Can the glasma affect jet quenching in AA-collisions at RHIC and LHC?

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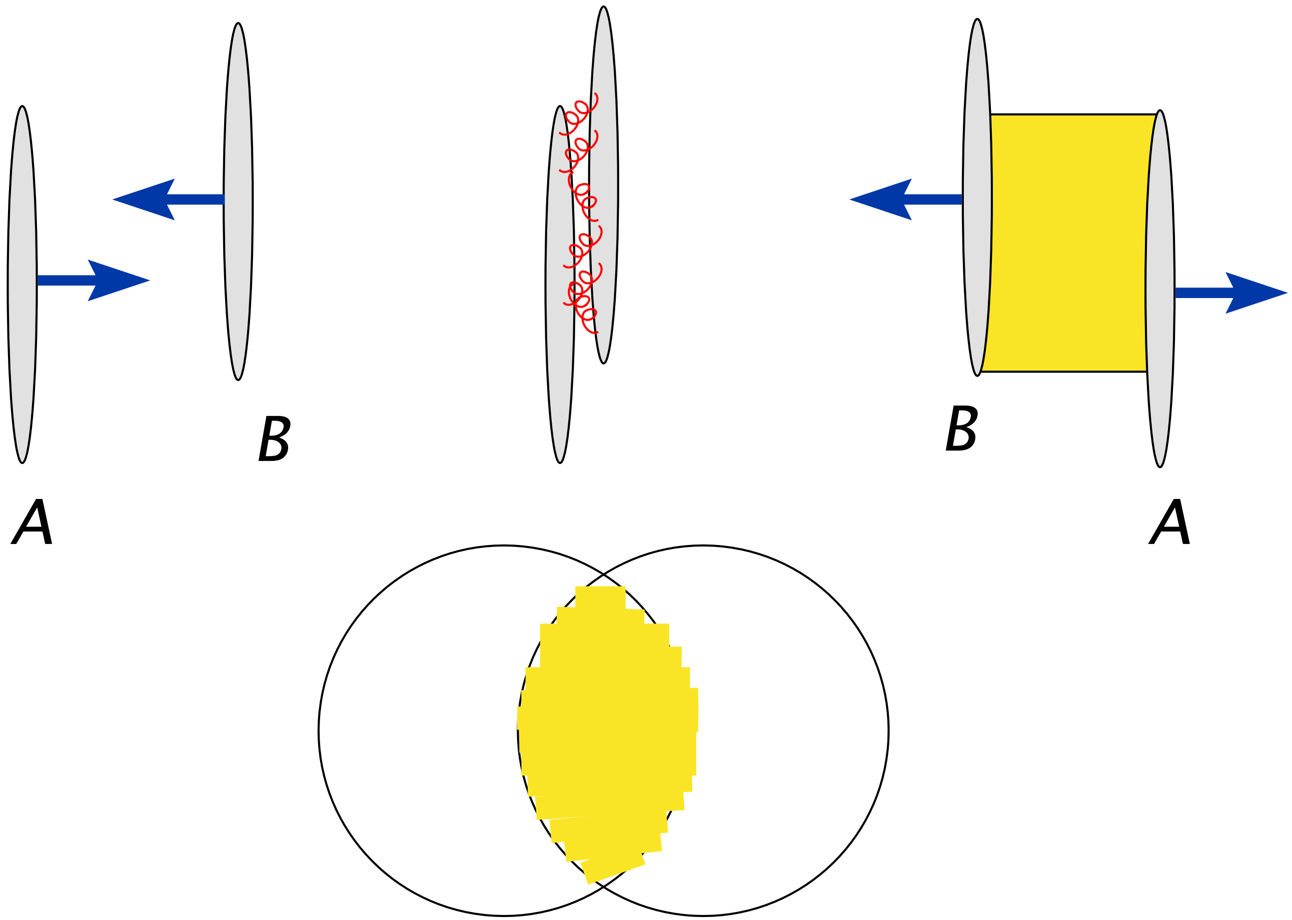
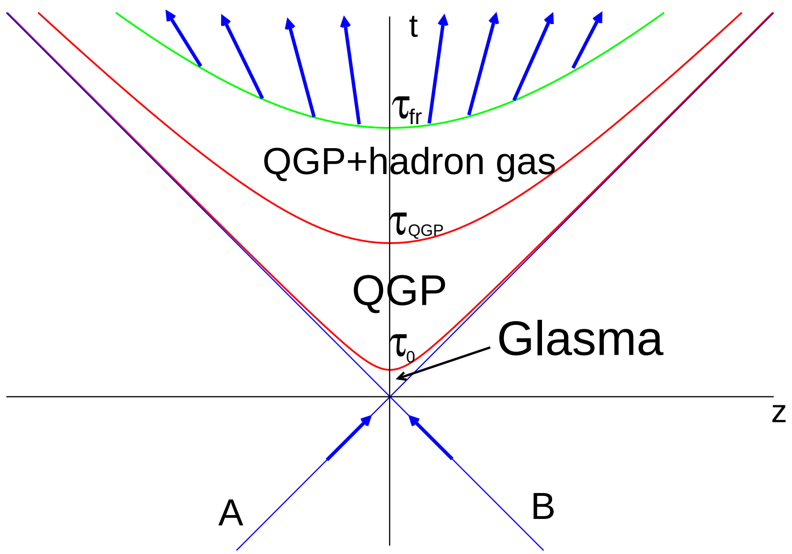
Outline:

- Introduction: Creation of hot QCD matter and jet quenching in AA-collisions at RHIC-LHC.
- Quasiclassical approach to the synchrotron-like gluon emission in a finite-size classical color fields.
- Numerical results for RHIC and LHC conditions.

Creation of hot QCD matter in AA-collisions



Bjorken expansion



In the Bjorken $1+1$ expansion $T^3\tau = T_0^3\tau_0$ in the QGP phase. $n(\tau) \approx n_0(\tau_0/\tau)$ in the whole range of t. For the QGP creation time $\tau_0 = 0.5$ fm from $dS/dy \Big/ dN_{ch}/d\eta \approx 7.67$ [B. Mueller and K. Rajagopal (2005)]

$\Rightarrow \langle T_0 \rangle \approx 300$ MeV (central Au+Au, $\sqrt{s} = 200$ GeV), $\langle T_0 \rangle \approx 400$ MeV (central Pb+Pb, $\sqrt{s} = 2.76$ TeV).

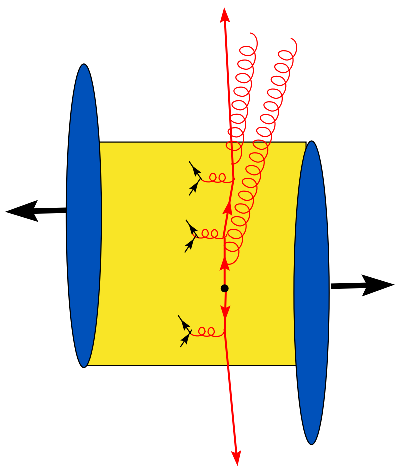
Glasma is the preequilibrium phase with strong coherent color fields predicted in the CGC model [T. Lappi and L.D. McLerran, Nucl. Phys. A772, 200 (2006)].

