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The changing scenario of the atomic nucleus—from nucleons and mesons to quarks and gluons

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Abstract. The general characteristics of the transition from hadronic matter of nucleons, three quark bags, mesons of quark antiquark pairs to quark gluon plasma is discussed. The phenomenological approach essentially guided by the MIT bag model and general thermodynamic criteria of first-order phase transition is elaborated. The more realistic calculations using the QCD lattice renormalization quark are touched upon. Possible signals of quark-gluon plasma are discussed. The central issue of deciphering plasma signals from the signals of hot hadronic matter is discussed in detail.

The signals of the quark-gluon plasma, a subject of considerable interest in contemporary literature are focussed only on (i) dileptons (ii) photon photon pairs and (iii) J/Ψ suppression (with special emphasis on CERN experiments). The lingering shadow of "EMC" effect is also mentioned.

Relics of the very early universe microseconds after the big bang in today's universe (~ 15 billion years later) are discussed. Finally, the outlook of this very exciting field is presented, a purely personal viewpoint, generalized eventually to poetic signals of the creation of the universe.

Keywords. Atomic nucleus; hadrons; chromoplasma quarks; gluons; quantum chromodynamics; nucleus-nucleus collision; heavy ion collision; dileptons; early universe relics.

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Preamble

Some years ago, as a graduate student in England, struggling through the finer details of nuclear structure physics, the present author came across a reference—a rather esoteric transformation from particles to holes. The method was powerful, I was told. The reference turned out to be a S P Pandya—a distinctly Indian name.

I was naturally impressed, if not proud. I had, in my mind's eye, image of a rather serious theoretical physicist, somber, aloof and reserved. Imagine my surprise therefore when I first met Sudhir Pandya, antithesis of my mind's image, eminently approachable, lucid, warm and above all humane. I salute you sir; in today's world of Indian science you are a rare person; the Pandya School you have so tenaciously and tenderly developed is probably the best tribute to you. I am sure, you shall guide us, inspire us and motivate us for years to come in your inimitable tradition.

1. Introduction

In this "colour singlet" world of hadrons (neutrons, protons and mesons are referred to as hadrons) that we live in, quark confinement is assured by colourlessness—thus, to create a chromoplasma of quarks, we imply deconfinement of quarks and gluons within an extended volume.

The associated symmetry is the chiral symmetry (Satz 1984). In an environment of very high temperature, the chiral symmetry is restored. In normal circumstances, chiral symmetry is just broken, giving rise to the finite mass of the Goldstone bosons, the π-mesons. As the chiral symmetry is restored the pions lose their mass.

One can therefore think rather broadly two sets of temperature, a chiral symmetry restoration temperature T_{Ch} and a deconfinement temperature T_d. The conventional consensus dictates (Cleymans *et al* 1986) that the two temperatures are identical. Thus, with deconfinement, the quarks become massless since the chiral symmetry is already restored.

There already exists a number of analogies in other branches of physics intimately connected with the restoration of chiral symmetry. In a ferromagnet, for example, rotational symmetry is broken in the normal world; rotational symmetry, in a similar fashion, is restored at a high temperature for a ferromagnet.

Similarly the analogy between colour charge and electric charge, up to a point is also rather instructive. As we go from the colour singlet state of hadronic matter of a normal world to a chromoplasma of quarks gluons (QGP) we see an equivalent "Mott transition" of the QCD sector; like in QED, there is a phase transition from an insulator of no electric conductivity to a metal which is an electric conductor.

What are the very essential criteria of a so-called phase transition from a hadronic world of QGP? Indeed what are the basic properties of this phase transition?

Clearly, the energy density "scale" indicates an elementary but instructive requirements for such a phase transition—normal matter corresponds to an energy density $\varepsilon_N = m_N \rho_0 = 0{\cdot}15\,\mathrm{GeV/fm^3}$ for a typical hadronic density of $\rho_0 = 0{\cdot}145\,\mathrm{fm}$; where for the hadron energy density we have to deal with

$$\varepsilon_p \cong \left(\frac{4\pi}{3} r_p^3\right)^{-1} \sim 1{\cdot}5\,\mathrm{GeV/fm^3}$$

for a typical radius $r_p \sim 0{\cdot}8\,\mathrm{fm}$, guided by (say) MIT bag model (Thomas 1988) Therefore, a phase transition from a dilute hadron gas ($\rho_0 \sim 0{\cdot}145\,\mathrm{fm^3}$; one nucleon per $7\,\mathrm{fm^3}$!) to a chromoplasma of QGP one has to deposit an extra energy density approximately of the order of a factor five. Curiously, even this rather naive approach leaves adequate room for considerable uncertainty. For example what does one choose for the hadron bag radius r_p? If one looks upon a nucleon rather like a pea pod in a soup of meson it is almost certain that the radius of the baryon will change; if one follows the tradition and the wisdom of the East (Long Island, Stony Brook) the radius r_p will decrease, resulting a substantial increase in ε_p.

Indeed, the previous discussion is a hint related to the problems associated with the studies of the phase transition from the hadron gas to QGP; in the following therefore, we shall perform a simple but "colourful" nevertheless, tour de force, on this rapidly "expanding" field.

2. Phase transition from hadron gas (HG) to a chromoplasma quarks and gluons (QGP)

As an introduction to the state-of-the-art let us consider a hadron world of the typical thermodynamic variables T^h, μ^h and P^h, the temperature, baryon chemical potential, ensuring baryon conservation and pressure respectively; similarly for the QGP, we have T^q, μ^q and P^q. Let us further consider two extreme cases (corresponding to two furthest points) $\mu^h \to 0$ but $T^h \neq 0$ and $\mu^h \neq 0$ but $T^h \to 0$, the former corresponding to a no-baryon region and a rather high temperature and the latter corresponding to a rather dense baryonic regime but a cold system. For collision of nuclei at ultra-relativistic energies, to be discussed later, one can envisage a central rapidity regime for the first case whereas the fragmentation regime corresponds to the relatively cold but high baryon density regions. For the early universe saga to be discussed later, the baryons were initially absent but the temperature was extremely high for an almost infinite deposition of energy density.

