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# ADS/CFT AND HEAVY ION COLLISIONS*

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I describe the parton picture at strong coupling as emerging from the gauge/gravity duality and its consequences for high-energy scattering and for the hard probes of a strongly coupled plasma, as potentially relevant for heavy collisions at RHIC and LHC. I emphasize the differences with respect to the corresponding picture in perturbative QCD.

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## 1. Introduction: from RHIC to lattice QCD

Some of the experimental discoveries at RHIC, notably the unexpectedly large medium effects known as elliptic flow and jet quenching, led to the suggestion that the deconfined hadronic matter produced in the intermediate stages of a heavy ion collision is a nearly perfect fluid, so like a strongly coupled plasma. The coupling constant $\alpha_{\rm s} = g^2/4\pi$ in QCD can never become large, because of asymptotic freedom, but it can be of order one at scales of order $\Lambda_{\rm QCD}$, and this might lead to an effectively strong-coupling behavior. It is notoriously difficult to do reliable estimates in QCD when $\alpha_{\rm s} \simeq 1$, so it has become common practice to look to the $\mathcal{N} = 4$ supersymmetric Yang–Mills (SYM) theory, whose strong coupling regime can be addressed within the AdS/CFT correspondence [1], for guidance as to general properties of strongly coupled plasmas (see the review papers [2–4]). Since conformal symmetry is an essential property of $\mathcal{N} = 4$ SYM, this theory cannot be used as a model for the dynamics in QCD in the vicinity of the deconfinement phase transition ($T \simeq T_{\rm c} \simeq 170$ MeV), but only for larger temperatures, within the range $2T_{\rm c} \lesssim T \lesssim 5T_{\rm c}$, where QCD itself is known (from lattice studies [5]) to be nearly conformal. This is precisely the temperature range which is relevant for the phenomenology of heavy ion collisions at RHIC and LHC.

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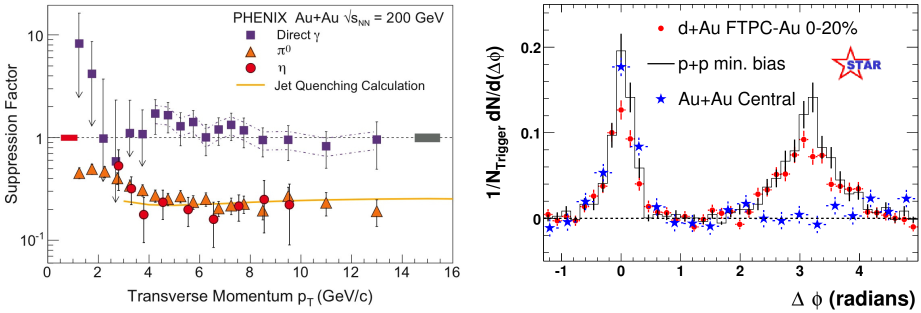


Fig. 1. Left: The ratio $R_{AA}$ of measured *versus* expected yield of various particles $(\pi^0, \eta, \gamma)$ in Au+Au collisions at $\sqrt{s_{NN}} = 200\,\mathrm{GeV}$ as function of the transverse momentum $p_{\mathrm{T}}$ (RHIC, PHENIX). Right: Azimuthal correlations for jet measurements at RHIC (STAR) in $p+p$, $d+$ Au, and Au+Au collisions.

Among the most intriguing RHIC data are those referring to *jet quenching*, *i.e.*, the energy loss and the transverse momentum broadening of a relatively "hard" probe (a heavy and/or energetic quark or lepton, with transverse momentum of a few GeV), for which one would naively expect a weak-coupling behavior, by asymptotic freedom. Yet, perturbative QCD seems to be unable to explain the strong suppression of particle production in Au+Au collisions, as shown in Fig. 1 (left). Namely, the ratio $R_{AA}$ between the particle yield in Au+Au collisions and that in $p+p$ collisions rescaled by the number of participants would be one in the absence of medium effects. But the RHIC measurements yield a very small value $R_{AA} \approx 0.2$ (for hadron production) which suggests that, after being produced through a hard scattering, the partonic jets are absorbed by the surrounding medium.

