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Quark gluon plasma

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Abstract. Recent trends in the research of quark gluon plasma (QGP) are surveyed and the current experimental and theoretical status regarding the properties and signals of QGP is reported. We hope that the experiments commencing at relativistic heavy-ion collider (RHIC) in 2000 will provide a glimpse of the QGP formation.

Keywords. Quark gluon plasma; heavy ion collisions; equation of state; hadron gas; early universe.

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1. Introduction

Intense experimental and theoretical activity has been continuing at present to explore the mechanisms of quark confinement as well as the properties of the vacuum state of quantum chromodynamics (QCD). Novel experimental tool employing relativistic heavy-ion collisions has been developed in the past decade to form an entirely new form of matter called as the quark gluon plasma (QGP) in which freely propagating quarks and gluons exist in the (colour charge) plasma. The QGP state formed in the nuclear collisions is a transient rearrangement of the correlations among quarks and gluons contained in the incident nucleons into a larger but globally still colour neutral system [1–3]. However, deconfined quarks and gluons of a QGP are not directly observable because of the confining property of the physical QCD vacuum. The observables indicating a QGP formation are hadrons and leptons. However, these are all indirect probes [4,5].

The energy density inside a proton is of the order of 0.6 GeV/fm^3 (for a proton radius $r_p \simeq 0.7$ fm) and we find that at such energy densities, quarks inside the proton behave almost as free particles. So we want to create a QGP at an energy density $\epsilon > 1$ GeV/fm^3. Such a large energy density can be created in two ways, either by heating the nuclear matter so that the kinetic energy of the particles becomes higher, or by compressing the matter by pumping more hadrons into the region so that the baryon density becomes very large. In both the situations, the quarks of the hadrons come closer to each other so that the long range force is Debye screened and a plasma of quarks and gluons results. The phase diagram of hadron gas and QGP has been shown in figure 1 where the precise determination of the critical line has been done by using Maxwell's construction of the first-order equilibrium phase transition.

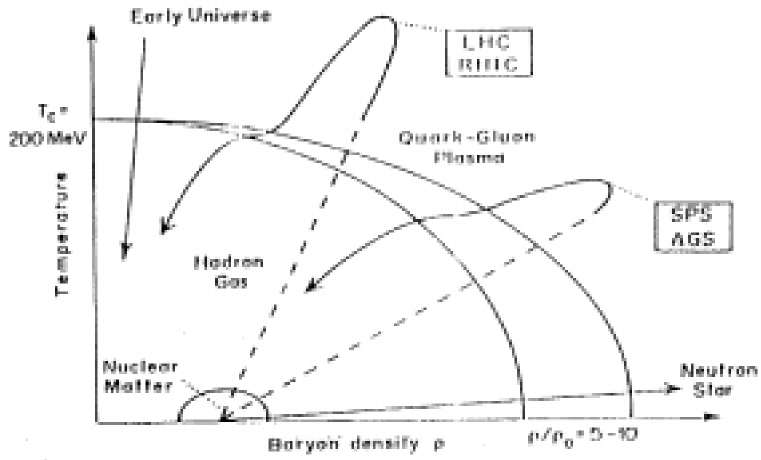

Figure 1. The phase diagram of strongly-interacting matter showing the hadronic phase at low temperature and baryon density, the interaction region (mixed phase), and the QGP phase. The solid line illustrates trajectories followed in Big-Bang evolution, and possibly in heavy-ion reactions at present and future accelerators.

QGP is a new and interdisciplinary field arising from the traditional domains of particle physics and nuclear physics. In combining methods and concepts from both these areas, the study of heavy-ion collisions at ultra relativistic energies leads to a new and original approach of the new matter and the interactions existing there. The focal theme of this review talk is to describe the properties of QGP as well as QGP diagnostics. In case QGP is not formed then we get an extremely hot and dense hadron gas. In order to distinguish the properties and signals of QGP, we must understand properly the hadronic background and its equation of state (EOS).

