HEAVY ION PHYSICS AT VERY HIGH ENERGIES

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1 Introduction

The principal aim of high energy heavy ion experiments is the study of strong interaction thermodynamics. We want to understand the behaviour of bulk matter at densities so high that the interaction of its constituents is governed by QCD. Statistical QCD predicts that at sufficiently high density, there will be a transition from hadronic matter to a plasma of deconfined quarks and gluons—a transition which in the early universe took place in the inverse direction some $10^{-5}$ seconds after the big bang. Ultimately, we hope that heavy ion experiments will provide the tool to study in the laboratory both the transition to and the properties of the primordial quark-gluon plasma. We therefore have to determine the energies necessary to achieve this and find the observables that will give us the information needed to test the thermal nature and probe the primordial state of the system.

A further proposed aim for heavy ion physics at very high energies is the study of coherent $A - A$ collisions (photon-photon or pomeron-pomeron interactions) in a regime where these could lead to the production of Higgs bosons, supersymmetric particles, or $W$ pairs. In particular, coherent heavy ion collisions can provide a high luminosity $\gamma - \gamma$ source; this could allow the observation of specific final states without the strong hadronic contamination present in the corresponding gluon-gluon interaction in proton-proton collisions. We thus have to investigate for what ion beam luminosities such experiments become feasible.

In Table 1 we summarize the present status and the future plans for heavy ion experimentation; it shows three distinct stages. At present, existing sources and existing accelerators are used to provide rather light ion beams (up to $^{28}S$) and $^{32}S$). The main aims at this stage are to establish the feasibility of high energy ion-ion experiments with their very abundant secondary hadron production, to show that there is a chance to obtain high densities, and to look for the onset of new, collective phenomena. It is generally agreed that the experiments carried out so far have achieved these objectives to a considerable degree and provide sufficient justification to continue.

In a second phase, to begin in about three years, both the BNL-AGS and the CERN-SPS will have new injectors able to accommodate really heavy ions, all the way up to uranium. This should make it possible to determine to what extent we can actually get into
a regime of thermodynamic behaviour. Moreover, if present estimates are correct, there should be a chance to obtain more conclusive evidence for the onset of quark deconfinement.

The third stage, expected to start around 1998, will bring us into the true high energy heavy ion regime. We should now get average energy densities well above the deconfinement threshold, so that a study of the properties of the quark-gluon plasma should become possible. Now we can also produce systems of nearly vanishing baryon number density (similar to the state of the early universe); it should thus become possible to study critical behaviour in a wide range of the temperature-density phase diagram of QCD matter. And it appears within reach to get even to energy densities for which the quark-gluon plasma is approaching its asymptotically free “ideal gas” form.

2 QCD Thermodynamics

2.1 Results from Statistical QCD

The advent of the quark structure of elementary particles led rather directly to the prediction of quark matter. Since then, both the new state and the transition between hadronic matter and quark-gluon plasma have been studied extensively. We expect deconfinement in dense matter, because the presence of many colour charges will screen the confining potential between the members of a given $q\bar{q}$ or $qqq$ system. The density of the system can be increased either by “compression” (an increase in baryon number density or baryonic chemical potential $\mu_B$), or by “heating” (an increase in the initial energy density $\epsilon$). This leads to a critical transition curve in the phase diagram of strongly interacting matter, as shown in Fig. 1.

The crucial features for our present considerations are the values of the transition parameters (temperature, density, energy density, screening length). These questions have been addressed in statistical QCD both in the lattice formulation and in various effective Lagrangian models.

In the computer simulation of statistical QCD on the lattice [1], one tries to calculate the relevant quantities from first principles, without any simplifying physical assumptions. The quantitative reliability of the results is at present, however, still somewhat limited by technical restrictions (memory size, operating speed of available supercomputers). Nevertheless, the results of lattice QCD today do give us a reasonably good general
In Fig. 2, we show the behaviour of energy density $\epsilon$ and pressure $P$ for QCD matter with light quarks of two flavours ($u$ and $d$), as calculated on the lattice [2]. We note that at a critical temperature $T_c$, the energy density undergoes a rapid transition from low values, corresponding to a hadron gas, to much higher values, corresponding to a quark-gluon plasma. For ideal (i.e., non-interacting) systems, the ratio of the energy densities of pion gas to quark-gluon plasma is given simply by the corresponding degrees of freedom; for $N_f = 2$, this means $\epsilon_x/\epsilon_Q = 1/10$. The energy density of the ideal plasma, including finite lattice corrections [3], is also shown in Fig. 2, and at high temperatures, the calculations appear to approach this value. In an ideal gas, however, energy density and pressure
are related by \((\epsilon - 3P)/T^4 = 0\), and we see that this condition is certainly not fulfilled for \(T < 1.5 T_c\): below \(1.5 T_c\), the energy density overshoots the Stefan-Boltzmann limit, the pressure falls much below it. For a pure \(SU(3)\) gauge system (i.e., for \(N_f = 0\)), these deviations from ideal gas behaviour were recently studied in detail [4]; the result is shown in Fig. 3 and indicates that the system may not become ideal until even higher temperatures.

The order of the transition is at present under intense investigation by lattice studies [1]. One finds a first order deconfinement transition for \(N_f = 0\), and a first order transition corresponding to chiral symmetry restoration and deconfinement for \(N_f \geq 3\), in the limit of massless quarks. For \(N_f = 2\), the transition appears to be continuous for present lattice sizes and quark masses; however, for two light and one heavy quark species (corresponding to the actual \(u, d, \bar{s}\) quarks), it becomes first order when the strange quark mass reaches a certain values [5]. More detailed quantitative studies are needed here. All results at vanishing baryon number density agree, however, on the same transition point for deconfinement and chiral symmetry restoration.

The physical value of the transition temperature is for present, not yet asymptotic lattice sizes best determined by calculating both \(T_c\) and the hadron masses in units of the lattice spacing; the ratio then gives us \(T_c\) in terms of meson or baryon masses. In Fig. 4 we show the result for \(T_c\) as determined from \(m_{\rho}\) for different \(N_f\). The most reasonable value for the actual physical case is \(T_c \approx 150\) MeV; we must keep in mind, however, that in present calculations the ratio of \(\pi\) to \(\rho\) mass has not yet reached its physical value, and hence quantitative results are not yet final. To be on the safe side, we shall consider the critical temperatures to lie in the range \(T_c = 150 - 200\) MeV; the corresponding critical values of the energy density necessary for deconfinement are \(\epsilon_e \approx 1 - 3\) GeV/fm\(^3\).

Finally we want to note the result of lattice calculations of the screening length. Since bound states will "melt" in dense matter when the screening radius becomes significantly smaller than the binding radius, the temperature dependence of the screening radius \(r_D(T)\) gives us some idea of when specific bound states will disappear. From Fig. 5 we see that above \(T \approx 1.2 T_c\), even the tightly bound \(c\bar{c}\) state \(J/\psi\) will become deconfined.