Additional evidence comes from direct studies of jets, *cf.* Fig. 1 (right). A high-energy proton–proton collision generally produces a pair of partons whose subsequent evolution via fragmentation and hadronisation leaves two jets of hadrons which propagate back-to-back in the center of mass frame (see Fig. 2 (left)). Hence, the distribution of the final state radiation w.r.t. the azimuthal angle $\Delta\Phi$ shows two pronounced peaks, at $\Delta\Phi = 0$ and $\pi$ — the curve denoted as "$p+p$ min. bias" in Fig. 1 (right). A similar distribution is seen in deuteron–gold collisions at RHIC, but not in central Au+Au collisions, where the peak at $\Delta\Phi = \pi$ has disappeared, as visible in Fig. 1 (right). It is then natural to imagine that the hard scattering producing the jets has occurred near the edge of the interaction region, so that the near side jet has escaped to the detector, while the away side jet has been absorbed within the medium (see Fig. 2 (right)).

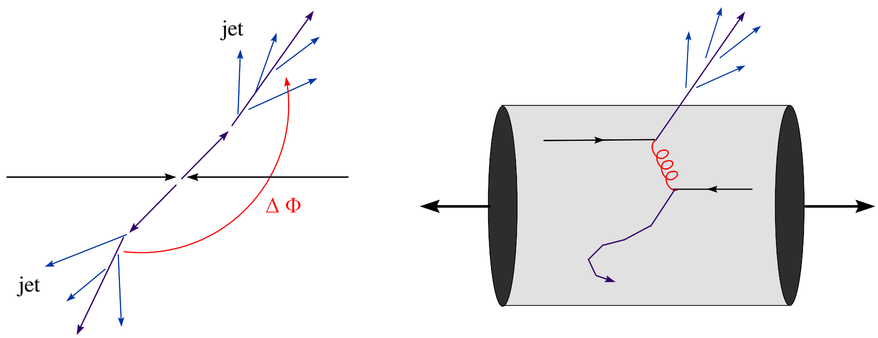


Fig. 2. Jet production in high-energy scattering. Left: a proton–proton collision: the leading partons fragment into two back-to-back hadronic jets. Right: a nucleus–nucleus collision: one of the leading partons escapes the interaction region and enters the detector, but the other one is absorbed in the surrounding matter.

These results at RHIC show that the matter produced in the intermediate stages of a heavy ion collision is *opaque*, which may well mean that this matter is dense, or strongly-coupled, or both. It is then natural to ask whether one can discriminate between weak coupling and strong coupling behavior on the basis of the lattice results for the QCD thermodynamics. Unfortunately however, the theoretical interpretation of these results is not completely free of ambiguities, as we now explain.

Figure 3 exhibits the lattice results for the conformal anomaly $(\varepsilon - 3p)/T^4$ (left) and for the pressure $p$ (right) in units of $T^4$. The left figure confirms that the QCD plasma is nearly conformal for temperatures $T \gtrsim 2T_{\rm c}$: after showing a peak at the deconfinement phase transition, the relative conformal anomaly $(\varepsilon - 3p)/\varepsilon$ decreases very fast with increasing $T$ and becomes smaller than 10% when $T \gtrsim 2T_{\rm c}$. The right figure shows that, after a rapid jump at the phase transition, the pressure slowly approaches the respective ideal gas limit $p_0$, in such a way $p/p_0 \simeq 0.85$ at $T = 3T_{\rm c}$. This deviation from the ideal gas is rather small, of the order of the first perturbative correction $\mathcal{O}(\alpha_{\rm s})$, and indeed perturbation theory in QCD at finite temperature does a good job in reproducing the lattice results for temperatures $T \gtrsim 2.5T_{\rm c}$; this is shown by the "HTL" band in that figure [7]. On the other hand, this value $p/p_0 \simeq 0.85$ is not too far from the respective result in the strong-coupling limit of $\mathcal{N} = 4$ SYM (as computed within AdS/CFT), which is $p/p_0 = 0.75$ [1]. Thus, although consistent with weak coupling expectations, the lattice results for thermodynamics do not totally exclude a strong-coupling behavior in the QCD plasma in the considered range of temperatures.

