. This lower number is obtained with a fixed \( p_0 = 2 \) GeV, that we found to be necessary in order to reproduce \( pp \), \( p\bar{p} \) and lower energy \( BA \) data using the HIJING event generator\[14\].

The initial energy density reached in such collisions can be estimated using the Bjorken formula

\[
\epsilon(\tau_0) \approx \frac{1}{\pi R^2 \tau_0} \frac{dE_T}{dy}
\]  

For \( p_0 \approx 1 - 2 \) GeV, \( dE_T/dy \approx 400 - 2000 \) GeV, and therefore \( \epsilon(\tau_0 \sim 0.5 \) fm/c) \( > 10 \) GeV/fm\(^3\) should be easily reached at RHIC, well inside the the deconfinement phase of QCD. At SPS energies, on the other hand, our estimates indicate that nuclear collisions may just reach the transition region and cool below it in a very short time.

Recently CERN issued a press release\[16\] claiming that “We now have evidence of a new state of matter where quarks and gluons are not confined.” As discussed below, I disagree with the claim that deconfinement has been observed. What I find compelling is that some form of matter, much denser than ever studied before, was created. Inferences about the role of quark and gluon degrees of freedom cannot be drawn because at the relatively low momentum scales accessible at SPS energies, the quarks and gluon degrees of freedom in the dense matter are mostly not resolvable. Even at the highest transverse momentum measured, the pion spectra were shown to be very sensitive to nonperturbative model assumptions regarding the role of intrinsic momenta and soft initial state interactions\[17\]. The dynamics of the non-perturbative beam jet fragmentation and hadronic final state interactions simply cannot be disentangled at SPS energies\[18\]. While there are many
Figure 5. Charged particle rapidity density per participating baryon pair from [20] is shown versus the cm energy per baryon. The PHOBOS data [19] (filled triangles) for the 6% most central Au+Au are compared to pp and p̅p data (open symbols) and the NA49 Pb + Pb (central 5%) data HIJING1.35 (solid) and EKRT (dot-dashed) predictions are also shown. HIJING upper (lower) solid curves are without (with) jet quenching. The right panel shows the characteristic difference between the predicted dependence of the charged multiplicity on the number of participant baryons for HIJING (solid) and EKRT (dot-dashed).

interesting signatures showing that dense matter was formed at the SPS (through the non-linear in dependence of several observables on multiplicity or $A$), the bottom line is that those data have said nothing about whether the QCD predictions in Figs 1-3 are correct or not. We need higher resolution, i.e. energy, to see the quarks and gluons in action in the new Yukawa phase.

3. First RHIC Data on Multiplicity

The first published data [19] from RHIC are shown in Fig. 5. Amazingly, both HIJING and EKRT predictions are seen to be consistent with the central multiplicity density data in Au+Au at $\sqrt{s} = 65, 130$ AGeV as measured by PHOBOS [19]. However, as emphasized in [20] the scaling as a function of centrality will soon be able to differentiate between these two very different pQCD based estimates of RHIC initial conditions. In addition, the transverse energy systematics will provide an independent critical test of the hypothetical saturation picture [22] as we discuss below. While it is obviously premature to draw any physics conclusion from these first data, the consistency of the increased activity per baryon predicted due to the onset of mini-jet activity beyond $\sqrt{s} > 100$ AGeV is
reasuring. A critical test of the pQCD framework used above will be the upcoming RHIC data at $\sqrt{s} = 200$ AGeV. HIJING predicts sensitivity to jet quenching at that energy.

4. Barometric Measures of Collective Dynamics

The simplest global barometer of collectivity in nuclear reactions is the $A$ and energy dependence of the transverse energy and charged particle rapidity density. At RHIC energies, HIJING predicts that initial transverse energy density in central $A + A$ collisions scales nonlinearly with $A$

$$\frac{dE_\perp}{dy} \approx 1 \text{ GeV } A^{1.3} \left( 1 + \log\frac{\sqrt{s}}{200} \right)$$

(11)