Jet quenching in AA-collisions

The parton energy loss suppresses high-p_T hadrons. It is described by the nuclear modification factor

$$R_{AA}(b) = \frac{dN(A+A \to h+X,\vec{b})/d\vec{p}_T dy}{T_{AA}(b) d\sigma(N+N \to h+X)/d\vec{p}_T dy},$$

$T_{AA}(b) = \int d\vec{\rho} T_A(\vec{\rho}) T_A(\vec{\rho}-\vec{b})$, $T_A(\vec{\rho}) = \int dz \rho_A(\vec{\rho},z)$ is the nucleus profile function.
The suppression is strong $R_{AA} \sim 0.1-0.3$ for RHIC-LHC.
It is widely believed that the JQ is due to radiative (Bethe-Heitler) and collisional (Bjorken) energy losses in the QGP



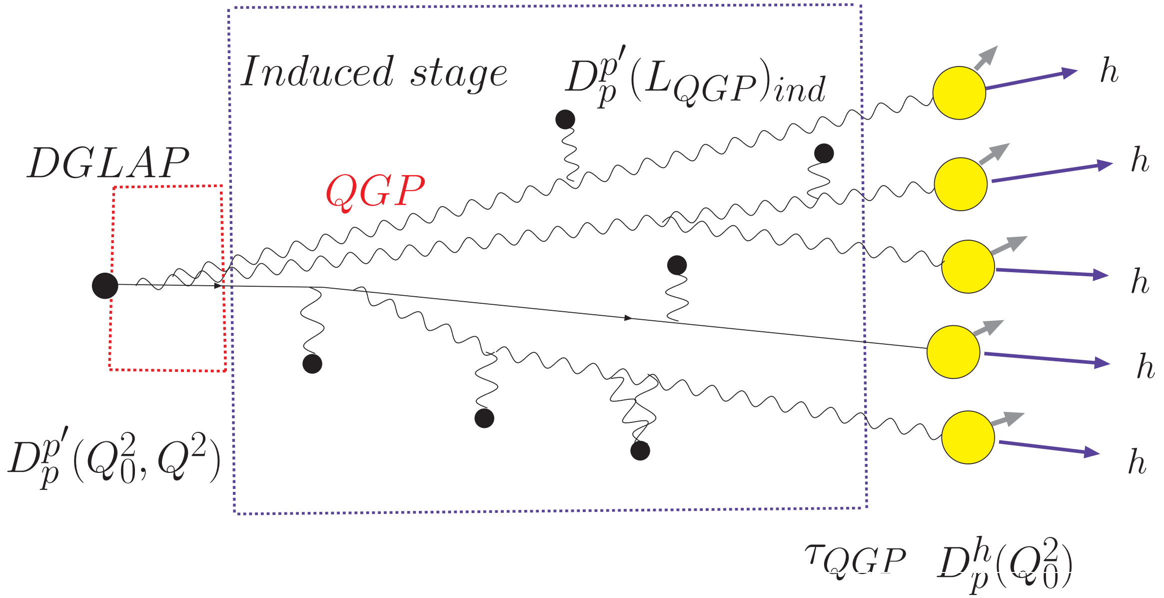
The radiative mechanism dominates since $\Delta E_{coll}/\Delta E_{rad} \sim 0.2-0.3$ [BGZ (2007)].
The theoretical uncertainties in R_{AA} are large (about a factor of $1.5-2$) but variation of R_{AA} from RHIC to LHC is more robust [BGZ (2011)].

The space-time pattern of jet distortion

The formation length for gluon emission

$$l_f \sim \frac{2x(1-x)E_q}{q_T^2+\epsilon^2}\,, \quad \epsilon^2 = m_q^2 x^2 + m_g^2(1-x)$$

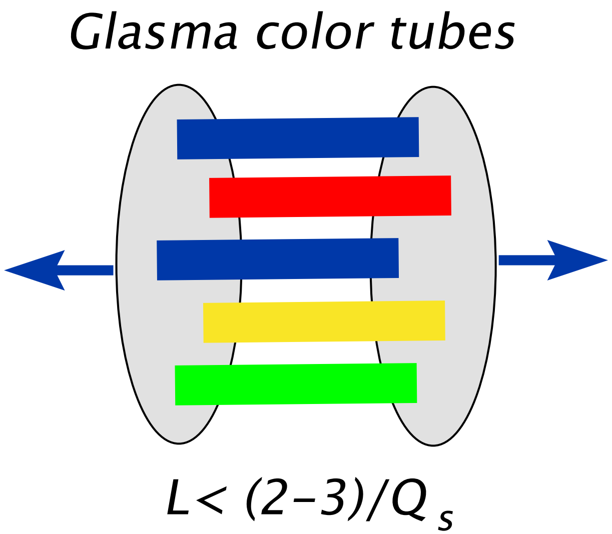
For the DGLAP $\bar{l}_F \sim 0.3-1$ fm for $E \lesssim 100$ GeV (if $m_q \sim 0.3$ GeV and $m_g \sim 0.75$ GeV). $\Rightarrow$We can neglect the overlapping of the DGLAP and induced emission stages [BGZ (2008)].



The glasma can modifies the DGLAP stage. How strong is this modification?

Glasma structure

The classical Yang-Mills fields in AA-collisions can be evaluated within the CGC model [McLerran and Venugopalan (1994)]. Just after the collision of the Lorentz contracted nuclei a system of the color flux tubes with the longitudinal boost invariant color electric and color magnetic fields (with $|E_z| \approx |B_z|$) should be produced.



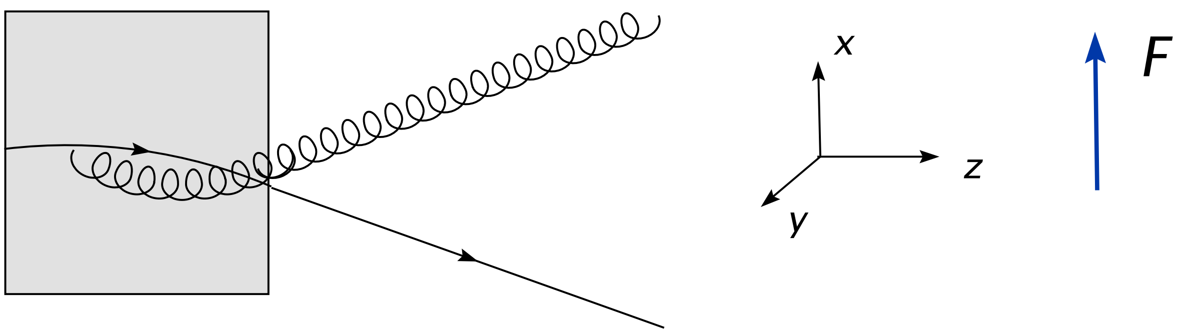
The transverse coherence length in this phase is $\sim 1/Q_s$, where Q_s ($\sim 1 - 1.5$ GeV for RHIC and LHC conditions) is the saturation scale of the nuclear parton distributions. At $\tau \sim 1/Q_s$ $gE_z \sim gB_z \sim Q_s^2$. The transverse fields (absent at $\tau = 0$) are also generated, at $\tau \gtrsim 1/Q_s$ they are close to the longitudinal ones. At $\tau \gtrsim 1/Q_s$ the energy density $\varepsilon = (E^2 + B^2)/2 \propto 1/\tau$ [Lappi (2006)].
The thermalization goes probably via instabilities rising quickly at $\tau \gtrsim 1/Q_s$ [Romatschke and Venugopalan (2006), Fujii and Itakura (2008), Iwazaki (2009)]. They should lead to a fast randomization of the color fields at $\tau \sim (2-3)/Q_s$.