An alternative approach to statistical QCD is offered by the study of effective La-
The basic idea here is to construct a model Lagrangian which incorporates as much as possible of the known low density hadron physics, and then check what it predicts at higher densities. Rather detailed studies in the framework of chiral perturbation theory [6] reproduce the known pion physics at zero temperature, and predict chiral symmetry restoration at $T_c \approx 190$ MeV, in accord with lattice results. Another effective Lagrangian study [7] has also been extended to non-zero baryon number density; it leads to the interesting phase diagram shown in Fig. 6, with a continuous transition at zero baryonic chemical potential $\mu_B$, which then turns into a first order transition at some tricritical point for $\mu_B > 0$. This behaviour illustrates that new features can still be expected in the region of the phase diagram not yet accessible to lattice studies.

Figure 4: Deconfinement temperature $T_c$ for QCD matter with $N_f$ flavours of light quarks, as function of the lattice size $N_T$; from [1].

Figure 5: Colour screening radius $r_D$ as function of the temperature $T$, for QCD matter with $N_f$ light quarks; from [1].
2.2 Conditions in Nuclear Collisions

In statistical QCD, we study the equilibrium thermodynamics of strongly interacting matter. We would like to test the results experimentally in high energy heavy ion collisions. This leads us immediately to the basic question which such studies have to face: do heavy ion collisions lead to systems dense enough, large enough, and long-lived enough to treat them by the equilibrium thermodynamics based on QCD? Let us first consider the density regime attainable.

The basic observable for an estimate of the initial energy density is multiplicity of the secondary hadrons emitted in the collision. The total multiplicity (1.5 times the observed charged multiplicity) per unit central rapidity interval can be parametrised in high energy proton-proton collisions by

\[
(dN/dy)_p = 0.8 \ln \sqrt{s};
\]

this form describes well all data from SPS to Tevatron energies [8], with \((dN/dy)_p\) growing from 2.4 at 20 GeV to about 6 at 1.8 TeV. This is extrapolated to central \(A-A\) collisions by

\[
(dN/dy)_A = A^\alpha (dN/dy)_p,
\]

with \(\alpha \geq 1\). For \(\alpha = 1\), we simply have a superposition of \(A\) independent \(p-p\) collisions; if there is rescattering between the different nucleons and/or the produced secondaries, we will have \(\alpha > 1\). Present data from nuclear collisions give \(\alpha \geq 1.1\) [9], and this leads for the multiplicity per unit central rapidity in \(Pb-Pb\) interactions (\(A=208\)) to the range of values:

<table>
<thead>
<tr>
<th>Energy</th>
<th>Multiplicity</th>
</tr>
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<tbody>
<tr>
<td>SPS (17 GeV)</td>
<td>500 - 800</td>
</tr>
<tr>
<td>RHIC (200 GeV)</td>
<td>900 - 1500</td>
</tr>
<tr>
<td>LHC (6300 GeV)</td>
<td>1500 - 2500</td>
</tr>
</tbody>
</table>

The lower number always corresponds to \(\alpha = 1\), the higher to \(\alpha = 1.1\). Large as these numbers may seem, it should be noted that fixed target experiments at SPS and AGS, which cover several units of rapidity, have shown that multiplicities in the hundreds can indeed be handled.
If the observed secondaries have come from the initial interaction region in “free flow”,
then the initial energy density $\varepsilon$ in a central $A - A$ collision is given by [10]

$$\varepsilon = \frac{\langle dN/dy \rangle_{A} m_{T}}{\pi R_{A}^{2} \tau}, \quad (4)$$

where $m_{T} = (p_{T}^{2} + m^{2})$ denotes the transverse energy of the secondary, $R_{A} \simeq 1.2A^{1/3}$
the nuclear radius, and $\tau$ the formation time or longitudinal extension of the equilibrium
system. With the estimates $m_{T} \simeq 0.5 \, \text{GeV}$ and $\tau \simeq 1\, \text{fm}$, we obtain from eqs.(1) and (2)

$$\varepsilon = 0.09A^{2/3} \ln \sqrt{s}, \quad (5)$$

and from this for central $Pb - Pb$ collisions:

<table>
<thead>
<tr>
<th></th>
<th></th>
<th>0.01</th>
<th>0.1</th>
<th>1.0</th>
<th>10</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>SPS</td>
<td>1.5 -</td>
<td>2.5</td>
<td>GeV/fm$^3$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>RHIC</td>
<td>2.8 -</td>
<td>4.7</td>
<td>GeV/fm$^3$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>LHC</td>
<td>4.6 -</td>
<td>7.8</td>
<td>GeV/fm$^3$</td>
<td>(6)</td>
</tr>
</tbody>
</table>

The values we have used for $m_{T}$ and $\tau$ are “conservative” — the transverse mass
increases somewhat with $\sqrt{s}$, and enhanced nuclear stopping would reduce $\tau$; there are
models [11] which give $\tau \sim A^{-1/6}$. On the other hand, a diffuse edge of the interaction
region would result in a larger value for $R_{A}$, which could offset these effects. We should
further note that the values in eq.(6) give the expected average initial energy density; it
should be possible to trigger on higher values. But even the average values give a gain by
a factor of 4–7 in comparison to the average energy density in $p - p$ collisions.

![Figure 7: The expected range of the average initial energy density $\bar{\varepsilon}$ versus incident CMS energy $\sqrt{s}$; also shown is the predicted range of the critical energy density $\varepsilon_c$ for deconfinement.](image)

In Fig. 7 we show the variation of the average energy density in central $Pb - Pb$
collisions, together with the critical energy density $\varepsilon_c$ determined in section 1.1. We note
that at the LHC we are well above this value and are in fact getting into the range in
which $\varepsilon$ becomes high enough to produce an ideal quark-gluon plasma.

We now want to consider briefly the question of how the energy density $\varepsilon$ can in practice
be varied for a given collider. In $A - B$ collisions, with $A \ll B$, we have different values
of $\epsilon$ for events at different impact parameter $b$, i.e., for events of different multiplicity or different transverse energy $E_T$. Going from peripheral to central collisions, we increase $\epsilon$: the effective transversal overlap area remains constant, as the smaller projectile after complete "immersion" hits the larger target at smaller and smaller $b$; the effective number of participants increases, however, since the target contains more nucleons in the center, at $b = 0$, than at the edges. In "symmetric" $A - A$ collisions, with large $A$, this is no longer the case, since over most of the impact parameter range, the number of participants and the overlap area are essentially proportional. The result [12] is illustrated in Fig. 8, where we show $\epsilon$ as function of $E_T(b)/E_T(0)$. For an infinite nucleus, $\epsilon$ becomes independent of $E_T(b)/E_T(0)$; deviations from constancy are thus of purely geometric origin. In particular, the drop of $\epsilon$ as $E_T(b)/E_T(0) \rightarrow 0$ is extremely dependent on the specific nuclear edge structure; it is more a surface than a volume effect and hence not useful to study the $\epsilon$ dependence of any observables. For $A - A$ collisions, we thus have to find some other way to change $\epsilon$.