This leads to about 0.6 TeV per unit rapidity at the present $\sqrt{s} = 130$ AGeV energy. In EKRT\cite{10}, on the other hand, the gluon saturation hypothesis predicts an approximate linear $A$ dependence of $E_\perp \sim A^{1.04}$, and a value several times that of HIJING. In Fig.(4), the initial gluon density however grows less than linear $A^{0.92}$ in that model. The $A^{1.3}$ scaling of HIJING with the fixed mini-jet scale $p_0 = 2$ GeV scale. This results from the combined increase of the number of binary interactions as well as the increased fraction of the energy that originated from the mini-jets. The initial energy density in HIJING varies approximately as

$$\epsilon_0 \approx 0.6 \text{ GeV/fm}^3 A^{0.63} \left( 1 + \log\frac{\sqrt{s}}{200} \right).$$

(12)

EKRT\cite{10} on the other hand predict $\epsilon(\tau = 1/p_{sat}) \approx 0.1A^{0.5}s^{0.38}$.

If local equilibrium is achieved and maintained, then hydrodynamics predicts that longitudinal boost invariant expansion together with $pdV$ work done by the plasma pushing matter down the beam pipe will cool the plasma and convert some its random transverse energy into collective longitudinal kinetic energy. For an equation of state, $p = c_s^2 \epsilon$, this cooling and expansion causes the energy density to decrease with proper time as

$$\epsilon(\tau) = \epsilon(\tau_0) \left( \frac{\tau_0}{\tau} \right)^{1+c_s^2}$$

(13)

Entropy conservation leads to a conservation of $\tau n(\tau)$, where $n$ is the proper parton density. At least if the matter is initially deep in the plasma phase, then (as seen in Fig.3) longitudinal will be done with $c_s^2 \approx 1/3$. Consequently, the transverse energy per particle should decrease by a factor 2-3 before freeze-out as\cite{21}

$$e_\perp(\tau) = \frac{dE_\perp}{dN} = e_\perp(\tau_0) \left( \frac{\tau_0}{\tau} \right)^{c_s^2}$$

(14)

However, dissipative effects due to rapid expansion and finite mean free paths reduce considerably the effective pressure\cite{22}. For the Bjorken expansion, the relaxation time, $\tau_c = 1/(\sigma T n) \propto \tau$, then increases with time as $n$ decreases. Numerical solution of 3+1 D transport equations with pQCD cross sections, $\sigma_T \sim 2$ mb, indicate that dissipation reduces the transverse energy loss for HIJING initial conditions rather significantly (see detailed comparisons in\cite{23,24}).
One of the important experimental observations at the SPS is that at $\sqrt{s} \sim 20$ AGeV, $dE_{\perp}/dy$ as well as $dN/dy$ scale nearly linearly with wounded nucleon or participant number ($\sim A^{1.07}$) [25]. Simple Glauber wounded nucleon models reproduce very well the nearly linear correlation between $E_{\perp}$ and the veto calorimeter (spectator) energy observed in all experiments at SPS[26]. The physics implications of this scaling depends on what is assumed for the $A$ dependence of the initial conditions. One view [10] is that the initial $e_{\perp}$ scales in just the right way that after hydrodynamic expansion the final $E_T$ and $e_{\perp}$ always scale linearly with $A$. My view is that at the SPS the linear dependence arises from additive nature of soft beam fragmentation together with the absence of $pdV$ work at early times. If the QGP transition region is just barely reached, as I believe, then the softness of the QCD equation of state with $c_s^2 \ll 1/3$ seen in Fig.(3) and dissipation conspire to prevent the dense matter from performing longitudinal work. However, it is impossible to tell from the data whether this observed null effect is due to a low pressure Hagedorn resonance gas of hadrons or to a low pressure lazy “plasma” with $c_s \ll 1$. At RHIC energies the plasma starts deep in the $c_s \approx 1/\sqrt{3}$ Yukawa regime, and longitudinal work should become observable in the $E_T$ systematics.