Another important property of QGP Lagrangian is the chiral symmetry. If we define right handed and left handed spinors as $\gamma_5\psi_R = \psi_R$ and $\gamma_5\psi_L = -\psi_L$, then $\psi_{R,L} = \{(1 \pm \gamma_5)/2\}\psi$ and $\bar{\psi}_{R,L} = \bar{\psi}\{(1 \pm \gamma_5)/2\}$. The term $\bar{\psi}\psi$ in the QCD Lagrangian can be written as the sum of $\bar{\psi}_L\psi_R + \bar{\psi}_R\psi_L$. So if the QCD vacuum is chirally symmetric, the term $\bar{\psi}\psi$ in the Lagrangian should be absent. Thus the presence of a mass term in the QCD Lagrangian breaks chiral symmetry explicitly. Chiral group SU(2) × SU(2) is a subgroup of global flavour rotation group U(N_f) × U(N_f). It is a transformation under combined γ_5 and flavour rotations acting on the quark fields i.e., $\psi' = \exp[i\gamma_5 \cdot (\tau \cdot \theta)/2]\psi$ where τ are the isospin matrices. The two Noether currents corresponding to SU(2) × SU(2) chiral group are isospin vector current $J^a_\mu = \bar{\psi}\gamma_\mu\{(\tau_a)/2\}\psi$ and axial current $A^a_\mu = \psi\gamma_\mu\gamma_5\{(\tau_a)/2\}\psi$. First is exactly conserved but the second is conserved only partially. $\partial^\mu A^a_\mu = 0$ only when pion mass is zero or quark mass is zero. The hadronic phase is a chiral symmetry broken phase because quarks acquire huge constituent mass. In the QGP phase quarks are free and have current quark mass (which is small) and hence chiral symmetry is restored in this phase. So in going from hardons to quark gluon phase we have chiral symmetry restoring phase transition. We still do not know fully well whether the critical line for chiral symmetry restoring phase transition should coincide with the one for the deconfining phase transition [1].

QCD has two vacua, one corresponding to nonperturbative low temperature phase of hadrons and the other corresponds to perturbative, high temperature phase of QGP in which quarks and gluons are weakly interacting. The value of vacuum pressure B in the bag model represents the energy density arising due to interaction effects which are responsible for a change in the vacuum structure between the low and high temperature phases. Minimizing the energy of a spherical bag, one gets the equilibrium energy density inside a proton $\epsilon_0 = 4B$ which is also the latent heat density required for the phase transition. If the pressure of a quark matter inside the bag is increased by increasing the temperature or kinetic energy of the quarks, then the outward pressure becomes comparable or larger than the inward vacuum pressure B. This results into a deconfined phase. Thus the bag model is the simplest QCD motivated model which indicates the formation of QGP under extreme conditions of temperature and density.

2. Lattice QCD simulations

Lattice Monte Carlo simulations are the best ways of performing first – principle non-perturbative computations on the QCD matter and the theory is formulated on a discrete lattice of space-time points. The discretization of the space-time continuum provides a natural way to regularize the ultra-violet divergences which correspond to short distance or large momentum behaviour. This is useful since in nonperturbative QCD, perturbative renormalization in which infinities are removed order by order by using the redefinition of observables (e.g. mass and coupling constant) fails in this case. The shortest distance scale in the theory is the spacing between the nearest lattice points and it automatically sets the maximum momentum scale for the problem in question and gives a momentum cut-off for the calculation. Another advantage of the lattice gauge theory is that by discretizing the space-time and going to the domain of imaginary time, the partition function can be evaluated in the path integral formalism and thus the equilibrium state of the system can be described by a Monte Carlo simulation. In practice, one cannot take a lattice of very large number of points because of the limitation of the computation memory and speed. It is, therefore, necessary to perform calculation with different number of lattice points (or different lattice spacing) until a scaling behaviour is reached [1].

In lattice formulation, the order of QCD phase transition has been studied for different number of quark flavours and for a wide variation of quark mass m_q. The order parameter for the deconfining phase transition is the expectation value of the Polyakov loop $\langle L \rangle = \exp[-F/T]$ where F is the free energy of the quark. Naturally $\langle L \rangle$ changes discontinuously from zero to some finite value as phase transition takes place from confinement to deconfinement region and thus one expects a first order phase transition. Lattice calculations indeed indicate a first order phase transition in the pure gauge theory calculation. Similarly for the chiral symmetry restoring phase transition, one finds that the order parameter $\langle \bar{\psi}\psi \rangle$ again indicates a first order phase transition in the pure SU(3) gauge calculations. When there are quarks in the theory, there is a big difference in physics for $N_f = 0$, 1, 2 and 3 flavours, respectively. Order of the phase transition is again a first order for $N_f \geq 3$ massless quarks. However, for $N_f = 3$ with u, d massless but $m_s > 0$, we obtain a second order transition. Obviously lattice simulations for $n_B \neq 0$ involve some technical troubles and the calculations are not unambiguous ones.

3. Perturbation theory

One can calculate the properties of QGP using perturbative field theory based on the asymptotically free QCD. The main problem one faces is due to infrared divergences. Resummation techniques for hard thermal loops have been found to be useful in this regard [6]. Recent calculations of perturbative theory up to $O(g^5)$ to the pressure have been done. However, the use of perturbation theory even at energy density $\epsilon \simeq 20$ GeV/fm^3 or $T = 600$ MeV is questionable because $g \simeq 2$ and one finds $gT \ll T$ is not satisfied. Therefore, nonperturbative results are more significant even at these temperatures and one finds much variation when one goes from $O(g^2)$ term to $O(g^3)$ term.