In principle, the best way would be to vary the incident energy $\sqrt{s}$ for given $A$, since this would change $\epsilon$ at constant volume. In practice, this requires large variations of the beam energy, since $\epsilon \sim \ln \sqrt{s}$. Moreover, a reduction of $\sqrt{s}$ at a given collider is in general accompanied by a considerable luminosity drop [13][14] and hence for many experiments not very useful.

This leaves us with $A - A$ collisions at several $A$ and fixed $\sqrt{s}$ as the only viable road to different $\epsilon$, even though this changes the associated volume as well. Going from $U - U$ to $S - S$ collisions reduces $\epsilon$ by a factor two or more; to achieve such a change by varying $\sqrt{s}$ would require going from the peak LHC energy of 6.3 TeV down to one-hundredth of this value, which is expected to decrease the luminosity by almost two orders of magnitude. In contrast, reducing $A$ will generally increase the luminosity, so that most rates should not be too much affected. For these reasons, it is very important to include the capability to run at varying $A$ from the beginning among the essential requirements in planning the heavy ion mode of any collider.

Figure 8: Energy density $\epsilon$ vs. transverse energy $E_T$ as function of the impact parameter $b$, for $O - Pb$ and $Pb - Pb$ collisions; the dashed line shows the behaviour for $A - A$ collisions in the limit of infinite nuclear size; from [12].
Finally we should note that in view of these considerations, the study of QCD thermodynamics at different accelerators, but for the same or similar $A$, plays an important complementary role.

What temperatures do the values of $\epsilon$ in eq.(6) correspond to? For an ideal plasma with three flavours of massless quarks we have

$$\epsilon = (47.5\pi^2/30)8T^4, \quad (7)$$

which leads to an initial temperature $T = \epsilon/(1953)^{1/4}$, with $\epsilon$ in GeV/fm$^3$ and $T$ in GeV. We can now use the energy density estimates of eq.(6) to get the corresponding temperatures. Before we list these, we want to note an alternative way to estimate $T$. If the initial bubble of matter undergoes longitudinal hydrodynamic expansion to attain the observed final state, rather than the free flow assumed in eq.(4), then the entropy is conserved, not the energy: part of the initial energy goes into work against the pressure of the vacuum on the system [16]. The initial entropy density $s$ in a central $A - A$ collision is obtained from

$$s = 3.6(dN/dy)A/(\pi R_A^2 \tau), \quad (8)$$

which for an ideal plasma leads to $T = (s/2605)^{1/3}$, with $T$ in GeV and $s$ in fm$^{-3}$. As already noted, this leads to somewhat higher initial temperatures, and hence also to higher initial energy densities. The temperature values for $Pb - Pb$ are (in MeV):

<table>
<thead>
<tr>
<th>Accelerator</th>
<th>Lower</th>
<th>Upper</th>
</tr>
</thead>
<tbody>
<tr>
<td>SPS</td>
<td>170 - 190</td>
<td>160 - 190</td>
</tr>
<tr>
<td>RHIC</td>
<td>200 - 220</td>
<td>200 - 240</td>
</tr>
<tr>
<td>LHC</td>
<td>220 - 250</td>
<td>230 - 280</td>
</tr>
</tbody>
</table>

(9)

The first column corresponds to free flow, the second to isentropic expansion. With increasing collision energy, the difference between the two temperature estimates (and hence between the corresponding energy densities) increases. For $Pb - Pb$, we obtain for free flow $\epsilon_{Bj} \simeq 0.52 \ln \sqrt{s}$, for isentropic expansion $\epsilon_{S} \simeq 0.32(\ln \sqrt{s})^{4/3}$. The resulting behaviour of $\epsilon$ is illustrated in Fig. 9. We note that only at LHC energies the two have

Figure 9: Initial energy density $\epsilon$ vs. incident CMS energy $\sqrt{s}$ for free flow [10] and hydrodynamic flow [16].
become really distinct, so that then the effects of longitudinal hydrodynamic expansion should become evident even for average quantities.

In section 1.1, we had seen that so far the most reliable calculations are available for systems of vanishing baryon number density; this is also the situation which presumably existed at the end of the quark-gluon phase of the early universe. When can we expect similar conditions in nuclear collisions? From $p - A$ collisions at $A \approx 200$, we know that in passing through the nuclear target, the projectile proton looses approximately two units in rapidity. The corresponding baryon number distribution is then centered at $y - \delta y$, where $Y \simeq \ln(\sqrt{s}/m_p)$ denotes the maximum rapidity and $\delta y$ the rapidity shift of a nucleon passing a nuclear target; the distribution vanishes at $Y$ and $Y - 2\delta y$. The overall baryon-free region in rapidity thus becomes

$$\left(\Delta y\right)_0 \simeq 2(Y - 2\delta y), \quad (10)$$

target. If $\delta y$ is the same in $A - A$ collisions as for $p - A$, then we get:

<table>
<thead>
<tr>
<th></th>
<th>$2Y$</th>
<th></th>
</tr>
</thead>
<tbody>
<tr>
<td>SPS</td>
<td>0</td>
<td>5.8</td>
</tr>
<tr>
<td>RHIC</td>
<td>2.7</td>
<td>10.7</td>
</tr>
<tr>
<td>LHC</td>
<td>9.6</td>
<td>17.6</td>
</tr>
</tbody>
</table>

Eq. (11) is consistent with recent $p - \bar{p}$ data [17], which show at $\sqrt{s} = 1800$ GeV, where $2Y \approx 15$, a baryon-free region of at least 6.5 units; even with $\delta y = 2$ we expect seven units. It is not clear, however, if $A - A$ collisions do not give enhanced stopping, i.e., $\delta y > 2$, and hence a smaller baryon-free region. This has to be studied by event generators for nuclear collisions, and we shall return to such studies shortly. In Fig. 10 we show the baryon-free regions at the different accelerator energies. Even if there is more stopping than expected from $p - A$ collisions, however, the LHC still provides an ample safety margin.

What can we say about the baryon-rich fragmentation region at very high energies? The kinematic compression experienced by target and projectile is essentially determined by the rapidity shift $\delta y$ for nucleons [18]; if this is not dependent on $\sqrt{s}$, then the kinematic compression will not increase by going to higher incident energies. The energy deposited in target or fragmentation regions is also expected to depend only very weakly on the incident

Figure 10: Baryon number density distribution in rapidity $y$ vs. incident energy $\sqrt{s}$. 

197
energy [9]. This leads to the conclusion that a baryon-rich state of strongly interacting matter is better studied at a fixed target machine with high $A$ beams; there appears to be no particular reason to consider this regime at high energy colliders, where it is not so easily accessible.