For $\mu^h \to 0$, considering massless pions only (for simplicity)

$$P^h(T, \mu^h \to 0) = (\pi^2/90)\cdot 3 \cdot T^4 = \tfrac{1}{3}[\varepsilon_H(T, \mu^h \to 0)] \tag{1}$$

and for QGP

$$P^q(T, \mu^q \to 0) = \left(\frac{\pi^2}{90}\right)\left[\overbrace{2 \times 8}^{\text{Gluons}} + \overbrace{\tfrac{7}{8} \times 2 \times 2 \times 2 \times 3}^{\text{Quarks}}\right] T^4 - B, \tag{2}$$

where the "bag" pressure B ensures phenomenologically, confinement of quarks, usually determined experimentally in the spirit of the MIT bag model (Thomas 1988) (see later, however, the interpretation from the QCD standpoint). We have chosen eight colour degrees of freedom for gluons and three colours for the quarks, each corresponding to two spin states. The equation of state for QGP is

$$P_Q = \tfrac{1}{3}(\varepsilon - 4B). \tag{3}$$

The important point to realize is the rapid change of total degrees of freedom as one transits from HG system to QGP, an increase $g^h \cong 3$ to $g^q \approx 42$ by a factor of fourteen. Clearly, the quenching of degrees of freedom and the associated entropy is a clear indication of a phase transition to a more ordered system with QGP $\to$ HG, the drop in the entropy being used up to give a finite mass attributed to the hadrons ($m_{\text{proton}} \sim 1$ GeV). The primordial saga of the nuclear synthesis, enacted in the first three minutes, of the creation of the universe is intimately connected with the quenching of the entropy (see later). Now, using the Gibbs criteria and a suitable bag parameter value of $B^{1/4} \sim 190$ MeV, we get for the critical temperature of transition

$$P^h(T_c, 0) = P^q(T_c, 0).$$

Therefore

$$\begin{aligned} T_c &= [(45/17\pi^2)B]^{1/4} \\ &\approx 0{\cdot}72 B^{1/4} \\ &= 140\,\text{MeV} \quad \text{for } B^{1/4} = 190\,\text{MeV}. \end{aligned} \tag{4}$$

Similarly, for a cold but strongly interacting system of density packed baryons we get

$$n_c \approx 0{\cdot}85 B^{3/4} \sim (4 \sim 5) n_0 : T \to 0, \tag{5}$$

where n_c is the critical density for a cold baryon dense regime in terms of the normal

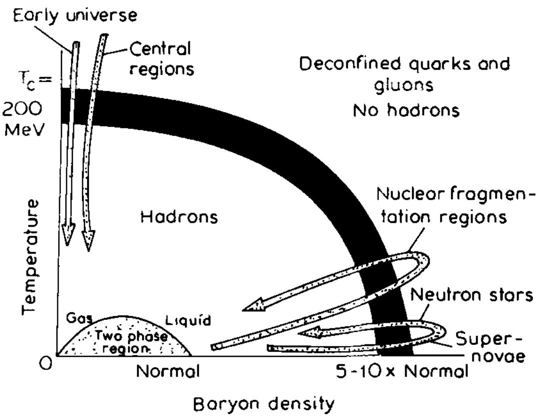

Figure 1. Baryon density.

nuclear density n_0. The last scenario corresponds to neutron star, the dense pulsar, remnant of the supernova explosion (see later).

The universal phase diagram in an approximate form is shown in figure 1.

2.1. *QCD Sector*

The elementary properties of the phase transition from HG to QGP in the sprit of the phenomenological model guided by the "bag" were discussed so far; we shall now briefly disucss the present state-of-the-art from the QCD sector using the much used lattice renormalization group (Kogut *et al* 1983).

The Tsars of the QCD kingdom, Helmut Satz (Celik *et al* 1983a, b) and John Kogut (Kogt 1984) and other workers (Fukugita and Ukawa 1986) have given us a rather impressive, if not glamorous, results about the phase transition. They have managed to enter or at least peep into the perturbative vacuum inside a baryon, the eye glass evidently is firmly stationed in the normal vacuum (non perturbative).

Much of quantum field theory can be understood in terms of the Feynman path integral formula for a thermodynamic partition function

$$Z = \int [\mathrm{d}A_\mu][\mathrm{d}\psi][\mathrm{d}\bar{\psi}]\exp\{-S(A_\mu, \psi, \bar{\psi}\},$$

where for the QCD sector S is the Euclidean action written in terms of the glue field A_μ and the quark fields ψ and $\bar{\psi}$, the partition function is the convolution integral over all space-time configuration of the fields. To ensure the intermediate results to be finite one introduces a space-time lattice with sites and an associated lattice spacing, a. Introducing the inverse temperature $\beta \approx T^{-1}$, and asserting $\beta = N_t a$ the awful mess and the hardest struggle are to evaluate the expectation value of an operator θ in the prescribed thermodynamic ensemble, the general definition being

$$\langle \theta \rangle = (Z^{-1}) \int [\mathrm{d}A_\mu][\mathrm{d}\psi][\mathrm{d}\bar{\psi}]\exp(-S\theta),$$

where θ is some function of the field variables, a simple depiction of the space-time lattice being given in figure 2.

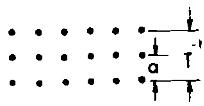

Figure 2. The space-time lattice.