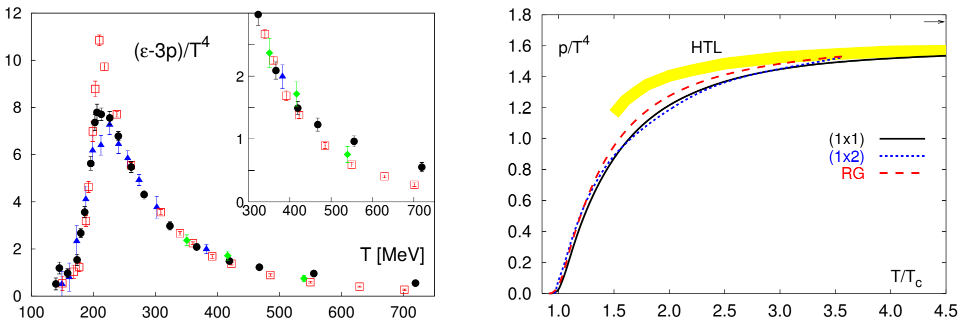


Fig. 3. Lattice results for the QCD trace anomaly, $T^\mu_\mu = \varepsilon - 3p$ (left) [5] and for the pressure of the SU(3) gauge theory (right) [6]. In the right figure, different lines correspond to different gauge actions, whereas the upper band denoted as "HTL" represents the results of a parameter-free resummation of perturbation theory [7]. The small arrow in the upper right corner indicates the pressure of an ideal gas.

## 2. A lattice test of the coupling strength in QCD

If the lattice results for the QCD thermodynamics cannot convincingly discriminate between weak and strong coupling behavior, is there any other test of the strength of the coupling in lattice QCD at finite temperature? In what follows, we shall describe a recent proposal in that sense [8], which involves the lattice measurement of leading-twist operators. These are the operators with spin $n$, classical dimension $d = n+2$, and twist $t = d-n = 2$, which in the weak coupling regime control the operator product expansion (OPE) of deep inelastic scattering (DIS) at large photon virtuality $Q^2$ [9]. There are two infinite sequences of leading-twist operators — fermionic and gluonic — among which we show here only those with $n = 2$:

$$\mathcal{O}_f^{\mu\nu} \equiv \tfrac{1}{2}\bar{q}\big(\gamma^\mu iD^\nu + \gamma^\nu iD^\mu\big)q\,, \tag{1}$$

(with an implicit sum over quark flavors and zero quark masses) and

$$\mathcal{O}_g^{\mu\nu} \equiv -F_a^{\mu\alpha}F_\alpha^{\nu,a} + \tfrac{1}{4}\,g^{\mu\nu}F_a^{\alpha\beta}F^a_{\alpha\beta}\,. \tag{2}$$

These are recognized as the energy-momentum tensors for quarks and gluons, respectively. The hadron expectation values of the spin-$n$ leading-twist operators measure the $(n-1)$-th moment of the longitudinal momentum fraction carried by the quark and gluon constituents of that hadron.

Being composite, these operators are well defined only with a renormalization prescription, and thus implicitly depend upon the renormalization scale $Q^2$. Physically, this dependence expresses the fact that quantum fields can radiate and their internal structure in terms of "bare" quanta depends upon the resolution scale $Q^2$ at which one probes this structure.

For instance, the success of the valence parton picture for high-energy scattering in QCD is deeply related to asymptotic freedom which guarantees that parton branching at $Q^2 \gg \Lambda^2_{\rm QCD}$ is controlled by weak coupling. This proceeds via bremsstrahlung, which favors the emission of "soft" and "collinear" quanta, *i.e.* quanta which carries only a small fraction $x$ of the longitudinal momenta of their parent partons and relatively small transverse momenta. Hence, although there are many small-$x$ gluons in the proton wavefunction at high energy, most of the proton longitudinal momentum is still carried by the point-like valence quarks.

By contrast, at strong coupling one expects parton branching to be fast and "quasi-democratic": the energy of the parent parton is quasi equally shared by the daughter partons. Through successive branchings, all partons should cascade down to small-$x$ constituents [10–12]. Hence, a strongly-coupled hadron or plasma cannot involve point-like constituents which carry a significant fraction of the total energy.

