The quark-gluon plasma at RHIC*

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I present a theory-guided review of RHIC data, arguing that they provide strong evidence for formation of a thermalized quark-gluon plasma at RHIC. Strong radial flow reflects high thermal pressure in the reaction zone. Large elliptic flow proves that the pressure builds up quickly and the system thermalizes on a very short time scale of < 1 fm/c. The observed hadron yields are consistent with statistical hadron formation from a quark-gluon plasma, followed by immediate chemical decoupling due to strong radial expansion. The observed suppression of jets appears to confirm the predicted large energy loss suffered by hard partons moving through a quark-gluon plasma; more work is required to quantitatively understand this effect. Source size measurements using two-particle correlations do not seem to fit into this picture; the origin of this discrepancy (“HBT puzzle”) is presently not understood.

1. THE LITTLE BANG: QGP AT RHIC

A quark-gluon plasma (QGP) is a thermalized system of deconfined quarks, antiquarks, and gluons. As such it has thermodynamic pressure \( P = \varepsilon c_s^2 \) where \( \varepsilon \) is the energy density and \( c_s \) is the sound velocity. Perturbative QCD gives \( c_s^2 = 1/3 \) at leading order, and lattice QCD confirms [1] that for \( T > 2T_c \) (where \( T_c \approx 170 \text{ MeV} \) is the critical temperature for color deconfinement in QCD [2]) about 80-85% of this value is reached. In heavy ion collisions the reaction zone is surrounded by vacuum with zero pressure; thus, if the collision fireball contains a QGP, the pressure gradient near the surface will lead to collective expansion (“flow”), in particular transverse to the beam in which direction the nuclear matter was initially at rest. Therefore, collective transverse flow is an unavoidable consequence of QGP formation, and the data must show it if a QGP has been created.

The converse is not necessarily true unless the observed flow is so strong that only a QGP could have generated sufficient pressure over a sufficiently long time to create it. In order to assess whether the latter is the case one exploits three facts:

(i) Due to incomplete stopping of the two nuclei at high collision energies, the reaction zone expands quickly in the beam direction, thereby rapidly cooling and diluting the matter inside; hence the pressure has only a limited time interval available to generate flow in the transverse directions. This time is the shorter the lower the initial energy density; one can therefore establish an upper limit for the amount of “radial” transverse flow that can be generated if the initial energy density (at the time of approximate thermalization) was never significantly above the critical value for color deconfinement (about 1 GeV/fm³).

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(ii) By energy conservation, the initial energy density is related to the final transversely emitted energy per unit rapidity, $dE_T/dy$, via geometric factors and the time of thermalization [2]: the shorter the thermalization time, the higher the initial energy density at fixed (measured) final $dE_T/dy$. The Bjorken formula [2] $\varepsilon(\tau_{\text{therm}}) = \frac{1}{\tau_{\text{therm}} \pi (R^2) dE_T/dy}$ is actually an underestimate of the initial energy density since it neglects longitudinal work done by the pressure during the expansion [3]. The measured values for $dE_T/dy$ at RHIC [7] are consistent with subcritical initial energy densities only if one assumes that thermalization takes at least $\tau_{\text{therm}} \geq 5 \text{ fm}/c$.

(iii) On the other hand, this thermalization time scale can be constrained by the measured elliptic flow in noncentral collisions [1]. In such collisions, the nuclear reaction zone is initially deformed in coordinate space: a cut transverse to the beam direction looks like an almond whose longer side points perpendicular to the reaction plane. For a given collision centrality, extracted e.g. from the measured total multiplicity $dN/dy$, this deformation can be calculated from the overlap geometry [5]. Thermalization then leads to pressure gradients which are anisotropic and larger in the short direction of the almond, causing a faster growth of the transverse flow into the reaction plane than perpendicular to it. This mechanism leads to an anisotropy in the final transverse momentum distribution, making it flatter in the direction of the impact parameter vector than perpendicular to it [6]. The elliptic flow coefficient $v_2$, defined as the second harmonic coefficient in an azimuthal Fourier decomposition of the measured transverse momentum spectrum, quantifies this anisotropy. Without reinteractions among the produced particles, no such momentum anisotropy would arise; rescattering in the hot fireball matter is required to transfer the initial spatial anisotropy into momentum space. Microscopic calculations [8] show that, for a given initial spatial deformation, the generated amount of $v_2$ is a monotonic function of the mean free path (or the product of density and scattering cross section) in the fireball; the maximal value is reached in the hydrodynamic limit of zero mean free path [9]. If thermalization is delayed, the initial spatial deformation disappears spontaneously by transverse free-streaming [10], thereby reducing the maximal possible momentum anisotropy. Thus, the measured elliptic flow at given collision centrality provides immediately an upper limit for the thermalization time scale, with larger $v_2$ values corresponding to smaller values for $\tau_{\text{therm}}$. In the hydrodynamic limit the extracted $\tau_{\text{therm}}$ depends parametrically on the stiffness of the equation of state (velocity of sound $c_s^2$), with smaller sound velocities resulting in shorter thermalization times at fixed $v_2$ [10]. The weakest upper limit for $\tau_{\text{therm}}$ is thus extracted from the $v_2$ data by taking the largest reasonable value for the speed of sound, $c_s^2 = \frac{1}{3}$, corresponding to an ideal gas of massless quarks and gluons.