4. Statistical thermodynamics

Statistical concepts can be used in nuclear collisions if we have very large multiplicities or number of particles and a sufficiently large size of the fireball ($\gg 1$ fm). Moreover, local thermal equilibrium can only be reached if the lifetime of the dense collision zone is larger than the typical relaxation time due to strong interactions which is of the order of 1 fm/c so that each particle can suffer many collisions. The lifetime of the fireball increases, if we can increase the energy density inside the fireball so that it takes a longer time in hadronization. Thus if heavy-ions collide at a very large energy and the sizes of the nuclei are also large, we can perhaps use the statistical concepts for the fireball formed in the collisions. The statistical description of QGP can easily be exploited in determining the critical parameters of phase transition provided the system is in the thermal and chemical equilibrium. Thermalization is defined as the transition of the system to a state in which it consists entirely of bosons and fermions distributed according to the equilibrium distributions. Chemical equilibrium means that different constituent species are present according to their relative thermodynamical weights.

In using the thermodynamical concepts, the equations of state (EOS) for the phases i.e. hadron gas and QGP are the most crucial factors. Use of ideal gas pictures for QGP and hadron gas separately can throw some light on the phase transition and its critical parameters. However, it is much too unrealistic. For example, unless one uses QGP contained in a big bag with a bag pressure B, one cannot construct a first order phase transition. Similarly in a hot and dense hadron gas, the interactions are intense and strong in magnitude.

The space-time development of the matter produced in the ultra relativistic heavy-ion collisions is described by the hydrodynamics. It gives an integrated picture of all the stages of the expansion from a QGP to a HG through the hadronization until a freeze-out into final state hadrons occur. The effect of a phase transition is significant on the hydrodynamical flow of the fluid and depends on the type of the phase transition dynamics. If the transition is of first order close to equilibrium, the effect slows down the expansion. Consequently the discontinuities in the thermodynamical quantities result in shock-like discontinuities of the fluid. If transition is of first order and involves a supercooled phase or a superheated hadron gas, deflagrations or detonations can occur and they can further cause fluctuations in the multiplicity distributions.

5. Non-equilibrium models

In the ultra relativistic heavy-ion collisions, nuclear matter is compressed and heated up and a maximum compression stage is reached. This stage is characterized by the conversion of the beam kinetic energy into the internal excitation energy of the newly formed matter and therefore, thermalization of the deposited energy starts. Non-equilibrium effects are certain to arise from the rapid time-dependence of the system, finite size effects, inhomogeneity etc. Transport theory is the basic tool to address such problems. Hydrodynamical model assumes that the initial condition in heavy-ion collision corresponds to a state in local thermal equilibrium and there local equilibrium is maintained during evolution. The initial condition in nuclear collisions is a coherent state $|AB\rangle$ of two nuclear systems. A non-equilibrium quantum evolution of $|AB\rangle$ introduces complex high order Fock-state components. Transport theory approximation to the evolution equations can then reveal if local equilibrium is achieved. However, the evolution again refers to an initial state $\rho(\tau_o)$. It should be stressed that we can use neither equilibrium nor near-equilibrium thermodynamics, nor can the nonequilibrium system be represented by an equation of state [3].

Parton cascade models are usually used at very high collider energies when nucleons of the colliding nuclei form a high density quark and gluon gas [7]. One can obtain initial parton distributions from the information on nucleon structure functions. One can use there the perturbative QCD to describe parton interactions and one can study the evolution of parton gas by inserting multiple scatterings. Finally partons are recombined or converted via string formation into final hadron states.

Hadronic transport models involve binary collisions of mesons, baryons or strings etc. One does not include partonic degrees of freedom. However, some models include the idea of colour ropes, breaking of multiple strings which clearly go beyond the hadronic physics [8].