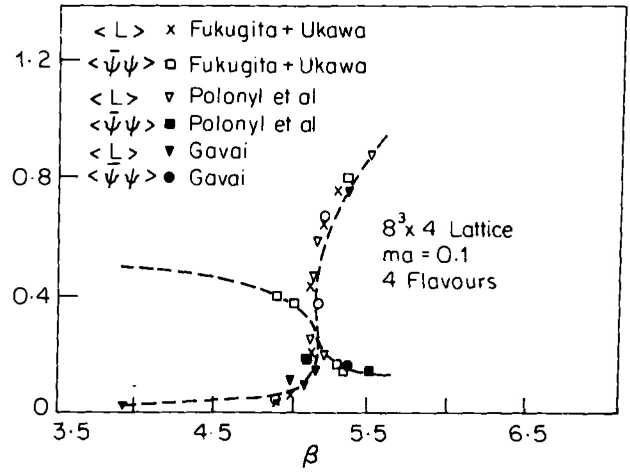

Figure 3. $\langle \Psi\Psi \rangle$ and $\langle L \rangle$ as a function of the coupling. All the data are for 4 flavour QCD with a quark mass 0·1 (Fukugita and Ukawa 1986).

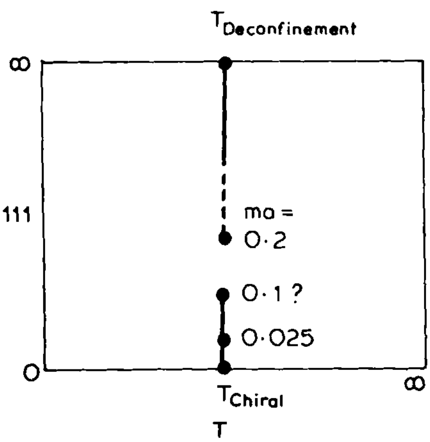

Figure 4. Schematic phase diagram of 3 flavour QCD on an $8^3 \times 4$ lattice.

The detailed technology is not within the scope of this paper as extensive reviews are available (Celik *et al* 1983a, b; Kogut 1984) and we shall quote here only the essential results.

For a pure Yang-Mills field of gluons for finite temperature massless quarks, the task appears to be relatively easy and the order of the phase transition is clearly defined viz first-order. The problem appears with fermions, the quarks and the gluons—the action S, defined above turns complex. Much work has recently been done in this field, the synopsis of which is presented in figures 3 and 4—for a finite and reasonable value for constituent mass of the light quarks, u and d, the scenario still remains clouded. (Figure 3)

2.2. *Nucleus-nucleus collision*

It is generally accepted now that ultra-relativistic heavy ion collision (typically, $50\,\mathrm{GeV} \leqslant E \leqslant 1\,\mathrm{TeV}$) may provide a scenario where hadronic matter is put under extreme conditions of pressure and temperature/density. If the energy density deposited as a result of collision is beyond a critical value, the new phase of QGP can replace the hadrons.

The central question therefore is how to decipher experimental signatures of such a QGP? The QGP, if at all formed, because of extreme pressure exerting outward will induce expansion and thus cooling of the system. The initial QGP, due to cooling will now make a phase transition to hadron bubbles. The mixed phase of QGP and HG after further expansion will freeze out to a typical temperature T_f (typically pion mass $\sim 140\,\mathrm{MeV}$) when the entire system will consist of non-interacting hadrons. The problem therefore is to seek out signals of the initial plasma state from the space-time integrated production rates of any species one cares to measure. The detectors can pick up only the convoluted integral over space-time $\int \mathrm{d}^4x\, f(x,\tau)\ldots$.

What can be the guiding scenario? What set of equations can one use to predict the space-time evolution of the thermodynamic variables already mentioned?

One important assumption is that, at least for the very central collisions, the quarks and gluons thermalize rather quickly within a time period of say $\tau_0 \sim 1\,\mathrm{fm}/c$. With this scenario, one can then apply the standard Landau hydrodynamics (Landau and Lifshiftz 1982). This automatically leads to well-defined global observables, characterizing the state of the matter, produced during the collision.

An obvious simplification occurs if one restricts oneself to central collisions of identical/or near-identical nuclei—cylindrical symmetry is the guiding symmetry under those boundary conditions and the hydrodynamical equation assumes a simple form.

The initial conditions which we shall consider are invariant under Lorentz boost in the longitudinal direction (Blaizot and Ollitrault 1986). The hydrodynamical evolution can then be understood in terms of the superposition of two elementary motions: a uniform longitudinal expansion, accompanied by rapid cooling and a transverse rarefaction. As the system cools, both as a result of longitudinal cooling and transverse expansion, the initial fireball will eventually undergo a transition from the QGP to the hadronic phase. In this paper we shall tacitly assume a thermodynamic equilibrium, in which QGP goes through a smooth phase transition with the formation of an intermediate mixed phase. More violent scenarios, precipitated by super cooling (Guylasay *et al* 1984; Kajantie and Kurki-Suonio 1986) have also been considered. We shall not however consider these cases here.

2.3. *Central heavy ion collisions*

We shall consider central collisions of two identical ions to demonstrate how ideas developed above can be used to compute production rates of signals from a hot fireball, formed by collisions, expanding through space and time.

Changing the variables to $\tau^2 = t^2 - z^2$ (Blaizot and Ollitrault 1986) and rapidity

$$y = \tfrac{1}{2}\ln[(t+z)/(t-z)],$$

the initial conditions of the hydrodynamic evolution can be specified in terms of a predetermined initial time and temperature. One can immediately calculate the entropy per unit rapidity in terms of initial entropy

$$dS/dy = \tau_0 \int_0^{R_0} d^2 r S(r,\tau) = \pi R_0^2 S_0 \tau_0, \tag{6}$$

where we have assumed that the initial entropy density remains uniform in the transverse direction

$$S(\tau_0, r) = S_0 \theta(R_0 - r) \tag{7}$$

for a cylinder of radius R_0.

It is well known (McLerran and Toimela 1985; Kajantie *et al* 1986) that the total entropy per unit rapidity is preserved, acting rather like a window connecting the initial times to final times. This scenario allows us to relate the initial entropy to the final multiplicity of particles.

$$(dS/dy) \approx (dS\pi/dy) \approx 3{\cdot}6(dN\pi/dy) \tag{8}$$

assuming only massless pions in the final state, a direct measure of entropy.