Experimentally, the strength of the radial flow can be extracted from the transverse mass ($m_\perp = \sqrt{m^2 + p_T^2}$) spectra of a variety of different mass hadrons by fitting them to a common flow spectrum [11]. At high $m_\perp \gg m$, the transverse flow flattens the thermal spectra by a simple common blueshift factor, and all spectra approach the same asymptotic slope from which the decoupling temperature and flow velocity cannot be separated. At low $m_\perp < 2m$, however, the flow induces an even stronger flattening which increases with the rest mass of the particles. Heavier hadrons thus develop a visible shoulder at small $m_\perp$, and this feature can be used to determine the average transverse flow velocity uniquely. An excellent systematic study of the shapes of the $m_\perp$ spectra of
π±, K±, φ, p, ¯p, Λ, ¯Λ, Ξ, ¯Ξ, Ω, ¯Ω, and d from Pb+Pb collisions at beam energies of 40, 80 and 160 A GeV at the SPS was recently presented by M. van Leeuwen at the Quark Matter 2002 conference [12], showing perfect consistency of all these spectra with a two-parameter flow fit [11] with kinetic freeze-out temperature $T_f = 120 - 130$ MeV and average transverse flow velocity $\langle \beta_\perp \rangle = 0.4 - 0.5$ ($T_f$ and $\langle \beta_\perp \rangle$ are anticorrelated in these fits). Even the $\Omega$ and $\bar{\Omega}$ show a clear flow shoulder and are consistent with these low freeze-out temperature and large flow values. This contradicts earlier conclusions drawn from simple exponential fits to $\Omega$ spectra measured at higher $m_\perp$ values by the WA97 Collaboration [13] which did not show clear evidence for the flow shoulder and which led to the suggestion [14] that $\Omega$ and $\bar{\Omega}$ are too heavy and too weakly coupled to the expanding pion fluid to pick up much of the transverse flow created during the late hadronic stages of the collision, allowing them to decouple earlier, i.e. at higher temperature and smaller flow. The new data [12] show that the $\Omega$ and $\bar{\Omega}$ fully participate in the flow and essentially decouple together with the rest of the hadron fluid. They also demonstrate the fallacies connected with characterising the spectra by a single slope parameter instead of comparing their entire non-exponential shape with a proper flow parametrization.

The hadron data at RHIC are already almost as rich as at the SPS, but not yet all published in final form. Available flow analyses of pion, kaon and (anti)proton spectra [14, 15, 27] give a similar range of freeze-out temperatures and average flow velocities as at the SPS, but in a different combination (either $T_f$ or $\langle \beta_\perp \rangle$ is about 5-10% higher than at the SPS). This results in about 5% flatter pion spectra while the effect on the proton slope at small $p_\perp$ is much larger and at least 25%. Again the flow fits describe all available spectra very well up to transverse momenta of about 2-3 GeV (which covers more than 99% of all hadrons), and only at higher $p_\perp$ one begins to see evidence for the power-law tails expected from hard QCD processes. Also at RHIC, preliminary $\Omega$ and $\bar{\Omega}$ spectra [18] are consistent with the $\Omega$ fully participating in the hadronic flow. A comparison with hydrodynamic predictions published in [19] works well if the $\Omega$ is assumed to decouple at $T_f \simeq 135 - 140$ MeV, i.e. only slightly before the pions freeze out at $T_f = 125 - 130$ MeV [21], but not if one assumes decoupling already at hadronization, $T_f = T_c = 170$ MeV, where the hydrodynamic model has not yet developed enough transverse flow.