These considerations are encoded in the renormalization group equations describing the evolution of the leading-twist operators with the resolution scale $\mu^2$. Up to operator mixing issues to which we shall return in a moment, these equations read (for a generic spin-$n$ operator $\mathcal{O}^{(n)}$)

$$\mu^2 \frac{d}{d\mu^2}\mathcal{O}^{(n)} = \gamma^{(n)}\mathcal{O}^{(n)} \implies \frac{\mathcal{O}^{(n)}(Q^2)}{\mathcal{O}^{(n)}(\mu_0^2)} = \exp\left\{\int_{\mu_0^2}^{Q^2} \frac{d\mu^2}{\mu^2}\gamma^{(n)}(\mu^2)\right\}, \qquad (3)$$

where $\gamma^{(n)}$ is the corresponding anomalous dimension and is strictly negative — meaning that the evolution increases the number of partons with small longitudinal momentum fraction $x$ while decreasing that of the partons with large $x$ — except for the total energy-momentum operator

$$T^{\mu\nu} = \mathcal{O}_f^{\mu\nu} + \mathcal{O}_g^{\mu\nu}\,, \qquad (4)$$

which has zero anomalous dimension since it is a conserved quantity (and hence it does not depend upon the resolution scale $Q^2$). Hence, in the continuum limit $Q^2 \to 0$, the expectation values of all the leading-twist operator except for $T$ must vanish. But the rate of this evolution is very different at weak and respectively strong coupling.

*(i) Weak coupling:* To lowest order in the running coupling, one has

$$\gamma^{(n)}(\mu^2) = -a^{(n)}\frac{\alpha_{\rm s}(\mu^2)}{4\pi} \implies \frac{\mathcal{O}^{(n)}(Q^2)}{\mathcal{O}^{(n)}(\mu_0^2)} = \left[\frac{\ln(\mu_0^2/\Lambda^2)}{\ln(Q^2/\Lambda^2)}\right]^{a^{(n)}/b_0}, \qquad (5)$$

with $\alpha_{\rm s}(\mu^2) = 4\pi/[b_0 \ln(\mu^2/\Lambda^2)]$, $\Lambda \approx 200\,$MeV, $b_0 = (11N_c - 2N_f)/3$, and $a^{(n)} > 0$. Thus, the perturbative evolution is merely logarithmic.

*(ii) Strong coupling and conformal field theory:* At strong coupling, direct calculations in QCD are not possible anymore, but we shall use the corresponding results for $\mathcal{N}=4$ SYM for some qualitative insight. In a conformal field theory, $\gamma^{(n)}$ is scale-independent and negative (with the exception of $T^{\mu\nu}$, of course), so the evolution is power-like:

$$\frac{\mathcal{O}^{(n)}(Q^2)}{\mathcal{O}^{(n)}(\mu_0^2)} = \left[\frac{\mu_0^2}{Q^2}\right]^{|\gamma^{(n)}|} . \tag{6}$$

Moreover, AdS/CFT predicts that, at strong coupling $\lambda \equiv g^2 N_c \gg 1$, all the non-zero anomalous dimensions are very large $|\gamma^{(n)}| \sim \mathcal{O}(\lambda^{1/4})$ [13], so the leading-twist operators rapidly die away with increasing $Q^2$ (meaning that all partons have fallen down to small values of $x$).

These results suggest that a natural way to measure the strength of the coupling in QCD at finite temperature is to compute thermal expectation values of leading-twist operators in lattice QCD [8]. These operators evolve from the natural physical scale $T$ up to the resolution scale $Q$ set by the lattice spacing: $Q = a^{-1}$. In practice, the ratio $Q/T = aT$ is not very large, $Q/T \lesssim 10$, so if the evolution is perturbative, *cf.* (5), the expectation value of an unprotected operator is only slightly reduced — by a few percent. On the other hand, if the plasma is effectively strongly coupled at the scale $T$, then at least the early stages of the evolution should involve a large negative anomalous dimension, leading to a strong suppression in the expectation value measured at the final scale $Q$.