5. Transverse Collective Flow

An entirely different measure of barometric collectivity is afforded by the study of the triple differential distributions, $dN/dyd^2p_{\perp}$. Already at sub-luminal Bevalac energies (< 1 AGeV), azimuthally asymmetric collective directed and elliptic flow were discovered long ago. For non central collisions, $b \neq 0$, the asymmetric transverse coordinate profile of the reaction region leads to different gradients of the pressure as a function of the azimuthal angle relative to beam axis. This leads to a “bounce” off of projectile and target fragments in the reaction plane and to azimuthally asymmetric transverse momentum dependence of particles with short mean free at mid rapidities. This phenomenon has now been observed at both AGS and SPS energies as well [27]. It has been predicted to be there also at RHIC and LHC energies [28]. In Fig.(3) the first STAR data [29] on the centrality dependence of the second Fourier component, $v_2$, of the azimuthal flow patterns is shown:

$$\frac{dN}{dyd^2p_{\perp}} \propto 1 + 2v_1\cos(\phi - \phi_R) + 2v_2\cos(2(\phi - \phi_R)) + \cdots . \quad (15)$$

These data show that the azimuthal asymmetry is about twice as large at RHIC than at the SPS. Furthermore, the differential azimuthal flow, $v_2(p_{\perp})$, rises linearly with $p_{\perp}$ up to 2 GeV/c.

The important question is how this type of barometer could serve to help search for evidence of the QCD transition. In [30] the idea was proposed that one could use $v_2$ systematics to search for the predicted softening of the QCD equation of state. Hydrodynamics calculations lead to a factor of two smaller $v_2$ for an equation of state with a soft critical point as in Fig 3 versus one in which the speed of sound remains $1/\sqrt{3}$. Searches for anomalous $v_2$ dependence as well as $v_1$ are underway [31]. As with the global barometer, dissipation can of course also simulate a soft equation of state. In ref[32] we studied the dependence of $v_2$ on the transport parton cross section for RHIC conditions. The results are rather sensitive to the cross sections and initial densities as shown in Fig(6).
Figure 6. (left) First STAR data [29] on the centrality dependence of the azimuthal flow moment $v_2$ for $Au + Au$ at 130 AGeV. Estimates of the corresponding spatial asymmetry are indicated by boxes. (Right) Predicted dependence of $v_2$ at mid rapidity $Au + Au$ at RHIC from parton cascade calculations ZPC[32] with HIJING initial conditions for $\sigma_{gg} = 3$ and 10 mb. (Note $v_2'$ is the $p_\perp$ weighed form of $v_2$)

For HIJING initial conditions, dissipation leads to a significant reduction of $v_2$ relative to hydrodynamics.

6. The QGP Stall and Interferometry

Hadron interferometry has been developed over the last several decades in heavy ion collisions into a precise tool to image the space-time region of the decoupling 4-volume. In [33] it was proposed that a possible signature of the QGP phase transition would be a time delay associated with very slow hadronization. The plasma "burns" into hadronic ashes along deflagration front that moves very slowly if the entropy drop across the transition is large[34]. In [35] we calculated the 2+1D evolution of a Bjorken cylinder with time and transverse coordinate using hydrodynamics with different equations of state. For an initial energy density $\sim 20$ GeV/fm$^3$ a time delay up to a factor of two was found even for a continuous transition as long as the entropy jump was $\sim 10$. This QGP "stall" of the transverse expansion is due to the small speed of sound in the mixed phase. It should be readily observable in pion interferometry by comparing $R_{out}$ and $R_{side}$ radii.

The time delay is a robust generic signature of a rapid cross over transition of the entropy density. In particular the "stall" is expected even for a smooth cross-over transition as long as the width $\Delta T/T_c < 0.1$. However, its magnitude also depends strongly on the entropy drop across that region. Unfortunately, as noted before lattice QCD has not yet been able to resolved the hadronic world below $T_c$ due to technical problems. If the
entropy jump in nature is less than a factor of three, then the stall signature would be much more difficult to observe.

High statistics measurements of pion and kaon interferometry searches for time delay at AGS and SPS have come up empty handed thus far. No time delay has ever been observed in any nuclear reactions thus far. This could be due to (a) the absence of a large rapid entropy drop in real QCD, or (b) to unfavorable kinematic conditions at AGS and SPS energies. From the hydrodynamic calculations in [35], it was found that a large time delay signal requires a favorable initial condition with initial temperature \( \sim 2T_c \).