So far, we have addressed general aspects of nuclear collisions, extrapolating from known features of $p - p$ and $p - A$ interactions. This is done in more detail in specific event generators, which are obtained by assuming some particular form of interaction in the course of the collision [19]. In PYTHIA, the interaction is simply taken to be a superposition of individual nucleon-nucleon collisions, without any secondary interactions; particular attention is paid, however, to obtain the correct high energy behaviour of the $p - p$ interaction, including minijet effects. In VENUS, on the other hand, the interaction is assumed to be string fragmentation, with strings formed also between nucleons and secondaries, as well as between different secondaries. This model has been used so far more at lower energies, and it probably has to be complemented for high energy features (minijets). It does give us some idea of rescattering effects already now. In addition, studies using FRITIOF and the dual parton model are under way. In all cases, the event generator should be tuned to account correctly for the general features of $p - p$ data at all available energies. For $A - A$ collisions, the different codes will then show us the role of different interaction schemes corresponding to different extrapolations.

Let us look at some first results. In Fig. 11, we see that both PYTHIA and VENUS describe correctly the energy dependence of the average multiplicity per unit central rapidity in $p - p$ collisions, and that this dependence is in fact well reproduced by eq. (1). Going to $Pb - Pb$ collisions at LHC energy, PYTHIA recovers the result listed in eq.(2), with $(dN/d\eta)_{Pb} = 1500$. On the other hand VENUS, because it includes rescattering, gives a considerably larger value, $(dN/d\eta)_{Pb} = 2500$. First estimates from the dual parton model give much larger values still. It is thus possible that our "conservative" estimates of energy densities and temperatures for RHIC and LHC are in fact too conservative.

In Fig. 12, we show the baryon number distributions from PYTHIA and VENUS. For $Pb - Pb$ at the LHC, PYTHIA gives about 14 units of baryon-free rapidity region. This
is more than the 10 units we had found in eq. (11) by extrapolating $p - A$ data, since PYTHIA contains no nuclear stopping. VENUS, with more secondary interactions, gets only 8 units, which, however, is still only half of the total LHC rapidity region.

These results are just the beginning of event generator studies in this field. Much further work is in progress and should give us information also on particle ratios and momentum spectra.

### 2.3 Volumes and Life-Times

Before we can use heavy ion collisions to study equilibrium thermodynamics, we must make sure that the volume and the life-time of the systems experimentally produced are sufficiently large. This becomes all the more crucial for the study of critical behaviour, which becomes really "critical" only in the infinite volume limit. How can we check volumes and life-times?

The initial interaction volume has already been introduced in eq. (3); for a central $A - A$ collision, it is $V_0 \approx \pi R_A^2 \tau$, where $R_A \approx 1.2 A^{1/3}$ denotes the effective transverse nuclear radius and $\tau \approx 1$ fm the initial longitudinal extension of the system which will later thermalize, or, equivalently, the formation time. This system now expands, and at the time of transition to hadronic matter, it has attained the size $V_c = V_0 (\varepsilon_o / \varepsilon_c)$, where we have used $\varepsilon_o$ to denote the initial energy density. This means that $V_c$ increases with $(dN/dy)$; at the LHC, the transition volume $V_c$ is about 5 - 8 times larger than the initial volume. For $Pb - Pb$, this gives at the transition point a volume of some $800 - 1200$ fm$^3$. The system now continues to expand until freeze-out. The freeze-out radius can be estimated by supposing that interactions stop when the energy density has dropped to that of an ideal pion gas at $T \approx T_c = 150$ MeV. For lead beams, this gives us a freeze-out radius

$$R_F^2 \approx 1.24 (dN/dy)^{1/3}. \quad (12)$$

Another possible estimate for $R_F$ is obtained if one supposes freeze-out to take place when the mean free path $\lambda$ of pions has reached the size of the system [20],[21]. With
\[ \lambda = V_F/(dN/dy)\sigma_\pi \text{ and } \sigma_\pi \simeq 20 \text{ mb}, \text{ this leads to} \]

\[ R_F^3 \simeq 0.69 \left( dN/dy \right)^{1/2}, \quad (13) \]

for the freeze-out radius. Introducing an energy dependence of \( \sigma_\pi \) [9,22] leads to yet another form, which grows as \( (dN/dy)^{5/12} \) and thus falls between the two cases we had discussed.

\[ \tau_c = \tau_0 (T_0/T_c)^3, \quad (14) \]

Experimentally, the freeze-out size can be determined by particle interferometry, based on the Hanbury-Brown-Twiss method to measure star sizes. In Fig. 13 we see that at present multiplicities, both forms (12) and (13) accommodate the data. It should be emphasized that these data lead to freeze-out radii which are almost a factor two larger than the radii of the projectiles. This supports the idea that nuclear collisions indeed produce bubbles of expanding strongly interacting matter and thus gives us a first hint that we are on the right track. From Fig. 14 we conclude that already a lead beam in the SPS should allow us to distinguish the two forms (12) and (13), and that at the LHC lead-lead collisions should produce volumes of some \( 10^4 \) to \( 10^5 \) fm\(^3\). Thus nuclear collisions can indeed be expected to lead to volumes which are very much larger than the few fm\(^3\) typical of nucleon-nucleon interactions.

The use of like-particle interferometry to determine the freeze-out radii will, however, encounter some difficulties at very high energies. The linear size of the system is inversely proportional to the momentum resolution needed to study the interference between partners, and this may at LHC energy surpass the feasible. Moreover, at a radius of some 20 fm the Coulomb interaction between like-sign charged particles becomes very strong and may mask the Bose-Einstein interference. Hence the proposal [23],[24] to study correlations in the longitudinal instead of the transverse dimension is very interesting. It would replace the transverse scale of nuclear size by the longitudinal scale of only one fermi, so that after expansion even by a factor five, the corresponding radii would still be less than 5–10 fermi.

The life-time of the produced bubble in a possible plasma phase can be estimated if we assume longitudinal hydrodynamic expansion; from entropy conservation we then get
Hence the chance for rescattering increases. It still has to be checked, however, how far at fixed $x/5$ and by increasing $\sqrt{s}$ at fixed $A$: in both cases, the number of secondaries and therefore want to move towards thermalisation by increasing the effective $A$ compared to the quoted $p-p$ value; central $S-W$ collisions give about 0.2 at this energy of GeV, central $S-S$ collisions [28] give 0.15 ± 0.03 at midrapidity for the $K^+/\pi^+$ ratio, to be checked how far nuclear collisions bring us on the road to equilibrium, and the observation of $K^+/\pi^+ \approx 0.19 \pm 0.03$ in $Si-Au$ collisions at the AGS [27] certainly points towards an onset of thermal behaviour. Data from the SPS are in agreement with this. At $\sqrt{s} = 20$ GeV, central $S-S$ collisions [28] give 0.15 ± 0.03 at midrapidity for the $K/\pi$ ratio, to be compared to the quoted $p-p$ value; central $S-W$ collisions give about 0.2 at this energy [29]. In general, we expect to move towards thermalisation by increasing the effective $A$ at fixed $\sqrt{s}$ and by increasing $\sqrt{s}$ at fixed $A$: in both cases, the number of secondaries and hence the chance for rescattering increases. It still has to be checked, however, how far

$$\tau \approx 1 \text{ fm}$$ for the initial state formation time. This gives

$$\text{SPS} \quad \text{RHIC} \quad \text{LHC}$$

$$\tau \text{[fm]} \quad 1.2 - 2.0 \quad 2.4 - 3.2 \quad 3.6 - 6.5$$

for the plasma life-time at the various energies. A higher initial temperature, and even more so a first order transition, would prolong considerably the time until the system is completely hadronized.