6. Phenomenological EOS for QGP

Lattice QCD yields a suitable EOS for the quark matter. However, lattice approach is still not a suitable tool for studying the strongly interacting matter in the baryon-rich environment. In such situations, QCD-motivated phenomenological models are useful. Recently Leonidov *et al* have proposed a modified EOS for QGP with a temperature T and baryon chemical potential μ dependent bag constant which makes entropy per baryon S/B ratio continuous along the phase boundary. In our earlier work we have modified the μ and τ dependence of $B(\mu,T)$ by incorporating the QCD perturbative corrections for the weak interactions existing in the quarks and gluons of QGP and we have explored in detail the consequences of such a bag constant on the deconfining phase transition [9,10]. We have recently generalised [11] the expression of $B(\mu,T)$ for N_f flavours in the low baryon density region and have investigated further the effect of a massless as well as a massive s-quark on the critical behaviour of the phase transition. We get expression for $B(\mu,T)$ as follows in the limit of $\mu \to 0$ and large T limit:

$$B(\mu,T) = B_0 + (1/9)\mu^2 T^2 - \frac{S_{\mathrm{QGP}}}{S_M}(P_B(\mu,T) - P_B(T,\mu=0))$$

$$-\frac{1}{2}\left(\frac{S'_{\rm QGP}}{S_M} - S_{\rm QGP}\frac{S'_M}{S_M^2}\right)(P_B(T,\mu) - P_B(T,\mu=0))^2. \qquad (1)$$

Here $S'_{\rm QGP} = \partial S_{\rm QGP}/\partial T$ and $S'_M = \partial S_M/\partial T \cdot P_B$ represents the baryonic pressure, $S_{\rm QGP}$ is the entropy of QGP and S_M is the entropy of mesons only. B_0 represents the bag pressure in the $T = 0$ and $\mu = 0$ limit. Hadron gas consists of all strange as well as non-strange hadrons. Thus

$$P_{\rm HG} = \frac{T}{V}\ln Z_B(T,\mu) + \frac{T}{V}\ln Z_M(T,\mu). \qquad (2)$$

From the equations of state for QGP and HG, one can construct a first-order, isentropic and equilibrium phase transition with the help of Gibbs criteria i.e., $P_{\rm QGP}(T_c,\mu_c) = P_{\rm HG}(T_c,\mu_c)$ and thus determine T_c, μ_c for the phase transition. Similarly one gets another phase boundary by determining T'_c and μ'_c for the phase transition from a hadron gas to an ideal QGP by putting $B\ (\mu'_c, T'_c) = 0$.

In figure 2, we have shown the variation of $B(\mu,T)$ as obtained from eq. (1) when u, d, s quarks are in QGP. We found that the values for $B(\mu,T)$ differ much from the case of u, d massless quarks. We find that $B(\mu,T) = 0$ for $T = 210$ MeV and $\mu = 50$ MeV if the interactions among quarks and gluons are incorporated and it is far less than the value $T = 440$ MeV obtained in the case for three massless quarks included in QGP without any interactions between them. However, we find that the calculation shows little change when we take u, d, massless quarks and s as a massive quark ($m_s = 150$ MeV) if compared with the case of three massless quarks. In figure 3, we have plotted the variation of the quantity $(\epsilon - 3p)/T^4$ with temperature T for a QGP consisting of u, d, s massless quarks.

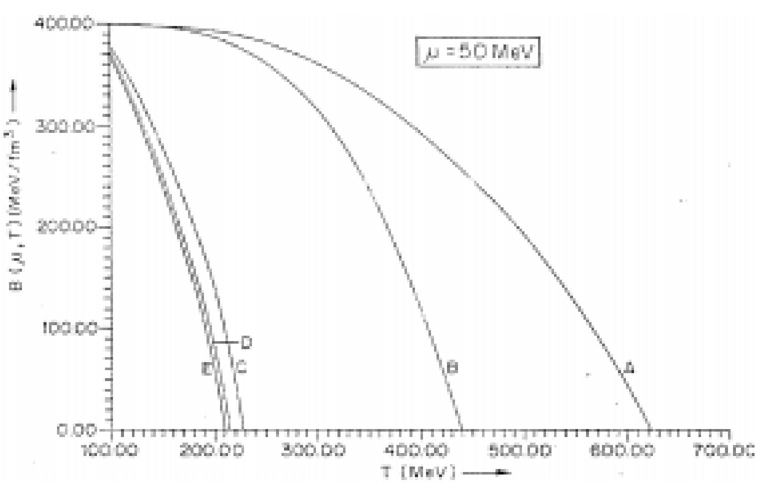

Figure 2. Variation of bag constant $B(\mu,T)$ with temperature (T) at a baryon chemical potential $\mu = 50$ MeV. Curve A represents the free QGP EOS consisting of u, d massless flavours. Curve B is for the free u, d, s massless flavours. Curve C stands for u, d massless and interacting flavours with QCD scale parameter $\Lambda = 100$ MeV. Curve D represents interacting QGP EOS but s-quark is massive ($m_s = 150$ MeV). Curve E is same as D but s-quark is massless. In all these curves, we have used $B_0^{1/4} = 235$ MeV.