Synchrotron-like gluon emission in glasma

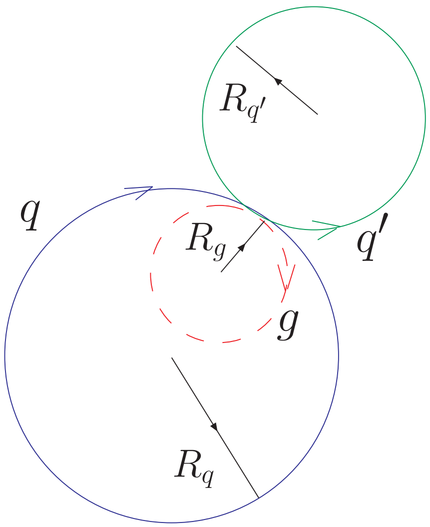
For RHIC-LHC the typical Lorentz force acting on a fast parton in the glasma $\sim Q_s^2 \sim 5 - 10$ GeV/fm. It is about 10–20 times that for the Debye screened color center in the QGP $\sim \alpha_s m_D^2$ (we take $\alpha_s \sim 0.3$ and $m_D \sim 0.5$ GeV).

The transverse momentum which a fast parton gets in the glasma should come mostly from interaction with the color field of the first crossed color flux tube. The random transverse momentum kicks at later times should should give a small effect due to weakness of the fields and the destructive interference of the chaotic contributions from different color tubes.

$\Rightarrow$ As a first step in understanding the glasma effect it seems reasonable to consider a model with a uniform time-dependent color field which acts only for a limited range of τ ($\tau = z$) about 2–3 units of $1/Q_s$.



Gluon synchrotron radiation in uniform field



$$\omega \gg m_g$$

$$E, E' \gg m_q$$

$$L_f \ll min(R_q, R_{q'}, R_g)$$

The small angle approximation is applicable at the scale $L \sim L_f$ ($L_f \sim$ min(L_1, L_2), where $L_1 = 2x(1-x)E/\epsilon^2$ and $L_2 = (24x(1-x)E/\vec{f}^{\,2})^{1/3}$).

$\Rightarrow$ $dN_s/d\omega dL$ can be calculated for a slab with thickness $R_{g,q} \gg L \gg L_f$ [BGZ JETP Lett. 88, 475 (2008)].

The results agree with recent calculations in the Schwinger method [A. Dbeyssi, D.A. Dirani and H. Zaraket, Phys.Rev. D84, 105033 (2011).], and disagree with previous calculations by Shuryak and Zahed [Phys. Rev. D67, 054025 (2003)].

Unfortunately, the Schwinger method is inapplicable for a finite-size slab. We use a quasiclassical approach in the small angle approximation.

Synchrotron emission in small angle approximation

- For $SU(3)$ it is enough to consider chromomagnetic field with color components $a = 3$ and $a = 8$. For radiated gluons we use the color states $Q = (Q_A, Q_B)$ with definite color isospin, Q_A, and color hypercharge, Q_B.
  There are 2 neutral gluons $A = G_3$ and $B = G_8$, and 6 charged gluons $X, Y, Z, \bar{X}, \bar{Y}, \bar{Z}$ given by

$$
\begin{aligned}
X &= (G_1 + iG_2)/\sqrt{2}, \quad Q = (-1, 0), \\
Y &= (G_4 + iG_5)/\sqrt{2}, \quad Q = (-1/2, -\sqrt{3}/2), \\
Z &= (G_6 + iG_7)/\sqrt{2}, \quad Q = (1/2, -\sqrt{3}/2).
\end{aligned}
$$

- The S-matrix element of the $q \to gq'$ synchrotron transition can be written as

$$
\langle gq'|\hat{S}|q\rangle = -ig \int dy \bar{\psi}_{q'}(y)\gamma^{\mu} G_{\mu}^{*}(y)\psi_q(y)\,.
$$

  We write each quark wave function in the form

$$
\psi_i(y) = \exp[-iE_i(t - z)]\hat{u}_{\lambda}\phi_i(z, \vec{\rho})/\sqrt{2E_i}\,,
$$

  where λ is quark helicity, $\hat{u}_{\lambda}$ is the Dirac spinor operator.

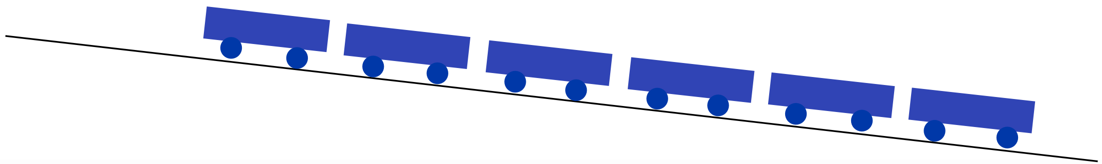
The z-dependence of the transverse wave functions ϕ_i for a parton with color vector $Q = (Q_A, Q_B)$ is governed by the two-dimensional Schrödinger equation

$$i\frac{\partial \phi_i(z,\vec{\rho})}{\partial z} = \left\{\frac{(\vec{p} - gQ_n\vec{G}_n)^2 + m_q^2}{2E_i} + U_i(z,\vec{\rho})\right\}\phi_i(z,\vec{\rho})$$

with the potential $U_i(z,\vec{\rho}) = gQ_n^i[G_n^0(z,\vec{\rho}) - G_n^3(z,\vec{\rho})]$. (the superscripts are the Lorentz indexes and $n = A, B$). The gluon wave function can be represented in a similar way.