### 2.4 The Onset of Thermalisation

How can we test whether the system produced in a heavy ion collision has reached thermal equilibrium? The starting point of the collision is evidently the interaction of individual nucleons in the projectile with those in the target. The extreme opposite to thermalisation is thus a superposition of nucleon-nucleon collisions, with each projectile nucleon interacting with just one target nucleon. On the way to thermalisation, we have “rescattering”: a given nucleon will interact with more than one other nucleon, it will interact with secondaries produced in previous collisions, and these secondaries will interact with each other. Particle ratios provide us with an effective tool to check whether rescattering has occurred and brought us closer to thermal equilibrium. Consider as example the production ratio $K^+/\pi^+$. In $p-p$ interactions at $\sqrt{s} = 20$ GeV, it is found to be $0.07 \pm 0.02$ [25]. Rescattering will increase this ratio, and after “enough” rescattering, there will be equilibrium. The $K^+/\pi^+$ ratio attains a value of about 0.20 in an equilibrium hadron gas at $T = 150$ MeV and vanishing baryon number density [26].
particle ratios such as the considered $K/\pi$ ratio increase in event generator codes which include only one or two rescatterings and thus are probably not yet thermal. Furthermore,

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{figure15.png}
\caption{The $K/\pi$ ratio in $p - p$ and $p - \bar{p}$ collisions as function of the incident energy square root of $s$, from [25]; also shown are the hadron gas limits for two temperatures [26].}
\end{figure}

it should be noted that in $p - p$ and $p - \bar{p}$ interactions, the $K/\pi$ ratio increases with increasing $\sqrt{s}$ (Fig. 15) [25] as well as with increasing multiplicity (Fig. 16) [17]. This could be due to rescattering among the produced secondaries, but as well to enhanced strangeness production in minijets. Both can again be checked by event generator studies.

The onset of thermalisation can thus be tested by studying the evolution of strangeness away from the $p - p$ level, for kaons as well as hyperons. We have here implied that a gas of non-interacting hadrons is the equilibrated end-phase; as well, one could consider the ratios obtained from an equilibrium quark-gluon plasma, which expands and subsequently hadronises. It is found, however, that for most ratios the two scenarios agree [30], even though it may take longer to attain equilibrium by hadron interactions alone [31].

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{figure16.png}
\caption{The $K/\pi$ ratio vs. central multiplicity, from [17].}
\end{figure}
may be cases, however, for which hadron gas and quark gluon plasma do lead to very different ratios, and hence the study of strangeness evolution has also been proposed as a way to obtain information about the primordial state [31],[32]; we shall return to this point later on.

2.5 Primordial Features

There are two ways to test whether the system produced in a high energy heavy ion collision was in its early "primordial" history in a deconfined state.

We can look for signals which are produced at such early times and are not affected by the subsequent hadronisation. Possible signals of this type are thermal dileptons or direct photons, which are emitted by the plasma and then escape [33],[34],[35]. In this spirit, one may also study the effect of the produced medium on the observed production rates of heavy quark bound states [36],[37] or hard jets [38]; their initial production is non-thermal, presumably occurs very early in the collision, and can be understood reasonably well in terms of perturbative QCD.

Another approach is to look for primordial remnants in the observed hadronic features. Possible candidates considered in this vein are discontinuities in the momentum distribution of the secondaries, reflecting a first order phase transition [39]; particle ratios which are significantly different for a hadron gas and a hadronising quark-gluon plasma [31],[32]; and droplets of strange matter, baryonic states of very low charge to mass ratio [40].

2.5.1 Thermal Dileptons and Direct Photons

There are several sources for dilepton production in nuclear collisions. They are produced in the decay of the low mass vector mesons \(\rho,\omega\) and \(\phi\). Thermal dileptons are emitted when \(\pi^+\pi^-\) or \(q\bar{q}\) pairs annihilate in a pion or quark gas. Finally, there are comparatively rare "hard" interactions between incident partons at a very early stage of the collision, leading to Drell-Yan production or to the production of heavy (cc or bb) vector mesons which subsequently decay into lepton pairs. The distinction between dileptons from low mass and high mass vector mesons, as well as that between thermal dileptons from \(q\bar{q}\) annihilation and Drell-Yan pairs is somewhat arbitrary; it is mainly motivated by the fact that the production mechanism of high mass dileptons seems to be comparatively well understood in terms of perturbative QCD. In the typical "soft" hadronic regime around the \(\rho,\omega\) and \(\phi\) this is not the case. However, it may well be that below the \(\rho\) region, for very soft dileptons, finite temperature perturbation theory may become applicable; this region has been the subject of much recent theoretical work [41],[42],[43].

Thermal dileptons and Drell-Yan pairs moreover lead to different functional behaviour in the dilepton mass \(M\). We have

\[
\frac{d^2\sigma}{dM^2 dy}|_{y=0} \sim \exp\left(-\frac{M}{T}\right) \quad (16)
\]

for thermal pairs emitted from a system at temperature \(T\), and

\[
\frac{d^2\sigma}{dM^2 dy}|_{y=0} \sim M^{-4} f(\tau) \quad (17)
\]

for Drell-Yan pairs, with \(\tau \equiv M^2/s\). Eq. (17) does not include contributions from scaling violation terms; these are, however, expected to contribute more at large \(P_T\) and not affect greatly the integrated cross-section. Since the mass distribution of thermal dileptons contains directly the temperature of the emitting system, it has been proposed as a "thermometer" for strongly interacting matter [34]. Note here that if used in this way, it does
not tell us whether the system is deconfined or not, since a pion gas as well as a quark-gluon plasma would emit thermal dileptons. In addition, however, we have to ask under what conditions thermal dileptons are at all observable: is there some window between the “hadronic” region around the low mass vector mesons and the high mass Drell-Yan regime, in which thermal pairs should dominate? Theoretical studies [44] based on what we consider today extreme temperatures ($T = 500 - 800$ MeV) led to significant thermal production in the mass region above 2.5 GeV, dominating the Drell-Yan distribution. To check this more quantitatively, we need calculations of the thermal spectrum, based on LHC conditions as upper limit, as well as the corresponding Drell-Yan calculations for $Pb - Pb$ collisions, with $3 \leq M \leq 10$ GeV, including higher order QCD contributions. For the latter task, we encounter some further problems: since at the LHC even $M = 10$ GeV leads to $\tau \approx 10^{-6}$, we need quark and gluon structure functions at very small $x$, where they are not yet known. Moreover, it is known that at small $x$ there is “shadowing” in nuclear collisions, i.e., the structure function in a nucleus is at small $x$ reduced in comparison to that in a nucleon. Detailed studies lead us to expect for Drell-Yan production at the LHC a reduction of up to 50% [45].

