Title
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Publication Date
1983-06-01

LBL-16295 c.2

Lawrence Berkeley Laboratory
UNIVERSITY OF CALIFORNIA

Presented at the Sixth International Balaton Topical Conference, Balatonfured, Hungary, June 6-11, 1983

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June 1983

NUCLEAR SCIENCE DIVISION

Prepared for the U.S. Department of Energy under Contract DE-AC03-76SF00098

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LBL-16295

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This work was supported by the Director, Office of Energy Research, Division of Nuclear Physics of the Office of High Energy and Nuclear Physics of the U.S. Department of Energy under Contract DE-AC03-76SF00098.

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Abstract: The production of high baryon density plasmas at 10 GeV/A and low baryon density plasmas at 1 TeV/A via nuclear collisions is discussed.

1. Introduction

There is mounting theoretical evidence[1] that QCD predicts a transition from hadronic to quark-glue matter at high energy densities. That transition appears to be first order for SU(3) and thus can be characterized by two numbers $\varepsilon_H(\mu)$ and $\varepsilon_Q(\mu)$ where μ is the chemical potential. When the energy density $\varepsilon < \varepsilon_H$, the system behaves as a complex hadronic gas. For $\varepsilon > \varepsilon_Q$, on the other hand, deconfinement is complete, and the system behaves as an ideal (Stefan-Boltzmann) quark-glue plasma. In between $\varepsilon_H < \varepsilon < \varepsilon_Q$ there is a mixed phase. Figure 1 illustrates the dependence of ε and the pressure p as a function of temperature.

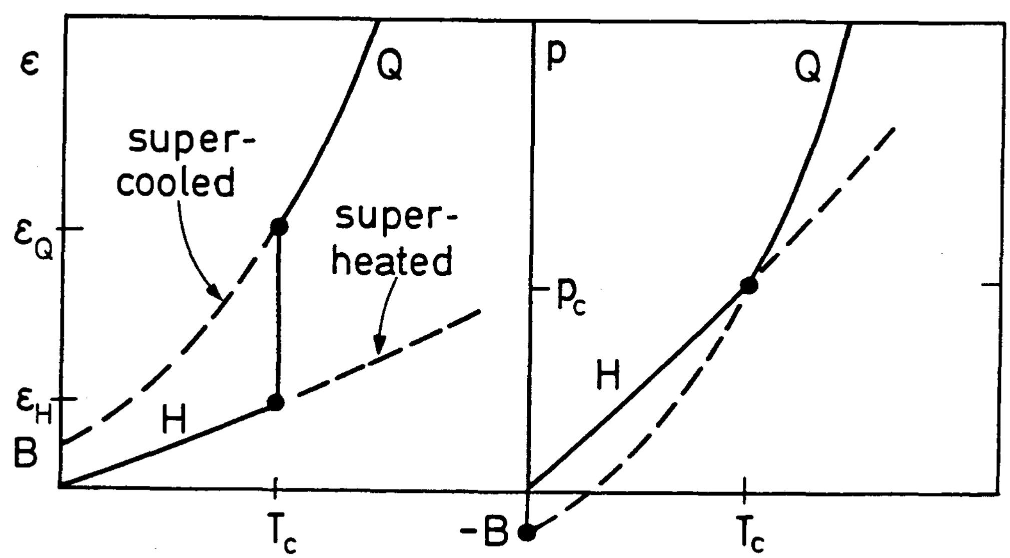

Fig. 1. The energy density ε and pressure p as a function of temperature. At $T_c \sim$ 150-200 MeV the pressure p_c in the hadronic and quark phases coincide. A latent heat per unit volume $\varepsilon_Q - \varepsilon_H$ must be supplied to complete the transition.

We also note that there may exist a superheated hadronic metastable phase and a supercooled plasma metastable phase indicated by the dashed lines. Virtually nothing is known about these metastable phases, but it is important to keep them in mind because in nuclear collisions the compressions and expansions occur so fast that an adiabatic evolution along the stable branches may not be possible.