The previous argument applies to the unprotected operators, which include all the higher-spin operators with $n \geq 4$. Unfortunately, however, it turns out that it is very difficult, if not impossible, to accurately measure on the lattice such high-spin operators. There is another possibility, though, which should be easier in practice: this is to measure the linear combination of the spin-2 operators in Eqs. (1), (2) which is *orthogonal* to $T^{\mu\nu}$ within the renormalization flow and therefore has a non-zero, and negative, anomalous dimension. In full QCD, one is unable to identify the relevant orthogonal operator except in perturbation theory. However, there is a simpler, but still non-trivial, version of the theory in which the answer to this question is known for *any* value of the coupling: this is *quenched* QCD, *i.e.* the theory obtained from QCD after removing all the quark loops. On the lattice, this is non-perturbatively defined by removing the fermionic determinant from the QCD action. As argued in Ref. [8], $\mathcal{O}_f^{\mu\nu}$ is the operator orthogonal to $T^{\mu\nu}$ in this framework: indeed, in quenched QCD, a quark can emit gluons, but the emitted gluons, as well as those from the thermal bath, are not allowed to emit quark–antiquark pairs. Hence, when the system is probed on a sufficiently hard scale, most of the total energy appears in the gluons.

To summarize, the proposal made in Ref. [8] is to measure the thermal expectation value $\langle \mathcal{O}_f^{00}(Q^2)\rangle_T$ of the quark energy density in lattice quenched QCD, for a temperature between $2T_{\rm c}$ and $5T_{\rm c}$. If the deviation from the corresponding result for the ideal Fermi–Dirac gas turns out to be relatively small, say $\lesssim 30\%$, then one can conclude that the QCD plasma is weakly coupled at the scale $T$. If on the other hand the lattice result turns out to be considerably smaller, then there must be a regime in $\mu$ around $T$ where QCD is effectively strongly coupled.

## 3. DIS and parton saturation at strong coupling

The previous discussion emphasized the importance of understanding parton evolution at strong coupling. At least for the $\mathcal{N}=4$ SYM theory, this can be done within the gauge/string duality. The simplest version of the formalism, known as the "supergravity approximation", is obtained by taking the large-$N_c$ limit, or, equivalently, the large 't Hooft coupling limit: $\lambda = g^2N_c \to \infty$ with $g$ fixed and small ($g \ll 1$). This is generally not a good limit to study a scattering process, since the corresponding amplitude is suppressed as $1/N_c^2$ [10, 11]. Yet, this is meaningful for processes taking place in a deconfined plasma, which involves $N_c^2$ degrees of freedom per unit volume, thus yielding finite amplitudes when $N_c \to \infty$. In this limit, the $\mathcal{N}=4$ SYM plasma at finite temperature is described as a "black-hole" (more properly, a black-brane) embedded in AdS$_5$ and the dynamics reduces to classical gravity in this curved space-time [1–4]. It should be emphasized that the black hole (BH) is homogeneous in the physical 4 dimensions[1] but it has an horizon in the radial, or "fifth", dimension of AdS$_5$, at a position which is determined by the temperature of the plasma.

The AdS$_5$ BH geometry is illustrated in Fig. 4, which also shows the supergravity process dual to DIS off the $\mathcal{N}=4$ SYM plasma: A space-like virtual photon, with 4-momentum $q^\mu = (\omega, 0, 0, q)$ and virtuality $Q^2 \equiv q^2 - \omega^2 \gg T^2$, acts as a perturbation on the Minkowski boundary of AdS$_5$ (at $\chi = 0$), thus inducing a massless, vector, supergravity field $A_m$ (with $m = \mu$ or $\chi$) which propagates within the bulk of AdS$_5$ ($\chi > 0$), according to the Maxwell equations in curved space-time:

$$\partial_m\big(\sqrt{-g}g^{mp}g^{nq}F_{pq}\big) = 0\,, \qquad \text{where} \qquad F_{mn} = \partial_m A_n - \partial_n A_m\,. \qquad (7)$$

Here, $g^{mn}$ is the metric tensor in 5-dimensions which in particular contains the information about the BH horizon at $\chi = 1/T$. Thus (7) describes the gravitational interaction between the Maxwell field $A_m$ and the BH. As

[1] More recently, a finite-length plasma "slice" has been considered too, as a model for a nucleus which admits a simple supergravity dual (a "shockwave") [15].