Both thermal [46] and Drell-Yan [47] production have now been calculated for LHC conditions; for the latter, the DFLM [48] structure functions were used, without any nuclear shadowing. The result is shown in Fig. 17 for two different initial temperature values, corresponding to our conservative density estimate for the lower and to about twice that density for the higher value. We conclude that for the lower density, there does not seem to be a clear-cut window for the observation of thermal dileptons; this, incidentally, is in accord with the observation that the mass distribution for $1.7 \leq M \leq 3.5$ GeV in present dilepton data from nuclear collisions [49] is in accord with the functional form found in $p - p$ Drell-Yan production. At the higher density, the situation improves somewhat. All in all, the usefulness of thermal dileptons as thermometer for strongly interacting matter is not a priori clear; if we can get to densities at or above the upper end of our estimated average, it appears possible.

For direct photons, the main competition at low momenta comes from the decay of hadrons, mainly $\pi^0$ and $\eta$ [50]. At high momenta, there are in addition direct photons from
2.5.2 The Spectral Analysis of $c\bar{c}$ and $b\bar{b}$ States.

In view of the problems encountered for thermal dileptons and photons as tools to probe the primordial features of strongly interacting matter, another means of analysis would certainly be very helpful. This may be provided by studying the spectra of heavy quark bound states ($c\bar{c}$ and $b\bar{b}$) produced in heavy ion collisions. A suppression of the $J/\psi$ signal relative to the Drell-Yan continuum had in fact been predicted as a signature for deconfinement [36] and was subsequently observed by the NA38 collaboration at the CERN-SPS [51]. The observed features have in the meantime also been accounted for by absorption in dense hadronic matter, coupled with initial state parton scattering [52][53]. The effect thus does seem to establish the production of dense, strongly interacting matter; to check whether this matter is already deconfined or still in a hadronic state, further experimental study is required. But we can use it as a starting point for a spectral analysis of QCD matter [12], very similar to the spectral analysis of stellar matter used in astrophysics. Stellar matter emits radiation containing spectral lines from the excitation and ionisation of various elements. The hotter the matter is, the lower is the intensity of the spectral lines from low-lying atomic excitation/ionisation states: with increasing temperature, these states become “suppressed” by thermal excitation. In much the same way, we expect “cold” QCD matter to show charmonium and bottomium signals at the same rate, relative to the Drell-Yan continuum, as in $p - p$ collisions; an increase in temperature should then lead to $J/\psi$ and $\psi'$ suppression, and still further increase to stronger suppression and to that of higher mass bound states (Fig. 19). Different suppression mechanisms (deconfinement, absorption) will also lead to different suppression patterns in energy density and in the

![Figure 18: Production rates for direct (thermal) photons [46] compared to photons from hadron decay [50] and Compton scattering.](image)
$P_T$ of the bound states. In Fig. 20, we show the predicted $P_T$ distributions for $J/\psi$ production in $S-U$ and $Pb-Pb$ collisions at the SPS, based on deconfinement, absorption, and absorption with initial state parton scattering as suppression mechanisms. We note that the present $S-U$ data are in accord with all three schemes; considerable differences in the predictions arise, however, already for the large $P_T$ behaviour in $Pb-Pb$ at the SPS. In Fig. 21, the suppression patterns of the $\psi$ at SPS, RHIC, and LHC energies are shown, as predicted by deconfinement. The difference between the suppression patterns by deconfinement and absorption is in Fig. 22 illustrated for $\psi$ and the $\Upsilon$ production at the LHC. The basis for the proposed analysis is certainly not yet complete, however: at TeV energies and large $P_T$, there is abundant $J/\psi$ production from $B$ decay [54], and much of the observable $\Upsilon$ production will come from $\chi_b$ decay [55]. Both these effects have to be taken into account in a realistic analysis of charmonium and bottomonium production in high energy heavy ion collisions.

Finally, let us consider for what energies and luminosities such a spectral analysis is at all possible. We measure the suppression of the different bound states relative to the Drell-Yan continuum, whose overall production rates are unaffected by the nuclear environment. Therefore we want to know under what conditions the Drell-Yan production of lepton pairs in the $J/\psi$ and the $\Upsilon$ range is feasible. To obtain an estimate of the Drell-Yan rates, we use a scaling cross-section form [56] for $p-p$ collisions

$$
(d^2 \sigma / dM^2 dy)_{p=0} = 3.75 \times 10^{-5} M^{-4} \exp(-15 \sqrt{T}).
$$

(18)

It fits all existing $p-p$ data; to obtain the rates for central $A-A$ collisions, we simply multiply it by $A^2$. Scaling violations are expected to increase the results somewhat [47], and shadowing will reduce it by up to 50% at high energies [45]. Neglecting these effects, we get the cross-sections and rates shown in Table 2. In general, there is a danger that the increase in cross-section with energy is compensated by the lower luminosity of high energy colliders; the need for high luminosities can thus not be emphasized enough. To get a rough estimate of the production rates for heavy quark resonances, assuming a mass resolution of about 100 MeV, we can multiply the Drell-Yan rates by a factor $10^2$ for $J/\psi$ and $\Upsilon$, by a factor 10 for $\psi'$ and $\Upsilon'$. On the other hand, the restriction to central collisions
Figure 20: $J/\psi$ suppression as function of transverse momentum, in central $S-U$ (a) and $Pb-Pb$ (b) collisions at the SPS. The behaviour for by deconfinement (QGP), absorption (HG), and absorption after initial state scattering (IPS) is calculated with parameters tuned for $S-U$ collisions [12]. Data in (a) is from [51].

Figure 21: $\psi$ suppression by deconfinement, as function of transverse momentum, for central collisions at SPS, RHIC, and LHC energies; from [12].

Figure 22: $\psi$ and $\Upsilon$ suppression for central collisions at LHC energy, by deconfinement (QGP) and absorption (HG); from [12].
and the experimental acceptance will reduce the rates by at least a factor 10, perhaps by as much as $10^2$. Thus a study of the $b\bar{b}$ regime may well be possible only at the LHC.