Numerically, estimates[1] give $\epsilon_H \sim$ 0.2-0.5 GeV/fm^3 and $\epsilon_Q \sim$ 1-2 GeV/fm^3. The latent heat $\epsilon_Q - \epsilon_H \sim 4B \sim 1 \pm 1/2$ GeV/fm^3 characterizes the amount of energy needed to melt the nonperturbative vacuum that confines the hadrons.

The question addressed in this lecture is whether high enough energy densities can be produced in nuclear collisions to produce the plasma phase. We show that high baryon density ($\mu > 0$) plasmas may be produced at surpris- ingly low energies ($E_{lab} \sim$ 5-10 GeV/A), while low baryon density plasmas ($\mu \approx 0$) could be produced only at very high energies ($E_{lab} \gtrsim$ 1 TeV/A).

2. Formation of High Baryon Density Plasmas

From the Bevalac, we have learned[2] that medium mass nuclei (A ~ 100) can stop each other at energies ~1 GeV/A. By stopping we mean the formation of a transient fireball system in the center of mass. A natural question is up to what energies can a fireball be produced. A simple estimate of that can be made if we recall[2] that in a typical hadron-hadron collision both hadrons suffer a rapidity shift $\Delta y \sim 1$ independent of the rapidity gap between them. This follows from the fact that the inelasticity $\eta \sim 1/2$ at all energies. After ν independent collisions the rapidity of a hadron is reduced by $\nu\Delta y$. Clearly, stopping in the c.m. occurs if $\nu > y_{cm}/\Delta y$. At "low" energies $\nu \sim R/\lambda \sim 0.65\ A^{0.3}$. Using this estimate for ν, we see that two uranium nuclei ($\nu \sim 3$-4) can stop each other if $E_{lab} \lesssim 2^{2\nu-1} m_N \sim 60$ GeV. However, the assumption of independent collisions at high energies breaks down because of longitudinal growth.[4] In simplest terms, a nucleon can respond (recoil, emit a pion) in its rest frame only on a time scale $\tau_0 \gtrsim 1$ fm/c. In a frame where that nucleon has rapidity y, the response time is dilated to $t \sim \tau_0 \cosh y$. Hence, the space-time scale for hadronic processes grows rapidly with energy. A corollary of this is that a projectile nucleon will not have completed its interaction with a target nucleon in the lab frame before a time $t \sim \tau_0 \cosh y$. A 10 GeV nucleon traverses $\Delta z \sim 10$ fm before it can suffer another independent collision. This effect is responsible for the observation[3,4] that $(dN/dy)_{hA} \sim (dN/dy)_{hh}$ near the projectile fragmentation regions at high energies. Even though $(dN/dy)_{hA} \sim \nu (dN/dy)_{hh}$ in the target fragmentation regions, this enhancement does not change significantly the inelasticity. Therefore, the rapidity shift of the projectile remains $\Delta y \sim 1$ nearly independent of the number of target nucleons hit during that time.

We can now estimate the minimum nuclear thickness necessary to stop a nucleon in the c.m., i.e., reduce its lab rapidity by one-half. The first collision requires $\Delta Z_1 \sim \tau_0 \mathrm{ch} y_0$, where y_0 is the incident rapidity. After that collision the nucleon has rapidity $y_0 - \Delta y$, so that the second collision requires only $\Delta Z_2 \sim \tau_0 \mathrm{ch}(y_0 - \Delta y) \sim \Delta Z_1 \exp(-\Delta y) \sim \Delta Z_1/2$. The rapidity after the n^{th} collision is $y_0 - n\Delta y$, and thus $\Delta Z_{n+1} = \tau_0 \mathrm{ch}(y_0 - n\Delta y) \sim \Delta Z_1/2^n$. This geometrical sum gives a total traversal length $\Delta Z = \Sigma \Delta Z_n \sim \tau_0 \mathrm{ch} y_0/(1-e^{-\Delta y}) \lesssim 2\tau_0 \mathrm{ch} y_0$. Therefore, the minimum thickness of a nucleus needed to stop a nucleon is

$$R_{min} \sim (1-2)\, \tau_0 \,\mathrm{ch} y_0 \quad . \quad (1)$$

Uranium should therefore be able to stop a 5-10 GeV nucleon even with longitudinal growth taken into account.