The discussions, just presented, tacitly assume throughout the space-time evolution that a steady thermodynamic equilibrium is maintained, a scenario which has an immediate consequence of the thermalization of initial species, presumably quarks and gluons. It is doubtful however that before $\tau_0 \sim 1$ fm/c the system could have been in thermal equilibrium. Rather recently Heinz and other workers have been trying to look at these issues using techniques of non-equilibrium thermodynamics (see Sinha 1987).

3. Diagnostic of quark-qluon plasma, dileptons, photons and J/Ψ suppression

The most significant channel for production of lepton pairs from a plasma of quarks and gluons is $q\bar{q} \to \mu^+\mu^- + X$. Clearly using the Fermi distribution for the q, $\bar{q}$ number density, one finds in the naive non-relativistic limit of the Boltzmann distribution $n(q)n(\bar{q}) \sim \exp(-2\varepsilon/T)$ which turns out to be independent of the baryonic chemical potential. We shall elucidate on this simple but somewhat revealing point later. Considering only 'up' and 'down' quarks (strange quarks to follow) the number of lepton pairs in a space-time volume element d^4x and of invariant mass M is given by (Sinha 1986)

$$\frac{dN}{d^4 \times dM^2} = \frac{5}{3}\frac{\sigma M^2 T^2}{(2\pi)^4} \int_0^\infty dx\,dy \frac{1}{[\exp(x-z)+1][\exp(x+z)+1]} \theta(xy - u^2/4), \tag{9}$$

where $u = M/T$, $z = \mu_q/T$, $\mu_q \approx \frac{1}{3}\mu$, μ being the baryonic chemical potential and is the elementary annihilation cross-section into lepton pairs of mass m_μ from quark-antiquark pairs at a temperature T

$$\sigma = \left(\frac{4\pi\alpha^2}{M^2}\right)\left(1 + \frac{2m_\mu^2}{M^2}\right)\left(1 - \frac{4m_\mu^2}{M^2}\right)^{1/2}, \tag{10}$$

$m_\mu \equiv \mu^+\mu^-$ in our case. After some algebraic manipulation, one can write without any approximation (Sinha 1986)

$$\frac{dN}{d^4x\,dM^2} = \frac{5}{3}\frac{\sigma M^2 T^2}{(2\pi)^4}\sum_{n=1}^{\infty}\frac{(-1)^{n+1}}{n}\exp(-nz)\int_0^\infty \frac{\exp(-mu^2/4x)}{[1+\exp(x-z)]}\,dx \quad (11)$$

It is important to point out that the elementary cross-section goes to zero at an invariant mass $M = 2m_\mu$.

Please note that unlike the u, d case the integrand for the strange quark is cut off at m_S/T. The total contribution therefore for dilepton production is given by (10) plus (11).

The dilepton production rate from the hadronic world has been subject of many discussions. The three most important channels for dilepton are

$$\pi^+\pi^- \to \mu^+\mu^- + X, \quad (12a)$$

$$\pi N \to \mu^+\mu^- + X, \quad (12b)$$

$$N\bar{N} \to \mu^+\mu^- + X. \quad (12c)$$

It has already been argued that the annihilation channel (12a) is the most important; especially for small $M < 500$ MeV; (12b) does not contribute significantly whereas the antibaryon density being small in this temperature regime (12c) contribution is rather small. The other important process for dilepton production is the Drell-Yan mechanism, which contributes mostly to hard lepton pairs as compared to those arising from the fireball. For small-invariant mass, a regime of our interest, in the present study, the most important channel for dilepton production is thus $\pi^+\pi^- \to \mu^+\mu^-$. The production rate for dileptons from the hadronic world is given by (Sinha 1983, 1986; Domokos and Goldman 1985)

$$\frac{dN}{d^4x\,dM^2} = \frac{1}{12}\left(\frac{\alpha^2}{2\pi^3}\right)|F_\pi(M^2)|^2\left(1-\frac{4M_\pi^2}{M^2}\right)\left(1+\frac{2m_\mu^2}{M^2}\right) \times\left(1-\frac{4m_\mu^2}{M^2}\right)^{1/2} MT^2G(M,\lambda), \quad (13)$$

where $W = M^2/2m^2 - 1$, $\lambda = m/T$ and

$$G(W,\lambda) = \int_\lambda^\infty dx(e^x-1)^{-1}\ln\left[\frac{1-\exp(-Wx-P(X^2-\lambda^2)^{1/2}}{1-\exp(-Wx+P(X^2-\lambda^2)^{1/2})}\right] \quad (14)$$

where $P = (W^2-1)^{1/2}$ and the pion form factor is given by

$$|F_\pi(M^2)|^2 = C_1/[(m_\rho^2 - M^2) + C_2(M^2 - 4M_\pi^2)/M^2], \quad (15)$$

where m_ρ (the ρ meson mass) $= 0{\cdot}7723$ GeV, $C_1 = 0{\cdot}3894$ and $C_2 = 0{\cdot}0469$.

Clearly, the production rate from the pions goes to zero at $M = 2m_\pi$, in other words, the production rate of dileptons from the hadronic world turns off at an invariant mass larger than $2m_\pi$ hinting therefore at the possibility of detecting signals of quark-gluon plasma in the invariant mass window $2m_\mu \leqslant M \leqslant 2m_\pi$. What is interesting, however, is that if the chiral symmetry restoration temperature be identical as the deconfinement temperature, the pion mass, at the point of transition will go to zero, leaving a world of massless quarks and gluons—the peak of dilepton production from the hadronic world at $M = 0{\cdot}775\,\mathrm{GeV}$ gets decoupled and the dilepton production for transverse mass turns independent of the invariant mass. Deconfinement and the restoration of the chiral symmetry can thus be diagnosed by the disappearance of the resonance at $M\rho$ along with a substantial signal of dilepton at small invariant mass $2m_\mu \leqslant M \leqslant 2M_\pi$.