For too high an initial temperature (as at LHC) the large transverse collective flow that develops prevents a stall from happening. For too low a temperature (as at SPS) the time delay is suppressed because the system spends too little time in the mixed phase as a result of longitudinal expansion. Therefore, RHIC is the most likely energy regime where a QGP stall may be observable. Recent work [36] has also shown that high \( p_\perp \) kaon interferometry is an especially sensitivity probe of time delay. If observed, a time delay signature would be one of the most powerful indicators of novel collective behavior that can only arise if the produced matter has an anomalous (softest point) equation of state.

7. J/\( \psi \) Suppression

In 1986, Matsui and Satz proposed an intriguing direct measure of the transmutation of the \( q\bar{q} \) forces shown in Fig. 1. The idea was that \( J/\psi \) can form in the vacuum because the confining Luscher potential can bind a \( c\bar{c} \) pair into that vector meson. If that pair were placed in a hot medium in which the potential is screened, then above the temperature where the screening length is smaller than the \( J/\psi \) radius, the \( c\bar{c} \) would become unbound. The charm quarks would then emerge from the reaction region as an open charm \( D\bar{D} \) pair.

\( J/\psi \) suppression was soon seen in 1987 in \( O+U \) reactions by NA38. Since then this smoking gun has (unfortunately) never stopped smoking! \( J/\psi \) suppression seems to be as ubiquitous as the Yukawa potential. It was subsequently observed in \( p+A \) reactions as well. High mass Drell-Yan pairs, on the other hand, formed via \( q\bar{q} \to \ell\bar{\ell} \) was observed to scale perfectly linearly with the number of binary collisions. This is because lepton pairs suffer no final state interactions and the quark initial state (Cronin) interactions are invisible in \( p_\perp \) integrated DY yields.

The “normal, conventional” \( J/\psi \) suppression leading to a \((AB)^{0.9} \) dependence, is presumed to be due to a mechanism that is independent of QGP production. In Pb+Pb, NA50 observed an excess 25% suppression of \( J/\psi \) above that normal expectation. NA50 therefore claimed that this enhanced suppression is finally the sought after smoking gun of QGP formation [37].

While the deviation from the empirical \((AB)^{0.9} \) scaling is convincing, the dynamical origin of the effect is far from settled in my opinion. As shown by the several curves in Fig. 4, a \( \sim 50\% \) drop in the \( J/\psi \) yield as a function of \( E_T \) is consistent with final state co-moving hadronic absorption. While the detailed wiggles are not reproduced, large theoretical uncertainties about several key dynamical ingredients preclude more precise comparisons at this time. It is important to emphasize that plasma scenario models suffer at least as large theoretical uncertainties as the so called “conventional hadronic” models.
Figure 7. “Evidence for deconfinement of quarks and gluons” claimed by NA50 in Ref. [37]. The curves on left show transport theory estimates for hadronic final state dissociation. The enhanced suppression relative to $\exp(-\sigma_N \rho_0 L)$ nuclear suppression observed in $p + A$ and $S + U$ is shown versus a very rough estimate of the unobserved initial energy density.

Key uncertain elements include (1) the proper QCD treatment of cold nuclear absorption responsible for the nonplasma $AB^{0.9}$ suppression observed in $p + A$, (2) the unknown hadronic $M + \psi \rightarrow D\bar{D}X$ reaction rates, and (3) the detailed density evolution, $\epsilon(\vec{x},t)$.

For example, only a schematic “octet model” of pre-hadronization $c\bar{c}$ exists at present to calculate the so called “normal” nuclear suppression. That this is model may not be sufficient to adequately account for that component can be seen in the work of Ref. [38]. The observed $J/\psi$ production cross section (ignoring final state interactions) may be expressed as

$$\sigma_{AB \rightarrow \psi X} = \int d\sigma_{AB \rightarrow c\bar{c}} F_{c\bar{c} \rightarrow \psi}(q^2)$$ \hspace{1cm} (16)$$