In terms of the rates just discussed, we can illustrate particularly well the effect of dense strongly interacting matter on charmonium and bottomium production. Without suppression, we should expect $20\ 000 - 100\ 000\ \psi$'s per month at the LHC, produced predominantly at small $P_T$; with suppression by deconfinement or absorption, there should be essentially none in the range $P_T \leq 3 - 4\ \text{GeV}$.

### 2.5.3 Jet Quenching and Jet Structure

The production of "hard" jets at large transverse momenta is, just as Drell-Yan or $c\bar{c}/b\bar{b}$ production, a process which in its early stage should be describable by perturbative QCD and hence be understood. We can therefore hope to learn something about dense strongly interacting systems by studying the effect they have on this process. In a $p - p$ collision, a hard parton will be emitted and subsequently hadronize by forming $q\bar{q}$ pairs as it passes through the vacuum. This process leads for the parton to a certain energy loss per unit rapidity. In dense matter, this loss will presumably increase ("jet quenching"), but it should be quite sensitive to the state of matter [38]. Recently, the energy loss for jet production in hadronic matter was compared to that in a quark-gluon plasma [57],[58]; the latter is found to give very little damping ("jet unquenching"), so that a change in the jet production could signal a change of state of the medium through which it has passed. More detailed studies comparing jet production in $p - p$ collisions to that in hadronic matter and in a plasma are certainly needed.

Another possible tool to probe the early history of heavy ion collisions may be the study of the fractal structure of multiparticle production. Jet cascades, again presumably describable by perturbative QCD in their early stages, would lead to intermittent hadron distributions, if these cascades were self-similar [59]. For most types of cascade, this results in a multifractal pattern, whereas it was shown that continuous phase transitions in spin systems lead to fractal behaviour [60]. It has therefore been proposed recently [61],[62] that the fractal dimension of the production process provides information about the thermal or non-thermal nature of the primordial state.

### 2.5.4 Primordial Remnants in Hadronic Observables.

Finally we want to consider the possibility that the observed hadronic final state in some specific features still reflects its early history.

If we produce a thermalised system which subsequently undergoes hydrodynamic expansion, then this should be reflected also in the momentum spectrum of the observed

<table>
<thead>
<tr>
<th>Energy</th>
<th>SPS (17 GeV)</th>
<th>RHIC (200 GeV)</th>
<th>LHC (6300 GeV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\mathcal{L}$ [cm$^{-2}\text{s}^{-1}$]</td>
<td>$10^{26}$</td>
<td>$10^{26}$</td>
<td>$10^{27}$</td>
</tr>
<tr>
<td>$(d\sigma/dM^2dy)_{\gamma=0}^{M=4}$</td>
<td>$1.9 \times 10^{-4}$</td>
<td>$4.8 \times 10^{-3}$</td>
<td>$6.5 \times 10^{-3}$</td>
</tr>
<tr>
<td>$(d\sigma/dM^2dy)_{\gamma=0}^{M=9}$</td>
<td>$8.7 \times 10^{-5}$</td>
<td>$1.3 \times 10^{-4}$</td>
<td>$2.4 \times 10^{-4}$</td>
</tr>
<tr>
<td>(events/month)$_{\gamma=0}^{M=4}$</td>
<td>$4.9 \times 10^3$</td>
<td>$1.2^2$</td>
<td>$1.7 \times 10^4$</td>
</tr>
<tr>
<td>(events/month)$_{\gamma=0}^{M=9}$</td>
<td>$2.3$</td>
<td>$3.4$</td>
<td>$6.2 \times 10^2$</td>
</tr>
</tbody>
</table>
particles: they should gain in momentum from the collective flow of the medium, and the
gain should be greater for heavier particles [63]. As already indicated, it is not so clear,
however, whether such flow can be detected at energies below the LHC. The effect of a
first order phase transition on such flow has been proposed [39] as a way to check if the
system has passed from a quark-gluon plasma to a hadron gas. If we increase the initial en-
ergy density of the collision, collective flow will increase the average transverse momentum
of the emitted secondaries. In case of a first order transition, there will be a coexistence
regime in energy density, for which the pressure (and hence flow) remains unchanged. This
predicts a flattening of the $p_T$ distribution of the secondaries when the multiplicity of the
reaction is increased. A further increase of multiplicity (and hence initial energy density)
should eventually bring us into the plasma regime, with an increasing pressure and hence
increasing $<p_T>$. Before we can judge whether a transverse momentum pattern showing
such behaviour indeed provides a indication for a phase transition, we must understand
more clearly the role of minijets in $p-p$ or $p-\bar{p}$ collisions. At the Tevatron [17], an increase
and subsequent flattening of $<p_T>$ of the type just discussed is observed for pions (Fig.
23). An increase is also obtained in event generator studies (Fig. 24), using PYTHIA
(with minijets and no rescattering) as well as VENUS (with rescattering but no minijets).
Before any interpretation of the energy density variation of $<p_T>$ in nuclear collisions
becomes possible, we clearly have to understand better the behaviour in hadron-hadron
interactions, in particular the role of jet and mini-jet production [64].

Specific features of strangeness production in heavy ion collisions may constitute an-
other hadronic observable capable of reflecting primordial properties, in particular whether
the different quark species were present according to their ideal gas ratios, i.e., were in
"chemical" equilibrium [32]. An overall increase of strangeness production in comparison
to the average value for $p-p$ collisions is, as already discussed, simply a consequence of
multiple rescattering and arises as well for a hadron gas as for a quark-gluon plasma [9][26].
Moreover, it is also observed for high multiplicities in $p-\bar{p}$ collisions at the Tevatron [17].
Nevertheless, for certain channels, hadron gas and quark-gluon plasma may lead to very
different predictions. Thus the $\phi$ remains on the hadronic level just a slightly more massive

![Figure 23: The average $P_T$ of $K^\pm$ and $\pi^\pm$ in $p-\bar{p}$ collisions, as function of the charged particle multiplicity; from [17].](image-url)
ground, became particularly enticing [67]. Today, after a number of rather detailed studies of non-vanishing baryon number, containing all three quark species in equal amounts and thus electrically neutral. The observation of baryonic states with $A >> 1$ and $Z/A \approx 0$ would certainly constitute a very striking signal for the production of quark-gluon matter.

\[ \omega; \text{ on the quark level, it is a } s\bar{s} \text{ bound state and thus could reflect an enhanced presence of strangeness in the early stages of the system. Similarly, the production of multistrange antibaryons should be enhanced, if the primordial state was a quark-gluon plasma [32]; in both cases, a comparison of the ratios for a hadron gas and for a quark-gluon plasma would be very interesting.} \]

Finally we note the possibility of producing and observing hitherto unknown forms of hadronic matter. If up, down and strange quarks are present in equal amounts in an initial state of chemical equilibrium, then bubbles of strange matter may arise as stable end products in the subsequent hadronisation [40]. These “strangelets” would be systems of non-vanishing baryon number, containing all three quark species in equal amounts and thus electrically neutral. The observation of baryonic states with $A >> 1$ and $Z/A \approx 0$ would certainly constitute a very striking signal for the production of quark-gluon matter.