Hard photons emanating from quark-gluon plasma can arise essentially from two processes, the so-called unusual compton process $gq \to \gamma q$, g being a gluon and secondly $q\bar{q} \to \gamma g$. Taking into account the three flavours u, d and s quarks, the elementary cross-section of the above processes are given by (Sinha 1985, 1986)

$$\sigma(gq \to \gamma q) = (16\pi\alpha\alpha_s/81s)\ln(s/4m_q^2)$$

and

$$\sigma(q\bar{q} \to \gamma g) = (4\pi\alpha\alpha_s/3s)\ln(s/4m_q^2), \tag{16}$$

where α_s is the running coupling constant and m is the mass of the quark. The production rate is given by (Sinha 1986)

$$\left(\frac{\mathrm{d}\sigma}{\mathrm{d}^4x}\right)^\gamma = \left[\frac{1}{36}\pi\alpha\alpha_s\ln(T/m_{u,d}) + \frac{1}{144}\pi\alpha\alpha_s\ln(T/m_s)\right]T^4. \tag{17}$$

For two colliding nuclei it is useful to follow the remedy prescribed by Kajantie and (Kajantie and Niettinen 1981) to shield the quark mass singularity due to finite temperature and finite size effects such that $\ln(T/m) \sim \ln(1/\alpha_s)$ so that

$$\left(\frac{\mathrm{d}\sigma}{\mathrm{d}^4x}\right)^\gamma \cong \frac{5}{144}\pi\alpha\alpha_s\ln(1/\alpha_s)T^4. \tag{18}$$

The ratio as observed from pp and $p\bar{p}$ data at the ISR experiments has this interesting property that for $P_T \sim 3\,\mathrm{GeV}/c$ the QCD hard photons, in other words (γ/π^0 at $P_T \sim 3\,\mathrm{GeV}$) from the QCD processes simply turn zero, dominated and eventually drowned by background photons ($\pi^0 \to 2\gamma$ etc.); exactly in the kinematic regime where the thermal photons from the plasma tend to be most relevant. This observation allows one to select a precise kinematic window for the plasma signature of (γ/π^0) at or around P_T selected clearly from the hadronic pollutant, yet not drowned by the QCD photons (Sinha 1988).

The result for (γ/π^0),

$$(\gamma/\pi^0)^{\text{plasma}} \cong g_\pi^q\alpha\alpha_s\ln(2/\alpha_s)(R_A + R_B)^2 4r_0/3\pi,$$

where one can choose conveniently $\alpha_s \sim 0{\cdot}6$; $r_0 \sim 1/m_\pi$; $g_\pi^q = 0{\cdot}2$.

In figure 5 we compare the results obtained for the thermal (γ/π^0) from the plasma with the hard QCD photons, as determined recently in the ISR experiments (Camilleri 1987). The experimental results were computed with the pp reaction, the close similarity of the results hint at the relative unimportant role of the QCD processes arising from the interaction due to antiquarks. Clearly, the background photons, arising from the decay of mesons tend to populate the low P_T regime, for P_T typically, $P_T \sim (2\text{–}3)\,\mathrm{GeV}/c$, the background is small, the hard QCD photon contribution remaining almost zero, figure 5. Thus, a sharp rise for (γ/π^0) cross-section around $P_T \sim (2\text{–}3)\,\mathrm{GeV}/c$ is a significant signature of the plasma, as long as the signal is

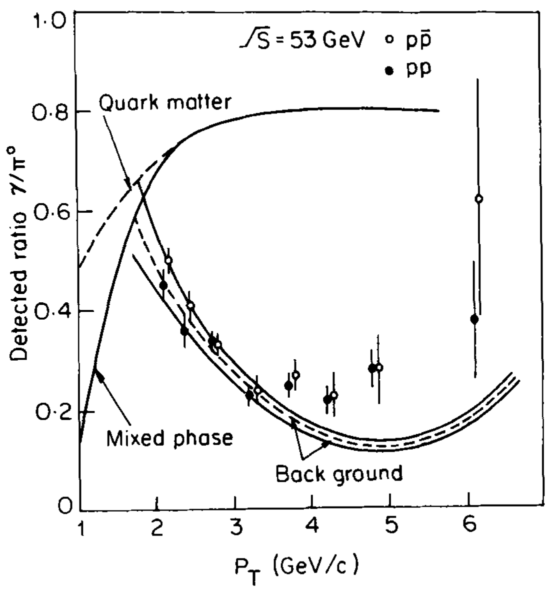

Figure 5. The thermal photons from the quark-gluon plasma phase compared to the background due to meson decay; also shown the recent ISR results for photon production essentially by QCD processes.

measured in conjunction with the dileptons, the dileptons acting rather like a filter for selecting the appropriate kinematic window.

The crucial importance of the size effect, as manifestly evident from the $(R_A + R_B)^2$ term, indicates a rise of (γ/π^0) cross-section by almost a factor of two going from the $^{16}O + ^{208}Pb$ system to $^{208}Pb + ^{208}Pb$ system.

The preceeding discussion is with the implicit assumption that preferential motion (transverse and/or longitudinal) of the fireball is ignored; in the real world the initial fireball in the early times is moving preferentially in the longitudinal direction; during the mixed phase and afterwards, with the space-time evolution, the longitudinal motion is replaced by a transverse motion, now unlike the initial time, dominating the space-time evolution. The preliminary results using the standard first-order phase transition code substantiates such an observation of (γ/π^0) turning insensitive to P_T. However, we already know that for the plasma signal (γ/π^0) insensitivity to P_T, can clearly be considered as a further complementary signature of the plasma. It is indeed rather remarkable that the very recent results of WA80 (Sinha 1988), as demonstrated in figure 6 precisely indicates such an observation. Even more interesting, tinged no doubt with a certain degree of cautious naiveity, that beyond 3 GeV/c, the (γ/π^0) signal turns independent of P_T, a claim, justified by the October (87) CERN experiments. Why should (γ/π^0) for very high P_T, from the plasma turn flat? A tentative argument goes as follows:

For very high P_T (ignoring QCD regime for the present) implies large temperature corresponding to very high initial temperature for very early times, much larger than T_c, the critical temperature. The plasma between $T > T_c$ upto $T \sim T_c$ remains in the form of non-interesting fermions, for the quark-gluon system, nothing changes (with no transverse motion) during that space-time scenario upto $T \sim T_c$ when the mixed phase is turned on. The flatenning of the (γ/π^0) for high P_T is thus a motion picture

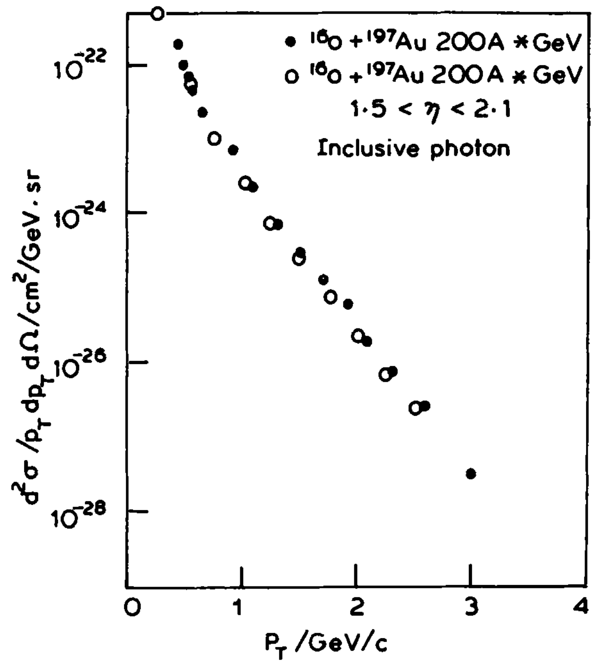

Figure 6. The recent WA80 results π^0 and γ superimposed as a function P_T; please note that γ/π^0 remains almost constant, as predicted by (Fukugita and Ukawa 1986).

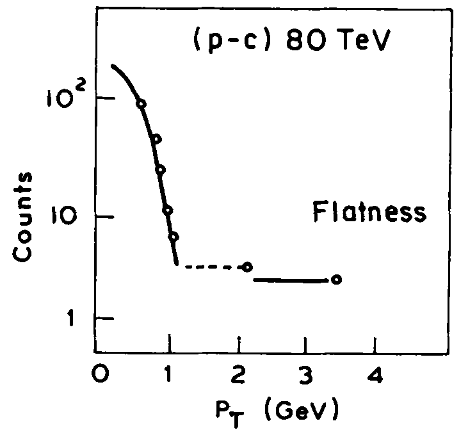

Figure 7. $(p-c)$ at 80 TeV JACEE experiments: The flatness of the spectrum with respect to P_T is the signal of the plasma.

of the plasma and plasma alone which remains unconcerned with the kinematic variables—the very high energy JACEE experimental results also seem to indicate the same outlook (figure 7).

4. J/Ψ suppression

A unique signal of QGP as suggested by Matsui and Satz (1986) can be the suppression of J/Ψ peak relative to the dimuon continuum (figure 8). Indeed, rather recently, the

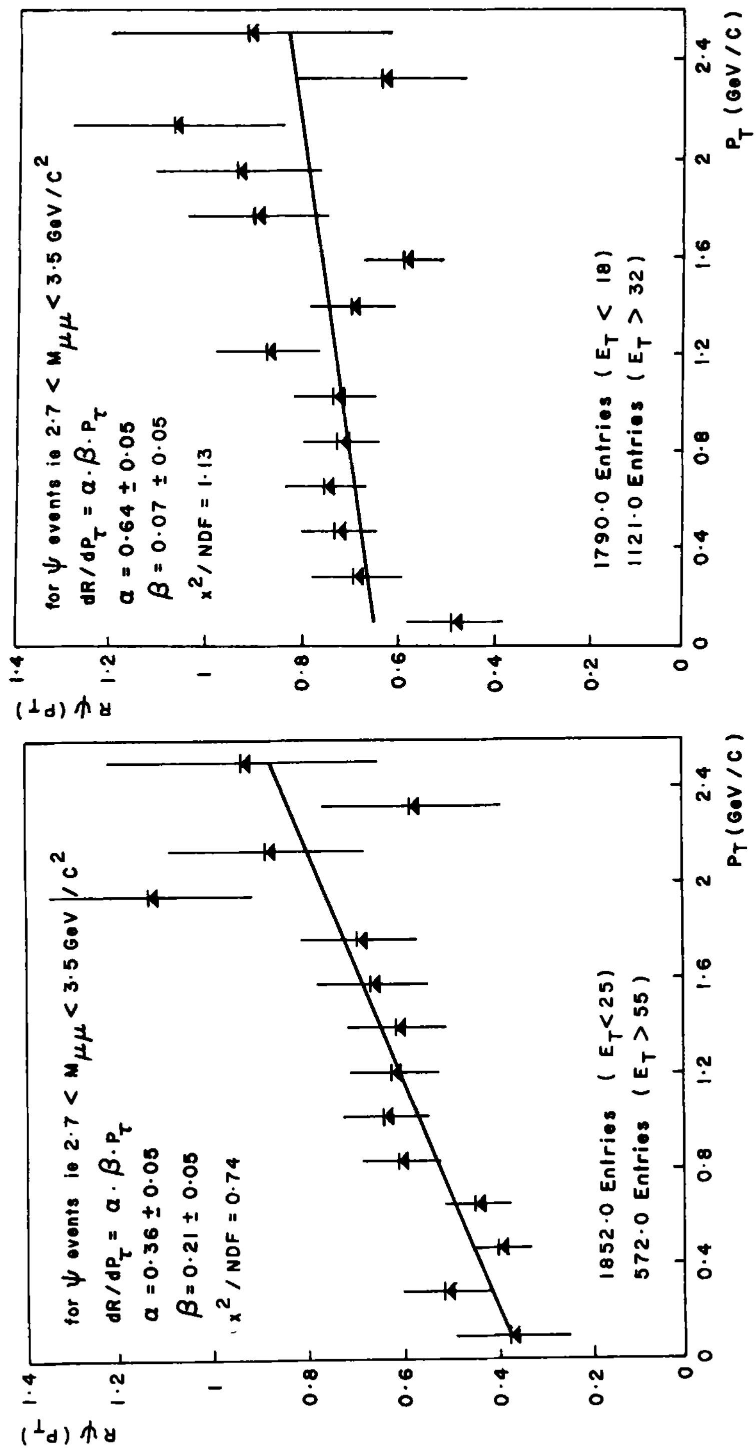

Figure 8. The ratio (Ψ high E_T/Ψ low E_T) as a function of P_T for $0 - U$ (a) $0 - \mathrm{Cu}$ (b) data.