We take the external vector potential in the form $G_n^0 = -\vec{\rho}\cdot\vec{E}_n$, $\vec{G}_n = 0$, and $G_n^3 = [\vec{H}_n \times \vec{\rho}]^3$, where $\vec{E}_n$ and $\vec{H}_n$ are the electric and magnetic fields. Then $U_i = -\vec{F}_i \cdot \vec{\rho}$, where $\vec{F}_i$ is the Lorentz force. The $\phi_i(z,\vec{\rho})$ can be taken in the form

$$\phi_i(z,\vec{\rho}) = \exp\left\{i\vec{p}_i(z)\vec{\rho} - \frac{i}{2E_i}\int_0^z dz'[\vec{p}_i^2(z') + m_q^2]\right\}, \qquad \frac{d\vec{p}_i}{dz} = \vec{F}_i(z)\,.$$



$$\langle gq'|\hat{S}|q\rangle = -ig(2\pi)^3\delta(E_g + E_{q'} - E_q)\int_{-\infty}^{\infty} dz V(z,\{\lambda\})\delta(\vec{p}_g(z) + \vec{p}_{q'}(z) - \vec{p}_q(z))$$

$$\times \exp\left\{ i\int_0^z dz' \left[\frac{\vec{p}_{q'}^{\,2}(z') + m_q^2}{2E_{q'}} + \frac{\vec{p}_g^{\,2}(z') + m_g^2}{2E_g} - \frac{\vec{p}_q^{\,2}(z') + m_q^2}{2E_q}\right]\right\},$$

where V is the spin vertex factor. Here $\vec{p}_g(z) + \vec{p}_{q'}(z) - \vec{p}_q(z) = 0$ (since $\vec{F}_q = \vec{F}_g + \vec{F}_{q'}$) $\Rightarrow$

$$\langle gq'|\hat{S}|q\rangle = -i(2\pi)^3\delta(\omega + E_{q'} - E_q)\delta(\vec{p}_g^{\,+} + \vec{p}_{q'}^{\,+} - \vec{p}_q^{\,+})T\,,$$

$$T = g\int_0^{\infty} dz V(\vec{q}(z),\{\lambda\})\exp\left\{ i\int_0^z dz' \left[\frac{\vec{q}^{\,2}(z') + \epsilon^2}{2M}\right]\right\}.$$

$\vec{p}_i^{\,+} = \vec{p}_i(z=\infty)$, $\vec{q}(z) = \vec{p}_g(z)(1-x) - \vec{p}_{q'}(z)x$, $x = \omega/E_q$, $\epsilon^2 = m_q^2x^2 + (1-x)m_g^2$ and $M = E_q x(1-x)$.

$$\frac{dN}{d\omega d\vec{q}} = \frac{|T|^2}{8(2\pi)^3E_q^3x(1-x)}\,.$$

Without an external field $\vec{q}(z) = const$. Then for $q \to gq$ in vacuum

$$T_v = gV(\vec{q}, \{\lambda\}) \frac{2iM}{\vec{q}^2 + \epsilon^2} \,,$$

$$\frac{dN_v}{d\omega d\vec{q}} = \frac{2C_F \alpha_s}{\pi^2 x E_q} \Big(1 - x + \frac{x^2}{2}\Big) \frac{\vec{q}^2}{(\vec{q}^2 + \epsilon^2)^2} \,.$$

For a nonzero external field

$$T = T_v + T_s \,,$$

$$T_s = g \int_0^L dz V(\vec{q}(z), \{\lambda\}) \exp\left\{ i \int_0^z dz' \left[\frac{\vec{q}^2(z') + \epsilon^2}{2M} \right] \right\} - (\vec{q}(z) \to \vec{q}) \,.$$

The synchrotron correction to the LO gluon spectrum reads

$$\frac{dN_s}{d\omega d\vec{q}} = \frac{2\text{Re}(T_v T_s^*) + |T_s|^2}{8(2\pi)^3 E_q^3 x(1-x)} \,.$$

$\vec{q}$-integrated spectrum

$dN_s/d\omega$ can be easily obtained within the LCPI approach [BGZ (1996)] formulated in the impact parameter space

$$\frac{dN_s}{d\omega} = 2\mathrm{Re}\int_0^\infty dz_1 \int_{z_1}^\infty dz_2 \hat{g}\left[\mathcal{K}(\vec{\rho}_2, z_2|\vec{\rho}_1, z_1) - \mathcal{K}_0(\vec{\rho}_2, z_2|\vec{\rho}_1, z_1)\right]\Big|_{\vec{\rho}_1=\vec{\rho}_2=0},$$

$$\begin{aligned}\hat{g} &= \frac{\alpha_s}{8E_q^3 x(1-x)} \sum_{\{\lambda\}} V(-i\partial/\partial\vec{\rho}_1, \{\lambda\}) V^*(-i\partial/\partial\vec{\rho}_2, \{\lambda\}) \\ &= \frac{|C|^2\alpha_s}{E_q^3 x^3 (1-x)^2}\left(1 - x + \frac{x^2}{2}\right)\frac{\partial}{\partial\vec{\rho}_1}\cdot\frac{\partial}{\partial\vec{\rho}_2},\end{aligned}$$

$\mathcal{K}$ is the Green's function of the Schrödinger equation with the Hamiltonian

$$\hat{H} = -\frac{1}{2M}\left(\frac{\partial}{\partial\vec{\rho}}\right)^2 - \vec{f}\cdot\vec{\rho} + \frac{\epsilon^2}{2M},$$

and $\mathcal{K}_0$ is the Green's function for the Hamiltonian with $\vec{f} = 0$.
The subtraction of the $\mathcal{K}_0$ term in the LCPI formula corresponds to subtraction of the vacuum spectrum.

The Green’s function is known explicitly

$$\mathcal{K}(\vec{\rho}_2, z_2|\vec{\rho}_1, z_1) = \frac{M}{2\pi i\Delta z}\exp\left[iS_{cl}\right], \tag{1}$$

where $\Delta z = z_2 - z_1$ and

$$S_{cl} = -\frac{\Delta z\epsilon^2}{2M} + \frac{M}{2\Delta z}\left[(\vec{\rho}_2 - \vec{\rho}_1)^2 + \frac{2}{M}\int_{z_1}^{z_2} dt\,\vec{\rho}_2\cdot\vec{f}(t)(t - z_1)\right.$$
$$\left. + \frac{2}{M}\int_{z_1}^{z_2} dt\,\vec{\rho}_1\cdot\vec{f}(t)(z_2 - t) - \frac{2}{M^2}\int_{z_1}^{z_2} dt\int_{z_1}^{t} ds\,\vec{f}(t)\cdot\vec{f}(s)(z_2 - t)(s - z_1)\right]. \tag{2}$$

For RHIC-LHC conditions the energy loss spectrum can not be calculated accurately at gluon energy $\omega \lesssim 5$ GeV. Numerical calculations show that in this region the large angle emission becomes important. $\Rightarrow$ The $\vec{q}$-integration can not be performed accurately using the formulas based on the small angle approximation. For this reason our results on the energy loss can only be treated as qualitative estimates, rather than quantitative predictions.

Gluon emission in a uniform field

$$\Rightarrow \quad \frac{dN_s}{dLdx} = \frac{iM}{2\pi}\int_{-\infty}^{\infty}\frac{d\tau}{\tau}\left[\frac{g_1}{M^2}\left(\epsilon^2+\frac{\vec{f}^{\,2}\tau^2}{2}\right)\right]\exp\left\{-i\left[\frac{\epsilon^2\tau}{2M}+\frac{\vec{f}^{\,2}\tau^3}{24M}\right]\right\}.$$

$g_1 = \alpha_s|\lambda^a_{fi}\chi^*_a/2|^2(1-x+x^2/2)/x$, Here τ has a small negative imaginary part. The integral around the lower semicircle near the pole at $\tau=0$ plays the role of the $\vec{f}=0$ subtraction term.

$$\frac{dN_s}{dLdx} = \frac{a}{\kappa}\mathrm{Ai}'(\kappa) + b\int_{\kappa}^{\infty} dy\mathrm{Ai}(y)\,,$$

where $a=-2\epsilon^2 g_1/M$, $b=-\epsilon^2 g_1/M$, $\kappa=\epsilon^2/(M^2\vec{f}^{\,2})^{1/3}$.
The effect of the field is only accumulated in $\vec{f}^{\,2}=\vec{F}_{q'}^{\,2}x_g^2-2\vec{F}_{q'}\vec{F}_g x_{q'}x_g+\vec{F}_g^2 x_q{}^2$.