where $F$ is the formation probability of the $\psi$ from a $c\bar{c}$ that emerges from the cold nuclear target with an invariant mass $q^2 < 4M_D^2$. If the pre-resonance $c\bar{c}$ pair multiply scatters in the nucleus, then $q^2 \rightarrow q^2 + \delta q^2(\sigma \rho L)$ increases linearly with nuclear thickness as shown by the curve G in Fig. (8). This leads to an approximate exponential suppression that can account for the approximate Glauber nuclear absorption factor ansatz, $\exp(-\sigma_{eff} \rho L)$. This Gaussian model can thus account for the observed $(AB)^{0.9}$ scaling for light projectiles. However, it was shown in [38] that if the power law tails due to induced radiation in the medium are included (resulting from the multiple Rutherford rescattering of the color octet $c\bar{c}$), then an additional nonlinear suppression in the nuclear thickness $L$ could result (curve P). This is because induced gluon radiation provides a way to increase the...
Figure 8. Sensitivity of “pre-resonance” $J/\psi$ nuclear absorption to details of color octet models of $c\bar{c}$ interactions in a nuclear medium from Qiu et al [38]. The curve marked P for power law accounts for the anomalous absorption seen in $Pb + Pb$ and deviated from Gaussian Glauber-like G expectations.

The sensitivity of $J/\psi$ nuclear absorption to details of color octet models of $c\bar{c}$ interactions in a nuclear medium from Qiu et al [38]. The curve marked P for power law accounts for the anomalous absorption seen in $Pb + Pb$ and deviated from Gaussian Glauber-like G expectations.

The claim of anomalous suppression cannot therefore rest merely on generic enhanced suppression in $Pb$. It rests on the possible existence of singular “step-like” structure of the suppression pattern. The evidence for “steps” is however the weakest link experimentally because a rigorous $\chi^2$ test including the substantial systematic errors in the $E_T$ scale has yet to be performed. Nevertheless, if the “step-like” suppression pattern survives further experimental scrutiny, it would certainly be the most dramatic nonlinearity observed at SPS. Experimentally, the claims would carry considerable more weight if similar ”step-wise” patterns were observed in other systems, e.g. $Xe + Xe$ suitably shifted in $E_T$ due to the expected smaller energy densities achieved. I would put such an experiment highest on the SPS priority list. At RHIC, the striking prediction by H. Satz at QM99 [43], was that under RHIC conditions the higher energy density should lead to the same step-wise pattern in the lighter nuclear $Cu + Cu$ interactions. PHENIX will provide a definitive test of this prediction in a few years.

One significant inconsistency with the present plasma scenario interpretation is its failure to account for the observed $E_T$ dependence of the $J/\psi$ $p_{\perp}$ spectra in Fig.(9). Standard Glauber multiple collisions lead to a random walk in transverse momentum
Figure 9. The mean transverse momentum of the $J/\psi$ shows clear evidence of multiple scattering in the nuclear target as expected from Glauber theory. Even in $Pb + Pb$ the increase of the $p_T$ is understood from the same mechanism in contradiction to predictions based on the plasma scenario, where at high $E_T$, the surviving $\psi$ are expected to have been produced only near the nuclear surface regions.

that are expected enhance the $J/\psi$ transverse momentum as

$$\langle p_{T}^2 \rangle_{AB} = \langle p_{T}^2 \rangle_{pp} + \frac{L}{\lambda} \delta p_{T}^2.$$ (17)

This is found to hold experimentally in all reaction including $Pb + Pb$. In contrast, in the plasma scenario, only those $\psi$ are expected to be observed which are produced near the surface, where the effective nuclear depth $L$ is small. Thus, the prediction as shown in Fig(9) was that the $\langle p_{T}^2 \rangle$ should begin to DECREASE with increasing $E_T$. This was not observed.

8. The High $p_T$ Frontier

One of the most exiting new physics areas that will become accessible at RHIC is the study of high transverse momentum (short wavelength) jet probes. The rates of jet production and its fragmentation in the vacuum are well understood. The new jet physics at RHIC will be the study of partonic interactions at extreme densities through the phenomenon of jet quenching. Final state interactions of a jet in a dense QGP are expected to induce a large radiative energy loss. In fact, BDMPS discovered that non-Abelian energy loss is in fact non-linear as a function of the thickness of the medium.
Figure 10. Jet quenching at SPS vs RHIC from ref.[17] compared to WA98 data. At RHIC the power law tail extends far enough above the nonperturbative “noise” to make jet quenching observable (long dashed in right panel). At SPS the steep high $p_\perp$ tails are too sensitive multiple soft scatterings.