The extracted flow velocities at or above half the speed of light demonstrate that the collision fireball indeed undergoes a violent explosion – the “Little Bang”. It appears to be impossible to obtain such large flow velocities without assuming initial energy densities and pressures well above the critical value for deconfinement. This is already true at the SPS [27] but, as I show next, much more convincingly so at RHIC.

A decisive measurement is the observation at RHIC, first made by STAR [22] and then confirmed by PHENIX [23, 24] and PHOBOS [25], that the elliptic flow $v_2$ in non-central collisions is large and, for transverse momenta below $p_\perp \simeq 2$ GeV, almost exhausts the upper limit provided by earlier hydrodynamic predictions [14, 15, 20] (see Fig. 1). The agreement between theory and data requires that the hydrodynamic evolution starts no later than about 1 fm/$c$ after nuclear impact [14, 19] (the successful predictions in [14, 20] use $\tau_{\text{therm}} = 0.6$ fm/$c$). Since the spatial deformation responsible for the creation of flow anisotropies quickly decreases, $v_2$ develops and saturates early in the collision [4, 10], long before hadrons decouple. Hence, the conclusion that the measured large $v_2$ implies early thermalization is not changed [27] if the somewhat unrealistic sharp “Cooper-Frye
The “freeze-out” used in the hydrodynamic approach is replaced by a proper kinetic treatment of the freeze-out stage [28]. The agreement of $v_2(p_{\perp})$ out to transverse momenta of about 2 GeV, including the predicted [19] splitting of $v_2$ with the hadron rest mass (see Fig. 1), shows that the bulk of the matter (i.e., > 99% of the emitted hadrons) behaves hydrodynamically. (One should note, however, that the measured rapidity dependence of $v_2$ cannot be reproduced by existing hydrodynamic models [29].) Microscopic models share this success only if they approach the hydrodynamic limit by invoking unusually strong rescattering among the fireball constituents [8, 30, 31], far above the level conventionally expected from perturbative QCD. It is important to realize that, given the observed final multiplicity and transverse energy per unit rapidity, the maximum energy density in the fireball at $\tau_{\text{therm}} \approx 0.6 \text{ fm}/c$ is about 25 GeV/fm$^3$ in central collisions. Even after averaging over the transverse plane this is still more than an order of magnitude above the critical value for deconfinement and corresponds to about twice the critical temperature! Unless a completely different mechanism for creating the elliptic flow can be found, the conclusion seems unavoidable that a thermalized QGP at $\varepsilon > 10 \varepsilon_c$ and $T \geq 2T_c$ is created at RHIC which lives for about 5-7 fm/c before becoming sufficiently dilute to form hadrons. Perturbative mechanisms [32] seem unable to explain the phenomenologically required very short thermalization time scale, pointing to strong non-perturbative dynamics in the QGP even at or above $2T_c$.

2. STATISTICAL HADRONIZATION: MEASURING $T_c$

The quark-hadron phase transition is arguably the most strongly coupled regime of QCD. Soft hadronization happens through a multitude of different channels and is therefore most efficiently described in a statistical approach, as realized by Hagedorn more than 35 years ago [33]. The microscopic processes are only constrained by local conservation laws (valid inside causally connected volume regions $\Delta V$) for energy, baryon number and strangeness. Maximizing the entropy subject to these constraints results in local thermal and chemical “equilibrium” distributions for the hadrons [34] whose local temperature $T_{\text{chem}}$ and chemical potentials $\mu_B, \mu_S$ arise from Lagrange multipliers and reflect the local energy, baryon and strangeness densities at hadronization. This “equilibrium” is not
achieved kinetically (by hadronic rescattering), but statistically (by interference of many channels). It does not require the prehadronic state to be thermalized, although very strong deviations from thermal equilibrium (e.g. at high $p_\perp$) may survive the hadronization process. The statistical approach is only expected to work for soft hadron production.