Our spectrum disagrees with that obtained by Shuryak and Zahed [Phys. Rev. D**67**, 054025 (2003)] in the soft gluon limit within the Schwinger's proper time method.

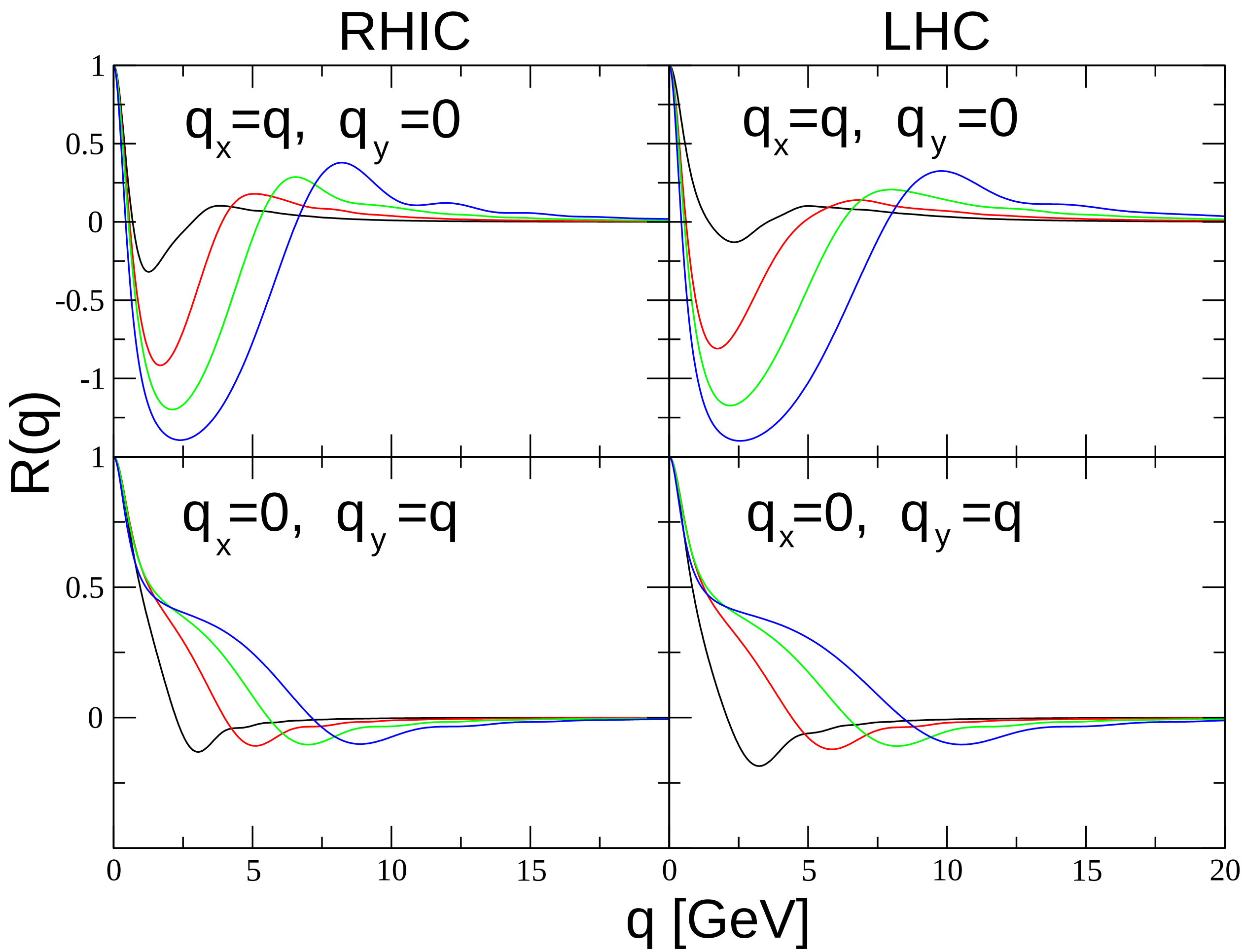
- In the SZ formula the argument of the exponential contains $\vec{F}_{q'}^{\,2}x_g^2+\vec{F}_g{}^2$.
- In the pre-exponential factor instead of $\vec{f}^2$ SZ have $\vec{F}_{q'}^{\,2}x_g^2$.

The SZ predictions are physically absurd: For $g\to gg$ the SZ spectrum has incorrect permutation properties. SZ results disagree with A. Dbeyssi, D.A. Dirani and H. Zaraket, Phys.Rev. D84, 105033 (2011).

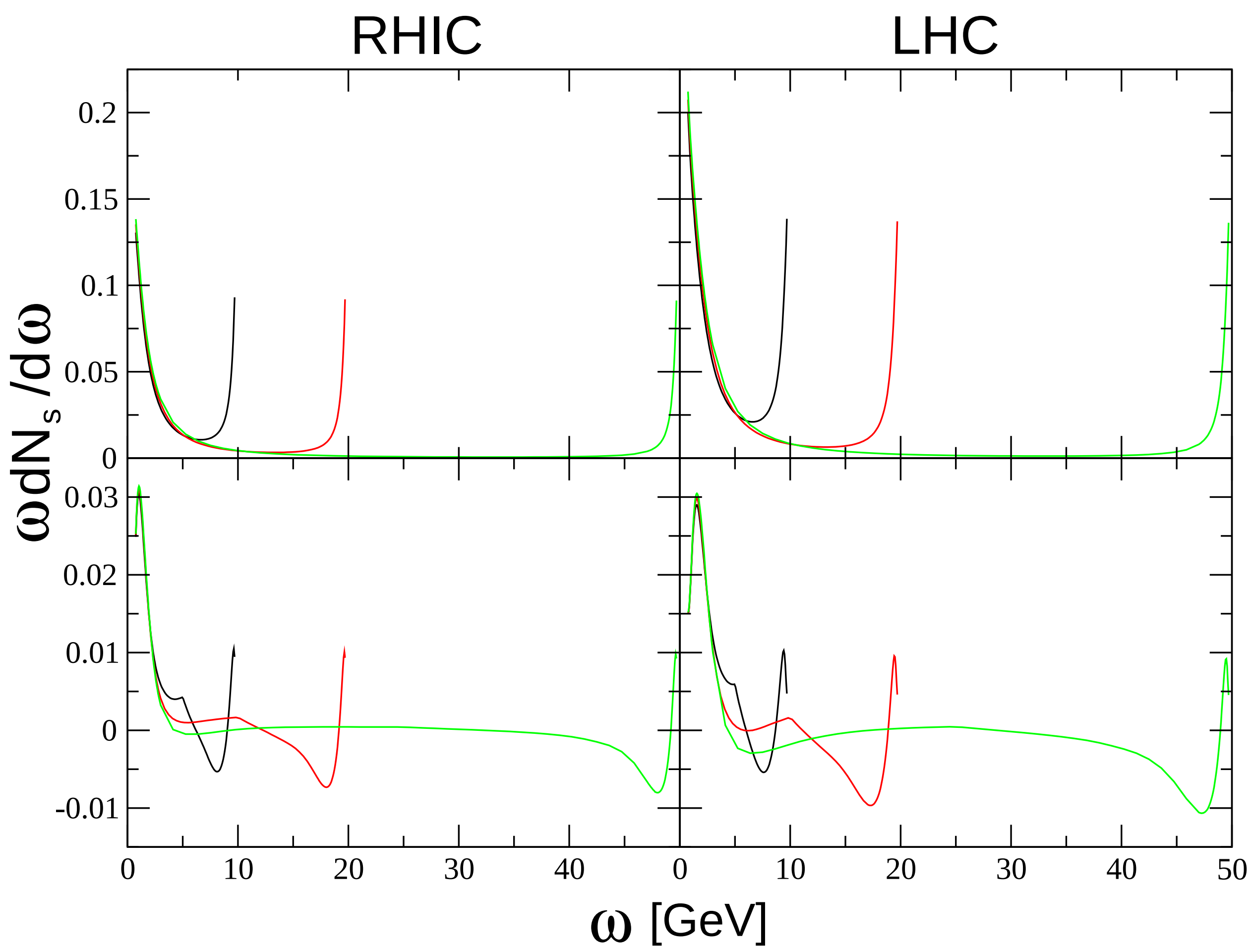
For numerical calculations we take $g = 2$ ($\alpha_s \approx 0.318$), $m_q = 0.3$ GeV, $m_g = 0.75$ GeV, and