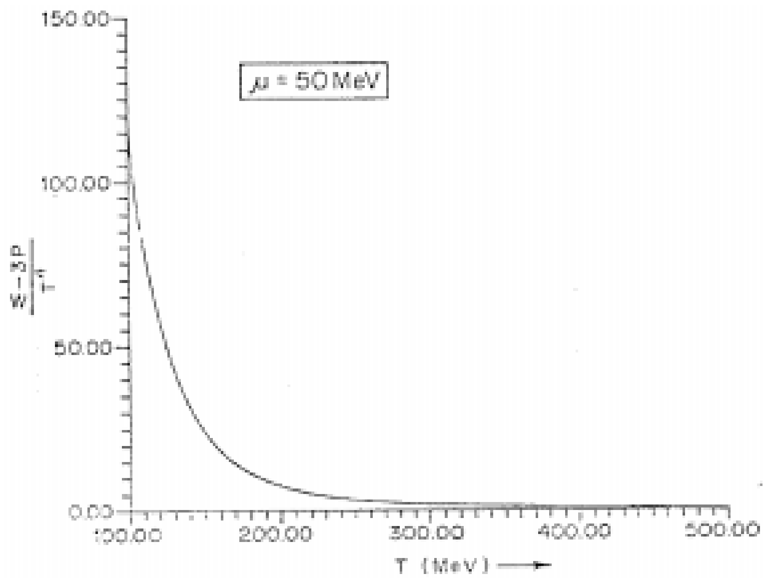

Figure 3. Variation of $(\epsilon - 3P)/T^4$ which is a measure of an ideal plasma behaviour with temperature (T) for $\mu = 50$ MeV and $B_0^{1/4} = 235$ MeV. Here ϵ is the energy density and P is the pressure of the QGP phase of three massless flavours.

Here ϵ represents the energy density and p is the pressure of the QGP. The quantity $(\epsilon - 3p)$ represents a measure of the ideal gas behaviour, since it vanishes for an ideal gas. We find that $(\epsilon - 3p)/T^4$ asymptotically vanishes. This curve compares well with the lattice gauge calculations. It gives us confidence in the QCD-motivated calculations for the EOS of a QGP matter.

7. EOS for hot and dense HG

We want to explore the properties of hadronic matter in unusual environments, in particular at large temperatures and/or densities. For this purpose we need an equation of state (EOS) for the hadron gas. However, very little empirical information is available in the high temperature and density regime on which we can base the theoretical construction of an EOS for the hadron gas. In a simple treatment of hadron gas, baryons were treated as non-interacting point-like objects. However, this gives an undesirable feature that at very large μ, the hadronic phase reappears as a stable configuration in the Gibbs construction of phase equilibrium between the HG and QGP phases. A simple remedy is provided by the inclusion of a hard-core repulsion among them at very high densities. Mean field-theoretical model and its phenomenological generalizations constitute another important approach in the construction of an EOS for the HG. The attractive interactions in such approach arises due to σ-exchange and repulsive interactions are due to ω-exchange which gives a mean potential energy of the hadron gas $U_B = (G^2 n_B)/m_\omega$. The potential energy U_B is proportional to n_B (net baryon density) and is zero when $n_B \to 0$. Thus the early Universe scenario $n_B \to 0$, although hadron density is extremely large. So hadrons effectively yield a point like behaviour due to vanishing of repulsive interaction and thus the mean field approach often provides an erroneous result in such cases. This underlines the importance of the excluded volume, geometrical approach in the EOS for a hot, dense HG. Many phenomenological EOS have been proposed in this direction. This gives rise to widely varying conclusions and ensuing controversies. The results for the same physical quantities from different models are found to be much different.

The EOS employed in excluded volume models suffer from two main and severe deficiencies. Firstly the EOS in these models is thermodynamically inconsistent because we do not have a well derived partition function or thermodynamical potential Ω such that the baryon density n_B can be obtained directly from it (i.e. $n_B \neq \partial\Omega/\partial\mu$). Second and more crucial deficiency is that these models violate causality, i.e. information travels at a velocity larger than the speed of light. In other words, sound velocity $v_s > 1$ in the unit of $c = 1$. Several proposals appeared in the literature which removed the first inconsistency [12,13]. But most of them still suffer from the second one. Recently we have presented a new excluded volume model which incorporated the excluded volume correction in a thermodynamically consistent way [13]. We obtained an EOS by directly evaluating the partition function of the grand canonical ensemble for a dense, hot HG consisting of baryons of kind i with finite size volume V_i as:

$$\ln Z_i^{\mathrm{ex}} = \frac{g_i}{6\pi^2 T}\int_{V_i}^{V-\Sigma_j N_j V_j} \mathrm{d}V \int_0^\infty \mathrm{d}k \frac{k^4}{(k^2+m_i^2)^{1/2}} \left[\lambda_i^{-1}\exp(E_i/T)+1\right]^{-1}, \tag{3}$$

