Abstract

This talk presents an overview of the theoretical contributions at the Hard Probes 2013 conference, held in Stellenbosch, South Africa, in November 2013.

Keywords: Heavy Ion Collisions, Quark Gluon Plasma, Quantum Chromodynamics

1. Introduction

Although quantum chromodynamics (QCD) has been recognized as the microscopic theory of strong interactions for over 40 years, its application to the study of complex reactions such as heavy ion collisions remains very challenging. These reactions involve a complicated mix of various momentum scales, ranging from very hard scales down to the confinement scale. For this reason, the theory of heavy ion collisions is fragmented into many different tools, effective descriptions and models, as illustrated in the figure. Arguably, the closest to QCD are lattice QCD and perturbation theory, but both are seriously limited in what questions they can address. Somewhat further away from QCD, one encounters some effective theories that approximate the QCD dynamics in a certain regime (high gluon density, heavy quarks, ...). Calculations based on the AdS/CFT correspondence allow one to study the strong coupling limit of certain cousin gauge theories, like the SUSY $\mathcal{N} = 4$ Yang-Mills theory. Finally, transport descriptions can be used at larger distance scales where the microscopic details are not so important and the dynamics is largely controlled by conservation laws.

Another important aspect of heavy ion collisions is that successful descriptions must include some mundane facts that have no connection with the aspects of QCD that motivate studying heavy-ion collisions, such as the geometry of
nuclei, how the nucleons are distributed inside nuclei and how they fluctuate, etc... All these things matter in the interpretation of heavy ion data, before one can get at the underlying QCD phenomena. The number of participants gives a good idea of the geometry of a given collision. It is not directly measurable, but in nucleus-nucleus collisions it is well correlated to some measurable quantities, such as the transverse energy or the deposit in zero degree calorimeters. A worrisome lesson from HP2013 is the fact that our handle on the geometry of proton-nucleus collisions is not very good. In these collisions, there is a rather weak correlation between the multiplicity and the centrality of the collision, which makes using pA collisions as a reference for AA collisions rather problematic. Because of this, understanding pA collisions has become a serious challenge in itself, quite bad news if one wishes to use them as a reference.

With this caveat in mind, the general idea of hard probes is to measure objects characterized by some large momentum scale (e.g. jets, photons, bound states containing heavy quarks), and to infer properties of the medium in which they propagate by studying how some of their properties are modified compared to collisions where one does not expect the formation of a quark-gluon plasma (e.g. proton-nucleus and proton-proton collisions). The theory contributions to HP2013 can be divided in roughly four topics:

1. Initial state. (16 contributions) This encompasses the very early stages of a collision, while the system is still dense enough to be characterized by a semi-hard momentum scale. This departs a bit from the general concept of hard probes, because in these studies the medium itself is treated perturbatively.

2. Electromagnetic probes. (7 contributions) Photons and dileptons, even when their momenta are not very large, are usually considered as part of the “hard probes” because their weak (electromagnetic only) interactions give them a mean free path much larger than the system size. Their spectrum therefore reflects the state of the system at the time of their production.

3. Energy loss. (16 contributions) Modifications of inclusive (identified or not) hadron spectra have been observed long ago in heavy ion collisions. More recently, these measurements have been extended to jets, which opens up new possibilities for studying the medium modifications of radiation.

4. Heavy flavors. (11 contributions) Rather unexpectedly, heavy quarks tend to lose a significant amount of energy in a dense medium, like their lighter siblings. Moreover, compact bound states made of heavy quarks can probe the local temperature of the medium.

The number of contributions in each of these subjects can be viewed as a vague indication of where the field is going, and what topics are considered to be in need of more work, or on the contrary are getting closer to a satisfactory state of understanding.

2. Initial state

The first step in studying hard probes is to determine their initial yield. In the case of jets and heavy flavors (for quark masses that are large compared to the saturation momentum), this can be done in the DGLAP framework, using nuclear parton distribution functions (NPDFs). In addition to having some shadowing (i.e. they are not a simple superposition of nucleon PDFs), these distributions can also be made transverse position dependent in order to account for the fact that the nuclear modification may depend on the local density of nucleons (I. Helenius, [1]).

When the typical momentum scale of the process under consideration is comparable to the saturation momentum, recombination corrections become important and cannot be included in the DGLAP framework. However, these effects can somehow be mimicked by introducing a form of “final state saturation” via an infrared cutoff. This is the essence of the EKRT model, that has now been extended to include NLO partonic cross-sections (R. Paatelainen, [2]).