$$Q_s = 1 \text{ GeV for RHIC and} \quad Q_s = 1.4 \text{ GeV for LHC}.$$

We take $L = 2/Q_s$. This gives $L(\text{RHIC}) \approx 0.4$ fm and $L(\text{LHC}) \approx 0.28$ fm.
To fix the z-dependence of the Lorentz force we use the τ-dependence of the glasma energy density $\varepsilon = (E^2 + H^2)/2$ obtained in the lattice simulations by Lappi [T. Lappi, Phys. Lett. B643, 11 (2006)].
We need only the field components transverse to the initial parton momentum. At $\tau \ll 1/Q_s$, when the electric and magnetic fields are almost parallel to the AA-collision axis the Lorentz force acting on a fast parton is purely transverse to the parton momentum. At such times for a unit color charge $F^2 = 2g^2\varepsilon$.
However, this relation is invalid at $\tau \gtrsim 1/Q_s$ when the contribution of the transverse (to the AA-collision axis) components of the color fields to the energy density becomes approximately equal to that from the components along the beam axis.
Since only half of these transverse fields squared contribute to the Lorentz force squared which we need (transverse to the jet direction) we can write in this regime $F^2 = g^2 3\varepsilon/2$.



The ratio $R(q) = dN_s/d\omega d\vec{q} \Big/ dN_s/d\omega d\vec{q}\Big|_{\vec{q}=0}$ at $E_q = 50$ GeV for $\omega =$2 (black), 5 (red), 10 (green), and 25 (blue) GeV. x-axis in $\vec{q}$-plane is parallel to the Lorentz force $\vec{f}$, and y-axis is perpendicular to it.



$\omega dN_s/d\omega$ at $E_q = 10$ (black), 20 (red), and 50 (green) GeV without kinematical constraint on $\vec{q}$ (upper panels) and with the restriction $|\vec{q}| < \mathrm{min}(\omega, E - \omega)$ (lower panels).

Energy loss due to synchrotron radiation

Even without the kinematical constraint the synchrotron corrections turn out to be smaller by a factor $\sim 10-20$ than the contribution of the induced gluon emission in the QGP.
It is also seen from calculation of the total energy loss

$$\Delta E_s = \int_{\omega_{min}}^{\omega_{max}} d\omega \omega dN_s/d\omega \,.$$

Taking $\omega_{min} = m_g$ and $\omega_{max} = E_q/2$ without kinematical constraint we obtained

RHIC: $\Delta E_s \approx 184$ MeV (E=10 GeV), $\Delta E_s \approx 276$ MeV (E=50 GeV),

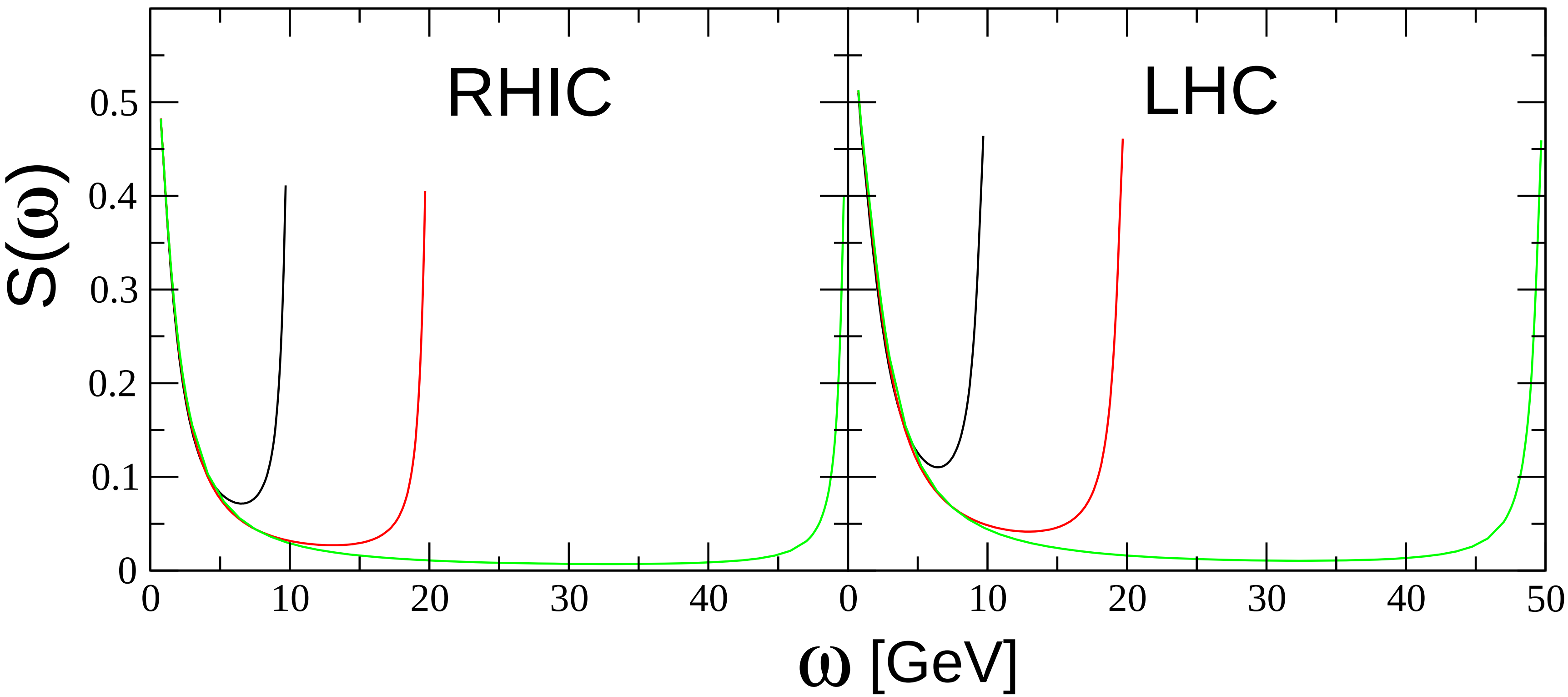
LHC: $\Delta E_s \approx 320$ MeV (E=10 GeV), $\Delta E_s \approx 495$ MeV (E=50 GeV),