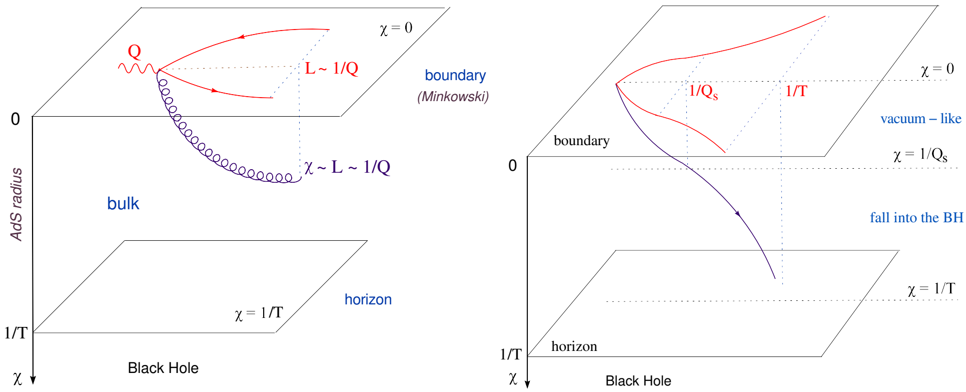


Fig. 4. Space-like current in the plasma: the trajectory of the wave packet in $AdS_5$ and its "shadow" on the boundary. Left: low energy — the Maxwell wave gets stuck near the boundary. Right: high energy — the wave falls into the BH.

usual, the strength of this interaction is proportional to the product $\omega^2 T^4$ between the energy densities in the two interacting systems. Interestingly, there is a threshold value for this quantity, of order $Q^6$, below which there is essentially no interaction [12]: so long as $\omega T^2 \ll Q^3$, the Maxwell wave is stuck within a distance $\chi \lesssim 1/Q \ll 1/T$ from the Minkowski boundary and does not "see" the BH (*cf.* Fig. 4 (left)). But for higher energies and/or temperatures, such that $\omega T^2 \gtrsim Q^3$, the wave gets attracted by the BH and eventually falls into the latter. Physically, this means that the energetic space-like photon is absorbed with probability one into the plasma — the "black disk limit" for DIS (*cf.* Fig. 4 (right)).

This critical value $Q_s \sim (\omega T^2)^{1/3}$, known as the *saturation momentum*, together with the physical picture of the scattering, can be understood with the help of the "UV/IR correspondance", which relates the 5th dimension of $AdS_5$ to the momenta (or typical sizes) of the quantum fluctuations which are implicitly integrated out in the boundary gauge theory. Namely, the radial penetration $\chi$ of the Maxwell wave packet in $AdS_5$ is proportional to the transverse size $L$ of the quantum fluctuation of the virtual photon in the dual gauge theory. By the uncertainty principle, we expect a highly energetic space-like photon with $\omega \simeq q \gg Q$ to fluctuate into a partonic system with transverse size $L \sim 1/Q$ — which indeed matches the radial penetration of the dual Maxwell field, as illustrated in Fig. 4 (left) — and a finite lifetime $\Delta t \sim \omega/Q^2$. In the vacuum, the space-like fluctuation cannot decay into on-shell partons, by energy-momentum conservation. But the situation can change in the presence of the plasma. Unlike the photon, which is color neutral, its partonic fluctuation has a dipolar color moment and thus can interact with the plasma. Via such interactions, the partons can acquire the energy and momentum necessary to get on-shell, and then the fluctuation decays: the space-like photon disappears (*cf.* Fig. 4 (right)).

Let us now return to the saturation momentum $Q_{\rm s} \sim (\omega T^2)^{1/3}$. The condition $Q \sim Q_{\rm s}$ can be rewritten as

$$Q \sim \frac{\omega}{Q^2}\, T^2\,, \tag{8}$$

which admits the following interpretation [12] : the scattering becomes strong when the lifetime $\Delta t \sim \omega/Q^2$ of the partonic fluctuation is large enough for the mechanical work $W = \Delta t \times F_T$ done by the *plasma force* $F_T \sim T^2$ acting on these partons to compensate for the energy deficit $Q$ of the space-like system. This plasma force $F_T \sim T^2$ represents the effect of the strongly-coupled plasma on color dipole fluctuations and can be viewed as a prediction of the AdS/CFT calculation. Introducing the Bjorken $x$ variable $x \equiv Q^2/(2\omega T)$ for DIS off the plasma (this represents the longitudinal momentum fraction of the plasma constituent which absorbs the virtual photon), one can rewrite the plasma *saturation line* as $Q_{\rm s}(x) = T/x$ or, alternatively, $x_{\rm s}(Q) = T/Q$.