where $\Sigma N_j V_j$ is the total excluded volume for j kinds of baryons, g_i is the spin–isospin degeneracy factor for i kind of baryon. We obtain finally the number density for i kind of baryons in Boltzmann approximation:

$$n_i = \frac{Q_i(1-\sum_{j\neq i} n_j V_j)}{\lambda_i V_i} \exp\left[\frac{1}{I_i V_i \lambda_i}\right], \tag{4}$$

where $Q_i = \int_0^{\lambda_i} \exp[-(1/I_i V_i \lambda_i)]\,\mathrm{d}\lambda_i$ and $\lambda_i = \exp(\mu_i/T)$ called the fugacity of the ith hadron, I_i is the momentum space integral

$$I_i = (g_i/2\pi^2) m_i^2 T K_2(m_i/T)$$

and K_2 is the modified Bessel function of second kind. Once we know baryon densities and $R = \Sigma_j n_j V_j$, we can easily get the pressure as

$$p^{\mathrm{ex}}(T,\mu) = (1-R)P^0(T,\mu) + P_m. \tag{5}$$

Here P^0 is the total pressure arising from the summation of the partial pressures of all kinds of point like baryons and P_m is the pressure due to mesons. We have used a HG with seven baryonic components $N, \Delta, \Lambda, \Sigma, \Xi, \Sigma^*, \Lambda^*$ and their antiparticles and $\pi, K, \eta, K^*, \rho, \omega, \phi$ mesons. We have taken an eigenvolume of each baryon component as $V_0 = (4/3)\pi r^3$ with a hard core radius $r = 0.8$ fm.

In figure 4, we have shown the variations of specific entropy i.e., entropy per particle with the temperature T of hadron gas for different values of μ. At large T and μ, we find that S/n becomes almost independent of temperature and chemical potential and the value of $S/n = 5.0$. This arises because free or available volume in an excluded volume model becomes small at very large T and/or μ. So entropy and baryon density separately saturated to a value for $T > 200$ MeV. In an ideal HG consisting of pions only $s_{\pi}/n_{\pi} \cong 4$. We have compared our calculation with those of Kuono and Takagi model and Rischke model. In figure 5, we have shown the variation of total pressure $P^{\rm ex}$ with respect to energy density at $S/n \simeq 5.0$. From the curve one can easily determine $v_s^2 = \partial P^{\rm ex}/\partial\epsilon$ at a constant S/n. We find that the value of $v_s = 0.7$ in our model. Thus our excluded volume model does not show any violation of causality. For an ideal gas of ultra relativistic particles, $v_s = (1/3)^{1/2} = 0.58$ and for a non-relativistic ideal gas $v_s = (2/3)^{1/2} = 0.8$. Our value lies close to the second value. Thus our excluded volume model gives a realistic equation of state for HG at a high T or μ [14].

In a recent work, it has been demonstrated that a classical Boltzmann gas with two particle interactions can be represented in a mean field model when the mean field is a function of temperature as well as density. The van der Waals' mean field term $U_{V{\rm d}W}(n,T)$ arises due to repulsive interaction between a pair of baryons or antibaryons and leads to an additional pressure term $P_{V{\rm d}W}(n,T)$ where n is the sum of baryons and antibaryons. This term thus incorporates the excluded volume effects in the Walecka mean-field formalism and the EOS in such a formalism can easily describe a HG which has densely packed baryons and antibaryons much below the $n \rightarrow (1/V_0)$ limit where V_0 is the eigenvolume