Since fireballs can be formed up to $E_{lab} \sim 10$ GeV/A, consider now the energy densities achieved. If a certain baryon density, ρ_B, can be generated, then trivial kinematics dictate that $\varepsilon = \gamma_{cm} m_N \rho_B$. For $E_{lab} = 10$ GeV, $\varepsilon = 2$ GeV/fm^3 already when $\rho_B \sim 6\rho_0$. To estimate the baryon density achieved, we first recall Goldhaber's estimate[3] for the minimum compression, $\rho_{min} = 2\gamma_{cm}\rho_0$, under the assumption that the stopping conditions are satisfied. Therefore, $\varepsilon_{min} = 2\gamma_{cm}^2 m_N \rho_0$. If shock waves are produced, then higher compressions $\rho_B = 4\gamma_{cm}\rho_0$ can in fact be reached.[5] These estimates indicate that $\varepsilon > 2$ GeV/fm^3 could easily be attained in the 10 GeV/A region at the price of high compression.

Thus far, we have considered only the problem of forming hadronic fireballs. To form a plasma we must pay the price of melting the nonperturbative vacuum. This condition implies that there is a minimum kinetic energy, E_{min}, below which plasma production is impossible. To determine E_{min}, we use the bag model equation of state[6,7]

$$\varepsilon = \varepsilon_{SB}(T,\mu) + B$$
$$P = \tfrac{1}{3}\varepsilon_{SB} - B \quad , \quad (2)$$

where ε_{SB} is the energy density of a Stefan-Boltzmann gas of quarks and gluons. The energy density of the perturbative vacuum is $B \sim 0.1$-0.5 GeV/fm^3 above the energy density of the nonperturbative vacuum. The latent heat in this model is 4B. If a plasma is formed in the c.m., then the properties of the plasma are constrained by the conservation of energy-momentum flux and baryon flux. Furthermore, positive entropy must be generated in the transition. This in turn means that the plasma temperature must be positive. A critical curve is determined by T = 0 in the plasma, at which point

$$\varepsilon = c\mu^4 + B \quad , \quad \rho_B = d\mu^3 \quad , \quad (3)$$

where $c = (54\pi^2)^{-1}$ and $d = 2/(9\pi)^2$. By energy conservation the energy per baryon in the plasma is $E/A = \gamma_{cm} m_N$, where $\gamma_{cm} = \operatorname{ch} y_{cm}$ is the c.m. gamma factor of the two colliding nuclei. The energy density of the plasma must therefore be $\varepsilon = \gamma_{cm} m_N \rho_B$. Setting this ε equal to the energy density of a zero temperature plasma gives γ_{crit}, below which no plasma can be produced. This critical c.m. gamma factor is thus[7]

$$\gamma_{crit} = a\left(\frac{\rho_B}{\rho_0}\right)^{1/3} + b\left(\frac{\rho_0}{\rho_B}\right) \quad , (4)$$

where a = 0.61, $b = B/\varepsilon_N$, and $\varepsilon_N = m_N \rho_0 \sim 0.15$ GeV/fm^3 is the energy density of nuclei. Eq. (4) is only an implicit equation in that ρ_B/ρ_0 depends on γ_{crit}. However, in the (ρ_B, γ_{cm}) plane (see Fig. 2), eq. (4) shows that the positive temperature plasma formation requires a kinetic minimum energy.