Tests of this and other aspects of non-Abelian multiple collision dynamical phenomena will soon be possible at RHIC[17].

At lower SPS energies, this physics is out of reach because nonperturbative multiparticle production physics (e.g., soft multiple $p_\perp$ kicks) dominates as shown in Fig.(10) from ref. [17]. HIJING happens to fit the WA98 data with or without jet quenching at SPS energies. While at the SPS no clean separation of soft and hard dynamics is kinematically possible, at RHIC energies, the high $p_\perp$ power law tails of the single inclusive distributions stick out far enough above the soft physics ”noise” that sensitivity to the form of the non-Abelian $dE/dx$ is expected as shown in Fig.(10). This problem is closely related also to the problem of pre-resonance $J/\psi$ absorption discussed previously[38]. Understanding jet quenching is also necessary to develop more accurate covariant parton transport theories[23,24].

9. Summary

The next Yukawa phase of QCD is awaiting discovery. The SPS data have provided many intriguing indirect hints that new physics is already operating near the transition region. Many claims and counter claims remain at SPS energies because both hadronic and partonic models have partial overlapping domains of validity. This “duality” is analogous to the problem of interpreting the R factor in $e^+e^-$ collisions below $\sqrt{s} < 10$ GeV. The
ratio of hadronic to leptonic production cross sections only reaches the magic 11/3 of pQCD above the threshold regions for $s\bar{s}, c\bar{c}, b\bar{b}$. Near the threshold regions conventional vector mesons models of nonperturbative hadronic physics provide an adequate “dual” description of the physics. The simpler continuum partonic description is only applicable far above that region. Similarly, reactions at SPS energies are near the threshold region where resonances become too broad due to multiple interactions. In nuclear reactions, the continuum parton QGP description can only be expected to apply far above that region as produced at collider energies.

At RHIC with a factor of ten increase in the initial energy density, only a continuum parton description of the initial conditions is tenable. The dense partonic matter so formed also will have much more time to evolve and produce collective signatures of its existence. Furthermore, a factor of ten smaller wavelength (jet) probes will finally allow experimentalists to resolve (i.e. see) individual quark and gluon degrees of freedom of that plasma. Direct observation of longitudinal work, transverse azimuthal collectivity, time delay, step-wise $J/\psi$ suppression in Cu+Cu, and jet quenching should, among other signatures, allow experimentalists to measure the properties of the partonic Yukawa phase at RHIC.