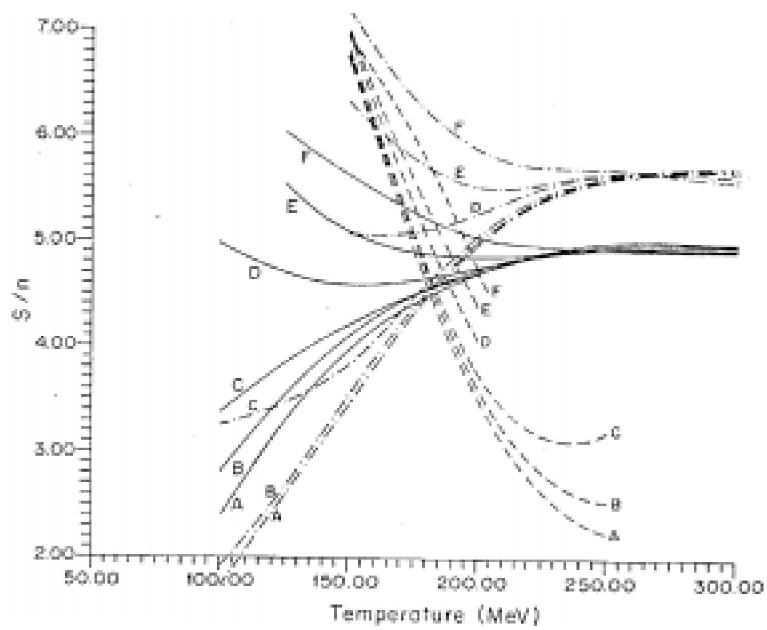

Figure 4. Variations of entropy per particle vs temperature at $\mu = 1400$ MeV, 1200 MeV, 1000 MeV, 750 MeV, 600 MeV and 450 MeV as calculated in our model are shown by the solid line curves A, B, C, D, E and F, respectively. The calculations in Rishke model are shown by dashed curves and those in Kuono–Tkagi model are shown by dashed-dotted cuves.

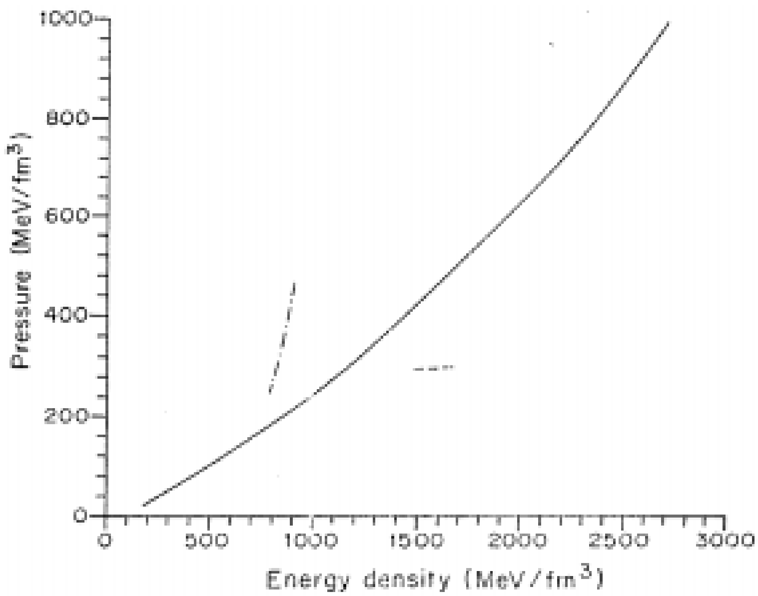

Figure 5. Variations of total pressure with respect to total energy density of the HG at $S/n = 5.0$ is shown in our model by the solid line. The calculations in Rishke model are shown by dashed-dotted curve and those in Kuono–Takagi model are shown by dashed curve.

of a baryon. Since nuclear mean-field theories are effective ones, baryons behave like quasi-particles in a medium. They interact with a constant background of classical scalar and vector mean-fields which are functions of T and n_B in the nuclear medium. Thus the masses of baryons become a variable quantity and it changes with T or n_B or both. One of the major shortcoming of the Walecka mean field model is in getting a large decrease in the effective mass of the nucleons in the nuclear medium $M_N^* < 0$ at $T > 150$ MeV. We have recently studied the mass variations of the nucleons and hyperons in the hot and dense medium in our model in which hard core repulsion has been incorporated into the mean field formalism [15–17].

These studies simply let us know how to achieve a realistic EOS of a hot, dense hadron gas and are essential in the QGP diagnostics.

8. QGP signals

One must identify signals to test whether or not the system produced in a high energy heavy ion collision was in a 'primordial' plasma phase. QGP after its formation subsequently expands and cools below the confining temperature T_c so that the hadrons are formed and at freeze-out, they decouple from the fireball. The hot and dense hadron gas forms a background in this case. The task in identifying QGP signals is rather difficult because one requires a precise knowledge of the HG under extreme conditions of temperature and density. The pictures for a HG we employ either involves an equilibrated statistical system or an ideal, non-interacting system or we consider nuclear collisions as multiple, coherent hadron–hadron collisions etc. We are not sure that such idealized pictures can describe the ultra relativistic heavy-ion collision in a realistic manner. The standard method used in the QGP diagnostics is to compare the predictions of heavy-ion collisions incorporating the presence of QGP with the predictions of models involving the dynamics of hot, dense HG. In case we find any anomalous difference between two types of the pictures, we can subscribe it to an exotic phenomenon like QGP formation.