A crude estimate for γ_{crit} can be made by locating $\partial\gamma_{crit}/\partial\rho_B = 0$. This gives

$$\gamma_{crit} \approx \frac{4}{3} a^{3/4} (3b)^{1/4} \approx 1.2\,(B/\varepsilon_N)^{1/4} \quad . (5)$$

The minimum laboratory kinetic energy is therefore $E_{min} = 2m_N(\gamma_{crit}^2 - 1) \sim$ 1-3 GeV/A for $B \sim (1-3)\varepsilon_N$. A more detailed calculation involves the determination of ρ_B via the Rankine-Hugoniot constraint giving[8]

$$\frac{\rho_B}{\rho_0} = \left[\frac{\varepsilon_2 + p_1}{\varepsilon_1 + p_2} \quad \frac{\varepsilon_2 + p_2}{\varepsilon_1 + p_1}\right]^{1/2} \quad (6)$$

$$= \frac{1}{2\gamma_{cm}}\left\{4\left(\gamma_{cm}^2 + \frac{B}{\varepsilon_N}\right) - 3 + \left[\left(4\left(\gamma_{cm}^2 + \frac{B}{\varepsilon_N}\right) - 3\right)^2 - 16\frac{B}{\varepsilon_N}\gamma_{cm}^2\right]^{1/2}\right\}$$

The actual excitation curve therefore depends on B. An example is shown in Fig. 2.

These curves can be continued until the thickness of the flame front separating the hadronic and plasma phases becomes too large. The minimum thickness of the flame front must grow linearly with due to longitudinal growth. In the c.m. the front becomes as thick as the Lorentz contracted nuclei when[3]

$$\gamma_{cm} \tau_0 = R/\gamma_{cm} \quad . (7)$$

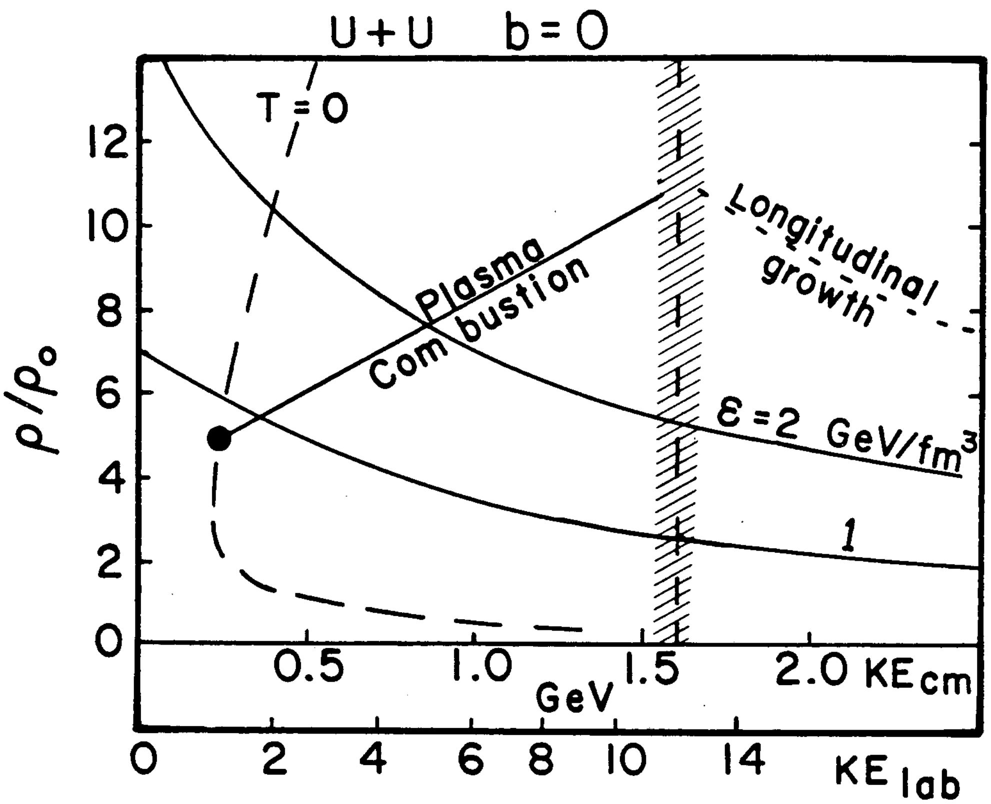

Fig. 2. The maximum baryon density in central U + U collisions versus beam energy. Solid contours show energy density. Dashed curve corresponds to zero temperature plasma. Plasma combustion curve, eq. (6), between $1 < KE_{lab} < 10$ GeV per nucleon is shown. Beyond the shaded boundary at $KE_{lab} \sim 10$ GeV shock formation is prevented by longitudinal growth.