At high transverse parton density, double parton scatterings can occur in a single collision. Ignoring quantum coherence effects, this can be modeled in terms of single parton scattering cross-sections [3]. However, the Color Glass Condensate is a more rigorous and systematic way to treat the physics of multi-parton interactions in QCD. It is based on the separation of the partons into “fast” partons that are described as light-cone color currents $J_{1,2}$ (one for each colliding projectile) and “slow” partons that are described by the usual quantum fields, glued together by an eikonal coupling: $$L = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} + A_{\mu} \cdot (J_{1}^{\mu} + J_{2}^{\mu}).$$ The currents $J_{1,2}$ reflect the instantaneous transverse positions and colors of the fast partons at the time of the collision. This can only be described probabilistically, with distributions
3. Electromagnetic probes

The production rate of photons and low mass dileptons by a hot plasma increases like $T^4$, which means that these probes are very sensitive to the temperature of the medium at the earliest stages. Moreover, their mean free path is much larger than the size of the system and therefore they can escape without any modification of their spectra. The only drawback of photons as probes of the quark-gluon plasma is that there are many other sources of photons:

\[ W[J_{1,2}] \] that depend solely on the nature of the projectile, and on the scale $\Lambda$ that separates the fast and slow partons. Like with conventional PDFs, the distributions $W[J_{1,2}]$ are non-perturbative, but their scale dependence is driven by a perturbative equation, the JIMWLK equation in the case of the CGC. Since the JIMWLK equation is a functional equation whose numerical solution is very time consuming, it is often replaced by a mean-field approximation, the BK equation, in phenomenological applications. The BK equation can be viewed as an approximation scheme that closes the equation for the dipole amplitude (that would otherwise depend on a 4-point function). More general truncation schemes have been proposed, such as the Gaussian truncation (A. Ramnath, [4], G. Jackson, [5]).

Due to the simplicity of their implementation, models based on the BK equation have been used to describe proton-nucleus collisions, where it is more legitimate to neglect final state effects. This can be done using $k_T$-factorization and BK evolution for the two projectiles (H. Mantysaari, [6]) or in a hybrid formulation where the BK evolution is applied only to the nucleus, while the proton is described with collinear factorization (B.-W. Xiao, [7]).

The BK equation had also the advantage of being well known at next-to-leading log accuracy. However, it is now becoming possible to include running coupling effects in the Langevin formulation of the JIMWLK equation (T. Laptev, [8]). Moreover, a full NLL JIMWLK equation has recently appeared on the horizon (S. Caron-Huot, [9] and M. Lublinsky, A. Kovner, Y. Mullan, [10]). Interestingly, simple improvements in the way the kinematics is treated in leading log evolution equations can account for a significant part of the NLL corrections (G. Beuf, [11]).

Another recent improvement to the JIMWLK equation is the possibility to tag some of the gluons that are radiated in the course of the evolution (E. Iancu, [12]), in order to compute multi-gluon correlations (this is necessary for instance when studying ridge-like correlations in proton-nucleus collisions). The practicality of such a generalized evolution equation is somewhat limited at the moment, since the required computational resources would scale as $S_\perp^{\# \text{gluons}}$, where $S_\perp$ is the number of lattice sites in the discretization of the transverse plane.

In nucleus-nucleus collisions, the expansion of the bulk of the matter is usually described by relativistic hydrodynamics, and models such as the CGC are used only to describe the pre-hydrodynamical stages of the collision, and to provide some initial data (energy density, pressure, flow velocity and stress tensor at some initial time). Besides the CGC, there are many models exist to compute these quantities with various degrees of sophistication (S. Jeon, [14]): Optical Glauber, MC Glauber, AMPT, NEXUS, MC-KLN, MC-rcBK, IP-GLasma (ordered here by the amount of QCD they contain). These models differ in the spatial scales at which they have fluctuations. In principle, by studying the event-by-event mapping from the initial shape harmonics $\tilde{e}_0$ to the final flow harmonics $v_n$, one could constrain the parameters that define these models, and at the same time determine transport coefficients such as the shear viscosity.

At Leading Order, the pressure tensor predicted by the CGC remains very anisotropic at all times. Since this is hard to accommodate within hydrodynamics, the matching to hydrodynamics usually uses only the CGC energy density and assumes an equilibrium equation of state to guess the corresponding pressure. However, it seems that resumming some higher order corrections in the CGC calculation has an important effect on the longitudinal pressure (T. Epelbaum, [15]). In addition, by a time comparable to the inverse saturation momentum, a non-zero flow velocity has developed (R. Fries, [16]), made of a term proportional to the gradient of the energy density, and a term specific to Yang-Mills fields that involves field commutators. This initial flow should certainly be part of the hydrodynamical initial conditions, if one wants the final results to be as insensitive as possible to the time at which it is initialized.