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REFERENCES

1. H. Yukawa, Proc. Phys.-Math. Soc. Jpn. Ser.3, 17 (1935) 48.
2. R. Machleidt, The Meson Theory of Nuclear Forces and Nuclear Structure, in Advances in Nuclear Physics, 19, eds. J.W. Negele and E. Vogt,(Plenum Press, New York, 1989), chap.2.
3. P. Debyee and E. Hückel, Z. Phys. 24 (1923) 185.
4. M. Lüscher, K. Symanzig, and P. Weise, NPB 173 (1980) 365.
5. M. Gao, Phys. Rev. D40 (1989) 2708.
6. O. Kaczmarek, F. Karsch, E. Laermann and M. Lutgemeier, [hep-lat/9907010].
7. A. K. Rebhan, Nucl. Phys. B430, 319 (1994).
8. C. Bernard et al. [MILC Collaboration], Phys. Rev. D55, 6861 (1997) [hep-lat/9612025].
9. L. McLerran and R. Venugopalan, Phys. Rev. D49, 2233 (1994) [hep-ph/9309289].
10. K. J. Eskola, K. Kajantie, P. V. Runskanen and K. Tuominen, Nucl. Phys. B570, 379 (2000) [hep-ph/9909456].
11. K. J. Eskola, K. Kajantie and J. Lindford, Nucl. Phys. B323, 37 (1989).
12. J. P. Blaizot and A. H. Mueller, Nucl. Phys. B289, 847 (1987).
13. M. Gyulassy and L. McLerran, Phys. Rev. C56, 2219 (1997) [nucl-th/9704034].
14. X. Wang and M. Gyulassy, Phys. Rev. D44, 3501 (1991).
15. X. Wang, Phys. Rept. 280, 287 (1997) [hep-ph/9605214].
16. CERN Press Release 2/8/00: http://press.web.cern.ch/Press/Releases00/
17. M. Gyulassy and P. Levai, Phys. Lett. B442, 1 (1998) [hep-ph/9807247].
18. S. A. Bass, M. Gyulassy, H. Stocker and W. Greiner, J. Phys. G G25, R1 (1999) [hep-ph/9810281].
19. B.B. Back et al., PHOBOS Collaboration, [hep-ex/0007036].
20. X. Wang and M. Gyulassy, [nucl-th/0008014].
21. G. C. Nayak, A. Dumitru, L. McLerran and W. Greiner, [hep-ph/0001202].
22. A. Dumitru and M. Gyulassy, [hep-ph/0006257].
23. M. Gyulassy, Y. Pang and B. Zhang, Prog. Theor. Phys. Suppl. 129 (1997) 21; Nucl. Phys. A626, 999 (1997) [nucl-th/9709023].
24. D. Mohnar, M. Gyulassy, e-Print Archive: [nucl-th/0005051].
25. H. Schlagheck [WA98 Collaboration], Nucl. Phys. A663&664, 725 (2000) [nucl-ex/9909003].
26. D. Kharzeev, C. Lourenco, M. Nardi and H. Satz, Z. Phys. C74, 307 (1997) [hep-ph/9612217].
27. M. Aggarwal et al WA98, Nucl. Phys. A610 (1996) 200c.
28. S. A. Voloshin and A. M. Poskanzer, Phys. Lett. B474, 27 (2000)
29. K. H. Ackermann et al. [STAR Collaboration], [nucl-ex/0009011].
30. J.-Y. Ollitrault, Phys. Rev. D 46 (1992) 229.
31. J.-Y. Ollitrault, Nucl. Phys. A590 (1995) 561c; W. Reisdorf and H.G. Ritter, Annu. Rev. Nucl. Part. Sci. 47 (1997) 663.
32. B. Zhang, M. Gyulassy and C. M. Ko, Phys. Lett. B455, 45 (1999) [nucl-th/9902016].
33. S. Pratt, Phys. Rev. C 49 (1994) 2722, Phys. Rev. D 33 (1986) 1314.
34. L. Van Hove, Z. Phys. C 21 (1983) 93.
35. D. H. Rischke and M. Gyulassy, Nucl. Phys. A608, 479 (1996) [nucl-th/9606039].
36. S. Bernard, D. H. Rischke, J. A. Maruhn and W. Greiner, Nucl. Phys. A625, 473 (1997) [nucl-th/9703017].
37. M. C. Abreu et al. [NA50 Collaboration], Phys. Lett. B477, 28 (2000).
38. J. Qiu, J. P. Vary and X. Zhang, [hep-ph/9809442].
39. O. Drapier, Thesis, Universit Lyon-I, 1998.
40. J. L. Nagle and M. J. Bennett, Phys. Lett. B465, 21 (1999) [nucl-th/9907004].
41. S. Gavin and M. Gyulassy, Phys. Lett. B214, 241 (1988).
42. D. Kharzeev, M. Nardi and H. Satz, Phys. Lett. B405, 14 (1997) [hep-ph/9702273].
43. S. A. Bass et al., Nucl. Phys. A661, 205 (1999) [nucl-th/9907090].
44. M. Gyulassy, M. Plümer, M.H. Thoma and X.-N. Wang, Nucl. Phys. A 538 (1992) 37c; X.-N. Wang and M. Gyulassy, Phys. Rev. Lett. 68 (1992) 1480.
45. M. Gyulassy and X.-N. Wang, Nucl. Phys. B 420 (1994) 583.
46. R. Baier, Yu.L. Dokshitzer, A.H. Mueller and D. Schiff, Nucl. Phys. B 531 (1998) 403 and refs therein.
47. M. Gyulassy, P. Lévi, I. Vitev, Nucl. Phys. B 571 (2000) 197; e-Print nucl-th/0005032 ; nucl-th/0006010.