results of the NA38 experiments with 200A GeV ^{16}O beam on heavy targets do indeed show such a suppression in high transverse energy events. One has to agree with Hwa (1988) that the continuum contribution at different transverse energy density and/or transverse momenta should be examined rather carefully before any conclusion is drawn about the supression. In this paper we shall only point out that nuclear effect on J/Ψ production is not insignificant.

To this effect, let us recapitulate with the essential idea that an empirical measure of the amount of coherence can be obtained from the observed power-low behaviour of the total cross-section as a function of the target mass number A_T.

$$\sigma_{h-A} = \sigma_0 A_T^{\beta},$$

where σ_0 is the basic cross-section of the process in question off a hydrogen target with the same beam. In dissipative (totally coherent) processes, $\beta = 2/3$, on the other hand, for high $-Q^2$ process where the sub-nucleonic degrees of freedom are expected to play a significant role, it is reasonable to expect $\beta \sim 1$.

There is thus a fundamental difference between resonances and dileptons in terms of their production mechanism in space and time. In the simplest version of Drell-Yan processes, a $Q\bar{Q}$ pair annihilates into a virtual photon which then decays into a dilepton pair—thus, the spatial size of the dilepton pair at the time of formation is thus limited to $\sim 1/M$. The scenario with resonances, such as J/Ψ is however different. Thus after the $C\bar{C}$ is produced in a spatial size of $M(J/\Psi)^{-1}$, the pair must move away from each other to reach the typical size of the J/Ψ radius. A degree of coherence is thus introduced in this process.

To this end, it is necessary to estimate to what extend a single nucleon in a large nucleus is affected by the surrounding nucleons; in short, how exactly the ground-state properties of the nucleus within the framework of the nuclear many-body problem accommodate this sub-nuclear phenomena.

Defining the sharp radius of the nucleus $R = r_0 A^{1/3}$, the radius parameter remains remarkably constant 1·2 fm. We may naively visualize that the nucleonic sharp $r_0 = r_{\text{int}} + r_N$ is a measure of the nucleonic interaction radius. It is well established on the other hand (Raha and Sinha 1987) that, $r_{\text{int}} = 0{\cdot}5\sqrt{J_d}$ (where J_d is the volume integral of the two-body effective interaction) and that J_d tends to decrease with A_T, a direct consequence of the weakening of the effective interaction with the increase of the density. The decrease in r_{int} with A_T and constancy of r_0 immediately implies that r_N and constancy of r_0 immediately implies that r_N must increase with A_T.

How does the bag radius r_N change from dileptons to J/Ψ? As argued before, $r_N \to r'_N = r_N + r_B$ for the case of resonances whereas $r_N \to r'_N = r_N + 1/M$ for dileptons (Raha and Sinha 1987). This corresponds to an effective decrease of the interaction volume per nucleon somewhat more for J/Ψ compared to dileptons, since for J/Ψ $r_B \gg 1/M$, corresponding to dileptons.

The details of the results can be found in Raha and Sinha (1987) and we would like to quote the results here

$$\beta(J/\Psi) \sim 0{\cdot}95 : \beta(\mu^+\mu^-) \sim 1$$

for $M_{\mu^+\mu^-} \gtrsim 3{\cdot}5$. Thus there exists a suppression of J/Ψ with respect to $\mu^+\mu^-$ even in hA process; certainly $|\lhd\beta| \cong 0{\cdot}5$ is not enough to explain the observed suppression of J/Ψ but the point to bear in mind is that nuclear many-body effects should be first eliminated for a reasonable assessment of the QGP signal.

The preceding discussion is for cold nucleus; what does happen to J/Ψ production scenario with hot hadronic matter? After all, to reach a critical temperature $T_c \sim 200$ MeV, the nuclear medium increases in temperature and density substantially; for a temperature somewhat less than the critical temperature the effect may be significant.

The radius of the bag, in zeroth approximation, is inversely proportional to the fourth root of the bag constant. One expects the bag constant, B, to decrease with increase in temperature. Following Shuryak (1980), the temperature of the bag constant can be utilized to obtain

$$r_N(T) \sim T_N[1-(T/T_c)^4]^{-1/4}$$

a symptom of QGP is for $r_N \to \infty$ as $T \to T_c$. Thus one immediately sees that with increase in temperature the effective volume available for (J/Ψ) production for hot/dense hadronic medium decreases, reducing the value of β even further. The change $|\triangleleft\beta|$ with temperature, given the present state-of-the-art, cannot be evaluated with conclusive accuracy. Suffice it to say these contributions should be examined rather carefully. Indeed there is already an avalanche of publications in this very rich field of research and other references therein; the exciting experimental results of NA38, in my opinion, however, is a remarkable experiment, quark gluon plasma or hadron gas!

5. Early universe relics of QGP

According to the standard model the universe was created some fifteen billion years ago in a "big bang"—the initial singularity in space and time, for temperature $T \to \infty$ as $t \to 0$; to date, all known laws of physics tend to become meaningless beyond the Planck mass scale ($M_P = 1/\sqrt{G} \sim 10^{19}$ GeV) without quantisation of the gravitational field.

Our scenario of the early universe starts much later in time of the order of a few microseconds, at a point of the space-time evolution when the temperature is of the order of ~ 200 MeV.