Somewhat more on the exploratory side, there have been some attempts to mimic a nucleus-nucleus collision by colliding two shockwaves in the strong coupling regime of the SUSY $N = 4$ Yang-Mills theory, using the AdS/CFT correspondence. The energy of the collision can be controlled by tuning the thickness of the shockwaves. By varying this parameter, a transition from a situation close to Landau hydrodynamics (at low energy) to Bjorken-like hydrodynamics (at high energy) is observed (J. Casalderrey-Solana, [17]).

3. Electromagnetic probes

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\[ ^1 \text{We were reminded by A. Kovner [13] of some mechanisms that can produce the azimuthal collimation of the ridge without invoking any flow.} \]
prompt photons produced at the impact of the two nuclei, before the formation of the QGP, photons produced by the hot hadronic gas that forms after the confinement transition, and a copious background of photons from the decay of unstable hadrons in the final state, e.g. π0's.

The calculation of the photon and dilepton production rates in a thermal medium can be done by calculating the relevant squared matrix elements, integrated over the phase-space of the unobserved particles with the appropriate statistical factors. Alternatively, the bookkeeping can be made more systematic by relating the photon production rate to the imaginary part of the photon polarization tensor, that can be calculated in QCD at finite temperature by using cutting rules.

The lowest order processes ($q\bar{q} \to \gamma^*$ for dileptons, $q\bar{q} \to \gamma g$ and $gg \to \gamma q$ for real photons) have first been calculated long ago. In the last two processes, the quark exchanged in the t-channel can become soft, which makes the rate formally infinite. This is cured by thermally generated quark masses (resummed via the Hard Thermal Loops effective theory), which turns the infinity into a logarithm.

However, it was realized shortly afterwards that the higher order bremsstrahlung graphs also have the same order of magnitude, due to a strong collinear singularity regularized by a quark thermal mass. Physically, this divergence arises when the photon formation time becomes larger than the quark mean free path, and the cure involves resumming multiple scattering corrections, in order to capture the Landau-Pomeranchuk-Migdal interference effect. The full leading order photon and dilepton rates have been obtained in [18] and [19].

Major improvements over these old results have been performed recently and were presented at HP2013. Firstly, the production rate of real photons has now been extended to NLO in the strong coupling constant (J. G. Multiple scattering corrections, in order to capture the Landau-Pomeranchuk-Migdal interference effect. The full leading order photon and dilepton rates have been obtained in [18] and [19].

The thermal dilepton rate has also been calculated at NLO (M. Laine, [21]) for pair invariant masses that are large enough so that the resummation of the multiple scattering corrections is not necessary. Here it was observed that the NLO corrections become very large compared to the LO contribution at small invariant masses, possibly due to new production channels opening up at NLO. Interestingly, this calculation shows a fairly good agreement with correlators evaluated on the lattice at zero spatial momentum.

Let us now digress on the collision kernel $C(q_{\perp})$ that enters in the calculation of photon rates and low mass dilepton rates. In short, $C(q_{\perp})$ is the squared scattering amplitude of the probe off a particle of the medium, integrated over the phase-space of the scattering center with the appropriate statistical distribution. The t-channel propagator in the amplitude must also be dressed for medium effects such as Debye screening. This object also appears in the discussion of the transverse momentum broadening undergone by a fast parton moving through a thermal medium, since $\Delta k_T^2$ per collision $= \int \frac{dq}{2\pi} q^2_{\perp} C(q_{\perp})$. Perturbatively, this collision kernel can be expressed in terms of in-medium gluon spectral functions:

$$C(q_{\perp}) = \int \frac{dq_0 dq_{\perp}}{(2\pi)^2} \frac{2\pi \delta(q_0 - q_{\perp})}{v_0 v_\gamma} \left( \rho_{\perp}(Q) + \rho_{\perp}(Q) \right).$$