We know from lattice QCD that a hadronic equilibrium distribution at $\varepsilon > \varepsilon_c$ is unstable against deconfinement, so statistical hadronization can only proceed once the energy density has dropped to $\varepsilon_{\text{had}} = \varepsilon_c \approx 1 \text{GeV/fm}^3$ \cite{33}. The resulting equilibrium distribution will thus be characterized by $T_{\text{chem}} = T_c$. Hadronic cascade models have shown \cite{36} that, at small net baryon density, the total particle density below $T_c$ is too small and the expansion of the collision fireball is too rapid to maintain chemical equilibrium by inelastic hadronic collisions. Hence, the hadron abundances decouple directly at hadronization, and the chemical freeze-out temperature extracted from a statistical model fit to the measured abundances reflects directly the (de)confinement temperature, $T_{\text{chem}} = T_c$ \cite{37}.

![Figure 2](image.png)

Figure 2. Particle ratios from 130A GeV Au+Au collisions, measured by the STAR Collaboration at RHIC \cite{18} and fitted by M. Kaneta to a thermal statistical model \cite{38,39}. Hadron ratios measured by the BRAHMS and PHENIX Collaborations agree with the shown STAR data.

Figure 2 \cite{18} shows that the RHIC data confirm this expectation, just as it was previously confirmed at the SPS \cite{40}: the chemical decoupling temperature extracted from the particle ratios is consistent with the critical temperature for deconfinement, $T_c \approx 170 \text{MeV}$, extracted from lattice QCD simulations \cite{1}. Fitting the maximum entropy parameters to some of the more abundant particle yields one is able to reproduce all measured particle ratios, including those involving the rare multistrange (anti)baryons; this is quite impressive. **Thermal freeze-out** of the momentum distributions is delayed by strong quasi-elastic scattering via hadron resonances, such as $\pi + N \rightarrow \Delta \rightarrow \pi + N$, $\pi + N \rightarrow \Delta \rightarrow \pi + N$, $\pi + K \rightarrow K^* \rightarrow \pi + K$, etc., which do not modify the total measured yields of pions, kaons, and nucleons but continue to adjust their momentum distributions to the falling temperature and growing collective transverse flow until also resonance scattering ceases. The momentum spectra thus reflect a lower freeze-out temperature $T_f < T_{\text{chem}}$, as seen in section \cite{4}. Note that only the total abundances (after all decays of unstable resonances are taken into account, e.g. $N_\pi^{\text{tot}} = N_\pi + 2N_\rho + N_\Delta + N_{K^*} + \ldots$) freeze out at $T_{\text{chem}}$ while the fraction of pions stored in resonances still changes: due to their strong coupling to
the cooling pion fluid, the resonance abundances keep readjusting to the decreasing temperature, and the larger their pionic decay width the later they decouple. This implies that the measured $K^*/K$, $\Delta/N$, $\Lambda^*/\Lambda$ etc. ratios should generally be smaller than their chemical equilibrium values at $T_{\text{chem}} = T_c$ and, when themselves translated into a decoupling temperature, reflect more closely the spectral temperature $T_f$ than $T_c$ [41]. First results on resonance/ground-state ratios were reported at Quark Matter 2002 [42], and this prediction will be checked soon.

3. JET QUENCHING: STRONG PARTON ENERGY LOSS IN THE QGP

In 1982 Bjorken [43] suggested that fast partons travelling through a QGP might lose large amounts of energy by elastic scattering with the plasma constituents, resulting in the suppression of jets from the interior of the collision fireball in relativistic heavy ion collisions. Although the energy loss mechanism envisaged by him turned out to be ineffective, his qualitative prediction appears to be impressively confirmed by recent RHIC data. A good review of the theory of parton energy loss was recently given by Baier [44], and I refer to his talk for references. In a series of papers in the early 1990’s, Gyulassy, Plümer and X.N. Wang identified radiative energy loss as the dominant jet suppression mechanism, and in 1995-1997 Baier, Dokshitzer, Mueller, Peigné and Schiff (BDMPS) realized that an important effect on the energy loss rate results from multiple color interactions of the radiated gluon with the colored plasma constituents [45]. In the limit of an optically thick quark-gluon plasma they found [45] that the energy loss $\Delta E$ of the fast parton increases quadratically with the distance $L$ travelled before escaping and hadronizing,

$$\Delta E \approx \frac{\alpha_s \mu^2}{2} \frac{L^2}{\lambda^2}, \quad \text{with} \quad \frac{\mu^2}{\lambda} = \rho \int dq_{\perp}^2 q_{\perp}^2 d\sigma dq_{\perp}^2,$$