The AdS/CFT results suggest the following partonic picture for the strongly-coupled plasma [12]. For $Q \gg Q_{\rm s}(x)$ (or, equivalently, $x \gg x_{\rm s}(Q)$), the scattering is negligible and the DIS structure function $F_2$ is exponentially small: $F_2(x,Q^2) \sim \exp\{-Q/Q_{\rm s}(x)\}$. This shows that there are no pointlike constituents in the strongly coupled plasma, in agreement with the OPE argument in Sec. 2. For $x \lesssim x_{\rm s}(Q)$, the scattering is strong and the structure function is found as $F_2(x,Q^2) \sim x N_c^2 Q^2$. This is in agreement with our physical expectation that partons must accumulate at small values of $x$, as a result of branching, and is moreover consistent with energy-momentum conservation, which requires the integral $\int_0^1 dx\, F_2(x,Q^2)$ to have a finite limit as $Q^2 \to \infty$ [9]. The previous results imply indeed

$$\int_0^1 dx\, F_2(x,Q^2) \simeq x_{\rm s} F_2(x_{\rm s},Q^2) \sim N_c^2 T^2\,, \tag{9}$$

where the integral is dominated by $x \simeq x_{\rm s}(Q)$: the energy and momentum of the plasma as probed on a "hard" resolution scale $Q^2 \gg T^2$ is fully carried by the partons "along the saturation line", *i.e.*, those having $x \simeq T/Q$. A similar picture holds for other hadronic targets so like a "glueball" [10,11] or a "nuclear" shockwave [15], but the respective saturation momentum rises slower with $1/x$ than for the infinite plasma: $Q_{\rm s}^2(x) \propto 1/x$ for a finite-size "hadron" as opposed to $Q_{\rm s}^2(x) \propto 1/x^2$ for the plasma. The additional factor $1/x$ in the case of the plasma comes from the lifetime $\Delta t \sim \omega/Q^2 \sim 1/xT$ of the partonic fluctuation: since the medium is infinite, the effects of the scattering accumulate all the way along the parton lifetime.

## 4. High-energy scattering and hard probes

The previously discussed parton picture at strong coupling has some striking physical consequences for the high-energy scattering problem which look very different from our experience with QCD. For instance, the rapid energy growth $Q_{\rm s}^2(x) \propto 1/x$ of the saturation momentum is much faster than the respective growth observed in the HERA data, namely $Q_{\rm s}^2 \sim 1/x^{\omega}$ with $\omega \simeq 0.2 \div 0.3$, and which is in fact well accounted for by perturbative QCD [16]. Also, the absence of large-$x$ partons in a hadronic wavefunction at strong coupling means that, in a hypothetical scattering between two such hadrons, there would be no particle production at either forward or backward rapidities: the two nuclei colliding with each other at strong coupling would fully stop each other [17]. This is in sharp contrast to the situation at RHIC, where the large-$x$ partons from the incoming nuclei are seen to emerge from the collision, as hadronic jets, along their original trajectories.

A related prediction of AdS/CFT is the absence of jets in electron–positron annihilation at strong coupling [12, 14]. Fig. 5 exhibits the typical, 2-jet, final state in $e^+e^-$ annihilation at weak coupling (left) together with what should be the corresponding state at strong coupling (right). In both cases, the final state is produced via the decay of a time-like photon into a pair of partons and the subsequent evolution of this pair. At weak coupling this evolution typically involves the emission of soft and collinear gluons, with the result that the leading partons get dressed into a pair of well-collimated jets of hadrons (*cf.* Fig. 5 (left)). Multi-jet ($n \geq 3$) events are possible as well, but they have a lower probability, as they require hard parton emissions in the final state, which are suppressed by asymptotic freedom [9]. At strong coupling, on the other hand, parton branching is much more efficient, as previously explained, and rapidly leads to a system of numerous and relatively soft quanta, with energies and momenta of the order of the soft, confinement, scale, which are isotropically distributed in space. Thus, the respective final state shown no sign of jets, but only an isotropic distribution of hadronic matter (*cf.* Fig. 5 (right)) [14].