From this expression and the Hard Thermal Loop results for the gluon spectral functions, one obtains the following LO expression [22]. $C(q_{\perp}) = m_D^2/(q_{\perp}^2(q_{\perp}^2 + m_D^2))$ (where $m_D$ is the Debye mass). In order to go beyond this result, one may use the following non-perturbative definition

$$e^{-t \tilde{C}(r_{\perp})} = \frac{1}{N_c} \text{Tr}(U(0_{\perp})U^+(r_{\perp}))$$

that relates the Fourier transform $\tilde{C}(r_{\perp})$ of the collision kernel to a correlator of Wilson lines along the light-cone. The NLO collision kernel has been calculated in [23]. Recently, this formula has been used in conjunction with the EQCD effective theory, obtained by dimensional reduction after integrating out all the modes with a non-zero Matsubara frequency, (M. Panero, [24]) in order to calculate $C(r_{\perp})$ on the lattice. This computation leads to an estimated value $\hat{q} \approx 6 \text{GeV}^2/\text{fm}$ for the jet quenching parameter at typical RHIC temperatures. Interestingly, if scaled appropriately by the Debye mass, this lattice result agrees quite well with the NLO perturbative result at small distance $r_{\perp}$.

In the context of heavy quarks energy loss, it was proposed by M. Djordjevic [25] to modify the LO collision...
in order to include phenomenologically a magnetic screening mass $m_\mu$ (lacking at LO). Although this modification seems to be helpful with the phenomenology of energy loss, one should keep in mind the fact that magnetic screening is intrinsically non-perturbative in QCD, and cannot be encompassed completely by this mere change in otherwise leading order formulas.

The local production photon and dilepton production rates are then folded into a hydrodynamical evolution in order to perform the space and time integrations that lead to the produced spectra. Basically, one just needs to boost the rates in the rest frame of each fluid cell by the flow velocity of that cell. When one fits the measured spectra by an exponential of the form \(\exp(-p_\perp/T)\), it is important to keep in mind that the slope parameter is not the initial temperature of the quark-gluon plasma (U. Heinz, [26]), due to the integration over the whole space-time history of the system.

When using viscous hydrodynamics in this procedure, one is really considering an out-of-equilibrium system: a non-zero viscous stress tensor \(\tau^{\mu\nu}\) implies that the distributions \(f(p)\) depart from the equilibrium ones. Consequently, the photon and dilepton rates themselves should be corrected to account for these modified quark and gluon distributions. This has been done for the \(2 \to 2\) production processes (C. Shen, [26]), and should be done also for the multiple scattering corrections to bremsstrahlung. It also seems that the photon $v_2$ exhibits some sensitivity to the parameters of the early flow (G. Vujanovic, [27]).

When one compares the photon spectra generated by these models to the data from nucleus-nucleus collisions, it appears that the predicted yield is a bit too small. But the largest discrepancy is in the $v_2$ of the photons, whose predicted value is too small by a factor \(~4\). The conjunction of these two facts suggests that the present models may underpredict photon production in the hadronic phase, where the flow is already developed (for a step in this direction, see W. Cassing, [28]).

### 4. Energy loss and jet quenching

The modifications of the propagation and fragmentation of a parton when it propagates through a dense medium is a natural way to probe the properties of this medium. These medium modifications depend on a wide range of scales, going from the initial energy of the probe to the Debye screening scale in the medium, which makes it rather complicated to describe in detail. This complexity has lead to several popular approaches (BDMPS, GLV, AMY,...), where the emphasis is put on different aspects of the problem. A recent trend has been to go beyond these original models, especially the BDMPS one, in order to include physics that was not in there originally (C. Salgado, [29], E. Iancu, [30]).

The common ground of all these models is soft and collinear gluon radiation in QCD, \(dP \sim \alpha_s(dx/x)d^2k_\perp/k_\perp^2\), which is enhanced by the momentum kicks that the probe receives while propagating through a medium. For independent scatterings, the transverse momentum of the probe increases as \(k_\perp^2 \sim \hat{q} t_f\) (in this discussion, \(\hat{q}\) is the only parameter characterizing the medium), while the gluon formation time is \(t_f \sim \omega/k_\perp^2\). Therefore, \(t_f \sim \sqrt{\omega/\hat{q}}\). For induced radiation to happen, \(t_f\) should be less than the medium size, which puts an upper limit \(\omega_c = \hat{q} L^2\) on the energy of the radiated gluons. The angular aperture of these emissions is given by \(\theta \sim k_\perp/\omega \sim (\hat{q}/\omega_3)^{1/4}\), which means that the hardest gluons are emitted at small angles. The mean energy lost in a single emission is \(\langle \omega \rangle \sim \alpha_s \omega_c \sim \alpha_s \hat{q} L^2\), i.e. it is dominated by \(\omega\) close to the upper limit \(\omega_c\) and it stays in a small cone around the trajectory of the probe. Therefore, this explains the energy loss of single hadrons, but it does not explain why jets are quenched.