(1)

where $\alpha_s$ is the strong coupling constant, $\rho$ is the plasma density, $\mu^2$ is the Debye screening mass for color electric fields in the plasma and $\lambda$ is the gluon mean free path. This was subsequently improved upon for plasmas with finite opacity by Gyulassy, Levai, and Vitev [46] and by Wiedemann [47], using an opacity expansion. For optically thin plasmas this leads to an important dependence of the energy loss $\Delta E$ on the energy $E$ of the fast parton [46] and to deviations from the quadratic path length dependence for small values of $L$. Both corrections reduce the predicted energy loss relative to the BDPMS result [41], in particular at “low” (SPS and RHIC) energies, but it remains still significantly larger in a QGP than in subcritical normal hadronic matter [44]. With these improved results one may nurture the hope that parton energy loss and jet quenching, if observed, can be used as a quantitative probe of the density and its early time evolution in quark-gluon plasmas created in heavy-ion collisions.

Two important observations at RHIC have moved this hope to the brink of reality. The first is the discovery of a suppression of high-$p_{\perp}$ particle production in central Au+Au collisions at RHIC. When studied as a function of collision centrality, the production rates of charged hadrons at high $p_{\perp} \gg 2$ GeV were found [48, 49, 50] not to follow the expected scaling with the number of binary collisions which characterizes hard QCD processes, but instead to continue to scale with $N_{\text{part}}$ (the number of wounded nucleons), just as in the soft low-$p_{\perp}$ regime where phase chereence and the Landau-Pomeranchuk-Migdal
effect are known to suppress particle production. (There seems to be an intermediate $p_\perp$ region, $1.5 \text{ GeV} < p_\perp < 3.5 \text{ GeV}$, where particle production grows a bit faster than $N_{\text{part}}$ but more slowly than the number of binary collisions.) From the Glauber model the number of binary $NN$ collisions is known to scale with $N_{\text{part}}^{4/3}$ while $N_{\text{part}}$ is proportional to the volume of nuclear matter participating in the collision. The observed suppression or “quenching” of hard particle production by a factor $N_{\text{part}}^{1/3}$ is thus proportional to the linear size (“radius”) of the nuclear overlap volume. It is qualitatively consistent with expectations from pQCD-based jet quenching calculations [51] and inconsistent with pQCD calculations without jet quenching. It is seen for charged and neutral pions [50] but apparently not for high-$p_\perp$ protons [52] (at least not in the presently accessible $p_\perp$ range for proton PID); to what extent this “lack of hard proton suppression” is simply a reflection of the observed increase of the $p/\pi$ ratio with increasing $p_\perp$ [53] which, at least up to $p_\perp \approx 2 \text{ GeV}$, is naturally explained by the stronger transverse flow effects on protons than on pions, is still unclear.

![Figure 3](image.png)

The second important observation, made by STAR [54] and (not yet as convincingly) by PHENIX [55], is that high-$p_\perp$ particle production is indeed due to jets, and that the observed overall suppression of high-$p_\perp$ particle production in central Au+Au collisions goes hand-in-hand with the complete suppression of the far-side partners of the observed jets. Perturbatively, hadron jets result from hard parton collisions leading to parton pairs with large, approximately balancing transverse momenta which fragment into pairs of back-to-back jets. In an angular distribution around the beam axis, jets appear in the two-particle correlation function as two peaks separated by $180^\circ$. By triggering on a high-$p_\perp$ particle and correlating it with other particles of $p_\perp > 2 \text{ GeV}$ (in order to suppress the uncorrelated soft background), these pairs of peaks were indeed identified in pp and peripheral Au+Au collisions at 200 A GeV at RHIC [54], but in central Au+Au collisions only the correlation peak at small relative angles is found. Already in peripheral Au+Au collisions the far-side peak at $180^\circ$ is much smaller than expected from a superposition of pp collisions [54, 55], and in central collisions it completely disappears after accounting for the collective azimuthal correlations generated by elliptic flow [54] (see Fig. 3).