and with the kinematical constraint $q <$ min$(\omega, E-\omega)$

RHIC: $\Delta Es \approx 48$ MeV (E=10 GeV), $\Delta E_s \approx 45$ MeV (E=50 GeV),

LHC: $\Delta E_s \approx 63$ MeV (E=10 GeV), $\Delta E_s \approx 39$ MeV (E=50 GeV)

$\Delta E_s \ll \Delta E_{ind}$. In QGP at $E_q \sim 10-50$ GeV $\Delta E_{ind} \sim 5-15$ GeV for RHIC, and $\Delta E_{ind} \sim 10-30$ GeV for LHC. The synchrotron radiation is weak mostly due to the finite-size effects since $L_f \gtrsim 1/Q_s$.



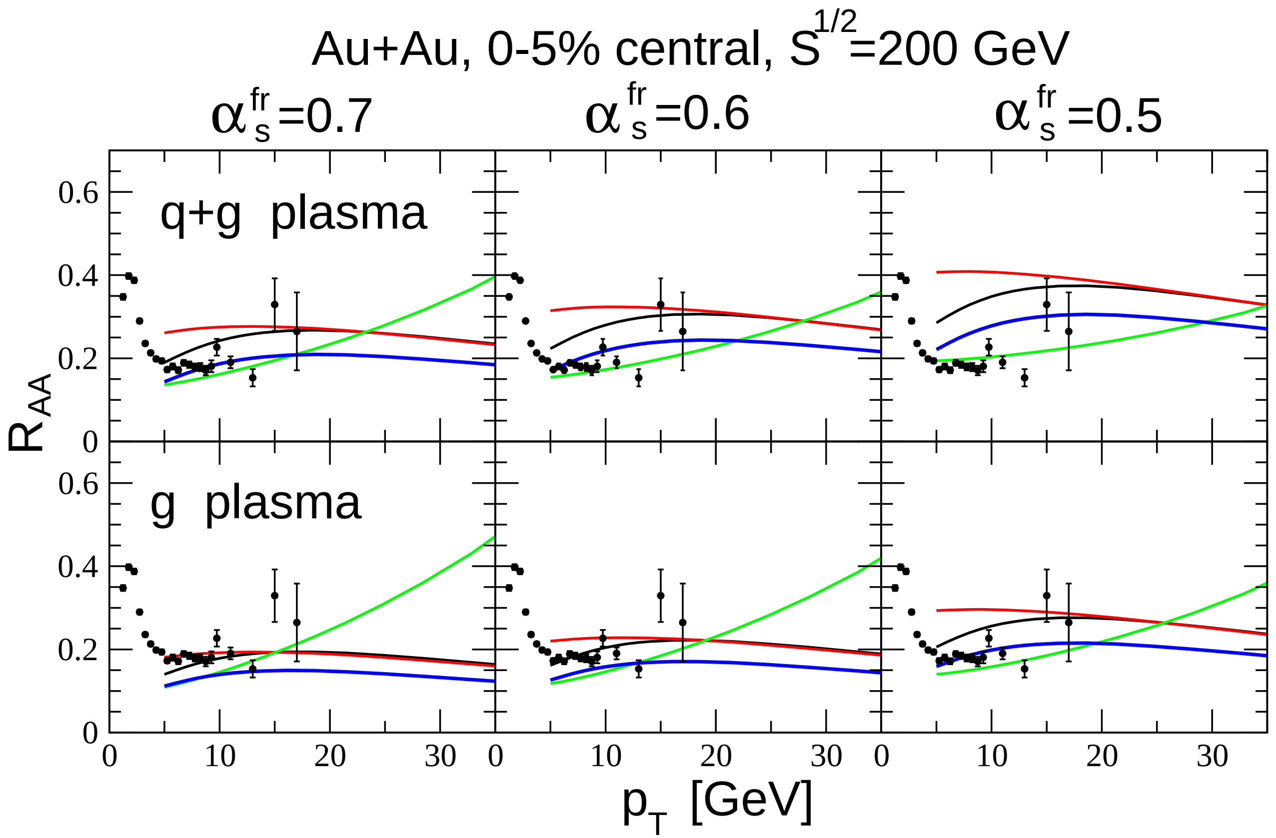
The finite-size suppression factor for the ω-distribution obtained without kinematical constraint on transverse momentum at $E_q = 10$ (black), 20 (red), and 50 (green) GeV.

Conclusions:

- We have developed a quasiclassical theory of the synchrotron-like gluon radiation in a finite-size external color field.
- We have found that the large angle gluon emission in glasma may be important for soft gluons with $\omega \lesssim 5$ GeV.
- Despite a huge Lorentz force acting on fast partons in the glasma, the effect of the glasma color tubes on parton energy loss turns out to be rather small. This smallness is due partly to strong finite-size effects.