The suppression of jets requires that energy be radiated at large angles (outside of the jet cone). This can happen for soft emissions, but many of them will be necessary to achieve a significant suppression. The color fields of the medium tend to randomize the colors of the emitters, which destroys the vacuum coherent antenna effects for out-of-cone emissions. Because of this, large angle emissions are not suppressed, and the energy radiated out-of-cone is rapidly degraded all the way down to the medium temperature scale. In-cone, the medium does not resolve the medium tends to randomize the colors of the emitters, which destroys the vacuum coherent antenna effects for out-of-cone emissions. Because of this, large angle emissions are not suppressed, and the energy radiated out-of-cone is rapidly degraded all the way down to the medium temperature scale. In-cone, the medium does not resolve the individual jet constituents, and the jet fragmentation proceeds as in vacuum.

Another effect is the correction to the parameter \(\hat{q}\) due to the fact that the gluon emissions contribute to the accumulated transverse momentum (Y. Mehtar-Tani, [31]). This effect is suppressed by \(\alpha_s\) but enhanced by logarithms,
and if resummed it introduces an anomalous $L$ dependence in $\hat{q} : \langle \omega \rangle \sim \alpha_s q_0 L^2 (q_0 L/m_t^2)^{2 - 2\alpha_s/\pi}$. Interestingly, this result is between the BDMPS result $L^2$ and the strong coupling result $L^3$. Several technical improvements were also presented, such as the complete single gluon radiation off a quark, including the in- and out-of-medium emission and their interference (L. Apolinario, [32]) and a study of the effect of the mass of heavy quarks on the destruction of the antenna coherence (M. Calvo, [33]).

Phenomenological models of energy loss that accommodate a range of behavior from elastic collisions only to the strong coupling result, embedded in viscous hydrodynamics, tend to prefer an energy loss that varies with a power law close to $L^2$, while stronger $L$ dependences are disfavored (B. Betz, [34]). This may not exclude totally strong coupling scenarios, since it appears that shooting string holographic models of jet quenching can also produce a range of powers of $L$ (A. Ficnar, [35]). Another way to circumvent the shortcomings of plain AdS/CFT is a hybrid strong/weak coupled approach to jet quenching (D. Fabbri).

The last step in more realistic studies is to embed a given microscopic model of energy loss in a model for the expansion of the fireball. For instance, the $v_2$ is quite sensitive to the details of the latter (D. Molnar, [36]). How one chooses the triggers is also important, as this can bias the event sample in several ways (chemical composition, kinematics, location of the initial hard scattering, etc.). This may be used to select specially interesting situations, but it also makes the comparison to data more complicated (T. Renk, [37]). A number of real world implementations were also presented at HP2013, that can be roughly categorized into hydro-based models (J. Xu, [38], Y. Tachibana, [39], R. Andrade, [40]) or cascade-like models (J. Tan, [41], Y. Zhu, [42], K. Tywoniuk, [43], K. Zapp, [44]).

5. Heavy flavors

An important subtopic in the field of heavy flavors is that of energy loss. Until recently, there was an apparent contradiction between the observation that heavy quarks seem to loose almost as much energy as lighter quarks, and the theoretical expectation that the dead-cone effect forbids radiation at angles $\theta < M/E$. This observation triggered a lot of interest for strong coupling models based on AdS/CFT, and a revived interest in collisional energy loss.

It now appears that, by carefully taking the partonic energy loss (both radiative and collisional, and by including running coupling as well as a phenomenological form of magnetic screening), the fragmentation functions and the subsequent decay of the D mesons, one can now obtain a reasonably good agreement with the suppression data for inclusive hadrons and for D mesons (M. Djordjevic, [45]). This improvement seems due to several smaller changes in various places, rather than having a unique cause.

The energy loss of heavy quarks can also be studied in cascade models (C. Greiner, J. Uphoff, [46]). In order to account for the radiative processes, it is necessary to include $2 \rightarrow 3$ and $3 \rightarrow 2$ processes as well as $2 \rightarrow 2$ processes, and the LPM interference is mimicked by vetoing rescatterings during the formation time of the emitted gluons. For inclusive hadrons, this leads to a reasonable $R_{AA}$, while for heavy quarks some tension persists between $R_{AA}$ and $v_2$.