To decide the evolution of the universe let us consider the standard Einstein field equation for a finite volume V and radius S, giving rise to

$$(1/V)(\mathrm{d}V/\mathrm{d}t) = (3/R)(\mathrm{d}R/\mathrm{d}t) \approx (24\pi G\rho)^{1/2}, \tag{19}$$

where the total mass-energy density will be given by (Wagonar 1973)

$$\rho = (\pi^2 N_f/30)T^4 \quad \hbar = c = k = 1, \tag{20}$$

where N_f is the effective number of degrees of freedom $N_f = N_B + \frac{7}{8}N_F$ for the baryons B and fermions F, since (20) gives $R^2 \propto t$ and $\rho \propto R^{-4}$, one gets immediately

$$T = (2{\cdot}42 \times 10^6 \times N_f)^{-1/2} t^{-1/2}. \tag{21}$$

For a typical time scale of $(2\text{–}3) \times 10^{-6}$ s, the temperature will be of the order $\sim(150\text{–}200)$ MeV, the approximate regime for the universe being in the state of quark-gluon plasma. This epoch is well before the regime of nuclear synthesis, much later in the scenario.

What conceivable signature can there be of that primordial epoch in today's universe?

Just prior to the few microseconds the universe consisted of quarks and gluons, hopefully in the form of plasma and a negligibly small baryons. The present-day wisdom puts a boundary condition $(n_b/n_\gamma) \sim (10^{-9}\text{–}10^{-10})$. From consideration of conservation of baryon number and total entropy (only approximately) during the evaluation of the universe, one immediately finds the ratio of baryons to entropy in the vicinity of an epoch, microsecond-old universe,

$$n_b/s \sim 10^{-10}$$

As the universe cools below T_c, it is more than likely that the plasma of quarks and gluons will supercool (Shuryak 1980 and Wagoner 1973) and go through a first-order phase transition (isentropic) to a hot hadronic state of mostly mesons and a small number of baryons.

Supercooling (Witten 1984) will almost inevitably introduce formation of small bubbles of hadronic matter surrounded in a hot quark soup. There can now be two distinct scenarios—fast nucleation followed by release of latent heat which brings back the temperature of the universe to T_c. With further expansion of the universe the low entropy hadron bubbles grow relative to the quark matter; eventually the hadron bubbles begin to percolate (Witten 1984) leading to breaking up of the quarks phase to small bubbles, the so-called "nuggets" of Witten (1984) surrounded by hot hadronic medium.

The other scenario can be rather a slow process of nucleation with a substantial release of latent heat leading to a highly turbulent state (Kajantie and Kurki-Suonio 1986). For the present, there is no a priori reason to decide between one course to the other and we shall concentrate on the first scenario of rapid nucleation.

The rapid nucleation scenario favours trapping of baryon number in the quark bubbles, yielding the ratio (Witten 1984)

$$R^{h,q} = \rho_b^h/\rho_b^q = 3(2m/\pi T_c)^{3/2} \exp(-m/T_c), \tag{22}$$

where m can be the mass of a nucleon in the ground state 1 GeV—the ratio $R^{h,q}$ can be as small 0·004 for $T_c = 100$ MeV even, 0·14 for $T_c = 200$ MeV.

Two processes can be important for cooling of these baryon-trapped quark nuggets—neutrino diffusion and surface emission of hadrons—leading to a decrease in their baryon number and eventual evaporation.

It is now understood (Applegate *et al* 1987; Farhi and Jaffe 1984; Baym 1988) that for quark nuggets to survive after fifteen billion years the initial size has to be almost of the size of our earth—rendering such as exciting signature of early universe an unlikely candidate—but, precisely for this reason, there can be yet another very exciting possibility and that is as follows.

The relic fluctuation of (Linde 1984; Linde and Weinberg 1983) and the consequent variations of local neutron/proton ratio arising from the evaporation of the baryons from the bubble can enhance the primodial concentration of entropy during the process of nuclear synthesis—tentative hint for at least a partial candidate for dark matter, in the form of quark matter.

The field evidently is very exciting and developing rather rapidly, and we end our comments on the early universe by mentioning that the very suggestion of dark matter resulted from the inflationary universe scenario and there is now at least a plausible candidate for solving the most famous riddle at least partly—the initial conditions of the universe consisting of quarks and gluons should be thought out more carefully as a part of the space-time evolution for the inflationary epoch, and even before. Irreversibility introduced at the very beginning of the creation tends to create pockets of high density quark baryon matter.

6. Outlook and conclusion

To seek out signals of quark-gluon plasma either by bombarding nuclei at ultra-relativistic energies or by a meaningful survey of the cosmological scenario, early universe or neutron stars, is one of the most exciting areas of research, recommended rather strongly to any research students. The fusion of statistical mechanics, particle physics, nuclear physics and early universe cosmology make the field a melting point of the most traditional yet the most original ideas.

The question remains however as to how efficiently one can decipher signals of quark-gluon plasma from the pollutants of the hot hadronic matter—so little is known about the physical properties of hot hadronic matter when the external energy input gets soaked up in creating the excited states of baryonic matter and the meson soup. The vexed question of the fate of the baryon bag, the conventional habitat of three-quark system, is yet to be settled. Further remains the daunting question of the nature of the phase transition and the sancitity of thermodynamic equilibrium. A lot of work has to be done in understanding the pre-equilibrium conditions of the quark-gluon system, a most likely scenario of the early times of nucleus-nucleus collision.

Finally the adventurous saga of the early universe, the counting of time starting from the Planck mass epoch. Why the universe was created and indeed what remains beyond the point of creation and how exactly was it created remains beyond the domain of science that we know today.

We end our paper on a poetic signal from the early epoch when physics melts into metaphysics as recorded by one of the greatest poet of our land only a week before his death.......

> The Primordial Sun, at the
> Moment of creation asked the question
> 'Who art thou'? There was no answer!
>
> The final sun, after millions of years,
> Fading across the Vast Ocean of time,
> Now dark and still, asked the question
> 'Who art thou?' There was no answer!
>
> Rabindra Nath Tagore
> "The Primordial Sun"
> (Transliteration, B S)

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