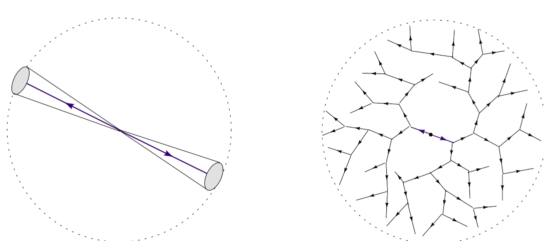

Fig. 5. Final state in $e^+e^-$ annihilation. Left: weak coupling, right: strong coupling.

Let us finally address the problem of interest for heavy ion physics, namely the propagation of a 'hard probe' through a strongly-coupled plasma. Consider *e.g.* the energy loss of a heavy quark: the respective AdS/CFT calculation has been given in [18], but the result can be also inferred from the present parton picture [20]. Among the virtual, space-like, quanta which are continuously emitted and reabsorbed by the heavy quark, only those can escape to the plasma which have a virtuality $Q$ lower than the plasma saturation momentum $Q_{\rm s}(x)$ for a value of $x$ set by the lifetime of the fluctuation: $1/x \sim T\Delta t$ with $\Delta t \sim \omega/Q^2$. Since $dE/dt \propto \omega/\Delta t \sim Q^2$, the energy loss is controlled by the fluctuations having the maximal possible value for the virtuality, that is, $Q \sim Q_{\rm s}(x)$ with $x$ set by the rapidity $\gamma$ of the heavy quark. Using $\gamma = \omega/Q$ and $Q_{\rm s} \sim T/x \sim \gamma T^2/Q_{\rm s}$, one finds $Q_{\rm s} \sim \sqrt{\gamma}\, T$ and

$$-\frac{dE}{dt} \simeq \sqrt{\lambda}\frac{\omega}{(\omega/Q_{\rm s}^2)} \simeq \sqrt{\lambda}\, Q_{\rm s}^2 \sim \sqrt{\lambda}\,\gamma\, T^2\,, \tag{10}$$

where the factor $\sqrt{\lambda}$ expresses the fact that, at strong coupling, the heavy quark does not radiate just a single quanta per time $\Delta t$, but rather a large number $\sim \sqrt{\lambda}$. Eq. (10) is parametrically consistent with the respective AdS/CFT result [18]. Note the strong enhancement of the medium effects at high energy, as expressed by the Lorentz $\gamma$ factor in the r.h.s. of (10): this is in qualitative agreement with the strong suppression of particle production seen in Au+Au collisions at RHIC, but one should be very careful before directly comparing such AdS/CFT results with the QCD phenomenology.

Consider similarly the momentum broadening: the $\sqrt{\lambda}$ quanta emitted during a time interval $\Delta t$ are uncorrelated with each other, so their transverse momenta are randomly oriented and the cumulating recoils increase the average *squared* transverse momentum of the heavy quark:

$$\frac{d\langle p_\perp^2\rangle}{dt} \sim \frac{\sqrt{\lambda}\, Q_{\rm s}^2}{(\omega/Q_{\rm s}^2)} \sim \sqrt{\lambda}\,\frac{Q_{\rm s}^4}{\gamma Q_{\rm s}} \sim \sqrt{\lambda}\,\sqrt{\gamma}\, T^3\,. \tag{11}$$

This is parametrically the same as the respective AdS/CFT results [19]. But this mechanism is different from the one at work in perturbative QCD, where the dominant contribution to momentum broadening rather comes from the medium rescattering.

To summarize, the strong-coupling picture of high-energy scattering appears to be very different from everything we know, theoretically and experimentally, about real-life QCD. There are no valence partons, the saturation momentum (and hence the cross-sections) grows much too fast with increasing energy, and there are no jets in the final state. This is not necessarily a surprise: within QCD, these high-energy phenomena are controlled by hard momentum exchanges and thus by weak coupling, because of asymptotic freedom. On the other hand, AdS/CFT might give us some qualitative insight in the semi-hard physics of particle production in heavy ion collisions, and also in some longstanding problems like thermalization and the emergence of hydrodynamics in the late stages of the collision [21].

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