Hybrid models mixing transport ideas and an underlying hydrodynamical background were also presented at HP2013 (M. Nahrgang, P.B. Gossiaux, [47]). In this particular model, EPOS was used for the description of the bulk, and the initial heavy quark distribution was taken from FONLL. A similar idea was presented by M. Nardini, [48]. Another, somewhat related, approach is to modify the Langevin equation that describes the momentum diffusion of heavy quarks by adding a force term that mimics the radiative energy loss (S. Cao, [49]).

In the strong coupling limit, the drag force that a medium exerts on a heavy quark has been known for a long time, $d\vec{p}/dt = (\sqrt{\alpha_s/2\pi})(\pi T)^2 B/\sqrt{1 - B^2}$, in the case of a static and isotropic quark-gluon plasma. A more sophisticated study in the case of a medium that has an anisotropic energy-momentum tensor, as would be the case for a longitudinally expanding quark-gluon plasma, indicates that the magnitude of the drag force is not dramatically modified, with only relatively small changes due to the anisotropy (M. Lekaveckas, [50]).

The second subtopic related to heavy quarks is that of quarkonia dissociation in a hot and dense medium. There is now a clear signal at the LHC of the suppression of the excited states of the $\Upsilon$, which is very suggestive of a sequential dissociation pattern, where the ground state melts at a higher temperature than the excited ones.

A lot of activity in lattice QCD is devoted to the study of quantities relevant to the physics of heavy quarks in the quark-gluon plasma. A notoriously difficult problem is the calculation of a spectral function $\sigma$. On the lattice, the
closest object one can compute is an Euclidean correlator \( G_z \), related to \( \sigma \) by

\[
G_z(\tau, p) = \int_0^{\infty} d\omega \sigma(\omega, p) \frac{\cosh(\omega(\tau - \beta/2))}{\sinh(\omega \beta/2)}.
\]

The inversion of this integral relationship is an ill-posed problem, unless one introduces some extra informations or assumptions about the spectral function, which can be done with the maximum entropy method. An additional difficulty is that the Euclidean correlator \( G_z \) is not very sensitive to changes of the spectral function (H.-T. Ding, [51]). This method can be used to extract the spectral function in various channels with heavy quarks, in order to directly visualize the disappearance of the bound state peaks as the temperature increases. In the case of the \( J/\Psi \), it appears that the peak is almost gone at \( T = 1.45T_c \). This method can also be used to compute transport coefficients such as the diffusion constant of heavy quarks. The same approach with light quarks can give access to the electrical conductivity of the quark–gluon plasma, and to the dilepton production rate.

The previous method is more and more difficult to apply as the temperature increases, because the extent of the imaginary time axis shrinks as \( 1/T \). This is especially problematic for extracting transport coefficients, that are obtained from the zero momentum limit of the spectral function. An alternative approach is to consider an Euclidean correlator in the spatial \( z \) direction, whose extent is not limited by the temperature,

\[
G(z) \equiv \int_0^\beta d\tau \int dx dy \langle J(\tau, x) J(0, 0) \rangle = \int_0^\infty dz e^{ipz} \int_0^{\infty} d\omega \sigma(\omega_s(0, 0, p_z)).
\]

This alternate method is more sensitive to changes of the spectral function (A. Mocsy, [52]), and high-temperature modifications become much more obvious.

One can also address the fate of quarkonia made of heavy quarks in effective theories, such as pNRQCD,

\[
\mathcal{L} = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} + \mathcal{Q}^a q \bar{q} + \frac{1}{2} q^\dagger S \left[ i\gamma_5 \partial_\mu - \frac{(\nabla_\mu)^2}{M} - V_s(r) \right] S + \cdots
\]

where \( S \) is an effective field for singlet \( Q \bar{Q} \) states. The potential \( V_s \) is generally complex, with real and imaginary parts respectively related to the shift of the peak and the broadening of the spectral function. In this effective description at lowest order, the field \( S \) obeys a non-relativistic Schrödinger equation with the potential \( V_s \). In the past, the singlet free energy was used in this equation, but it has no imaginary part. A more rigorous derivation is from the spectral function (P. Petreczky, [53]). First results indicate that the real part of the potential is between the zero temperature potential and the singlet free energy, with a non negligible imaginary part. When this potential is used in the Schrödinger equation, one finds that all the charmonium states disappear for \( T > 240 \text{ MeV} \). In the case of the Upsilon states, the \( \Upsilon(2s) \) and \( \Upsilon(3s) \) melt around \( T \sim 250 \text{ MeV} \), while the ground state \( \Upsilon(1s) \) survives until \( T \sim 450 \text{ MeV} \).