In high energy nuclear collisions, two beams of nucleons collide. The quarks and gluons are confined in the colliding nucleons. After the primary collisions, we expect multiple scatterings and hence entropy rapidly increases and the system quickly thermalizes. An important question which we investigate, is whether confinement survives this thermalization. In case confinement survives, we have hadrons in the system and if it does not, we then have a QGP in the fireball. Thus we primarily investigate in the QGP diagnostic studies whether deconfinement has really occurred in the heavy-ion collisions. The main proposals are:

(a) *Strangeness enhancement*: The hadrons resulting after QGP formation will confirm a very large number of strange mesons and antibaryons so that we will get a larger value for the ratios $K^+/\pi^+, \bar{\Lambda}/\Lambda, \bar{\Xi}/\Xi$ etc. Enhancement of strangeness also indicates (partial) restoration of chiral symmetry in dense, hot matter. We have recently found that by suitably formulating the EOS for QGP and HG phases separately, one can propose some unique signatures. For example, $\Phi/(\rho+\omega), \bar{\Lambda}/\Lambda, \bar{\Xi}/\Xi$ etc. are anomalously large when QGP is formed. However, recent experimental ratios can be explained on the basis of HG without QGP formation [18,19].

(b) *Photons and dileptons*: These electromagnetic signals for the QGP probe the earliest and hottest phase of the evolution of the fireball and are not affected by the final state interactions. Thus they are the thermometers and probably the best probe for the QGP formation. The recent calculations indicate that at the present CERN energies, the HG background contributions almost matches with the QGP contribution and one cannot draw any conclusive evidence regarding QGP formation.

(c) J/ψ *suppression*: The suppression of J/ψ production caused by the colour screening in the dense medium was originally proposed by Matsui and Satz. The ground state of a c$\bar{\text{c}}$ pair does not exist when the colour screening length λ_D is less than the bound state radius $\langle r^2_{j/\psi}\rangle^{1/2}$, since the Debye length $\lambda_D \simeq (1/gT)$, it becomes smaller at higher temperatures. So the above condition is satisfied when $T/T_c > 1.2$. Recent NA50 data reveal very large suppression of J/ψ in Pb–Pb collision and it gives the first indication for the deconfining phase transition.

Similarly many new directions have emerged in QGP diagnostics studies by explaining intermittency and fractal patterns in multiplicity distributions, HBT inteferometry, disoriented chiral condensates etc. In conclusion, there are many QGP signals proposed in the literature. The problem in the detection of QGP lies in the fact that such a state lives for a brief time and hadrons exist before and after the QGP phase. To what extent, the signals are affected by hadronization is not fully understood at present.

In conclusion, the understanding of the properties and signals of the deconfined phase i.e. QGP is one of the key questions at the end of this millenium. We hope the experiments at RHIC and LHC will unveil this novel phase of matter.

References

[1] C P Singh, Proc. of XII DAE Symp. on HEP, *Indian J. Phys.* **A72**, 601 (1998)
[2] C P Singh, *Int. J. Mod. Phys.* **A71**, 7185 (1992)
[3] S A Bass, M Gyulassy, H Stocker and W Greiner, *J. Phys.* (submitted)
[4] J W Harris and B Muller, *Ann. Rev. Nucl. Part. Sci.* **46**, 71 (1996)
[5] C P Singh, *Phys. Rep.* **236**, 147 (1993)
[6] E Braaten and R D Pisarski, *Nucl. Phys.* **B337**, 569 (1990)
[7] K Geiger, *Phys. Rep.* **258**, 237 (1995)
[8] S A Bass *et al*, *Prog. Part. Nucl. Phys.* **41**, 225 (1998)
[9] B K Patra and C P Singh, *Phys. Rev.* **D54**, 3551 (1996)
[10] B K Patra and C P Singh, *Z. Phys.* **C42**, 265 (1997)
[11] B K Patra, K K Singh and C P Singh, *J. Phys.* (submitted)
[12] D H Rischke *et al*, *Z. Phys.* **C51**, 485 (1991)
[13] C P Singh, B K Patra and K K Singh, *Phys. Lett.* **B387**, 680 (1996)
[14] N Prasad, K K Singh and C P Singh, *Phys. Rev.* (to be published)
[15] V K Tiwari, K K Singh, N Prasad and C P Singh, *Nucl. Phys.* **A637**, 159 (1998)
[16] V K Tiwari, N Prasad and C P Singh, *Phys. Rev.* **C58**, 439 (1998)
[17] N Prasad, V K Tiwari and C P Singh, *Phys. Rev.* **C59**, 2948 (1999)
[18] V K Tiwari and C P Singh, *Phys. Lett.* **B421**, 363 (1998)
[19] C P Singh, V K Tiwari and K K Singh, *Phys. Lett.* **B393**, 188 (1997)