This leads to a maximum laboratory kinetic energy

$$E_{max} \sim 2\frac{R}{\gamma_0}m_N \sim 10 \text{ GeV/A} \qquad (8)$$

Beyond that point the Rankine-Hugoniot combustion solution no longer makes sense. It is possible that a certain fraction of the nuclear material is still converted into a plasma at higher energies (see next section), but a uniform plasma cannot be produced. Furthermore, since the stopping distance increases with $\gamma_{cm}\gamma_0$, the baryon density must decrease[3] for $E > E_{max}$ as $\rho_B \propto \gamma_{cm}^{-1}$.

Therefore, we find a narrow window between 3 and 10 GeV/A where high baryon density plasmas could in principle be produced. It should be remarked, however, that the window could be closed by more careful consideration of the E_{min} and E_{max} estimates. For example, Stocker showed[6] that the first order perturbation correction to the

Stefan-Boltzmann law has the effect of multiplying the coefficient a in eq. (3) by $1+\alpha_s$. This could raise E_{min} by a factor of two.[6] Also, it is possible that the proper width of the flame front is larger than $\tau_0 \sim 1$ fm/c. This would have the effect of reducing E_{max}. Furthermore, at E_{min} only a T = 0 plasma could be produced. Such a state has zero entropy and there would be zero branching ratio to it. For $E > E_{min}$ one must compare the entropy generated in the plasma and hadronic phases to compute the branching ratio.[6] It could be that $E > E_{max}$ is required for the branching ratio to become significant to the plasma phase. However, within a plausible range of parameters a plasma window could exist at these "low" energies.

3. Fragmentation Region Asymptotics

At higher energies, nuclear transparency certainly sets in. If we concentrate on the fragmentation regions, then there remains a residual excitation and compression but at the price of substantial gradients [see M. Gyulassy in Ref. (1)]. Simple estimates show that the proper energy density increases linearly with the depth in the target.

For heavy nuclei, $\mathcal{E} \sim 2$ GeV/fm^3 could be reached at depths Z > 10 fm. However, detailed hydrodynamic calculations[9] only seem to reach $\mathcal{E} \sim 1$ GeV/fm^3 in the fragmentation regions. In any case, only a small fraction of the nucleus is likely to be in the plasma phase at best. A large fraction of the nucleus will be heated enough to be in the mixed phase, though. For observables such as strangeness production or dilepton production this means that there will be a convolution of signals from hadronic, mixed, and plasma phases across the nuclei in the fragmentation regions. It would probably be difficult to disentangle that convoluted signal. However, I suggest that the occurrence of a mixed phase in the middle of a nucleus could result in spectacular signatures in terms of fluctuations. Since $dp/d\varepsilon = 0$ in the mixed phase, density fluctuations are unstable. These unstable fluctuations could lead to enhanced rapidity density fluctuations in the fragmentation regions.

4. Central Region Asymptotics

I turn finally to the production of low baryon density plasmas at very high energies. The new observation[10] I report on here is that the initial conditions of the plasma are connected in a nonlinear manner with the final rapidity density. The physics is simple. If the plasma expands hydrodynamically, then each cell does work on expansion. Consequently, the energy density of each cell at breakup time is less than at formation time. Specifically, consider the scaling hydrodynamics[9] equations

$$\tau \frac{d\varepsilon}{d\tau} + (\varepsilon + p) = 0$$

$$\frac{d}{d\tau}(\tau s(\tau)) = 0 \quad , \quad (9)$$

where $s = (\varepsilon + p)/T$ is the entropy density of the ($\mu = 0$) plasma. From eq. (9) follows that $\varepsilon(\tau) = \varepsilon(\tau_0)(\tau_0/\tau)^{4/3}$ and that τs is a constant of motion. As Landau suggested the density of quanta in a fluid cell is proprtional to the entropy density