Summary This 6th edition of the Hard Probes conference has seen many new results, both in hardcore theory and in the modeling of nucleus-nucleus collisions. Some areas of the theory of heavy ion collisions are close to reaching a state where the corresponding data becomes understandable from a QCD perspective. This conference also shed some light on the fact that proton-nucleus collisions are not as simple as one may have thought originally, with many issues related to the centrality determination, making them more difficult to use reliably as a control experiment.

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References
[1] I. Helenius, K. J. Eskola, H. Paukkunen, Constraining nPDFs with inclusive pions and direct photons at forward rapidities in p+Pb collisions at the LHC. arXiv:1404.0470
[2] R. Paatelainen, K. Eskola, H. Holopainen, K. Tuominen, Multiplicities and \( p_T \) spectra in ultrarelativistic heavy ion collisions from a next-to-leading order improved perturbative QCD + saturation + hydrodynamics model, Phys.Rev. C87 (4) (2013) 044904. arXiv:1211.0461
[3] D. d’Enterria, A. Singirev, Double-parton scattering cross sections in proton-nucleus and nucleus-nucleus collisions at the LHC. arXiv:1403.2335

[4] A. Ramnath, These proceedings.

[5] G. Jackson, W. Horowitz, Predictions for the Spatial Distribution of Gluons in the Initial Nuclear State. arXiv:1404.4743

[6] T. Lappi, H. Mantysaari, Particle production from the Color Glass Condensate: proton-nucleus collisions in light of the HERA data. arXiv:1403.6994

[7] B.-W. Xiao, The Search for Gluon Saturation in pA Collisions. arXiv:1404.0179

[8] T. Lappi, H. Mantysaari, Proposal for a running coupling JIMWLK equation. arXiv:1407.7289

[9] S. Caron-Huot, When does the gluon reggeize? arXiv:1309.6521

[10] A. Kovner, M. Lublinsky, Y. Mulian, Jalilian-Marian, Iancu, McLerran, Weigert, Leonidov, Kovner evolution at next to leading order, Phys.Rev. D89 (2014) 061704. arXiv:1310.0378

[11] G. Beuf, Improving the kinematics for low-x QCD evolution equations in coordinate space. arXiv:1401.0313

[12] E. Iancu, D. Triantafyllopoulos, JIMWLK evolution for multi-particle production in Langevin form, JHEP 1311 (2013) 067. arXiv:1308.1959

[13] A. Kovner, M. Lublinsky, Angular and long range rapidity correlations in particle production at high energy, Int.J.Mod.Phys. E22 (2013) 1330001. arXiv:1211.1928

[14] S. Jeon, These proceedings.

[15] T. Epelbaum, F. Gelis, Pressure isotropization in high energy heavy ion collisions, Phys.Rev.Lett. 111 (2013) 232301. arXiv:1307.2214

[16] G. Chen, R. J. Fries, Global Flow of Glasma in High Energy Nuclear Collisions, Phys.Lett. B723 (2013) 417–421. arXiv:1303.2360

[17] J. Casalderrey-Solana, M. P. Heller, D. Mateos, W. van der Schee, From full stopping to transparency in a holographic model of heavy ion collisions, Phys. Rev. Lett. 111, 181601 (2013) 181601. arXiv:1306.4919

[18] P. B. Arnold, G. D. Moore, L. G. Vaffa, Photon emission from quark gluon plasma: Complete leading order results, JHEP 0112 (2001) 009. arXiv:hep-ph/0111107

[19] P. Aurenche, F. Gelis, G. Moore, H. Zaratek, Landau-Pomeranchuk-Migdal resummation for dilepton production, JHEP 0212 (2002) 006. arXiv:0204011

[20] J. Ghiglieri, J. Hong, A. Kukelka, E. Lu, G. D. Moore, et al., Next-to-leading order thermal photon production in a weakly coupled quark-gluon plasma, JHEP 1305 (2013) 010. arXiv:1302.5870

[21] M. Laine, NLO thermal dilepton rate at non-zero momentum, JHEP 1311 (2013) 120. arXiv:1310.0164

[22] P. Aurenche, F. Gelis, H. Zaratek, A Simple sum rule for the thermal gluon spectral function and applications, JHEP 0205 (2002) 043. arXiv:hep-ph/0204146

[23] S. Caron-Huot, O(g) plasma effects in jet quenching, Phys.Rev. D79 (2009) 065039. arXiv:0811.1650

[24] M. Panero, K. Rummukainen, A. Schaefel, A lattice study of the jet quenching parameter. arXiv:1307.5850

[25] M. Djordjevic, Heavy flavor suppression in a dynamical QCD medium with finite magnetic mass. arXiv:1211.0198