3 Coherent Heavy Ion Collisions

High energy heavy ion collisions in general lead to abundant production of hadronic secondaries; because of just this feature we hope to use such collisions for the study of statistical QCD. However, at large impact parameters $b > 2R_A$, where $R_A$ denotes the nuclear radius, the ion-ion interaction is predominantly electromagnetic. In coherent collisions, in which the ion acts as a whole, this interaction becomes very strong, since it is proportional to the nuclear charge: compared to $p - p$ collisions, we gain a factor $Z^4$, i.e., for $Pb - Pb$ a factor $5 \times 10^7$, in the photon-photon interaction cross-section. Clearly this factor will be reduced by coherence requirements; nevertheless, it stimulated considerable interest in the possibility to use coherent heavy ion collisions for the production of Higgs bosons and supersymmetric particles [65],[66]. In the mass range above about 80 GeV (the upper limit for LEP II) and below about 160 GeV (the threshold for $ZZ$ production), any Higgs production in $p - p$ colliders would be severely contaminated by hadronic background. Hence a possible alternative, using coherent heavy ion collisions without any such hadronic background, became particularly enticing [67]. Today, after a number of rather detailed studies
[68],[69],[70],[71] this possibility does not seem so promising any more. Let us rather briefly look at what difficulties arise.

The Higgs produced in a $\gamma\gamma$ interaction would subsequently decay into a $b\bar{b}$ pair; hence the signal would be two high $P_T$ $b\bar{b}$ jets, without any associated “debris” from hadronic $A - A$ interactions, to assure coherence. For a Higgs of 100 GeV mass, this leads for $Pb - Pb$ collisions at the LHC to a cross-section of about 100 pb [67]. The upper bound for the ion luminosity at the LHC is [14] $\mathcal{L} \simeq 10^{28}$cm$^{-2}$s$^{-1}$; this would give about 10 events per year. It was then pointed out, however, that even after coherent Higgs production, the ions could still interact, leading to hadron production; hence this part of the cross-section must be experimentally vetoed [68]. As a result, the accepted cross-section is reduced by a factor 2–5 [68] - [71], as shown in Fig. 25.

![Figure 25: Higgs production cross-section as functions of Higgs mass $M_H$, without corrections for final state ion-ion interactions (DEZ) and with such corrections (BF,CJ).](image)

Furthermore, the Higgs decay into a $b\bar{b}$ pair has a strong irreducible background of direct $b\bar{b}$ production through $\gamma\gamma$ interactions. Since the direct production is peaked along the $\gamma\gamma$ collision axis, a restriction to large $P_Tb$ jets increases the signal-to-background ratio. At best, however, it becomes about $1/5$ (see Fig. 26). For the 100 events necessary to provide a 4 standard deviation signal in this case, a luminosity of $10^{30}$ cm$^{-2}$s$^{-1}$ would be needed—well above the possible LHC limit [14].

In addition, there is a strong $c\bar{c}$ background from $\gamma\gamma \rightarrow c\bar{c}$; the $2/3$ charge of the $c$ makes the $c\bar{c}$ cross-section a factor 16 larger than that for the $b\bar{b}$. Hence an excellent $b$ identification is necessary.

Finally, we have cross-sections for hadronic $b\bar{b}$ production, such as $gg \rightarrow b\bar{b}$ or $\gamma g \rightarrow b\bar{b}$, which are by orders of magnitude larger than the photon induced production [72]. Hence an excellent veto against beam jets is absolutely essential.

After taking all these factors into account, we conclude [70] [71] that a Higgs search in $Pb - Pb$ collisions at the LHC, with a luminosity of $\mathcal{L} = 10^{28}$cm$^{-2}$s$^{-1}$ or less, is not feasible.

A possible search for supersymmetric particles or a production of $W$ pairs turned out to be of as little promise [71]. For new charged particles, the mass range accessible to the LHC falls below that for LEP II. As far as $W^+W^-$ production is concerned, the top LHC ion luminosity leads to some 50 – 100 pairs per year. The produced $W$’s decay
Fig. 26: Higgs production cross-section with transverse momentum cut $P_T/M_{bb} > 0.4$, compared to hadronic $b$-quark background with same cut [70].

predominantly into jets however, which would make it difficult to reach sufficient statistics for a systematic study of the process $\gamma\gamma \rightarrow W^+W^-$. Finally we want to comment on the possibility to use diffractive heavy ion interactions for a Higgs search at the LHC [71]. Although the cross-section for Pomeron-Pomeron induced Higgs production is a factor $10^3$ higher for $Pb - Pb$ than for $p - p$ collisions, the attainable luminosity in the ion-ion mode is a factor $10^{-6}$ lower, making a Higgs search with $A - A$ collisions uninteresting.

The detailed considerations of coherent $A - A$ collisions at the LHC, carried out over the past year, thus only reinforce our earlier statement: the raison d'être of high energy heavy ion physics is the study of strong interaction thermodynamics.

Conclusions

We have seen that high energy heavy ion colliders can provide us with the best possible tool for the experimental study of statistical QCD. The LHC is moreover unique in several aspects:

- the expected energy density is well above the critical value for deconfinement and may be high enough to produce an asymptotically free quark-gluon plasma;

- the expected energy density should be high enough to observe any collective flow of the expanding dense matter produced in the collision;

- over half of the available rapidity range will have essentially zero baryon number density, even if there is enhanced stopping in $A - A$ collisions;

- the design luminosity should allow a full spectral analysis of both $c$ and $b$ bound states.

In closing, let us emphasise once more the two requirements which at present appear crucial for an optimal use of the LHC in the study of QCD thermodynamics: the integrated,
as well as the design luminosity must be sufficiently large, and it should be possible to accelerate ions of different $A$, in order to be able to vary the initial energy density.

**Acknowledgements**

This report is a physics assessment carried out for the ECFA Working Group HEAVY ION PHYSICS AT THE LHC. As such, it is to a large extent based on the work done in the various subgroups of this Working Group; in the APPENDIX, a list is shown of all those who have contributed at various stages to the work of the project, and who have worked very hard to obtain the results which I have presented here. To all of them, I want to express my sincere appreciation for their help in preparing this assessment. Many results will be given in more detail in Volume III of these Proceedings, and a discussion of the experimental aspects will be given by H. Specht in this volume. — In addition, I want to express particular thanks to U. Heinz, E. Papageorgiu, V. Ruuskanen and D. Zeppenfeld for helpful comments and a critical reading of parts of the manuscript.

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214

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   U. Heinz, Report at this Workshop, Proceedings Vol. III


APPENDIX

ECFA Working Group
HEAVY ION PHYSICS AT THE LHC
Conveners: H. Satz and H. Specht

QCD Thermodynamics


Coherent Heavy Ion Collisions