While attempts to quantitatively understand these observations have only just begun [56], they suggest a very simple and intriguing picture similar to the one already proposed in 1982 by Bjorken [43]: fast partons formed inside the hot and dense collision zone
suffer severe energy loss and become part of the approximately thermalized, collectively expanding “soup” which emits particles at low $p_\perp < 2\,\text{GeV}$. Only if the primary hard collision happens near the surface of the reaction zone, the “outward moving” parton fragments to become a jet, very much like in $pp$ collisions (upper set of symbols in Fig. 3), whereas the “inward moving” parton is more and more efficiently absorbed by the medium as the nuclei collide more centrally and the reaction fireball volume increases (lower set of symbols in Fig. 3). For central Au+Au collisions and $p_\perp < 8 - 10\,\text{GeV}$ \cite{54}, the “inward moving” parton doesn’t make it to the other side with sufficient energy left to fragment into a recognizable jet. Jets are therefore only emitted from a relatively thin surface layer while the bulk of the fireball volume is opaque to jets.

This very qualitative picture may appear naive but it nicely explains as a surface/volume effect the missing factor $N_{\text{part}}^{1/3}$ in the observed scaling of high-$p_\perp$ particle production. It also allows for a simple geometric estimate of the elliptic flow $v_2$ at high $p_\perp$, resulting from the anisotropic emission of surface jets from the spatially deformed overlap region \cite{57}. As pointed out by Voloshin \cite{58}, this simple geometric estimate \cite{57} appears to be consistent with the STAR data for $v_2$ at high $p_\perp$. Of course, at very high $p_\perp$ one eventually expects the “inward moving” parton to emerge on the other side with sufficient energy left to form a jet. When this happens, its energy loss can be studied as a function of its path length through the medium, in order to check the energy loss formula (1) and its various published variants. Selecting events of fixed centrality (impact parameter), the path length can be controlled geometrically by the azimuthal angle between the jet and the reaction plane, and a new era of “jet tomography” \cite{59} will begin.

4. THE RHIC HBT PUZZLE

Unfortunately, space-time limitations do not allow for a detailed discussion of one piece of the puzzle which so far fails to seamlessly fit into the picture painted above: the two-particle Hanbury Brown – Twiss (HBT) correlation measurements. Since I reviewed these elsewhere (see \cite{20} and references therein), let me be brief here. HBT correlations can be used to probe the size, shape and dynamical state of the source at hadronic decoupling. STAR and PHENIX have produced mutually consistent results showing that the HBT radii in the sideward and outward directions (i.e. perpendicular and parallel to the transverse emission vector) depend strongly on the transverse pair momentum and are almost equal to each other. Otherwise successful dynamical models such as those discussed in Sec. 1 fail to reproduce this strong transverse momentum dependence and consistently give smaller sideward than outward radii, in contradiction with the data. They also overpredict the longitudinal radius. The longitudinal and outward radii can be decreased by reducing the total lifetime of the fireball, but this would require even faster thermalization to accommodate the observed transverse flow, and the sideward radius remains too small. Equal outward and sideward radii not only require a very opaque source (such as the hydrodynamic ones with unrealistic sharp Cooper-Frye freeze-out), but also a short lifetime and a breaking of longitudinal boost invariance, at least in the decoupling process. While these problems are generally believed to reflect insufficiencies in the description of the late decoupling stage which will not affect our understanding of the early collision (QGP) stage, we can’t be sure of this until a model is found that works.
5. OUTLOOK

The first two years of RHIC running have brought a rich harvest of hadron production data and, as I discussed, abundantly fulfilled our hopes and expectations of finding more and stronger evidence for the making of quark-gluon plasma in heavy-ion collisions. The observation of large elliptic flow, with its compelling interpretation in terms of fast thermalization, and the discovery of jet quenching are two important milestones which go significantly beyond what was earlier achieved at the lower SPS energy [60] and which bring us closer to a generally accepted “proof” of QGP creation. But much more is to come: only now, with RHIC finally running at full energy and luminosity (and, hopefully, for the full promised time per year) it is possible to address such hallmark measurements as thermal dilepton and direct photon emission and heavy quarkonium production, all of which play crucial roles in the early diagnostics of the QGP which we are apparently mass-producing at RHIC. While trying to solve the HBT puzzle and to quantitatively understand jet quenching, we are looking forward to these high-luminosity measurements and any surprises they may bring.

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