BACK UP SLIDES

Comparison with RHIC PHENIX data, π^0



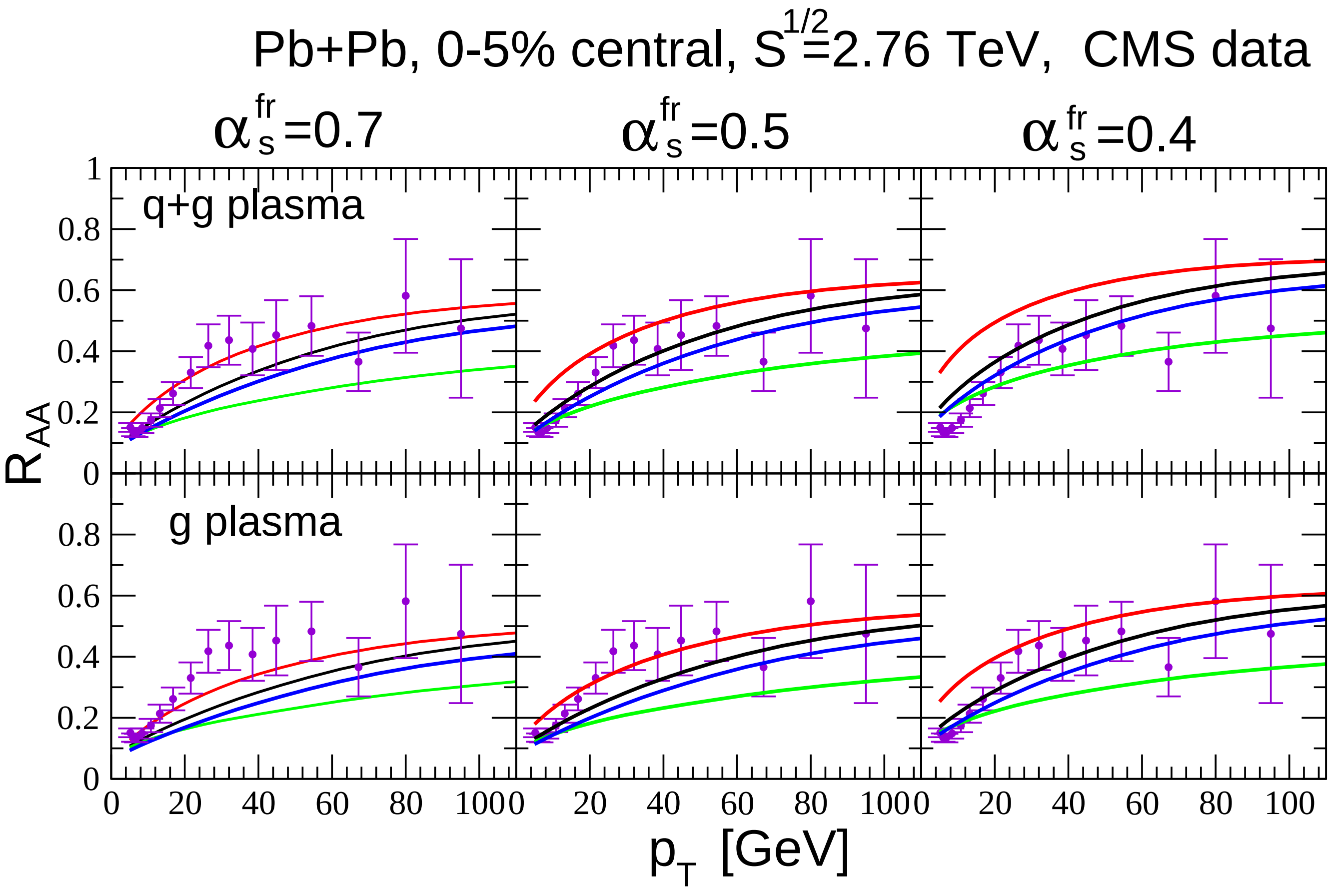
red: the radiative for quarks.

green: the radiative for gluons.

black: the radiative for quarks plus gluons.

blue: the total (quarks plus gluons) radiative plus collisional energy loss and the energy gain.

Comparison with CMS LHC data, $(h^+ + h^-)/2$



red: the radiative for quarks.

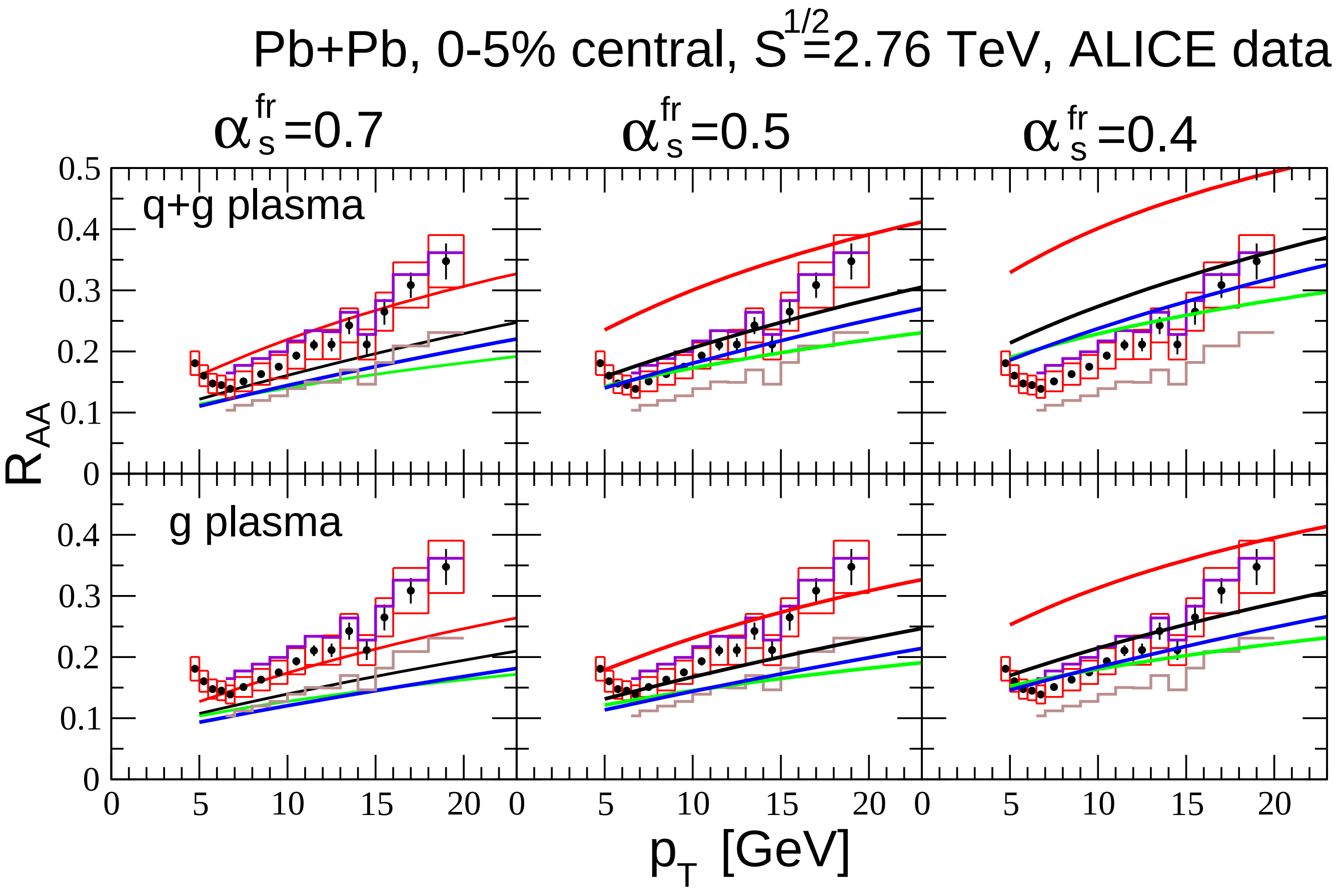
green: the radiative for gluons.

black: the radiative for quarks plus gluons.

blue: the total (quarks plus gluons) radiative plus collisional energy loss and the energy gain.

Experimental points: Preliminary CMS data [Y.J. Lee "Quark Matter 2011"]

Comparison with ALICE LHC data. $(h^+ + h^-)/2$



red: the radiative for quarks.

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blue: the total (quarks plus gluons) radiative plus collisional energy loss and the energy gain.

Experimental points: ALICE data [ALICE Collaboration, Phys. Lett. B696, 30 (2011)].