$$\rho = a s \qquad , (10)$$

where $a^{-1} \approx 4$. Therefore, the final multiplicity, N, obtained by integrating ρ along the breakup surface $\tau = \tau_f$ along which $\varepsilon(\tau_f) \sim m_\pi^4$ is given by

$$\begin{aligned} N &= \int d^4x \, \rho(x) \, \delta(\tau - \tau_f) \\ &= \int dy \, \tau_f \, d^2x_\perp \, \rho(\tau_f, y, x_\perp) \end{aligned} \qquad , (11)$$

where $y = \ln(t+z/\tau)$ is the rapidity variable. In the scaling domain, ρ depends only on τ and consequently

$$\frac{dN}{dy} = A_\perp \tau_f \, \rho(\tau_f) = A_\perp a \, \tau_0 \, s(\tau_0) \qquad , (12)$$

where we used eq. (10) and the conservation of τs. Substituting $s = k \varepsilon^{3/4}$ in to eq. (12) we find that

$$\varepsilon_0 = \varepsilon(\tau_0) = \left(\frac{1}{\tau_0 A_\perp a k} \frac{dN}{dy} \right)^{4/3} \qquad . (13)$$

This is the basic result showing the nonlinear connection between ε_0 and the final rapidity density. The constant of proportionality turns out[10] to be $(ak)^{-1} \approx 1.6$, where $A_\perp$ is the cross-section area of the plasma. Eq. (13) should be contrasted to the Bjorken estimate[11]

$$\varepsilon_{Bj}(\tau_0) = \frac{m_\perp}{\tau_0 A_\perp} \frac{dN}{dy} \qquad , (14)$$

which was obtained neglecting the work done on expansion. Numerically ε_0 turns out to be about twice as high with eq. (13) as with eq. (14). For the JACEE event[3] with dN/dy ~ 300, $A_\perp$ ~ 40 fm^2, ε_{Bj} ~ 2.5 GeV/fm^3 with $m_\perp \sim$ 0.35 GeV. Equation (13) gives, on the other hand, $\varepsilon \sim$ 5.5 GeV/fm^3 ~ $2\varepsilon_{Bj}$! This indicates that the JACEE event may be even more spectacular than thought before.[3] Roughly one-half of the energy per cell was used up as work. Even after taking finite energy corrections[9] [$\varepsilon_0 \rightarrow .75\,\varepsilon_0$] into account the initial energy density is well above current estimates for ε_Q in Fig. 1. We conclude that the conditions for plasma formation in the central region seem to be well satisfied already in "light" (Si+Ag) nuclear collisions in the TeV/A range.

Finally, I want to call attention to possible novel combustion phenomena[12] in expanding plasmas. Since the longitudinal expansion of the plasma is very rapid, $\varepsilon(\tau) \sim \tau^{-4/3}$, the plasma may be supercooled along the dashed curve of Fig. 1 (taken from Ref. 12). For typical initial conditions in the central region ε is reduced to ε_Q in a short time $\tau_Q = \tau_0(\varepsilon_0/\varepsilon_Q)^{3/4}$. For $\tau > \tau_Q$ a mixed phase would be produced in the system expanded adiabatically. However, no transition could take place before $\tau_0 \sim 1$ fm/c has elapsed in the supercooled phase. But by that time $\varepsilon(\tau_Q + \tau_0) \lesssim \varepsilon_0/2$! Thus considerable supercooling could occur. Under these conditions detonation or deflagration bubbles may form[12] in the plasma. These would in turn be observable as enhanced dN/dy fluctuations, medium range correlations, and large $dE_\perp/dy$.

This work was supported by the Director, Office of Energy Research, Division of Nuclear Physics of the Office of High Energy and Nuclear Physics of the U.S. Department of Energy under Contract DE-AC03-76SF00098.

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