[26] U. W. Heinz, J. Liu, C. Shen, Electromagnetic fingerprints of the Little Bang. arXiv:1403.8101

[27] G. Vujanovic, J.-F. Paquet, G. S. Denicol, M. Luzum, B. Schenke, et al., Probing the early-time dynamics of relativistic heavy-ion collisions with electromagnetic radiation. arXiv:1404.3714

[28] O. Linnyk, W. Cassing, E. Bratkovskaya, Centrality dependence of the direct photon yield and elliptic flow in heavy-ion collisions at sqrt(s)=200 GeV, Phys.Rev. C89 (2014) 034908. arXiv:1311.0279

[29] C. Salgado, These proceedings.

[30] E. Iancu, Di-jet asymmetry and wave turbulence. arXiv:1404.4566

[31] J.-P. Blaizot, Y. Mehtar-Tani, Renormalization of the jet-quenching parameter. arXiv:1403.2323

[32] L. Apolinario, N. Armesto, G. Milhano, C. A. Salgado, Medium-induced gluon radiation beyond the eikonal approximation. arXiv:1404.7079

[33] M. R. Calvo, M. R. Moldes, C. A. Salgado, Color coherence in a heavy quark antenna radiating gluons inside a QCD medium. arXiv:1403.4892

[34] B. Betz, M. Gyulassy, Azimuthal Jet Tomography at RHIC and LHC arXiv:1402.3419

[35] A. Ficnar, S. S. Gubser, M. Gyulassy, Holographic light quark jet quenching at RHIC and LHC via the shooting strings. arXiv:1404.0935

[36] D. Molnar, D. Sun, Realistic medium-averaging in radiative energy loss. arXiv:1209.2430

[37] T. Renk, Jet correlations - opportunities and pitfalls. arXiv:1404.0793

[38] J. Xu, A. Buzzatti, M. Gyulassy, The Tricky Azimuthal Dependence of Jet Quenching at RHIC and LHC via CUJET2.0. arXiv:1404.0384

[39] Y. Tachibana, T. Hirano, Di-jet asymmetric momentum transported by QGP fluid. arXiv:1404.1706

[40] R. P. G. Andrade, J. Noronha, G. S. Denicol, Jet quenching effects on the direct, elliptic, and triangular flow at RHIC. arXiv:1403.1789

[41] L. Tan, These proceedings.

[42] A.-N. Wang, Y. Zhu, Jet quenching and \(\gamma\)-jet correlation in high-energy heavy-ion collisions. arXiv:1404.0031

[43] Y. Mehtar-Tani, K. Tywoniuk, Jet (de)coherence in Pb-Pb collisions at the LHC. arXiv:1401.2893

[44] K. C. Zapp, Geometrical aspects of jet quenching in JEWEL. arXiv:1312.5536

[45] M. Djordjevic, Heavy flavor puzzle at LHC: a serendipitous interplay of jet suppression and fragmentation, Phys.Rev.Lett. 112 (2014) 042302. arXiv:1307.4702

[46] J. Uphoff, O. Fochler, Z. Xu, C. Greiner, Heavy vs. light flavor energy loss within a partonic transport model. arXiv:1310.3597

[47] M. Nahrgang, J. Aichelin, P. B. Gossiaux, K. Werner, Heavy-flavor azimuthal correlations of D mesons. arXiv:1310.2218

[48] W. Alberico, A. Beraudo, A. De Pace, A. Molinari, M. Monteno, et al., Heavy flavors in AA collisions: production, transport and final spectra, Eur.Phys.J. C73 (2013) 2481. arXiv:1305.7221

[49] S. Cao, G.-Y. Qin, S. A. Bass, Dynamical Evolution, Hadronization and Angular De-correlation of Heavy Flavor in a Hot and Dense QCD Medium. arXiv:1404.1081

[50] M. Lekaveckas, K. Rajagopal, Effects of Fluid Velocity Gradients on Heavy Quark Energy Loss, JHEP 1402 (2014) 068. arXiv:1311.5577
[51] H.-T. Ding, Hard and thermal probes of QGP from the perspective of Lattice QCD. arXiv:1404.5134
[52] A. Mocsy, P. Petreczky, M. Strickland, Quarkonia in the Quark Gluon Plasma, Int.J.Mod.Phys. A28 (2013) 1340012. arXiv:1302.2180
[53] A. Bazavov, Y. Burnier, P. Petreczky, Lattice calculation of the heavy quark potential at non-zero temperature. arXiv:1404.4267